Holographic Double Diffractive Production of Higgs and the AdS Graviton/Pomeron

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The holographic approach to double diffractive Higgs production is presented in terms of exchanging the AdS graviton/Pomeron. The goal is to provide a simple framework for central exclusive production from a dual strong coupling perspective.

1 Introduction

A promising production mechanism for Higgs meson at the LHC involves the forward proton-proton scattering \( pp \to pHp \). The protons scatter through very small angles with a large rapidity gaps separating the Higgs in the central region. Current phenomenological treatment for the diffractive Higgs production cross section have generally followed two approaches: perturbative (weak coupling) vs confining (strong coupling), or equivalently, in terms of the Regge language, often referred to as the “hard Pomeron” vs “soft Pomeron” methods. The Regge approach to high energy scattering, although well motivated phenomenologically, has suffered in the past by the lack of a precise theoretical underpinning. The advent of AdS/CFT has dramatically changed the situation. In a holographic approach, the Pomeron is a well-defined concept and it can be identified as the “AdS graviton” in the strong coupling \( \mathbb{I} \), or, simply the BPST Pomeron. We briefly review here the general properties of the BPST Pomeron and show how it can be used to describe central double-diffractive particle production, and, in particular, for Higgs.

The formulation of AdS/CFT for high energy diffractive collision has already a rather extensive literature to draw on \cite{2,3,4}. “Factorization in AdS” has emerged as a universal feature, applicable to scattering involving both particles and currents. For instance, for elastic scattering, the amplitude can be represented schematically in a factorizable form,

\[
A(s, t) = \Phi_{13} \ast \tilde{K}_P \ast \Phi_{24}.
\]  

(1)

where \( \Phi_{13} \) and \( \Phi_{24} \) represented two elastic vertex couplings, and \( \tilde{K}_P \) is an universal Pomeron kernel, with a characteristic power behavior at large \( s > |t|, \tilde{K}_P \sim s^{j_0} \), schematically represented by Fig. \( \mathbb{I} \). This “Pomeron intercept”, \( j_0 \), lies in the range \( 1 < j_0 < 2 \) and is a function of the ’t Hooft coupling, \( \lambda = g^2N_c \). The convolution in \( \mathbb{I} \), denoted by the \( \ast \)-operation, involves an integration over the AdS location in the
bulk. In moving from elastic to the small-$x$ limit for the deep inelastic scattering (DIS) \[6\], one simply replaces $\Phi_{13}$ in \[1\] by appropriate product of propagators for external currents \[5, 6\].

\[
A(s, s_1, s_2, t_1, t_2) = \Phi_{13} \ast \tilde{K}_P \ast V_H \ast \tilde{K}_P \ast \Phi_{24},
\]

schematically represented by Fig. \[1\]. A more refined analysis for Higgs production requires a careful treatment for that depicted in Fig. \[1\]. A new aspect, not addressed in \[7\], is the issue of scale invariance breaking. A proper accounting for a non-vanishing gluon condensate $\langle F^2 \rangle$ turns out to be a crucial ingredient in understanding the strength of diffractive Higgs production \[9, 11\].

The basic theoretical steps necessary in order to arrive at \[1\] and \[2\] are:

(a) **Diffractive Scattering and Large $N_c$ Limit**: In leading order of the $1/N_c$ expansion at fixed ’t Hooft coupling $\lambda = g^2 N_c$, diffraction is given perturbatively by the exchange of a network of gluons with the topology of a cylinder, corresponding in a confining theory to the t-channel exchange of a closed string for glueball states. Such a state can be identified with the Pomeron.

(b) **Strong Coupling and AdS-Graviton**: Prior to AdS/CFT, property of the Pomeron has been explored mostly from a perturbative approach. The advent of the AdS/CFT correspondence has provided a firmer foundation from which a non-perturbative treatment can now be carried out. For instance, for elastic scattering, in the extreme limit of $\lambda \rightarrow \infty$, the 2-to-2 amplitude can be represented by a Witten diagram of a single graviton propagating in $AdS_5$, schematically given in a factorizable form, $A(s, t) = \Phi_{13} \ast \tilde{K}_G \ast \Phi_{24}$, where $\Phi_{13}$ and $\Phi_{24}$ represented two elastic vertex couplings to the graviton and $\tilde{K}_G$ is dominated by the “$(+,-)$” component of the graviton propagator \[2\]. Since this corresponds to a spin-2 exchange, the dominant graviton kernel $\tilde{K}_G$ grows with an integral power, i.e., at fixed $t$, as $s^2$.

Similarly, double diffractive Higgs production will be dominated by a double-graviton exchange Witten diagram, leading to a similar factorizable expression for the production amplitude $A(s_1, s_2, t_1, t_2) = \Phi_{13} \ast \tilde{K}_G \ast V_H \ast \tilde{K}_G \ast \Phi_{24}$, where $V_H$ is a new Higgs production vertex. The central issue in a holographic description for diffractive Higgs production is the specification of this new vertex $V_H$.

(c) **From Graviton to Pomeron**: It has been shown in \[1\], for $\mathcal{N} = 4$ SUSY YM, the leading strong
coupling Pomeron \(1, 2, 3\) is at

\[ j_0 = 2 - \frac{2}{\sqrt{g^2 N_c}}. \]  

which is “lowered” from \(J = 2\) as one decreases \(\lambda = g^2 N_c\). In a realistic holographic approach to high energy scattering, one must work at \(\lambda\) large but finite in order to account for the Pomeron intercept of the order \(j_0 \approx 1.3\), leading to \(1\) and \(2\) for elastic and diffractive Higgs production respectively. Here the Pomeron kernel, \(\tilde{K}_P\), has hard components due to near conformality in the UV and soft Regge behavior in the IR.\(^1\)

(d) **Confinement**: Our discussion above has been formal since a CFT has no scale and one needs to be more precise in defining the Regge limit. With confinement deformation, both the dilaton and the transverse-traceless metric fluctuations become massive, leading to an infinite set of massive scalar and tensor glueballs respectively. In particular, each glueball state can be described by a normalizable wave function \(\Phi(z)\) in \(AdS\). The weight factor \(\Phi_{ij}\) in the respective factorized representation for the elastic and Higgs amplitudes, \(1\) and \(2\), is given by \(\Phi_{ij}(z) = e^{-2A(z)}\Phi_i(z)\Phi_j(z)\). In contrast, for amplitudes involving external currents, e.g., for DIS \(3, 4\), non-normalizable wave-functions will be used.

## 2 Pomeron-Pomeron Fusion Vertex

In a perturbative approach, often dubbed as “hard Pomeron”, Higgs production can be viewed as gluon fusion in the central rapidity region. A Higgs can be produced at central rapidity by the double Regge Higgs vertex through a heavy quark loop which in lowest order is a simple gluon fusion process, dominant for large parton \(x\) for the colliding gluons. A more elaborate picture emerges as one tries to go to the region of the softer (wee gluons) building up double Regge regime\(^2\). In the large \(N_c\) there are no quark loop in the bulk of \(AdS\) space and since the Higgs in the Standard Model only couples to quark via the Yukawa interactions there appears to be a problem with strong coupling Higgs production in leading \(1/N_c\). Fortunately the solution to this is to follow the standard procedure in Higgs phenomenology, which is to integrate out the quark field replacing the Higgs coupling to the gauge operator \(\text{Tr}[F^2]\).

Consider the Higgs coupling to quarks via a Yukawa coupling, and, for simplicity we will assume is dominated by the top quark. We will be more explicit in the next Section, and simply note here that, after taking advantage of the scale separations between the QCD scale, i.e., the Higgs mass and the top quark mass, \(\Lambda_{\text{QCD}} \ll m_H \ll 2m_t\), heavy quark decoupling allows one to replace the Yukawa coupling by an effective interaction, \(\mathcal{L} = \frac{\alpha_s}{24\pi M_W} F^a_{\mu
u} F^{a\mu\nu} \phi_H\), by evaluating the two gluon Higgs triangle graph in leading order \(O(M_H/m_t)\). Now the \(AdS/CFT\) dictionary simply requires that this be the source in the UV of the \(AdS\) dilaton field. It follows, effectively, for Higgs production, we are required to work with a five-point amplitudes, one of the external leg involves a scalar dilaton current coupling to \(\text{Tr}[F^2]\). For

\(^1\)Unlike the case of a graviton exchange in \(AdS\), this Pomeron kernel contains both real and imaginary parts. For more discussion, see \([1, 3, 11]\).

\(^2\)In addition to the Pomeron exchange contribution in these models must subsequently be reduced by large Sudakov correction at the Higgs vertex and by so called survival probability estimates for soft gluon emission, again reflecting the view that double diffraction Higgs production is intrinsically non-perturbative.
diffractive Higgs production, in the supergravity limit, the Higgs vertex \( V_H \) is given by a two-graviton-dilaton coupling, Fig. 1c.

We now must pause to realize that in any conformal theory there is no dimensional parameter to allow for such a dimensionful two-graviton-dilaton coupling, \( M^2 \phi h_{\mu \nu} h^{\mu \nu} \), emerging in an expansion of the AdS gravity action if scale invariance is maintained. However since QCD is not a conformal theory this is just one of many reasons to introduce conformal symmetry breaking. Many attempts have been made to supplement this phenomenological Lagrangian with other fields such as the gauge fields for the light quark Goldstone modes to provide a better holographic dual for QCD. In principle even at leading order of large \( N_c \) we should eventually require an infinite number of (higher spin) field in the bulk representation to correspond the yet undiscovered 2-d sigma model for the world-sheet string theory for QCD. Fortunately for the phenomenological level at high energy, these details are non-essential. To model an effective QCD background we will for the most part introduce two modifications of the pure AdS background: (1) an IR hardwall cut-off beyond \( z = 1/\Lambda_{\text{QCD}} \) to give confinement and linear static quark potential at large distances and (2) a slow deformation in the UV \( (z \to 0) \) to model the logarithmic running for asymptotic freedom. Both break conformal invariance, which as we will argue is required to couple the two gravitons to the dilaton and produce a Higgs in the central rapidity region.

After taking into account of finite \( \lambda \) correction, the leading order Higgs production diagram at large \( N_c \) can be schematically represented in Fig. 1c, with each of the left- and right-cylinder representing a BPST Pomeron. It should be pointed out, just as in the case of elastic scattering, it is necessary to consider higher order corrections, e.g., eikonal corrections. We will not do it here, but will address this issue in the conclusion section. In what follows, we shall focus on the Pomeron-Pomeron fusion vertex in the strong coupling limit.

Finally it should be noted that one critical missing ingredient of these ad hoc conformal breaking deformation of the AdS geometry in the UV and IR is the fact the spontaneous breaking of pure Yang Mills (and presumable QCD at large \( N_c \)), via “dimensional transmutation” eliminates the coupling, \( \lambda \), as a free parameter. It is fixed via the beta function in terms of a single integration constant (sometime called \( \Lambda_{\text{QCD}} \)) which provides the only mass scale. Thus the logarithmic scale violation in the UV are tied to the same parameter giving confinement in the IR. All holographic modes of QCD to date introduce two mass scales and thus neglect this constraint. The solution to this problem also presumably awaits the determination of the unique string theory for large \( N_c \) QCD.

We are now in a position to focus on the issue of double diffractive Higgs production from the perspective of String/Gauge duality, i.e., the Higgs vertex, \( V_H \). It is important to stress that our general discussion in moving from single-Pomeron exchange processes, (1), to double-Pomeron exchange, (2), applies equally well for both diffractive glueball production and for Higgs production. The difference lies in how to treat the new central vertex. For the production of a glueball, the vertex will be proportional to a normalizable AdS wave-function. There will also be an overall factor controlling the strength of coupling to the external states, e.g., the Pomeron-Pomeron-glueball couplings. For Higgs production, on the other hand, the central vertex, \( V_H \), involves a non-normalizable bulk-to-boundary propagator, appropriate for a scalar external current. This in turns leads to coupling to a Higgs scalar. The difference between these
two cases parallels the situation for four-point amplitudes in moving from proton-proton (p-p) elastic scattering to electron-proton deep-inelastic scattering (e-p DIS). In moving from p-p to DIS, one simply replaces one of the two pairs of normalizable proton wave-functions with a pair of non-normalizable counterparts appropriate for conserved external vector currents.

A Higgs scalar in the standard model couples exclusively to the quarks via Yukawa coupling, which for simplicity we will assume is dominated by the top quark, with \( \mathcal{L} = -\frac{g}{2M_W}m_t \bar{t}(x)t(x)\phi_H(x) \). Taking advantage of the scale separations between the QCD scale, the Higgs mass and the top quark mass, \( \Lambda_{qcd} \ll m_H \ll 2m_t \), heavy quark decoupling allows one to replace the Yukawa coupling by direct coupling of Higgs to gluons, which is treated as an external source in the AdS dictionary. Consequently, \( V_H \), in a coordinate representation, is replaced by the vertex for two AdS Pomeron fusing at \( (x_{\perp}^1, z_1^1) \) and \( (x_{\perp}^2, z_2^2) \) and propagating this disturbance to the \( \bar{t}(x)t(x) \) scalar current at the boundary of AdS. The double diffractive Higgs vertex \( V_H \) can then be obtained in a two-step process.

First, since the Yukawa Higgs quark coupling is proportional to the quark mass, it is dominated by the top quark. Assuming \( m_H \ll m_t \), this can be replaced by an effective interaction, \( (2) \), by evaluating the two gluon Higgs triangle graph in leading order \( O(M_H/m_t) \). Second, using the AdS/CFT dictionary, the external source for \( F_{\mu\nu}F_{\mu\nu}^a(x) \) is placed at the AdS boundary \( (z_0 \to 0) \) connecting to the Pomeron fusion vertex in the interior of AdS at \( b_H = (x_H^i’, z_H^i’) \), by a scalar bulk-to-boundary propagator, \( K(x_H^i’ - x_H^i, z_H^i, z_0) \).

We are finally in the position to put all the pieces together. Although we eventually want to go to a coordinate representation in order to perform eikonal unitarization, certain simplification can be achieved more easily in working with the momentum representation. The Higgs production amplitude, schematically given by \( (2) \), can then be written explicitly as

\[
A(s, s_1, s_2; t_1, t_2) \simeq \int dz_1 dz_2 \frac{1}{\sqrt{-g_1}} \frac{1}{\sqrt{-g_2}} \Phi_{13}(z_1) \times \bar{K}_P(s_1, t_1, z_1, z) V_H(q^2, z) \bar{K}_P(s_2, t_2, z, z_2) \Phi_{24}(z_2). \tag{4}
\]

where \( q^2 = -m_H^2 \). For this production vertex, we will keep it simple by expressing it as

\[
V_H(q^2, z) = V_{PP\phi}K(q^2, z)L_H. \tag{5}
\]

where \( K(q^2, z) \) is the conventionally normalized bulk to boundary propagator, \( V_{PP\phi} \) serves as an overall coupling from two-Pomeron to \( F^2 \), and \( L \) is the conversion factor from \( F^2 \) to Higgs, i.e., \( L_H = L(-m_H^2) \simeq \frac{\alpha_s^2}{24\pi^2 M_W^2} \). We shall treat the central vertex \( V_{PP\phi} \) as a constant, which follows from the super-gravity limit. This approximation gives an explicit factorizable form for Higgs production.

The current version for the holographic Higgs amplitude \( (1) \) involves 3 parameters: (1) the IR cut-off determined by the glueball mass, (2) the leading singularity in the \( J \)-plane determined \( \lambda \) by the ’t Hooft parameter \( \lambda \) and (3) the strength of the central vertex parameterized by the string coupling or Planck mass. A strategy must be provided in fixing these parameters.

\(^3\)In a true dual to QCD, there is no independent parameter for the strong coupling, because of “dimensional transmutation”, which fixes all dimensionful quantities relative to the a single mass scale \( \Lambda_{qcd} \), through the running coupling constant. For instance, the glueball mass in units of \( \Lambda_{qcd} \) is fixed and computed in lattice computations.
3 Strategy for Phenomenological Estimates

As a first step in this direction, we ask how the central vertex, \( V_H \), or equivalently, \( V_{PPφ} \), via (5), can be normalized, following the approach of Kharzeev and Levin [9] based on the analysis of trace anomaly. We also show how one can in principle use the elastic scattering to normalize the bare BPST Pomeron coupling to external protons and the ‘t Hooft coupling \( g^2 N_c \). As in the case of elastic scattering, it is pedagogically reasonable to begin by first treating the simplest case of double-Pomeron exchange for Higgs production, i.e., without absorptive correction. We discuss how phenomenologically reasonable simplifications can be made. This is followed by treating eikonal corrections in the next section, which provides a means of estimating the all-important survival probability.

(a) Continuation to Tensor Glueball Pole: Confinement deformation in AdS will lead to glueball states, e.g., the lowest tensor glueball state lying on the leading Pomeron trajectory [8]. There will also be scalar glueballs associated with the dilaton. With scalar invariance broken, this will also lead to non-vanishing couplings between a pair of tensor glueballs and scalar glueballs. In terms of the language of Witten diagram, corresponds to a non-vanishing graviton-graviton-dilaton coupling in the bulk, which in turn leads to \( V_H \neq 0 \).

Consider first the elastic amplitude. With confinement, each Pomeron kernel will contain a tensor glueball pole when \( t \) goes on-shell. Indeed, the propagator for our Pomeron kernel can be expressed as a discrete sum over pole contributions. That is, when \( t \approx m_0^2 \), where \( m_0 \) is the mass of the lightest tensor glueball, which lies on the leading Pomeron trajectory. In this limit, the elastic amplitude then takes on the expected pole-dominated form,

\[
A(s, t) \approx g_{13} \frac{s^2}{(t_1 - m_0^2)(t_2 - m_0^2)} g_{24}.
\]

Here \( \Gamma_{GGH} \) is the effective on-shell glueball-glueball-Higgs coupling, which can also be expressed as

\[
\Gamma_{GGH} = L_H F(-m_H^2), \quad F = \frac{\alpha_s g}{2\pi M_W^2} \quad \text{and} \quad F(q^2) = \langle G, ++, q | F_{\mu\nu} F_{\mu\nu} | G, --, q \rangle.
\]

which can again be expressed as an overlapping integral involving the bulk-boundary propagator and the glue ball wave function. What remains to be specified is the overall normalization, \( A(s, t) \) is dimensionless.

A similar analysis can also be carried out for the Higgs production amplitude, Eq. (4). Note that the Pomeron kernel now appears twice, \( \tilde{K}_P(s_1, t_1, z_1, z) \) and \( \tilde{K}_P(s_2, t_2, z_2, z) \). When nearing the respective tensor poles at \( t_1 \approx m_0^2 \) and \( t_2 \approx m_0^2 \), the amplitude can be expressed as

\[
A(s_1, s_2, t_1, t_2) \approx g_{13} \frac{\Gamma_{GGH} s^2}{(t_1 - m_0^2)(t_2 - m_0^2)} g_{24}.
\]

Here \( \Gamma_{GGH} \) is the effective on-shell glueball-glueball-Higgs coupling, which can also be expressed as

\[
\Gamma_{GGH} = L_H F(-m_H^2), \quad L_H = \frac{\alpha_s g}{2\pi M_W^2} \quad \text{and} \quad F = \text{a form factor } F(q^2) = \langle G, ++, q | F_{\mu\nu} F_{\mu\nu} | G, --, q \rangle.
\]

which can again be expressed as an overlapping integral involving the bulk-boundary propagator and the glue ball wave function. What remains to be specified is the overall normalization, \( F(0) \).

We next follow D. Kharzeev and E. M. Levin [9], who noted that, from the SYM side, \( F(q^2) \) at \( q^2 = 0 \), can be considered as the glueball condensate, leading directly to

\[
F(0) = \langle G | F_{\mu\nu} F_{\mu\nu} | G \rangle = -\frac{4\pi M_W^2}{3\beta}
\]

where \( \beta = -b_0/2\pi \), \( b = 11 - 2n_f/3 \), for \( N_c = 3 \). In what follows, we will use \( n_f = 3 \). Note that heavy quark contribution is not included in this limit. Since the conformal scale breaking is due the running
coupling constant in QCD, there is apparently a mapping between QCD scale breaking and breaking of the AdS background in the IR, which gives a finite mass to the glueball and to give a non-zero contribution to the gauge condensate.

(b) Extrapolation to the Near-Forward limit: To apply the above result to the physical region, one needs to extrapolate from $t$ near the tensor pole to the physical region where $t_1, t_2 \simeq 0$. In addition, exponential cutoff for $t_1$ and $t_2$ small also have to be taken into account. For elastic scattering, it is customary to express the amplitude in the near forward region as $A(s,t) \simeq e^{B_{\text{eff}}(s) t/2} A(s,0)$ where $B_{\text{eff}}(s)$ is a smoothly slowly increasing function of $s$, (we expect it to be logarithmic). Similarly, we also assume, for $t_1 < 0$, $t_2 < 0$ and small, the Higgs production amplitude is also strongly damped so that

$$A(s,s_1, s_2, t_1, t_2) \simeq e^{B'_{\text{eff}}(s_1) t_1/2} e^{B'_{\text{eff}}(s_2) t_2/2} A(s, s_1, s_2, t_1 \simeq 0, t_2 \simeq 0).$$

Although these steps are conceptually straightforward, considerable details have to be spelled out in arriving at a manageable expression, as done in [11]. One finds that

$$\frac{d\sigma}{dy_H} \simeq \left(1/\pi\right) \times C \times \left|\Gamma_{\text{GGH}}(0)/\tilde{m}^2\right|^2 \times \frac{\sigma(s)}{\sigma(m_H^2)} \times R_{el}^2(m_H \sqrt{2})$$

Here, $R_{el}(s) = \sigma_{el}/\sigma_{total} \simeq \frac{(1+\rho^2)\sigma_{total}(s)}{16m_{\text{eff}}B_{\text{eff}}(s)}$. In this expression above, both $C$ and $\tilde{m}^2$, like $m_{\text{eff}}^2$, are model dependent. It is nevertheless interesting to note that, since $\Gamma_{\text{GGH}}(0) \sim m_B^2$, the glueball mass scale also drops out, leaving a model-dependent ratio of order unity. In deriving the result above, we have replaced $B'_{\text{eff}}$ by $B_{\text{eff}}$ where the difference is unimportant at high energy. With $m_H \simeq 150$ GeV, $R_{el}$ can be taken to be in the range 0.1 to 0.2. For $C \simeq 1$, we find $\frac{d\sigma}{dy_H} \simeq 0.8 \sim 1.2$ pb. This is of the same order as estimated in [9]. However, as also pointed in [9], this should be considered as an overestimate. The major source of suppression will come from absorptive correction, which can lead to a central production cross section in the fb range. We turn to this next.

4 Discussion

We conclude by discussing how consideration of higher order contributions via an eikonal treatment leads to further suppression for the central Higgs production. Following by now established usage, the resulting production cross section can be expressed in terms of a “survival probability”.

For elastic scattering, the resulting eikonal sum leads to an impact representation for the 2-to-2 amplitude: $A(s, x^+ - x'^+) = -2is \int dz \ dz' \ P_{13}(z) P_{24}(z') \left[e^{ix(x,x'-x'^+,z,z')}-1\right]$. The eikonal $\chi$, as a function of $x_\perp - x'^\perp$, $z$, $z'$ and $s$, can be determined by matching the first order term in $\chi$ to the single-Pomeron contribution. This eikonal analysis can be extended directly to Higgs production. Since Higgs production is a small effect, we find the leading order Higgs production amplitude, to all order in $\chi$, becomes

$$A_H(s_1, s_2, x^+ - x_H^+, x'^+ - x_H^+, z_H) = 2s \int dz \ dz' \ P_{13}(z) P_{24}(z') \times \chi_H(s_1, s_2, x^+ - x_H^+, x'^+ - x_H^+, z, z', z_H) e^{ix(x,x'^+,z,z')}$$

7
where $\chi_H$ can be found by matching in the limit $\chi \to 0$ with the Higgs production amplitude, in an impact representation. The net effect of eikonal sum is to introduce a phase factor $e^{i\chi_{s,x},z'}$ into the production amplitude. Due to its absorptive part, $\text{Im} \chi > 0$, this eikonal factor provides a strong suppression for central Higgs production.

The effect of this suppression is often expressed in terms of a “Survival Probability”, $\langle S \rangle$. In a momentum representation, the cross section for Higgs production per unit of rapidity in the central region is

$$\frac{d\sigma_H(s,y_H)}{dy_H} = \frac{1}{\pi(16\pi \Lambda^2)^2} \int d^2q_1 d^2q_2 |A_H(s,y_H,q_{1\perp},q_{2\perp})|^2,$$

where $y_H$ is the rapidity of the produced Higgs, $q_{1\perp}$ and $q_{2\perp}$ are transverse momenta of two outgoing fast leading particle in the frame where the momenta of incoming particles are longitudinal. “Survival Probability” is conventionally defined by the ratio $\langle S \rangle \equiv \frac{\int d^2q_1 d^2q_2 |A_H(s,y_H,q_{1\perp},q_{2\perp})|^2}{\int d^2q_1 d^2q_2 |A_H^{(0)}(s,y_H,q_{1\perp},q_{2\perp})|^2}$, where $A_H^{(0)}$ is the corresponding amplitude before eikonal suppression, e.g., given by Eq. (4). For simplicity, we shall also focus on the mid-rapidity production, i.e., $y_H \simeq 0$ in the overall CM frame. In this case, $\langle S \rangle$ is a function of overall CM energy squared, $s$, or the equivalent total rapidity, $Y \simeq \log s$. Evaluating the survival probability as given by (4), though straightforward, is often tedious.

To gain a qualitative estimate, let us consider the local limit where $z \simeq \bar{z} \simeq z_0$ and $z' \simeq \bar{z}' \simeq z'_0$, with $z_0 \simeq z'_0 \simeq 1/\Lambda_{QCD}$. In this limit, one finds that this suppression factor reduces to $e^{-2\text{Im} \chi(s,x,z,z',0)}$, where $\text{Im} \chi > 0$ by unitarity. If follows that, in a super-gravity limit of strong coupling where the eikonal is strictly real, there will be no suppression and the survival probability is 1. Conversely, the fact that phenomenologically a small survival probability is required provides another evidence the need to work in an intermediate region where $1 < j_0 < 2$. In this more realistic limit, $\text{Im} \chi$ is large and cannot be neglected. In particular, it follows that the dominant region for diffractive Higgs production in pp scattering comes from the region where $\text{Im} \chi(s,x,z,z') = O(1)$, with $z \simeq z' = O(1/\Lambda_{QCD})$. Note that this is precisely the edge of the “disk region” for p-p scattering. In order to carry out a quantitative analysis, it is imperative that we learn the property of $\chi(s,\vec{b},z)$ for $|\vec{b}|$ large. From our experience with pp scattering, DIS at HERA, etc., we know that confinement will play a crucial role. In pp scattering, since $z \simeq z' = O(1/\Lambda_{QCD})$, we expect this condition is reached at relatively low energy, as is the case for total cross section. It therefore plays a dominant role in determining the magnitude of diffractive Higgs production at LHC. We will not discuss this issue here further; more pertinent discussions on how to determine $\chi(s,x_{1\perp},z,z')$ when confinement is important can be found in Ref. [6].

We have focussed here primarily on central exclusive Higgs production from a holographic perspective. In such a treatment, a non-perturbative Pomeron corresponds to exchanging a Reggeized Graviton in $AdS$. This approach can of course also be applied to central exclusive production of mesons [10]. Equally interesting is the possibility of further generalization of incorporation of Odderon exchanges [7]. For instance, a Pomeron-Odderon double-exchange can be used to discussion the production of “axions”, a subject which is under current investigation.
References

[1] R. C. Brower, J. Polchinski, M. J. Strassler and C. I. Tan, “The Pomeron and Gauge/String Duality,” JHEP 0712, 005 (2007).

[2] R. C. Brower, M. J. Strassler and C. I. Tan, “On the Eikonal Approximation in AdS Space,” JHEP 0903, 050 (2009).

[3] R. C. Brower, M. J. Strassler and C. I. Tan, “On The Pomeron at Large ’t Hooft Coupling,” JHEP 0903, 092 (2009).

[4] L. Cornalba, et al., “Eikonal approximation in AdS/CFT: Conformal partial waves and finite N four-point functions,” Nucl. Phys. B 767, 327 (2007); “Eikonal approximation in AdS/CFT: From shock waves to four-point functions,” JHEP 0708, 019 (2007); “Eikonal Approximation in AdS/CFT: Resumming the Gravitational Loop Expansion,” 0709, 037 (2007); “Eikonal Methods in AdS/CFT: BFKL Pomeron at Weak Coupling,” 0806, 048 (2008). “Saturation in Deep Inelastic Scattering from AdS/CFT,” Phys. Rev. D 78, 096010 (2008).

[5] J. Polchinski and M. J. Strassler, “Deep inelastic scattering and gauge/string duality,” JHEP 0305, 012 (2003) arXiv:hep-th/0209211.

[6] R. C. Brower, M. Djuric, I. Sarcevic and C. I. Tan, “String-Gauge Dual Description of Deep Inelastic Scattering at Small-x,” JHEP 1011, 051 (2010) arXiv:1007.2259 [hep-ph]; M. S. Costa and M. Djuric, “Deeply Virtual Compton Scattering from Gauge/Gravity Duality,” arXiv:1201.1307 [hep-th].

[7] C. P. Herzog, S. Paik, M. J. Strassler and E. G. Thompson, “Holographic Double Diffractive Scattering,” JHEP 0808, 010 (2008) arXiv:0806.0181 [hep-th].

[8] R. C. Brower, S. D. Mathur and C. I. Tan, “Glueball Spectrum for QCD from AdS Supergravity Duality,” Nucl. Phys. B 587, 249 (2000).

[9] D. Kharzeev and E. Levin, “Soft double-diffractive Higgs production at hadron colliders,” Phys. Rev. D 63, 073004 (2001) arXiv:hep-ph/0005311.

[10] N. Anderson, S. K. Domokos, J. A. Harvey and N. Mann, “Central Production of Eta and Eta-prime via Double Pomeron Exchange in the Sakai-Sugimoto Model,” arXiv:1406.7010 [hep-ph].

[11] R. C. Brower, M. Djuric and C. I. Tan, “Diffractive Higgs Production by AdS Pomeron Fusion,” JHEP 1209, 097 (2012) arXiv:1202.4953 [hep-ph].