Electroproduction of two light vector mesons in next-to-leading BFKL* 

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Abstract
We calculate the amplitude for the forward electroproduction of two light vector mesons in next-to-leading order BFKL. This amplitude represents the first next-to-leading order amplitude ever calculated for a collision process between strongly interacting colorless particles.

1 Introduction
It is believed that the “gold-plated” measurement for the possible realization of the BFKL dynamics [2] is the $\gamma^*\gamma^*$ total cross section. However, a prediction for this cross section is not yet available with next-to-leading accuracy since the calculation of one necessary ingredient, the $\gamma^*$ to $\gamma^*$ impact factor in the next-to-leading order, has not yet been completed after year-long efforts [3]. Here we propose as “silver-plated” measurement the differential cross section for the $\gamma^*\gamma^*$ to two light vector mesons process in the forward case, i.e. for minimum momentum transfer.

In the BFKL approach, both in the leading logarithmic approximation (LLA), which means resummation of leading energy logarithms, all terms $(\alpha_s \ln(s))^n$, and in the next-to-leading approximation (NLA), which means resummation of all terms $\alpha_s(\alpha_s \ln(s))^n$, the (imaginary part of the) amplitude for a large-s hard collision process can be written as the convolution of the Green’s function of two interacting Reggeized gluons with the impact factors of the colliding particles (see, for example, Fig. 1).

The NLA singlet BFKL Green’s function for the forward case of interest here is known since several years [4]. Moreover, the impact factor for the transition from a virtual photon $\gamma^*$ to a light neutral vector meson $V = \rho^0, \omega, \phi$ has been recently calculated in the NLA in the forward case, up to contributions suppressed as inverse powers of the photon virtuality [5]. Therefore we have all is needed to build the NLA amplitude for the $\gamma^*\gamma^* \to VV$ reaction.

The knowledge of this amplitude is interesting first of all for theoretical reasons, since it could shed light on the role and the optimal choice of the energy scales entering the BFKL

*The content of this contribution is based on Ref. [1], to which we refer for additional details and for a more exhaustive list of references.
approach. Moreover, it could be used as a test-ground for comparisons with approaches different from BFKL, such as DGLAP, and with possible next-to-leading order extensions of other approaches, such as color dipole and $k_t$-factorization. But it could be interesting also for the possible applications to the phenomenology. Indeed, the calculation of the $\gamma^* \to V$ impact factor is the first step towards the application of BFKL approach to the description of processes such as the vector meson electroproduction $\gamma^* p \to V p$, being carried out at the HERA collider, and the production of two mesons in the photon collision, $\gamma^* \gamma^* \to V V$ or $\gamma^* \gamma \to V J / \Psi$, which can be studied at high-energy $e^+ e^-$ and $e\gamma$ colliders.

The process considered here has been studied recently in the Born ($2$-gluon exchange) limit in [6] and in the LLA [7] for arbitrary transverse momentum. In Ref. [7] also an estimate of NLA effects has been given in the forward case.

2 The NLA amplitude

We consider the production of two light vector mesons ($V = \rho^0, \omega, \phi$) in the collision of two virtual photons,

$$\gamma^*(p) \gamma^*(p') \to V(p_1) V(p_2).$$

Here, $p_1$ and $p_2$ are taken as Sudakov vectors satisfying $p_1^2 = p_2^2 = 0$ and $2(p_1 p_2) = s$; the virtual photon momenta are instead

$$p = \alpha p_1 - \frac{Q_1^2}{\alpha s} p_2, \quad p' = \alpha' p_2 - \frac{Q_2^2}{\alpha' s} p_1,$$

so that the photon virtualities turn to be $p^2 = -Q_1^2$ and $(p')^2 = -Q_2^2$. We consider the kinematics when

$$s \gg Q_{1,2}^2 \gg \Lambda_{QCD}^2,$$

and

$$\alpha = 1 + \frac{Q_1^2}{s} + \mathcal{O}(s^{-2}), \quad \alpha' = 1 + \frac{Q_2^2}{s} + \mathcal{O}(s^{-2}).$$

In this case vector mesons are produced by longitudinally polarized photons in the longitudinally polarized state [5]. Other helicity amplitudes are power suppressed, with a suppression factor $\sim m_V / Q_{1,2}$. We will discuss here the amplitude of the forward scattering, i.e. when the transverse momenta of produced $V$ mesons are zero or when the variable $t = (p_1 - p)^2$ takes its maximal value $t_0 = -Q_1^2 Q_2^2 / s + \mathcal{O}(s^{-2})$.

The forward amplitude in the BFKL approach may be presented as follows

$$\mathcal{I} m_s (A) = \frac{s}{(2\pi)^2} \int \frac{d^2 \tilde{q}_1}{\tilde{q}_1^2} \Phi_1(\tilde{q}_1, s_0) \int \frac{d^2 \tilde{q}_2}{\tilde{q}_2^2} \Phi_2(\tilde{q}_2, s_0) \int \frac{d\omega}{2\pi i} \left( \frac{s}{s_0} \right)^\omega G_\omega(\tilde{q}_1, \tilde{q}_2).$$

This representation for the amplitude is valid with NLA accuracy. The scale $s_0$ is artificial. It is introduced in the BFKL approach to perform the Mellin transform from the s-space to the complex angular momentum plane and must disappear in the full expression for the amplitude at each fixed order of approximation. Using the result for the meson NLA impact factor such cancellation was demonstrated explicitly in Ref. [5] for the process in question.

In Eq. (5), $\Phi_1(\tilde{q}_1, s_0)$ and $\Phi_2(\tilde{q}_2, s_0)$ are the impact factors describing the transitions $\gamma^*(p) \to V(p_1)$ and $\gamma^*(p') \to V(p_2)$, respectively. They are presented as an expansion in $\alpha_s$

$$\Phi_{1,2}(\tilde{q}) = \alpha_s D_{1,2} \left[ C_{1,2}^{(0)}(\tilde{q}^2) + \bar{\alpha}_s C_{1,2}^{(1)}(\tilde{q}^2) \right], \quad D_{1,2} = -\frac{4\pi \alpha_f f_V}{N_c Q_{1,2}} \sqrt{N_c^2 - 1},$$

(6)
where $\bar{\alpha}_s = \alpha_s N_c / \pi$, $f_V$ is the meson dimensional coupling constant ($f_\rho \approx 200$ MeV) and $e_q$ should be replaced by $e / \sqrt{2}$, $e / (3 \sqrt{2})$ and $-e / 3$ for the case of $\rho^0$, $\omega$ and $\phi$ meson production, respectively. In the collinear factorization approach the meson transition impact factor is given as a convolution of the hard scattering amplitude for the production of a collinear quark–antiquark pair with the meson distribution amplitude (DA). The integration variable in this convolution is the fraction $z$ of the meson momentum carried by the quark ($\bar{z} \equiv 1 - z$ is the momentum fraction carried by the antiquark):

$$C_{1,2}^{(0)}(\vec{q}^2) = \int_0^1 dz \frac{\vec{q}^2}{\vec{q}^2 + z\bar{z}Q_{1,2}^2} \phi(\bar{z}) .$$

The NLA correction to the hard scattering amplitude, for a photon with virtuality equal to $Q^2$, is defined as follows

$$C_{1,2}^{(1)}(\vec{q}^2) = \frac{1}{4N_c} \int_0^1 dz \frac{\vec{q}^2}{\vec{q}^2 + z\bar{z}Q_{1,2}^2} [\tau(z) + \tau(1 - z)] \phi(\bar{z}) ,$$

with $\tau(z)$ given in the Eq. (75) of Ref. [5]. $C_{1,2}^{(1)}(\vec{q}^2)$ are given by the previous expression with $Q^2$ replaced everywhere in the integrand by $Q_{1,2}^2$, respectively. Below we will use the DA in the asymptotic form, $\phi^{as}_|| = 6z(1 - z)$, both for the simplicity of the presentation and because, according to QCD sum rules estimates [8], $\phi^{as}_||$ may be indeed a good approximation for the DA of light vector mesons.

The Green’s function in (5) is determined by the BFKL equation

$$\delta^2(\vec{q}_1 - \vec{q}_2) = \omega G_{\omega}(\vec{q}_1, \vec{q}_2) - \int d^2\vec{q} K(\vec{q}_1, \vec{q}) G_{\omega}(\vec{q}, \vec{q}_2) ,$$

where $K(\vec{q}_1, \vec{q}_2)$ is the BFKL kernel. In the transverse momentum representation (see [1] for details), it can be written as

$$\hat{K} = \bar{\alpha}_s \hat{K}^0 + \bar{\alpha}_s^2 \hat{K}^1 ,$$

where $\hat{K}^0$ is the BFKL kernel in the LLA, $\hat{K}^1$ represents the NLA correction.
We find that
\[ \hat{K}^0|\nu\rangle = \chi(\nu)|\nu\rangle, \quad \chi(\nu) = 2\psi(1) - \psi\left(\frac{1}{2} + i\nu\right) - \psi\left(\frac{1}{2} - i\nu\right), \]  
(11)
is given by the following set of functions:
\[ \langle \bar{q}|\nu\rangle = \frac{1}{\pi^{\frac{1}{2}}}(\bar{q}^2)^{\nu - \frac{1}{2}}, \quad \langle \nu'|\nu\rangle = \int \frac{d^2\bar{q}}{2\pi^2} (\bar{q}^2)^{\nu' - \nu - 1} = \delta(\nu - \nu'). \]  
(12)
The action of the full NLA BFKL kernel on these functions may be expressed as follows:
\[ \hat{K}|\nu\rangle = \bar{\alpha}_s(\mu_R)\chi(\nu)|\nu\rangle + \alpha_s^2(\mu_R)\left(\chi^{(1)}(\nu) + \frac{\beta_0}{4N_c}\chi(\nu)\ln(\mu_R^2)\right)|\nu\rangle \]
\[ + \bar{\alpha}_s^2(\mu_R)\frac{\beta_0}{4N_c}\chi(\nu) \left(i\frac{\partial}{\partial\nu}\right)|\nu\rangle, \]  
(13)
where the first term represents the action of LLA kernel, while the second and the third ones stand for the diagonal and the non-diagonal parts of the NLA kernel. We refer to Ref. [1] for the expression of \(\hat{\chi}(\nu)\).

We will need also the \(|\nu\rangle\) representation for the impact factors, which is defined by the following expressions
\[ \frac{C_1^{(0)}(\bar{q}^2)}{\bar{q}^2} = \int_{-\infty}^{+\infty} d\nu' c_1(\nu')(\nu'|\bar{q}), \quad \frac{C_2^{(0)}(\bar{q}^2)}{\bar{q}^2} = \int_{-\infty}^{+\infty} d\nu c_2(\nu) \langle \bar{q}|\nu\rangle, \]  
(14)
\[ c_1(\nu) = \int d^2q C_1^{(0)}(\bar{q}^2)(\bar{q}^2)^{\nu - \frac{3}{2}}\pi^{\frac{1}{2}}, \quad c_2(\nu) = \int d^2q C_2^{(0)}(\bar{q}^2)(\bar{q}^2)^{-\nu - \frac{3}{2}}\pi^{\frac{1}{2}}, \]  
(15)
and by similar equations for \(c_1^{(1)}(\nu)\) and \(c_2^{(1)}(\nu)\) from the NLA corrections to the impact factors, \(C_1^{(1)}(\bar{q}^2)\) and \(C_2^{(1)}(\bar{q}^2)\).

Using the above formulas one can derive, after some algebra, the following representation for the amplitude
\[ \frac{\mathcal{I}m_s(A)}{D_1D_2} = \frac{s}{(2\pi)^2} \int_{-\infty}^{+\infty} d\nu \left(\frac{s}{s_0}\right)^{\alpha_s(\mu_R)\chi(\nu)} \alpha_s^2(\mu_R)c_1(\nu)c_2(\nu) \left[1 + \bar{\alpha}_s(\mu_R) \left(\frac{c_1^{(1)}(\nu)}{c_1(\nu)} + \frac{c_2^{(1)}(\nu)}{c_2(\nu)}\right)\right] \]
\[ + \bar{\alpha}_s^2(\mu_R) \ln\left(\frac{s}{s_0}\right) \left(\hat{\chi}(\nu) + \frac{\beta_0}{8N_c}\chi(\nu) \left[-\chi(\nu) + \frac{10}{3} + i\frac{d\ln(c_2^{(1)}(\nu))}{d\nu} + 2\ln(\mu_R^2)\right]\right) \]  
(16)

We find that
\[ c_{1,2}(\nu) = \left(\frac{Q_{1,2}^2}{\sqrt{2}}\right)^{\pm\nu - \frac{1}{2}} \left[\frac{\Gamma[\frac{3}{2} + \pm i\nu]}{\Gamma[3 + 2i\nu]}\right]^2 \frac{6\pi}{\cosh(\pi\nu)}. \]  
(17)

Using Eq. (16) we construct the series representation for the amplitude
\[ \frac{Q_1Q_2 \mathcal{I}m_sA}{D_1D_2}s = \frac{1}{(2\pi)^2}\alpha_s(\mu_R)^2 \]
\[ \times \left[b_0 + \sum_{n=1}^{\infty} \bar{\alpha}_s(\mu_R)^n b_n \left(\ln\left(\frac{s}{s_0}\right)^n + d_n(s_0, \mu_R) \ln\left(\frac{s}{s_0}\right)^{n-1}\right)\right], \]  
(18)
where the coefficients
\[
\frac{b_n}{Q_1 Q_2} = \int_{-\infty}^{+\infty} d\nu c_1(\nu)c_2(\nu) \frac{\chi^n(\nu)}{n!},
\] (19)
are determined by the kernel and the impact factors in LLA. The expression for the coefficients \(d_n\) can be easily determined from Eq. (16) and is given in Ref. [1].

One should stress that both representations of the amplitude (18) and (16) are equivalent with NLA accuracy, since they differ only by next-to-NLA (NNLA) terms.

It can be easily shown that the amplitude (18) is independent in the NLA from the choice of energy and strong coupling scales. It is also possible to trace the contributions to each \(d_n\) coefficient coming from the NLA corrections to the BFKL kernel and from the NLA impact factors (see Ref. [1] for details).

3 Numerical results

In this section we present some numerical results for the amplitude given in Eq. (18) for the \(Q_1 = Q_2 \equiv Q\) kinematics, i.e. in the “pure” BFKL regime. The other interesting regime, \(Q_1 \gg Q_2\) or vice-versa, where collinear effects could come heavily into the game, will not be considered here. We will emphasize in particular the dependence on the renormalization scale \(\mu_R\) and \(s_0\) in the NLA result.

In all the forthcoming figures the quantity on the vertical axis is the L.H.S. of Eq. (18), \(\mathcal{I}_{\alpha_s}(A)Q^2/(s D_1 D_2)\). In the numerical analysis presented below we truncate the series in the R.H.S. of Eq. (18) to \(n = 20\), after having verified that this procedure gives a very good approximation of the infinite sum for the \(Y\) values \(Y \leq 10\). We use the two–loop running coupling corresponding to the value \(\alpha_s(M_Z) = 0.12\).

We have calculated numerically the \(b_n\) and \(d_n\) coefficients for \(n_f = 5\) and \(s_0 = Q^2 = \mu_R^2\), getting
\[
\begin{align*}
b_0 &= 17.0664 \quad b_1 = 34.5920 \quad b_2 = 40.7609 \quad b_3 = 33.0618 \quad b_4 = 20.7467 \\
b_5 &= 10.5698 \quad b_6 = 4.54792 \quad b_7 = 1.69128 \quad b_8 = 0.554475 \\
\end{align*}
\]
\[
\begin{align*}
d_1 &= -3.71087 \quad d_2 = -11.3057 \quad d_3 = -23.3879 \quad d_4 = -39.1123 \\
d_5 &= -59.207 \quad d_6 = -83.0365 \quad d_7 = -111.151 \quad d_8 = -143.06 \\
\end{align*}
\]
(20)

In this case contributions to the \(d_n\) coefficients originating from the NLA corrections to the impact factors are
\[
\begin{align*}
d_1^{\text{imp}} &= -3.71087 \quad d_2^{\text{imp}} = -8.4361 \quad d_3^{\text{imp}} = -13.1984 \quad d_4^{\text{imp}} = -18.0971 \\
d_5^{\text{imp}} &= -23.0235 \quad d_6^{\text{imp}} = -27.9877 \quad d_7^{\text{imp}} = -32.9676 \quad d_8^{\text{imp}} = -37.9618 .
\end{align*}
\]
(21)

Thus, comparing (20) and (21), we see that the contribution from the kernel starts to be larger than the impact factor one only for \(n \geq 4\).

These numbers make visible the effect of the NLA corrections: the \(d_n\) coefficients are negative and increasingly large in absolute values as the perturbative order increases. The NLA corrections turn to be very large. In this situation the optimization of perturbative expansion, in our case the choice of the renormalization scale \(\mu_R\) and of the energy scale \(s_0\), becomes an important issue. Below we will adopt the principle of minimal sensitivity (PMS) [9]. Usually PMS is used to fix the value of the renormalization scale for the strong coupling. We suggest to use this principle in a broader sense, requiring in our case the
minimal sensitivity of the predictions to the change of both the renormalization and the energy scales, $\mu_R$ and $s_0$.

We replace for convenience in (18) $\ln(s/s_0)$ with $Y - Y_0$, where $Y = \ln(s/Q^2)$ and $Y_0 = \ln(s_0/Q^2)$, and study the dependence of the amplitude on $Y_0$.

The next two figures illustrate the dependence on these parameters for $Q^2=24$ GeV$^2$ and $n_f = 5$. In Fig. 2 we show the dependence of amplitude on $Y_0$ for $\mu_R = 10Q$, when $Y$ takes the values 10, 8, 6, 4, 3.

We see that for each $Y$ the amplitude has an extremum in $Y_0$ near which it is not sensitive to the variation of $Y_0$, or $s_0$. Our choice of $\mu_R$ for this figure is motivated by the study of $\mu_R$ dependence. In Fig. 3 we present the $\mu_R$ dependence for $Y = 6$; the curves from above to below are for $Y_0=3, 2, 1, 0$.

Varying $\mu_R$ and $Y_0$ we found for each $Y$ quite large regions in $\mu_R$ and $Y_0$ where the amplitude is practically independent on $\mu_R$ and $Y_0$. We use this value as the NLA result for the amplitude at given $Y$. The resulting curve is compared with the curve obtained from the LLA prediction when the scales are chosen as $\mu_R = 10Q$ and $Y_0 = 2.2$, in order to make the LLA curve the closest possible (of course it is not an exact statement) to the NLA one in the given interval of $Y$. The two horizontal lines in Fig. 4 are the Born (2-gluon exchange) predictions calculated for $\mu_R = Q$ and $\mu_R = 10Q$.

We stress that one should take with care BFKL predictions for small values of $Y$, since in this region the contributions suppressed by powers of the energy should be taken into account. At the lowest order in $\alpha_s$ such contributions are given by diagrams with quark exchange in the $t$-channel and are proportional in our case to $\alpha_{EM}\alpha_s f^2_\gamma/Q^2$. At higher orders power suppressed contributions contain double logarithms, terms $\sim \alpha_s^n \ln^{2n} s$, which can lead to a significant enhancement. Such contributions were recently studied for the total cross section of $\gamma^*\gamma^*$ interactions [10].

If the NLA (and LLA) curves in Fig. 4 are compared with the Born (2-gluon exchange) results, one can conclude that the summation of BFKL series gives negative contribution to
the Born result for $Y < 6$ if one chooses for the scale of the strong coupling in the Born amplitude the value given by the kinematics, $\mu_R = Q$. We believe that our calculations show that one should at least accept with some caution the results obtained in the Born approximation, since they do not give necessarily an estimate of the observable from below.

Another important lesson from our calculation is the very large scale for $\alpha_s$ (and therefore the small $\alpha_s$ itself) we obtain using PMS. It appears to be much bigger than the kinematical scale and looks unnatural since there is no other scale for transverse momenta in the problem at question except $Q$. Moreover one can guess that at higher orders the typical transverse momenta are even smaller than $Q$ since they ”are shared” in the many-loop integrals and the strong coupling grows in the infrared. In our opinion the large values of $\mu_R$ we found is not an indication of the appearance of a new scale, but is rather a manifestation of the nature of the BFKL series. The fact is that NLA corrections are large and then, necessarily, since the exact amplitude should be renorm- and energy scale invariant, the NNLA terms should be large and of the opposite sign with respect to the NLA. We guess that if the NNLA corrections were known and we would apply PMS to the amplitude constructed as LLA + NLA-corrections + NNLA-corrections, we would obtain in such calculation more natural values of $\mu_R$.

In the last years strong efforts have been devoted to the improvement of the NLA BFKL kernel as a consequence of the analysis of collinear singularities of the NLA corrections and by the account of further collinear terms beyond NLA [11]. This strategy has something in common with ours, in the sense that it is also inspired by renormalization-group invariance and it also leads to the addition of terms beyond the NLA. These extra-terms are large and of opposite sign with respect to the NLA contribution, so that they partially compensate the NLA corrections. The findings of the present work suggest, however, that the corrections to the impact factors heavily contribute to the NLA amplitude, being even dominating in some interval of non-asymptotically high energies. Moreover, by inspection of the structure of the amplitude in the regime of strongly asymmetric photon virtualities, one can deduce
that also the impact factors generate collinear terms which add up to those arising from the kernel, see e.g. Eqs. (84) and (85) of Ref. [5]. This leads us to the conclusion that in the approaches based on kernel improvement the additional information coming from impact factors should somehow be taken into account when available. These issues certainly deserve further investigation and we believe that useful hints in this direction can be gained from the study of the $\gamma^*\gamma^* \to VV$ amplitude in the regime of strongly ordered photon virtualities.

We conclude this Section with a comment on the possible implications of our results for mesons electroproduction to the phenomenologically more important case of the $\gamma^*\gamma^*$ total cross section. By numerical inspection we have found that the ratios $b_n/b_0$ we got for the meson case agree for $n = 1 \div 10$ at $1 \div 2\%$ accuracy level with the analogous ratios for the longitudinal photon case and at $3.5 \div 30\%$ accuracy level with those for the transverse photon case. Should this similar behavior persist also in the NLA, our predictions could be easily translated to estimates of the $\gamma^*\gamma^*$ total cross section.

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