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Diffractive large transferred momentum photoproduction of vector mesons

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Abstract. The large t behavior of the helicity amplitudes of diffractive photoproduction is estimated relying on models of the photon and meson light-cone wave functions and on the double-logarithmic approximation to the exchanged gluon interaction. The role of large-size color dipole contributions to the photon–meson transition impact factor is discussed.

1 Introduction

Hard diffractive vector meson production is one of the topics considered in the analysis of HERA experiments [1]. Data on photoproduction are available [2–5] extending to relatively large t.

Diffractive photoproduction at relatively large momentum transfer is a particular example of semi-hard processes determined by two large scales, $s \gg -t \gg m_V^2$.

Because of the large momentum transfer one expects that an essential part of the interaction can be described by perturbative QCD; in particular the BFKL approach should be applicable for calculating the diffractive exchange. The intriguing question is whether the coupling of this perturbative exchange to the scattering particles, the photon–meson impact factor, is dominated by short distance configurations.

In the case of J/Ψ production at large t the heavy quark mass guarantees the applicability of the perturbatively calculated impact factor [6,7]. In the present paper we address the question to what extent a perturbative calculation can represent the photon–light meson diffractive transition, in particular, whether this process can be treated in the picture of small-size dipole interaction.

The case of light meson diffractive photoproduction at large t has been considered in a number of papers [8–12].

The small dipole contributions to the impact factor of all helicities have been calculated in [10] by using distribution amplitudes for both the photon and the vector meson. It has been suggested that the experimentally observed dominance of the transversely polarized meson production may be explained by a sizable chirally-odd contribution in the photon and meson light-cone wave functions. Calculating the exchange by the leading $\ln s$ BFKL equation allows one to describe the t-dependence of the photoproduction cross sections [11]. The BFKL formulation of the helicity amplitudes has been presented in [12], and its phenomenological consequences in the next article of the same authors [13].

The aim of the present paper is to emphasize the role of the large dipole size in the photon–meson transition impact factor. We adopt an ansatz for the meson light-cone wave function used in previous studies of diffractive electroproduction [14–16] and a similar ansatz for the photon light-cone wave function. We point out the contribution with large momentum transfer carried by both of the exchanged gluons, where the $q\bar{q}$ dipole size is not suppressed by the large t.

We include the leading effect of the exchanged gluon interaction by approximating the BFKL equation down to the lower level of double-logarithmic $\ln s \ \ln t$ accuracy [19]. In view of the complexity of the amplitudes constructed from the leading $\ln s$ BFKL solution as presented in [12], our approximate treatment is a reasonable simplification in order to study particular contributions. It allows us to demonstrate the main impact of the exchanged gluon interaction and to estimate the importance of the large dipole-size contributions.

2 Effective dipole scattering at large t

We recall the general approach to hard diffraction. The amplitude of the diffractive process $\gamma \to V$ can be represented as the integral over the transverse momenta of gluons in the t-channel (impact representation):

$$M^{\lambda_i \lambda_f}(s,q) = s \int_{-\mathrm{i}\infty}^{\mathrm{i}\infty} \frac{\mathrm{d}\omega}{2\pi \mathrm{i}} F^{\lambda_i \lambda_f}(\omega, q)$$

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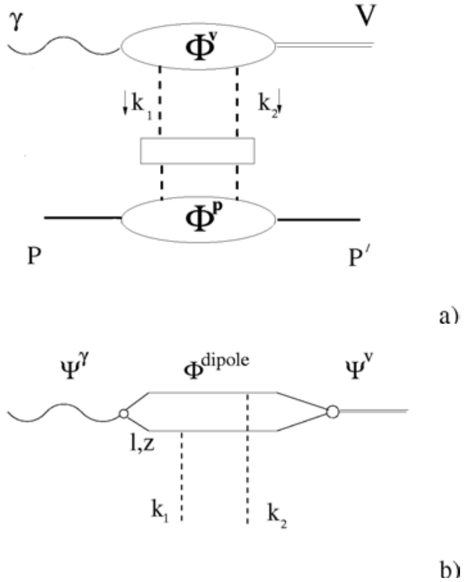

Fig. 1. a Impact factor form of the $\gamma \to V$ diffractive amplitude. **b** Contribution to the impact factor

$$\times \left[\left(\frac{s}{M^2(m,q)} \right)^{\omega} + \left(\frac{-s}{M^2(m,q)} \right)^{\omega} \right],$$

$$F^{\lambda_i \lambda_f}(\omega, q) \tag{2.1}$$
$$= \int \mathrm{d}^2\kappa \mathrm{d}^2\kappa' \Phi^{\lambda_i \lambda_f}(\kappa, q)\ \mathcal{G}(\kappa, \kappa', q, \omega)\ \Phi^P(\kappa', q).$$

Here q is the momentum transfer, κ, κ' the transversal momenta of the exchanged gluons, $\mathcal{G}$ the diffractive exchange (pomeron). $\Phi^{\lambda_i \lambda_f}$ and Φ^P are photon–meson and proton impact factors respectively.

The photon fluctuates into a $q\bar{q}$ pair a long time before, and this $q\bar{q}$ pair converts into the vector meson long time after the interaction with the proton. It is possible to represent the photon impact factor as the convolution of the impact factor for the $q\bar{q}$ dipole scattering with the light-cone wave functions of the incoming virtual photon and the outgoing vector meson (Fig. 1)

$$\Phi^{\lambda_i \lambda_f}(\kappa_1, \kappa_2)$$
$$= \int \mathrm{d}^2\ell_1 \mathrm{d}^2\ell_2 \mathrm{d}z \Psi^{(\gamma)\lambda_i}(\ell_1, z) \phi^{\mathrm{dip}}(\ell_1, \ell_2, \kappa_1, \kappa_2)$$
$$\times \Psi^{V^{\lambda_f *}}(\ell_2 - zq, z),$$
$$\phi^{\mathrm{dip}}(\ell_1, \ell_2, \kappa_1, \kappa_2)$$
$$= \alpha_{\mathrm{s}}[\delta^2(\ell_2 - \ell_1) + \delta^2(\ell_2 - \ell_1 + \kappa_1 + \kappa_2)$$
$$-\delta^2(\ell_2 - \ell_1 + \kappa_1) - \delta^2(\ell_2 - \ell_1 + \kappa_2)], \tag{2.2}$$

As the photon wave function we adopt the extrapolation of the virtual photon wave function. Whereas the latter is the result of a perturbative calculation, the extrapolation to $Q = 0$ is a model assumption:

$$\Psi^{(\gamma)\lambda}(\ell, z) = \frac{V^{\lambda}(\ell, z, 0) z\bar{z}}{|\ell|^2 + m_q^2},$$
$$V^{(+1)} = \frac{\ell^*}{z},\ V^{(-1)} = \frac{\ell}{\bar{z}}. \tag{2.3}$$

As vector meson wave function we can use

$$\Psi^{V\lambda}(\ell, z) = f_V \frac{V^{\lambda}(\ell, z, m_V)}{m_V^2} \exp\left[-\frac{|\ell|^2 + m_q^2}{z\bar{z} m_V^2} \right]. \tag{2.4}$$

This form has been used earlier [14–16]. It can be motivated by QCD sum rules, as resulting from the virtual photon wave function by a Borel transformation and by the substitution of the Borel variable by m_V^2. This wave function, being close to the one of the incoming photon, is a particular realization of the phenomenologically succesful concept of vector dominance.

Technically, the adopted forms of wave functions provide the advantage that the transverse momentum integration involved in the impact factor can easily be performed.

Unlike the case of electroproduction now the wave functions do not suppress the contributions from large dipole sizes. Such a suppression can result rather from the dipole impact factor involving the large momentum transfer. Transforming the dipole impact factor to coordinate representation we have

$$\int \mathrm{e}^{\mathrm{i}(\ell_1 r_1 - \ell_2 r_2)} \mathrm{d}\ell_1 \mathrm{d}\ell_2 \phi^{\mathrm{dip}}(\ell_1, \ell_2, \kappa_1, \kappa_2) \tag{2.5}$$
$$= \mathrm{e}^{\mathrm{i}zqr}\ (\mathrm{e}^{-\mathrm{i}\kappa_1 r} + \mathrm{e}^{-\mathrm{i}\kappa_2 r} - 1 - \mathrm{e}^{\mathrm{i}(\kappa_1 + \kappa_2) r})\ \delta^2(r_1 - r_2).$$

The large momentum transfer q leads to the dominance of small dipole sizes, $r_1 = r_2 = O(q^{-1})$, for generic values of the momenta κ_1, κ_2 ($\kappa_1 + \kappa_2 = -q$) of the exchanged gluons. This hard contribution to the photon–meson impact factor can be constructed with distribution amplitudes of both the photon and the vector meson. In [10] only this contribution has been considered. A particular feature is that one exchanged gluon carries a large, and the other a relatively small momentum, $\kappa_1 \ll q$ or $\kappa_2 \ll q$. Equation (2.5) shows that we have two further regions, where the dominant dipole size is not small. The first one is the vicinity of the end point, $z = 0, z = 1$. We have considered this contribution in the case of electroproduction [16]. The second one corresponds to small values of $\kappa_1 + zq$ or $\kappa_1 + \bar{z}q$. This means that here the large momentum transfer is shared by the two gluons. In this respect it is reminiscent of the Landshoff mechanism proposed for *pp* elastic scattering at large t [17]. We shall see that the photon–meson impact factor has extra terms which contribute to this region but are exponentially small outside of it.

3 Impact factor $\gamma_{\mathbf{real}} V$

First we want to consider the upper part of the diagram. We write down the photon impact factor in the following general form:

$$\Phi^{\lambda_i \lambda_f} = \int_0^1 \mathrm{d}z z\bar{z} \varphi_4^{\lambda_i \lambda_f}(z, \kappa, q), \quad (3.1)$$

$$\varphi_4(z, \kappa, q) = \varphi(z, \kappa, q) + \varphi(z, -\kappa - q, q) - \varphi(z, 0, q) - \varphi(z, -q, q), \quad (3.2)$$

where contributions of all four diagrams with different couplings of gluons to quarks are taken into account. Summing the diagrams with different momentum flow in effect we average over quark helecities in the quark loop.

For any contribution we have

$$\varphi(z, \kappa, q) = \frac{f_V}{m_V^2} \int \frac{\mathrm{d}^2\ell \langle V_i^{\lambda_i} V_f^{\lambda_f} \rangle}{\ell^2} \exp\left(-\frac{|\ell - (\kappa + zq)|^2}{m_V^2 z\bar{z}}\right) \quad (3.3)$$

The contractions of vertices for different helicities are

$$\langle V_i^1 V_f^1 \rangle = \ell^*(\ell - (\kappa + zq)) \left(\frac{1}{z^2} + \frac{1}{\bar{z}^2}\right), \quad (3.4)$$

$$\langle V_i^1 V_f^0 \rangle = \ell^* m_V \left(\frac{1}{z} - \frac{1}{\bar{z}}\right),$$

$$\langle V_i^1 V_f^{-1} \rangle = \ell^*(\ell - (\kappa - zq))^* \frac{2}{z\bar{z}}.$$

In (3.3) the integration over ℓ can be done without further approximation, e.g., in the case $\lambda_i = \lambda_f = 1$ it leads to

$$\int \mathrm{d}^2\widetilde{\ell} \frac{(\widetilde{\ell} - \widetilde{\kappa})\widehat{\ell}^*}{\widetilde{\ell}^2} \mathrm{e}^{-(\widetilde{\ell} - \widetilde{\kappa})^2} = \pi \mathrm{e}^{-\widetilde{\kappa}^2},$$

$$\widetilde{\ell} = \frac{\ell}{m_V \sqrt{z\bar{z}}}, \qquad \widetilde{\kappa} = \frac{\kappa + zq}{m_V \sqrt{z\bar{z}}}.$$

We obtain

$$\varphi^{1,1}(z, \kappa, q) = \pi f_V z\bar{z} \exp\left(-\frac{|\kappa + zq|^2}{m_V^2 z\bar{z}}\right) \left(\frac{1}{z^2} + \frac{1}{\bar{z}^2}\right) (3.5)$$

$$\varphi^{1,0}(z, \kappa, q) = 2\pi f_V m_V \frac{(\kappa + zq)^*}{|\kappa + zq|^2} \left(1 - \exp\left(-\frac{|\kappa + zq|^2}{m_V^2 z\bar{z}}\right)\right),$$

$$\varphi^{1,-1}(z, \kappa, q) = 2\pi f_V m_V \frac{(\kappa + zq)^{*2}}{|\kappa + zq|^2} \times \left(\left(1 + \frac{m_V^2 z\bar{z}}{|\kappa + zq|^2}\right) \exp\left(-\frac{|\kappa + zq|^2}{m_V^2 z\bar{z}}\right) - \frac{m_V^2 z\bar{z}}{|\kappa + zq|^2}\right).$$

Substituting into (3.3) we observe that there are terms contributing only in the vicinity of $|\kappa + zq| = 0$ or $|\kappa + \bar{z}q| = 0$ corresponding to Landshoff-type kinematics. We have hard contributions, $\kappa < q$, for $\lambda_f = 0$ and $\lambda_i = -\lambda_f$, which for $z = O(1)$ can be written as

$$z\bar{z}\varphi_4^{1,0}|_{\kappa \ll q} = \pi f_V \frac{\kappa}{q^2} \left(2 - \frac{1}{(z\bar{z})}\right), \quad (3.6)$$

$$z\bar{z}\varphi_4^{1,-1}|_{\kappa \ll q} = 2\pi f_V \frac{\kappa}{q^3} \left(-3 + \frac{1}{(z\bar{z})}\right).$$

The singularities at the end points are spurious. Actually the integration over z can be done with the result (3.5) without doing further approximations. However, the result can be represented approximately by a z-integral with (3.6) in the integrand and the range $\kappa/q < z < 1 - \kappa/q$. There is no hard chirally-even contribution to the impact factor $\lambda_i = \lambda_f = 1$ as the result of our particular choice of Ψ^γ. There is only the Landshoff-type contribution, at $\kappa' = \kappa + zq \ll q$:

$$z\bar{z}\varphi_4^{1,1}|_{\kappa \approx q} = \pi f_V (z^2 + \bar{z}^2) \exp\left(-\frac{|\kappa'|^2}{m_V z\bar{z}}\right), \quad (3.7)$$

and the analogous one at $\kappa + \bar{z}q \ll q$. There are extra Landshoff-type contributions to the other helicities. In the case of $\lambda_f = 0$ this results in a small contribution $O(q^{-5})$ to the amplitude and can be neglected.

4 BFKL in double-logarithmic approximation

Consider the BFKL equation [18] in the leading $\ln s$ approximation:

$$f(\omega, q, \kappa, \overline{\kappa}) = f_0 + \frac{g^2 N}{(2\pi)^3} \int \frac{\mathrm{d}^2\kappa'}{|\kappa'|^2 |q - \kappa|^2} K^0(\kappa', \kappa, q) f(\omega, q, \kappa', \overline{\kappa}) - \frac{g^2 N}{(2\pi)^3} [\alpha(\kappa) + \alpha(q - \kappa)] f(\omega, q, \kappa, \overline{\kappa}), \quad (4.1)$$

with the bare kernel

$$K^0(\kappa', \kappa, q) = \frac{\kappa_1' \kappa_1^* \kappa_2'^* \kappa_2 + \text{c.c.}}{|\kappa_1 - \kappa_1'|^2}. \quad (4.2)$$

We are going to simplify the equation in the double-log approximation, i.e., we shall approximate the transverse momentum integrals in $\ln t$ approximation [19]. In double-logarithmic approximation the gluon trajectory function can be written as

$$\alpha(\kappa) \approx \frac{g^2 N}{(2\pi)^3} \int_{\mu^2}^{|\kappa|^2} \frac{\mathrm{d}^2\kappa'}{|\kappa'|^2} = N \left(\frac{g^2}{4\pi}\right) \frac{1}{2\pi} \ln \frac{|\kappa|^2}{\mu^2}.$$

Taking also into account running of the coupling we get

$$\alpha(\kappa) \approx N \int_{\mu^2}^{|\kappa|^2} \frac{\mathrm{d}|\kappa'|^2}{|\kappa'|^2} \frac{\alpha_\mathrm{s}(|\kappa'|^2)}{2\pi} \equiv N\xi(\kappa). \quad (4.3)$$

According to Sect. 2 we want to consider two different kinematical cases.

(1) $\kappa_2' \approx \kappa_2 \approx q \gg \kappa_1' \gg \kappa_1$.

The evolution kernel in this region can be approximated as

$$K^0(\kappa', \kappa, q) = \frac{\kappa_1^*}{\kappa_1'^*} |q|^2 + \text{c.c.} \quad (4.4)$$

Replacing $f = \kappa \widetilde{f}$ the equation becomes

$$\widetilde{f} = \widetilde{f}_0 + \frac{g^2 N}{(2\pi)^3 \omega} \int_{\kappa^2}^{|\overline{\kappa}|^2} \frac{\mathrm{d}^2 \kappa'}{|\kappa'|^2} \widetilde{f}(\omega, \kappa'; \overline{\kappa}) - \dots$$

$$= \widetilde{f}_0 + N \int_{\xi(\kappa)}^{\xi(\overline{\kappa})} \mathrm{d}\xi' \widetilde{f}(\xi', \overline{\xi}) - \dots \qquad (4.5)$$

(2) $\kappa_1 \approx \kappa_1' \approx \kappa_2 \approx \kappa_2' \approx q \gg |\kappa_1 - \kappa_1'|$. We parameterize

$$\kappa_1 = zq + \widetilde{\kappa}; \quad \kappa_2 = \overline{z}q - \widetilde{\kappa}, \quad \kappa_1' = zq + \widetilde{\kappa}';$$
$$\kappa_2' = \overline{z}q - \widetilde{\kappa}' q \gg \widetilde{\kappa}' \gg \widetilde{\kappa}.$$

The kernel in this kinematics is approximately $K^0 = \frac{2|q|^4}{|\widetilde{\kappa}'|^2}$. Then the equation can be written as

$$f = f_0 + \frac{2g^2 N}{(2\pi)^3 \omega} \int_{|\mu|^2}^{|\overline{\kappa}|^2} \frac{\mathrm{d}^2 \widetilde{\kappa}'}{|\widetilde{\kappa}'|^2} f(\widetilde{\kappa}', \overline{\kappa}) - \dots$$

$$= f_0 + \frac{2N}{\omega} \int_{\xi(\kappa)}^{\xi(\overline{\kappa})} \mathrm{d}\xi' f(\xi', \overline{\xi}). \qquad (4.6)$$

In our approximation we can solve the evolution equations for both cases.

In the first case, substituting $\widetilde{f} \to \frac{\widetilde{f}}{\omega}$, we rewrite (4.5) as

$$(\omega + N\xi(\kappa) + N\xi(q))\, \widetilde{f}(\omega, q, \xi, \overline{\xi})$$
$$= \widetilde{f}_0 + N \int_{\xi(\kappa)}^{\xi(\overline{\kappa})} \mathrm{d}\xi' \widetilde{f}(\omega, \xi', \overline{\xi}),$$

and by another substitution,

$$\hat{f}(\omega, q, \xi, \overline{\xi}) \equiv (\omega + N\xi(\kappa) + N\xi(q))\, \widetilde{f}(\omega, q, \xi, \overline{\xi}),$$

we transform the equation into

$$\hat{f}(\omega, q, \xi, \overline{\xi}) = \widetilde{f}_0 + N \int_{\xi(\kappa)}^{\xi(\overline{\kappa})} \mathrm{d}\xi' \frac{\hat{f}(\omega, q, \xi', \overline{\xi})}{\omega + N\xi'(\kappa) + N\xi(q)}. \qquad (4.7)$$

The solution is

$$\widetilde{f}(\omega, q, \kappa, \overline{\kappa}) = \widetilde{f}_0 \left\{ \frac{1}{\omega + N\xi(\kappa) + N\xi(q)} + \frac{N\xi(\overline{\kappa}) - N\xi(\kappa)}{(\omega + N\xi(\kappa) + N\xi(q))^2} \right\}. \qquad (4.8)$$

Carrying out the Mellin transformation of this expression,

$$\mathcal{G}_{\mathrm{h}}(s, q, \kappa, \overline{\kappa}) = \int_{-\mathrm{i}\infty}^{\mathrm{i}\infty} \frac{\mathrm{d}\omega}{2\pi \mathrm{i}} \widetilde{f}(\omega, q, \kappa, \overline{\kappa}) \left(\frac{s}{|q|^2} \right)^{\omega},$$

we obtain the gluon exchange Green function of the scattering amplitude in double-logarithmic approximation:

$$\mathcal{G}_{\mathrm{h}}(s, q, \kappa, \overline{\kappa}) \qquad (4.9)$$

$$= \left(\frac{s}{|q|^2} \right)^{-N(\xi(\overline{\kappa}) + N\xi(\kappa))} \left(1 + N \ln \frac{s}{q^2} (\xi(\overline{\kappa}) - \xi(\kappa))\right)$$

In the other case of interest the evolution equation (4.6) is written as

$$f(\omega, q, \xi, \overline{\xi})$$
$$= \frac{f_0}{\omega} + \frac{2N}{\omega} \int_{\xi(\kappa)}^{\xi(\overline{\kappa})} \mathrm{d}\xi' f(\xi', \xi(\overline{\kappa})) - N(\xi(zq) + \xi(\overline{z}q))$$
$$\times f(\xi, \xi(\overline{\kappa})). \qquad (4.10)$$

Proceeding analogously to the previous case we get the differential equation

$$\frac{\mathrm{d}}{\mathrm{d}\xi} f(\omega, q, \xi, \overline{\xi}) = -\frac{2N f(\omega, q, \xi, \overline{\xi})}{\omega + N\xi(zq) + N\xi(\overline{z}q)}. \qquad (4.11)$$

The solution of this equation is

$$f(\omega, q, \xi(\kappa), \xi(\overline{\kappa})) = \frac{f_0}{(\omega + N\xi(zq) + N\xi(\overline{z}q))} \times \exp\left(\frac{2N(\xi(\overline{\kappa}) - \xi(\kappa))}{\omega + N\xi(zq) + N\xi(\overline{z}q)} \right). \qquad (4.12)$$

Carrying out the Mellin transformation we obtain the gluon exchange Green function:

$$\mathcal{G}_{\mathrm{L}}(s, q, \kappa, \overline{\kappa}) = \left(\frac{s}{|q|^2} \right)^{N(2\xi(\overline{\kappa}) - 2\xi(\kappa) - \xi(zq) - \xi(\overline{z}q))} \qquad (4.13)$$

5 Helicity amplitudes

At large t the diffractive exchange interacts with a single quark in the disintegrating proton. We write down the helicity amplitudes of diffractive scattering on a quark. In this case the proton impact factor reduces to a constant and the coupling of the exchange to the disintegrating proton does not influence the t-dependence. We obtain the diffractive γV amplitudes in terms of a sum of the hard and Landshoff-type contributions; each of them has the form

$$M^{\lambda_i \lambda_f} = \mathrm{i}s \int \mathrm{d}^2 \kappa \mathrm{d}z z \overline{z} \varphi_4^{\lambda_i \lambda_f}(\kappa, q, z) \frac{\mathcal{G}(s, q, \kappa, \overline{\kappa})}{|\kappa|^2 |\kappa - q|^2}. \qquad (5.1)$$

For comparison we consider first the amplitudes with simple two-gluon exchange, i.e. we substitute $\mathcal{G} = 1$, and we denote this by an additional subscript 1. The hard contribution to the cases $\lambda_f = 0$, and $\lambda_i = -\lambda_f$ can be calculated in this case without further approximation with the results

$$M_{1,h}^{1,0} = -2C m_V \pi^2 \left(4 - \frac{\pi^2}{3}\right) \frac{q}{t^2}, \qquad (5.2)$$

$$M_{1,h}^{1,-1} = C \frac{2\pi^2}{t^2} \left(\frac{2\pi^2}{3} - 8 \right).$$

The factor C denotes $C = \mathrm{i}s \frac{2}{3} \alpha_{\mathrm{s}}^2 e Q_q f_V \Phi_P$. The results for the hard contributions essentially coincide with the ones obtained in [10] for the corresponding chirally-even contributions.

The Landshoff-type contributions for the two-gluon case are

$$M_{1,\mathrm{L}}^{1,1} = \int \mathrm{d}^2\kappa' \int_0^1 \mathrm{d}z z\bar{z}\varphi_4^{11}(z,\kappa,q)\frac{1}{|\kappa|^2|q-\kappa|^2}$$
$$\approx C\int \mathrm{d}^2\kappa' \int_0^1 \mathrm{d}z\frac{\pi}{|q|^4}\exp\left(-\frac{|\kappa'|^2}{m_V^2 z\bar{z}}\right)\left(\frac{1}{z^2}+\frac{1}{\bar{z}^2}\right)$$
$$= C\frac{2\pi^2}{|q|^4}\int_{m_V/q}^{1-m_V/q}\left(\frac{1}{z\bar{z}}-2\right)\mathrm{d}z. \quad (5.3)$$

There is also a contribution to the double spin-flip amplitude which has to be added to the hard contribution:

$$M_{1,\mathrm{L}}^{1,-1} = C\frac{2\pi^2}{|q|^4}\int_{m_V/q}^{1-m_V/q}\left(\frac{1}{z\bar{z}}-3\right)dz, \quad (5.4)$$

and $M_{1,\mathrm{L}}^{1,0}$ is small compared to the hard contribution.

Now we evaluate the amplitudes including the double-log approximation to the BFKL evolution. We substitute $\mathcal{G}(s,k,q)$ by $\mathcal{G}_\mathrm{h}$, (4.9), for the hard contribution from the region $|\kappa| \ll |q|, |\kappa-q| \ll |q|$ and by $\mathcal{G}_\mathrm{L}$, (4.13), for the Landshoff-type contributions from the regions $|\kappa+zq| \ll |q|, |\kappa+\bar{z}q| \ll |q|$. The upper transverse momentum $\bar{\kappa}$ is of order q in $\mathcal{G}_\mathrm{h}$ and of order $\min(z,\bar{z})q$ in $\mathcal{G}_\mathrm{L}$. The integration is dominated by $k \approx m_V$. Therefore, the double-logarithmic interaction in the gluon exchange tends to suppress the hard contribution. There is no suppression in the Landshoff-type contribution for $z = O(1)$, but the end-point contributions are damped.

This affects the t-dependence of the amplitudes over a large t range. The combined effect of the end points and $\mathcal{G}_\mathrm{L}$ leads to a clear flattening of the t-dependence compared to the naive $1/|t|^2$ in the non-flip amplitude. In the small range $3 < |t| < 10\,\mathrm{GeV}^2$ the modification in the other amplitudes is small.

6 Numerical evaluation and discussion

The helicity amplitudes calculated above allow us to evaluate the t-dependence of the vector meson production cross section and of the angular-decay coefficients [20]. For the latter we use the relations

$$r_{00}^{04} \propto \frac{1}{N}|M^{10}|^2,$$
$$r_{10}^{04} \propto \frac{1}{2N}(M^{10*}M^{11}+M^{10*}M^{1-1}), \quad (6.1)$$
$$r_{1-1}^{04} \propto \frac{1}{N}(M^{1-1*}M^{11}),$$
$$N = |M^{10}|^2+|M^{11}|^2+|M^{1-1}|^2.$$

Our estimates are done for large t; therefore, it makes sense to extract the results for values of $|t|$ above $3\,\mathrm{GeV}^2$.

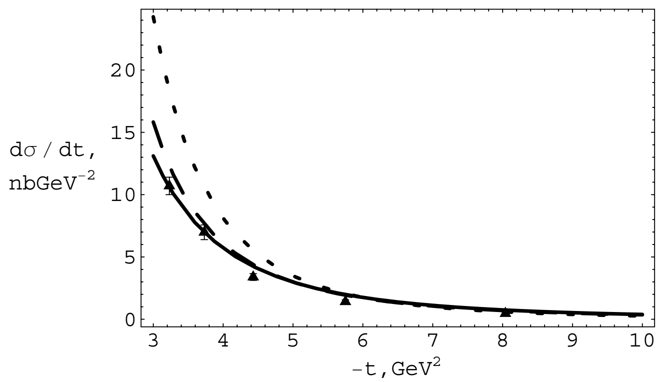

Fig. 2. The diffractive γV cross section. Two-gluon picture without double-log corrections (dense curve), with included double-logs (dashed curve), the result of [10] (dotted curve). Experimental points are due to ZEUS 2002

With our estimates we do not predict the normalization of the cross section. Besides of this, no parameters are fitted.

The predicted t-dependence of the cross section agrees reasonably with the data [2]. In Fig. 2 we show how the double-log interaction in the exchange improves the t-dependence.

We compare also with the t-dependence resulting from the amplitudes given in [10]; it deviates from the experimental behavior. We conclude that the combined effect of large-size dipole contributions and of exchange interaction improves the naive power behavior in t in agreement with experiment.

A more detailed description of the t-dependence has been achieved in [13], by including the leading ln s BFKL solution for the diffractive exchange and fitting some parameters. However, the angular-decay coefficients obtained there are not in full agreement with the data; in particular, r_{10}^{04} turns out to have the opposite sign. It has been pointed out that the right sign could be achieved by increasing the chirally-odd contribution parameterized there by the constituent quark mass, giving the latter an unreasonable large value.

In Fig. 3 we show our results on the angular-decay coefficients in comparison with the data and also with the predictions calculated from the results of [10]. It is clear that these coefficients are sensitive to the detailed structure of the amplitudes which is not resolved in the cross section.

We have seen that the relative magnitude of the amplitudes arising from the helicity dependence of the impact factors is influenced essentially by the end-point contributions and also by the interaction of the exchanged gluons. Neglecting these effects would result in the dominance of the longitudinal polarization of the produced vector meson. We obtain that in the experimentally accessible t range non-flip and single-flip are of approximately the same magnitude and t-dependence. The double-flip amplitude is about an order of magnitude smaller, due to partial cancellation of soft and hard contributions. The data on r_{00}^{04} suggest that the non-flip amplitude is even larger, exceeding the single-flip amplitude by a factor 3 to 4.

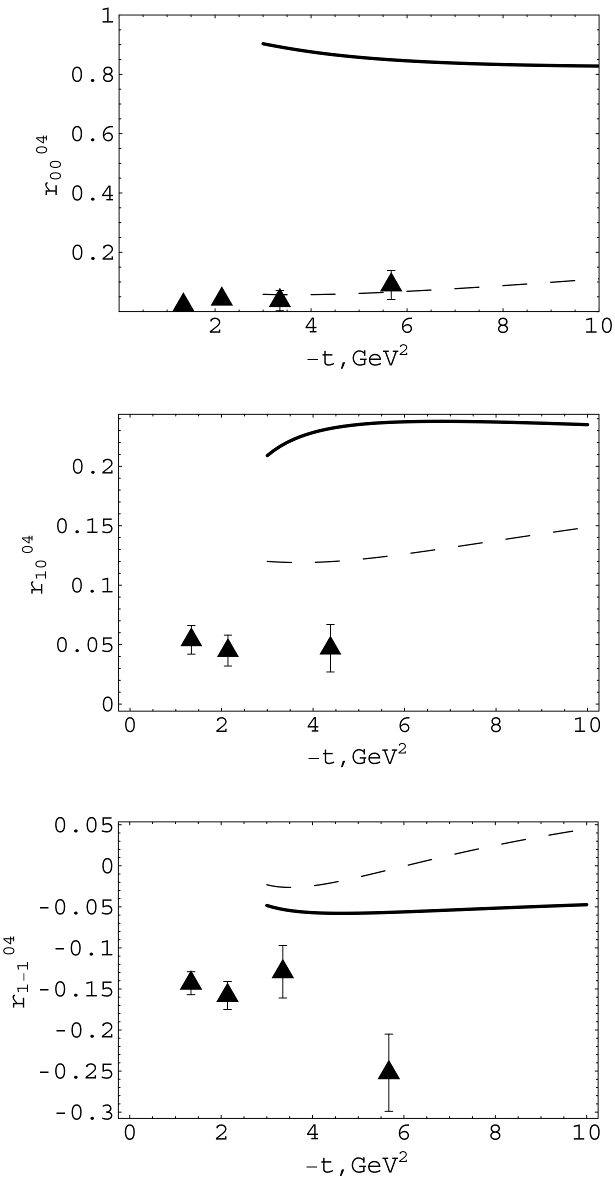

Fig. 3. The angular-decay coefficients r^{04}_{00}, r^{04}_{10}, r^{04}_{1-1}. The dotted curve is for results of [10], the dense curve presents the calculation. Experimental points are due to ZEUS 2002

Besides this our results on the angular-decay coefficient are in qualitative agreement with experiment and also with the results derived from [10]; in particular r^{04}_{10} agrees in sign with experiment. In the latter paper the inclusion of a sizable chirally-odd contribution was essential for the qualitative agreement with experiment, in particular for the observed dominance of the transverse polarization of the produced vector meson.

Our amplitudes do not include a chirally-odd contribution. In this way we have shown that the end-point and exchange interaction account partially for the correction achieved in [10] by the addition of a chirally-odd component.

In this paper we have emphasized the role of contributions with relatively large dipole size. We have estimated in the double-log approximation the dominant effect of the exchanged gluon interaction. Without including a chirally-odd contribution we obtain that the transverse and longitudinal production rates are of the same size in the t range of interest. The angular-decay distribution data suggest a stronger enhancement of the transverse production, leaving room for a sizable chirally-odd contribution to the photon and meson wave functions.

A lesson to be drawn from this study is that the large t transition impact factor has, besides the small dipole part, which can be parameterized by photon and meson distribution amplitudes, a part with relatively large dipole sizes. For the latter one has to introduce complete light-cone wave functions of the particles involved; their values in the vicinity of vanishing dipole size are not sufficient. The calculation of diffractive large t amplitudes requires the corresponding additional non-perturbative input.

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