September 1996 ULG-PNT-96-2-JRC

ELASTIC VECTOR-MESON PRODUCTION AT HERA

J.R. Cudell[1] and I. Royen[2]
Inst. de Physique, U. de Liège, Bât. B-5, Sart Tilman, B4000 Liège, Belgium

Abstract

We show that the lowest-order QCD calculation in a simple model of elastic vector-meson production does reproduce correctly the ratios of cross sections for ρ, ϕ and J/ψ, both in photoproduction and in high-Q^2 quasi-elastic scattering. The dependence of the slopes on the mass of the vector meson is reproduced as well. We examine the lower-energy data, and find that the energy dependence of the cross section does not depend on Q^2, but may depend on m_V.

[1]cudell@gw.unipc.ulg.ac.be
[2]royen@gw.unipc.ulg.ac.be

Elastic vector-meson production opens a precious window on the interface between perturbative QCD (pQCD) and non-perturbative hadronic physics, and is complementary to deep-inelastic scattering (DIS). Indeed, DIS has been until recently the triumph of perturbative ideas, whereas elastic processes were mostly treated through non-perturbative methods. These two processes now meet at HERA, where pQCD and Regge models have to be merged to obtain a full understanding of the data. Elastic vector-meson production has the extra advantage of containing by definition two scales, the mass of the produced vector meson, M_V, and the off-shellness of the photon, $Q^2 = -q^2$. We shall see that both dependences can be understood through a lowest-order calculation.

The data for $\gamma^* p \to Vp$, which both H1 [1, 2] and ZEUS [3] have obtained, for $V = \rho$, ϕ and J/ψ, exhibit the following main features: (i) the Q^2-distribution of the cross section is shallower for J/ψ than for ρ, the two cross sections becoming comparable around $Q^2 = 10$ GeV^2 and (ii) the t-slopes of the differential elastic cross sections depend on the mass of the produced meson, and become shallower as the mass increases.

Several models have been proposed to describe this process. Originally, Donnachie and Landshoff (DL) [4] extended their soft-pomeron model to predict ρ production at EMC [5]. This model works well there, but it has not been applied to predict the ratios of produced vector mesons. In photoproduction, DL preferred to resort to the quark counting rule to predict the ratios of cross sections [6], and reached the conclusion that it does not work perfectly, but argued that the violations were reasonable. The transition to perturbative QCD was first introduced by DL [7], who noticed the analogy between the pomeron expressions and two-gluon exchange, in the transverse case. They used a "constituent gluon" propagator and two-gluon exchange to model the pomeron. This analogy was pursued by one of us [8], who showed that such a model can give reasonable agreement with EMC data. This was later confirmed by NMC [5]. The final step to pQCD was performed by Ryskin [9] who observed that at high-Q^2 and high M_V, the effective intercept should be analogous to that found in $xg(x)$. This was later confirmed by Brodsky et al. [10]. So far, there has not been a model applied to the full range of masses and Q^2: this is the object of this letter.

A model for exclusive vector-meson production must include three sub-models, for which we shall adopt the simplest ones: the transition $\gamma^* \to V$ is described by a zero Fermi momentum wavefunction, the colour-singlet exchange is modeled à la Low-Nussinov [11], and we shall only consider the constituent quarks of the proton. There are in principle 72 diagrams contributing to the amplitude: the gluons can be hooked 4 different ways to the quarks of the vector meson, 9 different ways to the quarks of the proton, and both the direct and the $s \leftrightarrow u$ channels contribute to the amplitude in the high-s limit. As we shall explain however, the calculation of each part of the amplitude can be greatly simplified, so that one needs to calculate only the two diagrams shown in Fig. 1.

Let us first consider the kinematics of the process, which is spelled out in

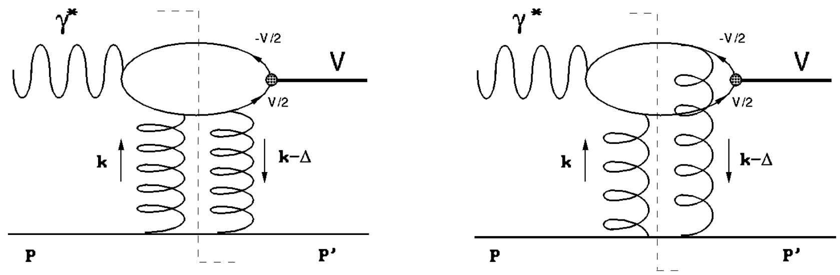

Figure 1: The two diagrams accounting for the transition $\gamma^* \to V$. The dashed line represents the cut which puts the intermediate state on-shell.

Fig. 1. Let P and q be respectively the 4-momenta of the proton and the photon, and $P - \Delta$ and $q + \Delta$ the momenta of the final-state proton and vector meson. We work in the high-w^2 limit, with $w^2 = (P+q)^2$. The on-shell condition for the proton and the vector meson imply that Δ is transverse to order $1/w^2$: $\Delta \approx \Delta_T$ with $\Delta_T \cdot P = \Delta_T \cdot q = 0$, and $|\Delta_T^2| \approx |\Delta^2| = -t$.

As it will turn out, the imaginary part of the amplitude is proportional to w^2. Crossing symmetry and analyticity then imply that the amplitude is purely imaginary, up to terms of order $1/w^2$, because the exchange is $C = +1$. Hence in the following, we shall calculate only the imaginary part, using Cutkovsky's rules, and putting intermediate states on-shell, as shown in Fig. 1. The quarks that make the vector meson are in the direction of q, whereas those that make the proton are parallel to P. The intermediate states come from the absorption of one gluon. In order for these states to be on-shell, the gluon must have vanishing components both in the direction of P and q, hence the gluon momenta are transverse to order $1/w^2$: $k \approx k_T$, with $k_T.P = k_T.q = 0$.

Hence the gluon momentum is essentially transverse, and furthermore we do not need to worry about the (purely real) crossed diagrams. The transversality of the gluons implies that the embedding of a process at the quark level into a proton is particularly simple. We follow here Gunion and Soper [12], and choose to represent the proton by a constituent model. This naturally leads to the quark counting rule, and has the extra advantage that the form factor of the proton is to a large extent measured. Indeed, when both gluons hit the same quark line, the form factor is given by the Dirac elastic form factor, measured in ep scattering:

$$\mathcal{E}_1(t = \Delta^2) \approx \frac{(3.53 - 2.79t)}{(3.53 - t)(1 - t/0.71)^2} \tag{1}$$

If the gluons hit different quark lines, then the form factor is not known, but we know that infinite wavelength gluons must average out the colour of the proton, hence we know that this form factor has to cancel the infrared singularity that would result from the pole in the gluon propagator, and therefore it must reduce to $\mathcal{E}_1(t)$ when either gluon goes on-shell. We choose to parametrise this form factor as:

$$\mathcal{E}_2(k, k-\Delta) = \mathcal{E}_1(k^2 + (k-\Delta)^2 + c\ k \cdot (k-\Delta)) \tag{2}$$

There are theoretical arguments [13] as well as phenomenological ones [14] which lead to the conclusion that $c \approx -1$.

The rule is then to calculate the process at the quark level. This leads to an integral over the transverse momenta of the gluons. One then introduces the difference of form factors $3(\mathcal{E}_1 - \mathcal{E}_2)$ into the integral to get the same process at the proton level, thereby reducing the number of diagrams by a factor 9. For the lower trace, and in the high-w^2 limit, we need to keep only the leading terms in p, the momentum of the quark, hence the trace along the quark line in the square of the amplitude, is given by $(4p^\alpha p^\beta) \times (4p^{\alpha'} p^{\beta'})$ (including a factor of 1/2 for spin averaging), with p the momentum of the quark inside the proton, and $\alpha^{(\prime)}$ and $\beta^{(\prime)}$ the indices of the gluons. Thus we can treat the process at the level of the amplitude, without the need to square it, provided that we write the contribution of the lower quark line as $4p^\alpha p^\beta$.

For the vector meson, we use a different model than for the proton, as we want to take into account the mass of the meson. The vector meson is modeled [15] by its lowest Fock state, with no Fermi momentum, which implies $m_q = m_V/2$, and the $V\bar{q}q$ vertex, including the two quark propagators, is given by:

$$\Phi_V = C'\ (m_q - \gamma.v)\ \gamma.e\ (m_q + \gamma.v) = C\ \gamma.e\ (m_V + \gamma.V) \tag{3}$$

where $v = \frac{V}{2}$ is the quark momentum within the vector meson and $C \equiv \sqrt{f_V/24}$ is the normalisation that reproduces the vector meson decay rate, with $f_V \approx 0.025(\mathrm{GeV})\ m_V$ the vector-meson decay constant squared. This effective vertex includes the propagators of the quark lines flowing into it. Reversing the direction of the quark current gives the same contribution, therefore we end up with only 2 diagrams to calculate - those shown in Fig. 1.

The traces corresponding to the upper bubble, dotted into $p_\alpha p_\beta$, are:

$$\mathcal{T}_1 = Tr[\not{e}\ (\gamma.V + m)\ \not{p}\ (\gamma.(q - v + k) + m_q)\ \not{p}\ (\gamma.(q - v) + m_q)\ \not{\epsilon}] \tag{4}$$

$$\mathcal{T}_2 = Tr[\not{e}\ (\gamma.V + m)\ \not{p}\ (\gamma.(q - v - k + \Delta) + m_q)\ \not{\epsilon} \times (\gamma.(-v - k + \Delta) + m_q)\ \not{p}] \tag{5}$$

with e and ϵ the polarisations of the vector meson and of the photon. The propagators of the off-shell quarks are: $p_1 = -(Q^2 + m_V^2 - t)/2$ and $p_2 = (m_V^2 + Q^2 - t - 4k.q)/2$. The answer is then proportional to: $\mathcal{T} = \mathcal{T}_1/p_1 + \mathcal{T}_2/p_2$. This answer is explicitly gauge invariant: substituting $\epsilon = q$ in $\mathcal{T}$ gives $0^\dagger$.

In the transverse case, the leading dot product is simply $\hat{w}^2 \equiv (p+q)^2 \approx 2p.q$. For longitudinal polarisation, further contributions appear: $\epsilon_L.p \approx -\hat{w}^2/(2\sqrt{Q^2})$

†Note that this was not the case for the expressions previously published in [8]. The difference at high-Q^2 amounts to a factor of 2 in the amplitude, presumably the discrepancy pointed out in [16].

and $e_L.p \approx -\hat{w}^2/(2m_V)$. Keeping track of these, the leading term is proportional to:

$$\begin{aligned}\mathcal{T} &= \frac{2m_V\ k\cdot(k-\Delta)}{t-m_V^2-Q^2+4\ k\cdot(k-\Delta)} \\ &\times\left(\frac{2\Delta.\epsilon\ e.p+2q.e\ \epsilon.p-\epsilon.e\ \hat{w}^2}{(m_V^2+Q^2-t)}-\frac{2}{\hat{w}^2}\ \epsilon.p\ e.p\right)\end{aligned} \tag{6}$$

Putting everything together, we obtain the following expression for the amplitude:

$$\mathcal{A} = -i\ \alpha_S^2\ g_V^{elm}\frac{2}{3}\sqrt{\frac{m_V\ f_V}{24}}\int\frac{d^2k_T}{k^2(k-\Delta)^2}\ 3[\mathcal{E}_1(t)-\mathcal{E}_2(k,k-\Delta)]\times 32\mathcal{T} \tag{7}$$

where $g_V^{elm} = \xi\sqrt{4\pi\alpha_{elm}}$ is the electromagnetic coupling of the different vector mesons: $\xi = \frac{1}{\sqrt{2}}$ for the ρ, $-1/3$ for the ϕ and $2/3$ for the J/ψ. For the various possible helicities, Eq. (7) gives:

$$\begin{aligned}\mathcal{A}(T\to T) &= i\hat{w}^2\frac{64\alpha_S^2\ g_V^{elm}m_V\sqrt{m_V\ f_V}}{\sqrt{6}} \\ &\times\int\frac{d^2k_T}{k^2(k-\Delta)^2}\frac{[\mathcal{E}_1(t)-\mathcal{E}_2(k,k-\Delta)]\ [k\cdot(k-\Delta)]}{(t-m_V^2-Q^2+4\ k\cdot(k-\Delta))\ \ (m_V^2+Q^2-t)}\end{aligned} \tag{8}$$

$$\mathcal{A}(L\to L) = \frac{\sqrt{Q^2}}{m_V}\times\mathcal{A}(T\to T) \tag{9}$$

The helicity violating amplitude $\mathcal{A}(L\to T)$ is suppressed by $1/\hat{w}^2$. As previously advertised, this answer is proportional to $\hat{w}^2$, therefore to w^2, and is thus purely imaginary. The resulting cross section is independent of w^2. Clearly this model cannot say anything about the energy dependence of the cross section. We shall assume that it comes in as a factor, R, and check whether the latter is mass- or Q^2-dependent. Note that we can only determine the value of that factor times α_S^2. In the following, we shall let α_S run with the off-shellness of the gluons, and freeze it at some value α_S^0. However, as we shall see later, the dominant contribution comes from gluons of small off-shellness, and the results we obtain are identical to fixed-coupling results for $\alpha_S = \alpha_S^0$.

The differential cross section is given by:

$$\frac{d\sigma}{dt} = \frac{d\sigma_T}{dt}+\varepsilon\frac{d\sigma_L}{dt} = \frac{R}{16\pi(\hat{w}^2)^2}\left[|A(T\to T)|^2+\varepsilon|A(L\to L)|^2\right] \tag{10}$$

with ε the polarisation of the photon beam: $\varepsilon \approx 1$ at HERA and $\varepsilon \approx 0.75$ at NMC.

We first give the results that we obtain for the various cross sections measured by ZEUS and H1. First of all, we show in Fig. 2 the dependence on Q^2 and m_V of

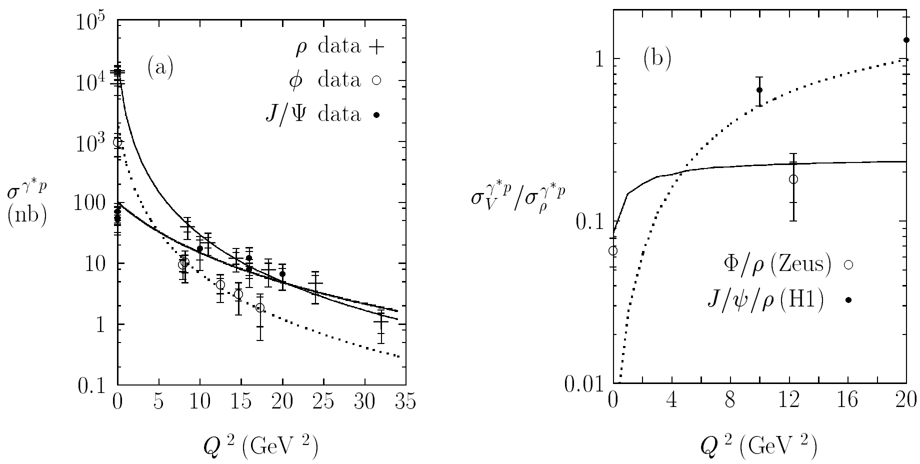

Figure 2: (a) Cross sections as functions of Q^2, compared with data from H1 [1, 2] and ZEUS [3] at $< w >\approx 100$ GeV, (b) Ratio of cross sections as functions of Q^2 at $< w >\approx 100$ GeV, compared with data from H1 [1, 2] and Zeus [3].

the integrated elastic cross section $\sigma(Q^2)$. We see that a common (Regge) factor is consistent with the data taken at HERA, as shown in Fig. 2(a). We insist on the fact that this factor is independent both of Q^2 and of m_V, as one would expect within Regge theory. Selecting only high-Q^2 data leads to a best value $R\alpha_S^2 = 0.6$, with a $\chi^2/d.o.f. = 0.39$. We do not find that $R\alpha_S^2$ varies significantly within the energy range of HERA. Although we see no reason why our model should work in photoproduction, it turns out that our curves do go through the photoproduction points. Including these in the fit brings $R\alpha_S^2$ to 0.55.

meson	A (nb)	n
ρ	5534	2.5
ϕ	1035.9	2.4
J/ψ	269.53	1.5

Table 1: Result of a fit of $\sigma(Q^2)$ to $A(Q^2)^{-n}$, for 5 GeV$^2 < Q^2 <$ 25 GeV^2.

We give in Table 1 the result of a fit to a power of Q^2 at large Q^2. Although asymptotically all cross sections behave like $(Q^2)^{-3}$, in agreement with [9, 10], we see that the data collected at HERA are not yet in that asymptotic regime. Note that our calculation holds only for $Q^2 << \hat{w}^2$, and hence it is not clear whether the asymptotic regime will be reachable at HERA. Our model is equivalent to those of refs [9, 10], with a gluon distribution corresponding to 2-gluon exchange $g(x) \sim 1/x$. We see that here the Q^2 dependence is already reproduced, hence there is no room for an extra dependence coming from the gluon distribution, and the use of the asymptotic formula [17] is misleading.

In Fig. 2(a), we show both the systematic and the statistical errors. In the ratio of cross sections, some of the systematic uncertainties cancel, and the reproduction of that ratio is a more stringent test of our model, especially as the normalisation then drops out of our prediction. We show in Fig. 2(b) the result of such a comparison. Again, we see that our model fares well, even in photoproduction.

Hence we can understand both the m_V- and the Q^2-dependence of the cross sections. One might object that this is because these are concentrated at low t, and argue that the t-dependence has to be wrong, as this is one of the well-known problems of perturbative calculations applied to diffractive scattering.

The behaviour of the slopes as a function of m_V and Q^2 can be understood as follows. We can approximate the proton form factor $\mathcal{E}_1(t) - \mathcal{E}_2(k, k-\Delta)$ as being proportional to $\mathcal{E}_1(t)[k\cdot(k-\Delta)]$, using the fact that F_1 is close to an exponential, and expanding for small $k\cdot(k-\Delta)$. The amplitude of Eq. (7) can then be written as $C(m_V, Q^2)\sqrt{R(t)}\mathcal{E}_1(t)\mathcal{F}\left(\frac{t}{Q^2+m_V^2}\right)$, with C a constant with respect to t, and $\mathcal{F}$ a calculable function. This means that the logarithmic derivative of $d\sigma/dt$ becomes:

$$b(t) = \frac{d\mathcal{E}_1^2}{\mathcal{E}_1^2 dt} + \frac{dR}{Rdt} + \frac{1}{(m_V^2+Q^2)}\frac{d\mathcal{F}^2}{\mathcal{F}^2 dx}\Big|_{x=\frac{t}{Q^2+m_V^2}} \tag{11}$$

Thus the slope is approximatively made of three terms: one corresponding to the proton response, one to the pomeron response, and one to the response of the loop which converts the photon into a vector meson. Only the latter depends on m_V and Q^2, and we see that it decreases rather fast with both of these factors. At large $Q^2+m_V^2$, it becomes negligible, and only the first two responses matter. This means that this kind of model predicts that all the slopes have to reach the same asymptotic value. This value is about 4 GeV^{-2}.

This variation of the slopes with $Q^2+m_V^2$ enables us to reproduce the measurements of HERA: in photoproduction, we obtain $1/<t_\rho>= 8.75$ GeV^{-2}, $1/<t_\phi>= 7.35$ GeV^{-2} and $1/<t_{J/\psi}>= 4.55$ GeV^{-2}, whereas the ρ slope is already down to 5 GeV^{-2} for $Q^2 = 5$ GeV^2. This is inconsistent with the data from H1, but agrees with those from Zeus. We want to point out that the experimental evaluation of the slope demands that the cross sections be exponential in t. We have quoted our results for $1/<t>$, which in the case of an exponential fall-off e^{bt} is equal to the logarithmic slope b. The fact that the differential cross section is not an exponential makes the comparison with data difficult. To illustrate the effect of the curvature, we compare in Fig. 3 our results with the data for ρ photoproduction in H1. We see that although the slope is supposed to be 10.9, our curve reproduces the data fairly well, with a $1/<t>$ of 8.75 GeV^{-2}. Hence the different definition of the slope leads to a possible systematic correction of about 2 GeV^{-2}. No matter which model is used, $d\sigma/dt$ is *not* an exponential, and we urge the experimentalists to quote a $<t>$ instead of a logarithmic slope. In fact, our model works although it assumes no shrinkage. There is room in the

ρ case to add a slope coming from the Regge factor R, which would contribute about 4.6 GeV^{-2} to b for a pomeron slope $\alpha' = 0.25$ GeV^{-2}. One would then expect a similar contribution in the J/ψ case. Hence if the H1 measurements of ρ slopes are correct, it is likely that the J/ψ numbers have a background problem: inelastic contributions would significantly reduce the slope.

The only problem at HERA seems to be the helicity structure of the cross section. The data support the prediction that helicity is conserved, but the ratio σ_L/σ_T does not follow the results of our model. This ratio has to behave as Q^2 for near-shell photons, as a consequence of gauge invariance. Our model fulfills this requirement, but it also predicts that this linear behaviour continues for all Q^2. The data, on the other hand, seem to indicate that the ratio reaches a plateau around 2 at high Q^2. This would indicate that our high-Q^2 transverse cross section is wrong. This is indeed possible: we have assumed that the meson wave function is dominated at high Q^2 by configurations in which both quarks have equal momenta. In fact, it is likely that further configurations exist [18] which would give additional contributions to the transverse cross section. This may account for the fact that this model does not reproduce the ration σ_L/σ_T measured at HERA. We plan to examine the role of Fermi momentum in a later paper. Another interesting possibility has recently been pointed out [17]: the form factor describing the recombination of quarks into a ρ meson is taken to be 1 once the quarks are restricted to the mass region of the ρ, and it is argued that the process will remain largely elastic. The ratio σ_T/σ_L is then shown, in the leading-log approximation, to be distorted, at $t = 0$, by the x and Q^2 dependence of the gluon distribution, in such a manner that it reproduces the data. This is the only model so far which manages to reproduce the observed ratio, but it can be used only for ρ mesons at $t = 0$, and heavily relies on the perturbative evolution of the gluon distribution.

Finally, we can now examine the w^2-dependence of the cross sections. We have seen that at HERA, the Regge factor does not seem to depend either on the meson mass or on Q^2. This is clearly reminiscent of the behaviour expected from a simple pole. We adopt the same philosophy when fitting to lower-energy cross sections from EMC and NMC [5], and write $R = s^{2\alpha_0 - 2}$, α_0 being the intercept of the pomeron. We show in Fig. 4 the curves which correspond to $\alpha_0 = 1.16$. This value of the soft intercept is rather high, but may not be entirely excluded by fits to total cross sections, especially once the effect of unitarisation is taken into account [19, 20]. We see that the fit is reasonable for the ρ and the ϕ, but fails for the J/ψ. The fact that the ρ data is somewhat high can be understood from the fact that there are contributions from lower trajectories to the ρ production cross section. The interference between a/f exchange and pomeron exchange could contribute as much as 20%, hence a high ρ cross section. In the J/ψ case, one has to realise that the calculation we have presented here is valid if $Q^2 + m_V^2 << \hat{w}^2$. This is not the case for the J/ψ at EMC. Besides, the data for J/ψ production from EMC is not subtracted for any diffractive background, and

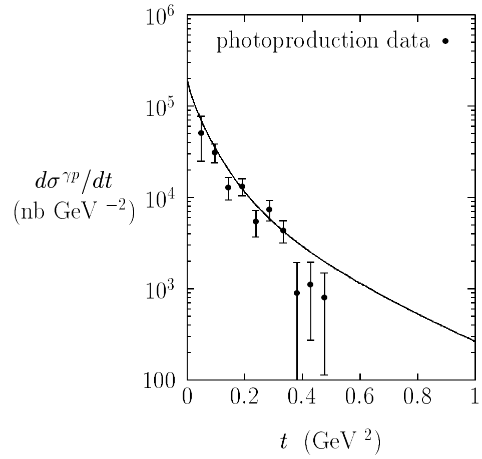

Figure 3: Photoproduction differential cross section $d\sigma/dt$, compared with H1 data [1].

can only constitute an upper bound on the elastic cross section. It is likely that there is a contamination of the order of a factor 2 from inelastic contributions, as the inelastic background is larger than in the ρ case, for which the EMC data were severely contaminated [5].

To check whether the behaviour is indeed that of a simple pole, we can fit the factor R separately for each meson. We give in Table 2 the best factors, and the corresponding intercepts, given the value of $R\alpha_S^2$ at HERA. We do not give the results for the J/ψ as it is impossible to obtain a good fit to the data. One sees that the ϕ and ρ data require somewhat different factors, This can be seen more clearly when one notices that the ratio r of cross sections has changed when going from HERA ($r = 0.18$) to NMC ($r = 0.12$). Hence there may be a problem with the smallness of the ϕ cross section at NMC, which seems to imply an intercept too large to be compatible with that of a soft pomeron. On the other hand, there is no sign in either ρ or ϕ data of a Q^2 dependence of the intercept.

meson	$R\alpha_S^2$	$\chi^2/d.o.f.$	energy (GeV)	intercept
ρ	0.4	0.6	$40 < \nu < 180$	1.10
ϕ	0.268	0.138	$40 < \nu < 180$	1.206
combined	0.32	1.19	$40 < \nu < 180$	1.161

Table 2: best values of $R\alpha_S^2$ at NMC.

Before concluding, we must mention that although the above looks like a successful perturbative calculation, most of the contribution to the total cross section comes from the infrared region. Imposing a cut-off $|k^2|, |(k-\Delta)^2| >$

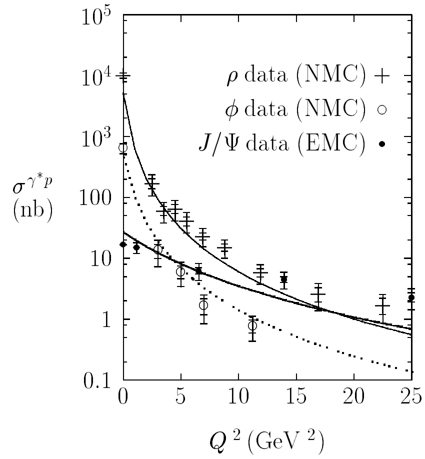

Figure 4: Cross sections compared to lower-energy EMC and NMC data [5], at $< w^2 > \approx 200$ GeV^2.

1 GeV^2 reduces the cross section by a factor 10. This dominance of the infrared region justifies a posteriori our choice of scale in α_S: we see no theoretical reason to make it run with either Q^2 or m_V^2, as these scales are unrelated to the off-shellnesses of the gluons entering the vertices. The dominance of the infrared region confirms the results of [21], where it is argued that the BFKL perturbative resummation is dominated by the non-perturbative region.

The simplest modification to the infrared region follows the ideas of Landshoff and Nachtmann [22], which have recently been further motivated by lattice studies [23], that the gluon propagator needs to be modified at low k^2, taming its behaviour to something softer than a pole. One of the main effects of that modification is that the differential cross section becomes much more linear, and that as a consequence $< t >$ becomes bigger. We have checked that such a model has all the other features of the one detailed above. The only major difference is that the slopes reach their asymptotic values much sooner, but it remain impossible to accommodate both the ρ shrinkage and the J/ψ data.

To sum up, we have shown that many features of elastic vector-meson production at HERA can be understood in a simple QCD model. We have also seen that comparison with lower-energy data proves somewhat problematic. The HERA data seem to indicate that the Regge factor does not depend either on Q^2 or m_V. The comparison with NMC ρ data points to a soft pomeron intercept, whereas the ϕ data would favor a larger one. We do not believe that much can be concluded from the J/ψ data, which seem to have an inelastic background. Following [14] we plan to model this background in a further publication.

Acknowledgments

We thank the H1 collaboration for making their data available to us, and in particular Pierre Marage, R. Roosen, P. Van Esch and B. Clerbaux for their hospitality and for numerous fruitful discussions. We also thank Peter Landshoff, Joseph Cugnon and Boris Kopeliovich for their comments and suggestions.

References

[1] H1 collaboration, S Aid et al., *Nucl. Phys.* **B468** (1996) 3; preprint DESY-96-037 (March 1996), e-Print Archive: hep-ex/9603005; *Nucl. Phys.* **B463** (1996) 3; H1 collaboration, T. Ahmed et al., *Phys. Lett.* **B338** (1994) 507.

[2] H1 collaboration, preliminary results : presented at the 28^{th} International Conference on High Energy Physics ICHEP'96, Warsaw, Poland, July 1996; also presented at the Topical Conference on Hard Diffractive Processes, Eilat, February 18-23 1996 and at the International Workshop on Deep Inelastic Scattering and related Phenomena, Roma, April 15-19, 1996. We are indebted to B. Clerbaux for making these data available to us.

[3] ZEUS collaboration, M. Derrick et al., *Phys. Lett.* **B356** (1995) 601; preprint DESY-96-067 (Apr 1996), e-Print Archive: hep-ex/9604008; *Phys. Lett.* **B350** (1995) 120; preprint DESY-96-002 (Jan. 1996), e-Print Archive: hep-ex/9601009; *Z. Phys.* **C69** (1995) 39.

[4] A. Donnachie and P.V. Landshoff, *Phys. Lett.* **185B** (1987) 403.

[5] A. Sandacz, preprint SINS-9-VI (April 1995); New Muon Collaboration, M. Arneodo et al., *Phys. Lett.* **B332** (1994) 195; *Phys. Lett.* **B332** (1994) 195; *Nucl. Phys.* **B429** (1994) 503; New Muon Collaboration, A. Sandacz et al., *Nucl. Phys.* **B429** (1994) 503.

[6] A. Donnachie and P.V. Landshoff, *Nucl. Phys.* **B311** (1989) 509.

[7] A. Donnachie and P.V. Landshoff, *Phys. Lett.* **B348** (1995) 213.

[8] J.R. Cudell, *Nucl. Phys.* **B336** (1990) 1.

[9] M.G. Ryskin, *Z. Phys.* **C57** (1993) 89; M.G. Ryskin, R.G. Roberts, A.D. Martin and E.M. Levin, preprint RAL-TR-95-065 (Nov. 1995), e-Print Archive: hep-ph/9511228.

[10] S.J. Brodsky, L. Frankfurt, J.F. Gunion, A.H. Mueller and M. Strikman, *Phys. Rev.* **D50** (1994) 3134.

[11] F. Low, *Phys. Rev.* **D12** (1975) 163; S. Nussinov, *Phys. Rev. Lett.* **34** (1975) 1286.

[12] J.F. Gunion and D. Soper, *Phys. Rev.* **D12** (1975) 163;

[13] J.R. Cudell and B.U. Nguyen, *Nucl. Phys.* **B420** (1994) 669. E.M. Levin and M.G. Ryskin, *Sov. J. Nucl. Phys.* **34** (1981) 619.

[14] J.R. Cudell and O.F. Hernández, *Nucl. Phys.* **B471** (1996) 471.

[15] A. Donnachie, R.R. Horgan, P.V. Landshoff, *Z. Phys.* **C10** (1981) 71 and *Phys. Lett.* **107B** (1981) 385; R.R. Horgan, P.V. Landshoff, D.M. Scott, *Phys. Lett.* **110B** (1982) 493.

[16] J.M. Laget and R. Mendez-Galain, *Nucl. Phys.* **A581**(1995) 397.

[17] A.D. Martin, M.G. Ryskin and T. Teubner, University of Durham preprint DTP/96/64, electronic archive hep-ph/9609448 (Sept. 1996).

[18] J. Nemchik, N.N. Nikolaev, E. Predazzi and B.G. Zakharov, preprint DFTT-71-95 (May 1996), e-Print Archive: hep-ph/9605231.

[19] J.R. Cudell, K. Kang and S.K. Kim, preprint SNUTP-95-103 (Jan. 1996), e-Print Archive: hep-ph/9601336, to appear in *Phys. Lett. B.*

[20] R.J.M. Covolan, J. Montanha and K. Goulianos, Rockefeller University preprint RU 96/E-34.

[21] J.R. Cudell, A. Donnachie and P.V. Landshoff, preprint ULG-PNT-96-1-JRC (Feb 1996), e-Print Archive: hep-ph/9602284, to appear in *Nucl. Phys. B.*

[22] P.V. Landshoff and O. Nachtmann, *Z. Phys.* **C35** (1987) 405.

[23] UKQCD Collaboration (D.S. Henty et al.), *Phys. Lett.* **B369** (1996) 130 and preprint EDINBURGH-96-12 (May 1996), e-Print Archive: hep-ph/9607258.