LUTP 92 - 12
March 1992

Diffractive J/Ψ-electroproduction in LLA QCD.

M. G. Ryskin

Department of Theoretical Physics, University of Lund
Sölvegatan 14A, S-22362 Lund, Sweden
and
St. Petersbourg Nuclear Physics Institute
188350, Gatchina, St. Petersbourg , Russia

Abstract: Cross section of diffractive J/Ψ-production in deep inelastic scattering in the Born and the leading log approximations of perturbative QCD are calculated.

1 Introduction

The process of J/Ψ -electroproduction arouses interest due to two reasons. First it can be calculated within the perturbative QCD and second its cross section is proportional to the gluon structure function. So it is a good way to study the gluon distribution inside a proton [1, 2].
In the reactions of heavy quark photoproduction $\gamma N \to c\bar{c}X$, a popular approach is the "photon-gluon fusion" mechanism [3, 1, 4, 5] based on the subprocess $\gamma g \to c\bar{c}$. The amplitude and cross section of inelastic J/Ψ-production via the same mechanism was calculated in ref. [6] and then discussed in ref.[7]. This approach has been called [5] diffractive J/Ψ-production, as (in the first approximation) the cross section does not depend on energy and there is no flavor exchange. Strictly speaking this is not a true diffractive process. There is a colour exchange in this case due to the colour of the gluon content in the target, and as a consequence, the inclusive J/Ψ cross section $\frac{d\sigma}{dz} \to const$ at $z \to 1$ instead of the $\delta(1-z)$ or $1/(1-z)$ behaviours that are usual for diffractive processes (z is the part of photon momenta carried away by the J/Ψ-meson).

The goal of this paper is to consider the exclusive (in some sense elastic) diffractive J/Ψ-electroproduction that is described by the exchange of a colourless two gluon system [1]; in the Born approximation by the diagrams in fig.1. In the leading log approximation (LLA) instead of the simple two gluon "Pomeron"[9] one has to use the whole system of LLA ladder diagrams that for $t=0$ reproduces exactly the gluon structure function $\bar{x}G(\bar{x},\bar{q}^2)$. Thus our amplitude is proportional to $\bar{x}G(\bar{x},\bar{q}^2)$ and the exclusive diffractive cross section - to the square of the gluon structure function.
Due to this fact the reaction $\gamma^* + N \to J/\Psi + N$ feels the variation of $\bar{x}G(\bar{x},\bar{q}^2)$ better than the inclusive J/Ψ cross section which depends on $\bar{x}G(\bar{x},\bar{q}^2)$ only linearly. Therefore this process is one of the best ways to measure the role of absorptive corrections (Pomeron cuts contributions) and to observe the saturation of gluon density predicted in the framework of perturbative QCD in ref.[10].
In Sect.2 we calculate the amplitude of diffractive J/Ψ-photoproduction. In Sect.3 we discuss the spin structure of this amplitude and correspondingly the distribution in azimuthal angle. In Sect.4 the numerical estimates of the single and double diffractive dissociation cross sections are given.

2 Amplitude of $\gamma^* + p \to J/\Psi + p$.

The Born amplitude of $\gamma^* + p \to J/\Psi + p$ reaction is described by the sum of the two diagrams in fig.1. As the binding energy of S-wave $c\bar{c}$ -quarks J/Ψ-system is small (much less than the charm quark mass $m_c \equiv m$) one can follow ref.[6] and use the nonrelativistic approximation writing the product of two propagators (k and k' in fig.1) and the J/Ψ-vertex (i.e. J/Ψ-wave function integrated over the relative momenta of $c\bar{c}$-quarks $\mathbf{k} = \mathbf{k}'$ in J/Ψ-rest frame system) in the form $g \cdot (\hat{k} + m)\gamma_\mu$. The constant g may be expressed in terms of the electronic width Γ^J_{ee} of $J/\Psi \to e^+e^-$ decay.

$$g^2 = \frac{3\Gamma^J_{ee} m_J}{64\pi\alpha^2_{e.m.}} \tag{1}$$

[1] The model for elastic and diffractive J/Ψ-production based on vector-meson dominance and Pomeron exchange was considered recently in ref. [8].

where m_J is the mass of J/Ψ-meson and $\alpha_{e.m.} = 1/137$.
Let us consider the large energy limit $s = (q+p)^2 \gg m_J^2$. In this case only the longitudinal polarization of gluons l and $l+Q$ is survived, i.e.

$$g_{\rho\sigma} = g^{\perp}_{\rho\sigma} + \frac{p'_\rho q'_\sigma + q'_\rho p'_\sigma}{(p'q')} \simeq \frac{p'_\rho q'_\sigma}{(p'q')}$$

$$\left(q_\mu = q'_\mu + \frac{q^2}{s}p'_\mu;\ p_\mu = p'_\mu + \frac{m_N^2}{s}q'_\mu \approx p_\mu;\ s = 2(p'q');\ p'^2 = q'^2 = 0\right)$$

(m_N is the nucleon mass)
The matrix element given by the Feynman graph in fig.1a takes the form

$$ImM^D_a = \frac{8}{s}\alpha_s{}^2\frac{2}{3}\int \frac{Sp[\gamma_\nu(\hat{k}+m)\gamma_\mu\hat{p}'(\hat{k}'+\hat{l}+m)\hat{p}'(\hat{r}+m)]}{(r^2-m^2)l^2(l+Q)^2} g \cdot e_c \Phi(l,Q) d^2 l_t \tag{2}$$

where $e_c = \frac{2}{3}\sqrt{\frac{4\pi}{137}}$ is the electric charge of c-quark, α_s is the QCD coupling constant, $\frac{2}{3}$ is the colour coefficient and the function $\Phi(l,Q)$ describes the emission of the gluon pair l and $l+Q$ by a proton.

As for diffractive processes a small momentum transfer $t = Q^2 \le 1/R_N^2$ (R_N is the nucleon radius) is essential we will below put $Q=0$. Then the function $\Phi(l,0) \equiv \Phi(l) \simeq 3$ for large $l \gg 1/R_N^2$ (each valence quark emits its own pair of gluons independently from the others quark-spectators), while at small $l \ll 1/R_N$ this quantity is $\Phi \propto l^2 R_N^2$.

Due to the even signature of the vacuum singularity (Pomeron) the real part of the amplitude M^D is small; it is equal to zero in the Born approximation. So the easiest way to get eq.(2) is to calculate the discontinuity shown by the crosses in fig.1. Note that in our nonrelativistic picture $k \simeq q^J/2$ and each c-quark (k and $-k'$) carries out $z = 1/2$ part of the photon momentum q; $k^2 = (k'+l)^2 = m^2$.
Adding the analogous formula for fig.1b graph and taking into account that there are two types of such diagrams (gluon l can interact with c-quark k or $\bar{c}$-antiquark k') one gets finally [2]

$$ImM^D(Q=0) = \frac{32\pi\alpha_s{}^2}{3} \cdot ge_c ms \int \Phi(l)\frac{dl^2}{l^4} \cdot$$
$$\cdot \left\{ \left[\frac{1}{(q-k)^2 - m^2} - \frac{1}{(q-k+l)^2 - m^2} \right] [sg_{\mu\nu} - 2q_\mu p_\nu - 2p_\mu q^J_\nu] - \frac{8}{s}\frac{p_\mu p_\nu (ql)}{[(q-k+l)^2 - m^2]} \right\} \tag{3}$$

Here: $(q-k)^2 - m^2 = r^2 - m^2 = -|q^2|/2 - 2m^2 - |q^J_t|^2/2$ and $(q-k+l)^2 - m^2 = r'^2 - m^2 = -(|q^2| + 4m^2)/2 - |\mathbf{q}^J - 2\mathbf{l}|_t^2/2$. (The longitudinal component of momentum l is negligible for large $s \gg 4m^2$ and $l^2 \simeq l_t^2$.)
After averaging over the azimuthal angles the expression in the curley brackets of eq.(3) is proportional to l^2 at small $l_t^2 \ll \bar{q}^2 = |q^2|/4 + m_t^2$, $(m_t^2 = m^2 + |k_t^2| = (m_J^2 + |q^J_t|^2)/4)$. Besides this at very small $l \ll 1/R_N$ the amplitude of two gluon emission $\Phi(l) \propto l^2$. So the main logarithmically large contribution to the integral over dl^2 comes from the region $\frac{1}{R_N^2} \ll |l_t^2| \ll \bar{q}^2$. This $log\bar{q}^2$represents the first loop logarithmic integration of the gluon

[2] Terms proportional to q_ν (or q^J_μ) that multiplied by the photon (or J/Ψ) polarization vector gives zero result are omitted here.

structure function and in the general case [3] in the LLA QCD one obtains

$$ImM^D = 16\pi^2\alpha_s \cdot ge_c m F_N^{2G}(t)\bar{x}G(\bar{x},\bar{q}^2)\frac{2\bar{q}^2 - |q_t^J|^2}{(2\bar{q}^2)^3}\{sg_{\mu\nu} - 2q_\mu p_\nu - 2p_\mu q_\nu^J + 4\frac{|2\bar{q}^2|}{s}p_\mu p_\nu\} \tag{4}$$

where: $\bar{q}^2 = (|q^2| + m_{Jt}^2)/4$ and $\bar{x} = 4\bar{q}^2/s$.
The new two gluon form factor $F_N^{2G}(t)$ that is introduced in eq.(4) in comparison with eq.(3) takes into account the t-dependence of the amplitude and characterizes the correlations between two gluons inside a proton, $(F_N^{2G}(t=0) = 1)$. For a rough estimate one can put $F_N^{2G} \approx F_N^{e.m.}(t)$- the electromagnetic proton form factor, but strictly speaking this is a new function that should be measured at the experiment.

3 Spin structure of the amplitude.

Let us discuss in more detail the structure of the last curley bracket in eq.(4). At large energies in the proton rest frame the polarization vectors $e_\nu^{(\lambda)} = (e_0, e_x, e_z, e_y)$ of photons with definite helicities-λ take the form: $e_\nu^{(\pm1)} = \frac{1}{\sqrt{2}}(0,0,1,\pm i)$ and $e_\nu^{(0)} = \frac{1}{\sqrt{|q^2|}}(E + \frac{|q^2|}{2E}, E, 0, 0)$. The same vectors for J/Ψ-meson are: $e_\mu^{(\pm1)} \simeq \frac{1}{\sqrt{2}}(0,0,1,\pm i)$ and $e^{(0)} \simeq \frac{1}{m_J}(E - \frac{m_J^2}{2E}, E, 0, 0)$, where E is the photon energy. For transverse polarizations $(i, j = \pm1)$ one gets $A_{i,j} = \delta_{i,j}$. It is easy to see from eq.(4) that a longitudinal zero helicity heavy photon only produces a longitudinal J/Ψ-meson. In particular

$$A_{0,0} = A_{1,1}\frac{\sqrt{|q^2|}}{m_J} \tag{5}$$

but at $t \neq 0$ zero helicity J/Ψ-meson can be produced by transverse photon also $A_{\pm1,0} = A_{1,1}\frac{|q_t^J|}{\sqrt{2m_J^2}}$. However the last amplitude is small in comparison to $A_{1,1}$, as for typical $q_t^J \sim 0.3$GeV the ratio $q_t^J/\sqrt{2m_J^2} \leq 0.07$. So we conclude that the amplitude in eq.(4) approximately conserves the helicity of the incoming photon and for large $|q^2| \gg m_J^2$ the longitidinal cross section $\sigma^L \gg \sigma^T$ (see eq.(5)). Within this accuracy there are no azimuthal correlations between initial leptons and the momentum of the J/Ψ-meson. The distribution of $\mu^+\mu^-$ coming from J/Ψ-decay is also flat for zero helicity J/Ψ-meson and photon (i.e. for longitudinal part of deep inelastic cross section σ^L). Only for transverse photons one gets nontrivial azimuthal distribution

$$1 - \sin^2\theta \cdot \cos^2\varphi$$

in the J/Ψ-rest frame, or

$$1 - \frac{|Q_t^\mu|}{m}\cos^2\varphi$$

in the proton rest frame for large $s_{\gamma^* p} \gg m_J^2$;
where Q_t^μ is the transverse momentum of muon from J/Ψ-decay and φ is the azimuthal angle between the vector $\mathbf{Q}_t^\mu$ and the plane of initial leptons $(e \to e')$ scattering.

[3]Summing up all the leading log contributions of the diagrams with the emission of an arbitrary number of additional partons (shown by short-dashed line on fig.1).

4 Concluding remarks.

1. The differential cross section of a transverse photon $\to J/\Psi$ diffractive dissociation given by the amplitude eq.(4)

$$\frac{d\sigma^T(\gamma p \to J/\Psi + p)}{dt} = \frac{|M|^2}{16\pi s^2} = [F_N^{2G}(t)]^2\frac{\alpha_s{}^2\Gamma_{ee}^J m_J^3}{3\alpha_{e.m.}}\pi^3\left[\bar{x}G(\bar{x},\bar{q}^2)\frac{2\bar{q}^2 - |q_t^J|^2}{(2\bar{q}^2)^3}\right]^2 \tag{6}$$

is of the same order (both from the numerical and parametrical points of views) as the inelastic J/Ψ production [6] cross section. Putting $\alpha_s = 1/4$ and $\bar{x}G(\bar{x},\bar{q}^2) = 2.5$ one gets for real photon dissociation

$$\frac{d\sigma^T(t=0)}{dt} = 58\frac{nb}{GeV^2}$$

in agreement with the experimental value

$$\frac{d\sigma(\gamma N \to J/\Psi N)}{dt}|_{t=0} = 52 \pm 5 \pm 10\frac{nb}{GeV^2}$$

extracted from 280GeV/c muon - iron coherent interactions [11]. The width of the essential t-region is about 0.2 to 0.25 GeV^2. So integrated over t the cross section will be $\sigma(\gamma N \to J/\Psi N) \simeq 13$nb, in agreement with the data [12].

2. In the framework of the vector dominance model[13] (see fig.2) one can use the value of $\frac{d\sigma(\gamma N \to J/\Psi N)}{dt}|_{t=0}$ to estimate the total cross section of J/Ψ-nucleon interaction. The result [11] $\sigma_{\Psi N}^{VDM} = 1.26 \pm 0.31$mb is too small in comparison to the data coming from the measurement of the A dependence of J/Ψ photoproduction ($\sigma_{\Psi N} = 3.5 \pm 0.8$mb) [14] which is more close to perturbative QCD calculations.
In the Born approximation the two gluon exchange model [9] of the Pomeron gives $\sigma_{\Psi N} = 4 - 7$ mb[15] for a reasonable value of J/Ψ radius R_Ψ.
The reason that the σ^{VDM} is so small is the following. In perturbative QCD the cross section is determined by the meson radius $\sigma_{\Psi N} \propto \langle R_\Psi^2 \rangle$. However in our reaction J/Ψ-meson is produced by a point like photon mostly through the small R_Ψ component of J/Ψ-wave function[4]. Therefore just after the photoproduction the radius $R_\Psi < \sqrt{\langle R_\Psi^2 \rangle}$ and the cross section $\sigma(\Psi N) \propto R_\Psi^2$ is smaller than its normal value $\sigma(\Psi N) \propto \langle R_\Psi^2 \rangle$.

3. The cross section of diffractive J/Ψ-electroproduction is also not extreamely small. At the HERA energies even for large q^2 one gets integrating over the region of $q^2 > 20GeV^2$ the value of $\sigma^T(ep \to e'J/\Psi + p) \simeq 20$pb and $\sigma^T(ep \to e'J/\Psi + p) \simeq 55$pb for $q^2 > 10GeV^2$, that I hope can be measured at HERA or LHC experiments. To estimate this cross sections the values of the gluon structure functions $\bar{x}G(\bar{x},\bar{q}^2)$=16 (at $\bar{x} = 7 \cdot 10^{-4}$, $q^2 > 20GeV^2$) and $\bar{x}G(\bar{x},\bar{q}^2)$=13 (at $\bar{x} = 4.5 \cdot 10^{-4}$, $q^2 > 10GeV^2$) from ref. [16] were used .
Note that due to the $\alpha_s ln\frac{1}{x}$ contributions the gluon density increases as $\bar{x}G(\bar{x},\bar{q}^2) \propto x^{\omega_0} \sim$

[4]Indeed, only due to the propagators $\frac{1}{r^2 - m'^2}$ or $\frac{1}{r'^2 - m'^2}$ the quarks get some possibilities to fly out from the initial point a (see fig.1) in the impact parameter plane. Without them (or in the case when the mass m' in these propagators would be very large) only the point-like components of the J/Ψ wave function contributes to the amplitudes fig.1 and one gets a zero result $d\sigma^T(\gamma \to J/\Psi)/dt \propto 1/m'^2 \to 0$. This is a reflection of the fact that the point-like colourless $c\bar{c}$-system is totally neutral. It does not emit the gluons $(l, l + Q)$ and thus does not interact at all.

$\frac{1}{\sqrt{x}}$[17, 18] up to the saturation boundary where the absorptive corrections stop the further growth of $\bar{x}G(\bar{x},\bar{q}^2)$. If for $\bar{q}^2 \simeq 10GeV^2$ there is no saturation at $x \geq 10^{-3}$ then from $\bar{x}G(\bar{x},\bar{q}^2)$=2.5 for $x \simeq$0.04 one gets $\bar{x}G(\bar{x},\bar{q}^2)\simeq$16 at $x = 10^{-3}$.
Based on the semihard phenomenology [19] one can expect to discover noticeable absorptive effects (the begining of gluon density saturation) at $x \leq 2\cdot 10^{-3}$ for $\bar{q}^2 = 10GeV^2$ (or even at $x \leq 10^{-2}$ for $\bar{q}^2 = 4GeV^2$). On the other hand if one prefers to choose the "week screening" parametrization of ref.[18] with $R = 5GeV^{-1}$ then the absorptive corrections will be almost negligible at the whole HERA kinematical range.
Thus the diffractive J/Ψ-production provides us with a good chance to find the real position of the saturation boundary and hence to answer this question.

4. Finally if the recoil proton will not be detected one can study the double diffractive process $\gamma^* + p \to J/\Psi + M^2$ (see fig.3). Unfortunately the triple-Reggeon vertex G_{3R} lies in the range of rather small virtualities $q'^2 \leq 1GeV^2$ and one can't use perturbative QCD to calculate it. As is known from the old triple-Reggeon analysis[20] the effective triple-Pomeron vertex G_{3P} is (20 - 40) times smaller than the elastic proton -Pomeron vertex g_{el} . However this smallness is compensated by the large phase volume of integration in longitudinal ($\int dz/(1-z)$) and transverse momentum spaces. Due to the small radius of the G_{3P}-vertex the integral over t in a double diffractive dissociation reaction is 4 - 5 times larger ($\int dt \sim 1GeV^2$) than in single one. Adding the excitation of nucleon resonances one gets the cross section of double dissociation to be of the same order as for the single dissociation. So the total cross section of diffractive J/Ψ-electroproduction may be 2 times larger than our previous estimate. All this cross section will be concentrated in a narrow peak at $z \to 1$ i.e. in a very small interval of $z > 0.99$.

5 Acknowledgements.

I am very grateful to G.A.Schuler for initial discussion. Mainly due to his initiative this work was done.I would like to acknowledge the hospitality of Department of Theoretical Physics University of LUND where this work was done and especially of Bo Andersson and Yuri Dokshitzer. I thank them for discussions, support and interest in this work.

References

[1] T. Weiler: Phys. Rev. Lett. 44 (1980) 304.

[2] D. Allasia *et al.*: Phys. Lett. B258 (1991) 493.

[3] H. Fritzsch and K. H. Streng: Phys. Lett. B72 (1978) 385.

[4] L. Jones and W. Wyld: Phys. Rev. D17 (1978) 759; M. Glück and E. Reya: Phys. Lett. B79 (1978) 453, B83 (1979) 98; J. P. Leveille and T. Weiler: Nucl. Phys. B147 (1979) 147; V. Barger, W. Y. Keung and R. J. N. Phillips: Phys. Rev. D20 (1979) 630, Phys. Lett. B91 (1980) 253, B92 (1980) 179.

[5] J. P. Leveille and T. Weiler: Phys. Lett. B86 (1979) 377, Phys. Rev. D24 (1981) 1789.

[6] E. L. Berger and D. Jones: Phys. Rev. D23 (1981) 1521.

[7] R. Baier and R. Rücl: Nucl. Phys. B201 (1982) 1.

[8] G. A. Schuler and J. Terron: preprint CERN-TH. 6403/92; DESY-92-17.

[9] F. E. Low: Phys. Rev. D12 (1975) 163; S. Nussinov: Phys. Rev. Lett. 34 (1975) 1268.

[10] L. V. Gribov, E. M. Levin and M. G. Ryskin: Phys. Rep. 100 (1983) 1.

[11] J. J. Aubert *et al.*: Phys. Lett. B89 (1980) 267.

[12] B. H. Denby *et al.*: Phys. Rev. Lett. 52 (1984) 795; M. Binkley *et al.*: Phys. Rev. Lett. 48 (1982) 73.

[13] J. Pumplin and W. Repko: Phys. Rev. D12 (1975) 1376; T. H. Bauer, R. D. Spital, D. R. Yennie and F. M. Pipkin: Rev. Mod. Phys. 50 (1978) 261.

[14] R. L. Anderson *et al.*: Phys. Rev. Lett. 38 (1977) 263.

[15] J. F. Gunion and H. Soper: Phys. Rev. D15 (1977) 2617; E. M. Levin and M. G. Ryskin: Sov. J. Nucl. Phys. 34 (1981) 619; P. W. Ensing and J. Rand: Phys. Rev. D28 (1983) 1796.

[16] V. T. Kim and M. G. Ryskin: preprint DESY 91-064 and Sov. J. Nucl. Phys. (in print).

[17] E. A. Kuraev, L. N. Lipatov and V. S. Fadin: Sov. Phys. JETP 45 (1977) 199.

[18] J. Kwiecinski; Proc. of DESY Topical Meeting on the small-X. Nucl. Phys. B. (Proc. Suppl.) 18C (1990) 133;
J. Kwiecinski, A. D. Martin, W. J. Stirling and R. G. Roberts: Phys. Rev. D42 (1990) 3645;
J. Kwiecinski, A. D. Martin and P. J. Sutton: Phys. Rev.D44 (1991) 2640.

[19] L. V. Gribov, E. M. Levin and M. G. Ryskin: Phys. Rep. 100 (1983) 1; E. M. Levin and M. G. Ryskin: Phys. Rep. 189 (1990) 267.

[20] A. B. Kaidalov, Yu. F. Pirogov, N. L. Ter-Isaakyan and V. A. Khoze: Sov. Phys. JETP Lett. 17 (1973) 440; A. B.Kaidalov and K. A. Ter-Martirosyan: Nucl. Phys. B75 (1974) 471; R. D. Field and G. C. Fox: Nucl.Phys. B80 (1974) 367.

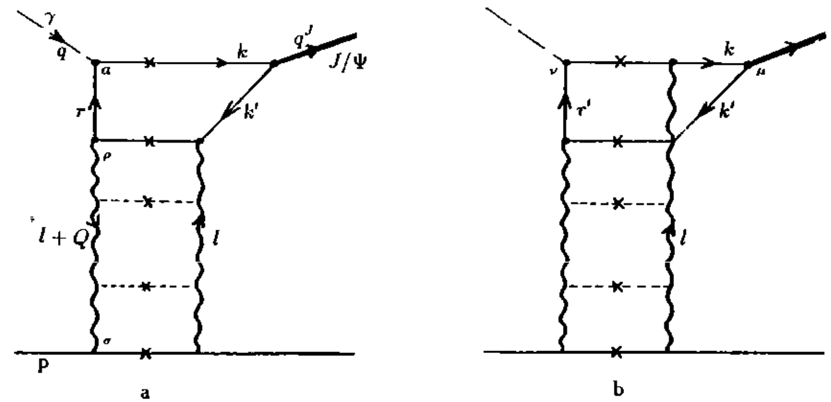

Fig.1. *Feynman diagrams for diffractive J/Ψ-production.*

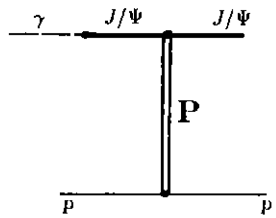

Fig.2. *Vector dominance model for J/Ψ-photoproduction.*

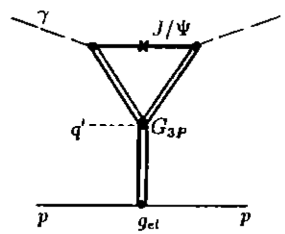

Fig.3. *Triple-Pomeron double diffractive dissociation* $\gamma^* + p \to J/\Psi + M^2$.