AZIMUTHAL DEPENDENCE OF DEEP INELASTIC HEAVY RESONANCE PRODUCTION

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The dependence on the azimuthal angle between the lepton-scattering plane and the plane defined by the virtual photon and the J/ψ is calculated in a photon–gluon fusion model in which the J/ψ is represented by the appropriate quark wave function. Our results differ substantially from models based on $c\bar{c}$ production below charmed particle threshold thus providing a way to distinguish phenomenologically between these two models.

The azimuthal dependence of the production of J/ψ resonances by muon beams has been the object of recent theoretical and experimental investigations [1–3]. It has been emphasized [1,2] that this dependence provides a sensitive test of the photon–gluon fusion model [4,5] and, beyond this, of the spin–parity assignment of the gluon. Because the photon–gluon fusion model has been criticized recently [6] (the $c\bar{c}$ is not in a color-singlet state and does not have $J^P = 1^-$) we have reanalysed the azimuthal dependence in a model, first proposed by Berger and Jones [6] for the photoproduction case, in which the J/ψ is represented by the appropriate quark wave function. Our results for inelastic J/ψ production by muon beams differ substantially from those of ref. [1], thus providing a way of distinguishing phenomenologically between the two models.

The general form of the cross section we are interested in is (see fig. 1):

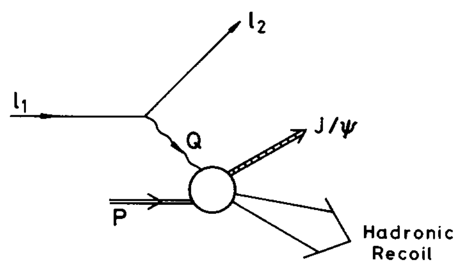

Fig. 1. Deep inelastic leptoproduction of J/ψ resonance.

$$\frac{d\sigma(\mu^- + p \to \mu^- + J/\psi + ...)}{d\nu\, dQ^2\, dp_\perp^2\, d\Phi\, dz} = A(1 + B\cos\Phi + C\cos 2\Phi)\,, \quad (1)$$

where Φ is the angle between the lepton-scattering plane and the plane defined by the virtual photon and the J/ψ (see eq. (10) below), Z is, in the laboratory frame, the energy of the J/ψ divided by the energy of the exchanged virtual photon and $p_\perp^2$ is the transverse momentum squared of the heavy resonance.

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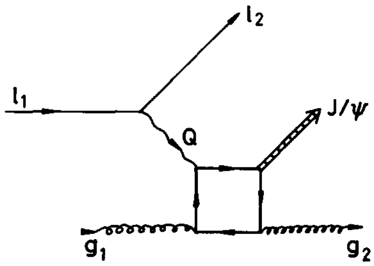

Fig. 2. One of the diagrams contributing to deep inelastic leptoproduction of the J/ψ resonance. There are five more diagrams corresponding to all possible ways of attaching the gluon lines to the quark loop.

The model we are considering is described by the diagram of fig. 2. It has been shown [7,8] that this model provides a good description of the t-dependence and of the Q^2-dependence (see ref. [3]). In the model the cross-section (1) is given explicitly by:

$$\frac{d\sigma(\mu^- + p \to \mu^- + J/\psi + ...)}{d\nu\, dQ^2\, dp_\perp^2\, d\Phi\, dz} = \frac{\pi G(\xi)\Sigma|\mathcal{M}(\mu^- + g_1 \to \mu^- + g_2 + J/\psi)|^2}{64(2\pi)^5 s_T^2 z(-t)}, \quad (2)$$

where s_T is the total incoming energy $s_T = (l_1 + p)^2$ and t is given by:

$$-t = p_\perp^2/z - Q^2(1-z) + M_{J/\psi}^2(1-z)/z\,. \quad (3)$$

$G(\xi)$ denotes the gluon distribution function inside the target particle, with

$$\xi = -t/2\nu(1-z)\,. \quad (4)$$

The summation in (2) refers to spins and colours. The matrix element in (2) has the following general form (see fig. 2) [‡1]:

$$\mathcal{M}(\mu^- + g_1 \to \mu^- + g_2 + J/\psi) = (e/Q^2)\bar{u}(l_2)\gamma_\mu u(l_1) T_{\mu\nu}\epsilon_\nu(J/\psi)\,. \quad (5)$$

When calculating the modulus squared of (4) we first take out constant factors corresponding to coupling constants and wave function at the origin:

[‡1] Unless stated otherwise we use the notations of ref. [9].

$$\sum_{\substack{\text{spins}\\ \text{colour}}} |\mathcal{M}(\mu^- + g_1 \to \mu^- + g_2 + J/\psi)|^2 = (e^2/Q^4) C L_{\mu\nu} H_{\mu\nu\mu'\nu'}(-g_{\mu'\nu'} + p_{J\mu'}p_{J\nu'}/M_J^2)\,, \quad (6)$$

where

$$C = (2\alpha_s^2/3\alpha)(64\pi)^2\Gamma(J \to \mu^+\mu^-)M_J^3\,,$$

and s, t and u correspond to the Mandelstam variables of the subprocess in fig. 2 and

$$H_{\mu\nu\mu'\nu'} = (64 m_c)^{-2} C_{\mu\nu\alpha\beta} C_{\mu'\nu'\alpha\beta}$$

in the notation of ref. [9]. $L_{\mu\nu}$ is the tensor arising from the leptonic part of the amplitude:

$$L_{\mu\nu} = l_{1\mu}l_{2\nu} + l_{1\mu}l_{2\nu} + \tfrac{1}{2}Q^2 g_{\mu\nu} = \tfrac{1}{2}(L_\mu L_\nu - Q_\mu Q_\nu + Q^2 g_{\mu\nu})\,, \quad (7)$$

with $L = l_1 + l_2$. To calculate the specifically hadronic part in (5) we introduce the following gauge invariant decomposition,

$$H_{\mu\nu\mu'\nu'}(-g_{\mu'\nu'} + P_{J\mu'}P_{J\nu'}/M_J^2) \quad (8)$$

$$= \mathcal{H}_1\hat{g}_{\mu\nu} + \mathcal{H}_2\hat{g}_{1\mu}\hat{g}_{1\nu} + \mathcal{H}_3\hat{g}_{2\mu}\hat{g}_{2\nu} + \mathcal{H}_4(\hat{g}_{1\mu}\hat{g}_{2\nu} + \hat{g}_{1\nu}\hat{g}_{2\mu})\,, \quad (8)$$

where the hatted quantities denote the following gauge invariant combinations:

$$\hat{g}_{\mu\nu} = g_{\mu\nu} - Q_\mu Q_\nu/Q^2, \quad \hat{g}_{i\mu} = g_{i\mu} - [(Qg_i)/Q^2]Q_\mu\,. \quad (9a,b)$$

The explicit expression for $\mathcal{H}_1$, $\mathcal{H}_2$, $\mathcal{H}_3$ and $\mathcal{H}_4$ are determined from the diagrams in fig. 2 as being

$$\mathcal{H}_1 = -(1/8D)\{[t^2 + (s - M_J^2)(M_J^2 - u)]^2 + Q^2 t[2M_J^2 t + (3M_J^2 - Q^2)(s - M_J^2)(M_J^2 - u)/M_J^2]\}\,, \quad (10a)$$

$$\mathcal{H}_2 = -(1/4D)[(u - Q^2)^2(M_J^2 - Q^2) + 2Q^2t^2]\,, \quad (10b)$$

$$\mathcal{H}_3 = -(1/4D)[(s - Q^2)^2(M_J^2 - Q^2) + 2Q^2t^2]\,, \quad (10c)$$

$$\mathcal{H}_4 = -(1/4D)(M_J^2 - Q^2) \times [(s - M_J^2)(M_J^2 - u)(M_J^2 - Q^2)/M_J^2 + M_J^2 t]\,, \quad (10d)$$

with

$$D \equiv (s - M_J^2)^2 (M_J^2 - u)^2 (t - M_J^2 + Q^2)^2 .$$

In order to exhibit the azimuthal dependence, we choose, in the virtual photon–gluon center of mass system, the momenta of the particles involved as follows:

$$g_1 = [(s - Q^2)/2\sqrt{s}](1, 0, 0, -1) , \quad (11a)$$

$$g_2 = [(s - M_J^2)/2\sqrt{s}](1, -\sin\theta_J, 0, -\cos\theta_J) , \quad (11b)$$

$$Q = \left(\frac{s + Q^2}{2\sqrt{s}}, 0, 0, \frac{s - Q^2}{2\sqrt{s}}\right), \quad (11c)$$

$$P_J = \left(\frac{s + M_J^2}{2\sqrt{s}}, \frac{s - M_J^2}{2\sqrt{s}} \sin\theta_J, 0, \frac{s - M_J^2}{2\sqrt{s}} \cos\theta_J\right), \quad (11d)$$

$$L = \left(\frac{s - Q^2}{2\sqrt{s}} \left(\frac{1+\epsilon}{1-\epsilon}\right)^{1/2}, \left(-2Q^2 \frac{\epsilon}{1-\epsilon}\right)^{1/2} \cos\Phi , - \left(-2Q^2 \frac{\epsilon}{1-\epsilon}\right)^{1/2} \sin\Phi, \frac{s + Q^2}{2\sqrt{s}} \left(\frac{1+\epsilon}{1-\epsilon}\right)^{1/2}\right) , \quad (11e)$$

where ϵ is the usual depolarization factor.

Combining eqs. (6)–(9) and (11) we obtain for the dependence of the modulus squared of the matrix element:

$$\sum_{\substack{\text{spins}\\ \text{colour}}} |\mathcal{M}(\mu^- + g_1 \to \mu^- + g_2 + J/\psi)|^2 = \frac{e^2 C(-Q^2)}{Q^4(1-\epsilon)}$$

$$\times \{\Sigma_u + \epsilon\Sigma_L + [2\epsilon(1+\epsilon)]^{1/2} \Sigma_I \cos\Phi + \epsilon\Sigma_T \cos 2\Phi\}, \quad (12)$$

where $\Sigma_u, ..., \Sigma_T$ are expressed in terms of the quantities $\mathcal{H}_1, ..., \mathcal{H}_4$ introduced in eq. (8) in the following way:

$$\Sigma_u = -\mathcal{H}_1 + [(s - M_J^2)^2/8s] \sin^2\theta_J \mathcal{H}_3 , \quad (13a)$$

$$\Sigma_L = \mathcal{H}_1 - [(s - Q^2)^2/4Q^2] \mathcal{H}_2 - \tfrac{1}{4} Q^{-2} [Q^2 - u + 2Q^2 t/(s - Q^2)]^2 \mathcal{H}_3 , \quad (13b)$$

$$\Sigma_I = -[(s - M_J^2)/4(-Q^2 s)^{1/2}] \sin\theta_J \times \{[Q^2 - u + 2Q^2 t/(s - Q^2)] \mathcal{H}_3 + (s - Q^2) \mathcal{H}_4\} , \quad (13c)$$

$$\Sigma_T = \tfrac{1}{8} \sin^2\theta_J [(s - M_J^2)^2/s] \mathcal{H}_3 . \quad (13d)$$

In eqs. (13b) and (13c) the singularities $1/Q^2$ and $1/\sqrt{-Q^2}$ are actually artifacts of the gauge invariant decomposition of eq. (8) and after eqs. (10) are inserted into eqs. (13b) and (13c), one can see that the artificial singularities cancel so that finally one obtains

$$\Sigma_L \sim Q^2 , \quad \Sigma_I \sim (-Q^2)^{1/2} . \quad (14)$$

These factors are of purely kinematic origin. This point is more easily seen in the helicity approach, which is described in detail in ref. [9]. Integrating eq. (16) of ref. [9] over the variables corresponding to the decay products of the heavy resonance one finds:

$$\Sigma_u = \frac{1}{8^4 M_J^2} \sum_{\lambda_2, J, \lambda_1} |H_{\lambda_2 J; \lambda_1 +}|^2 , \quad (15a)$$

$$\Sigma_L = \frac{1}{8^4 M_J^2} \sum_{\lambda_2, J, \lambda_1} |H_{\lambda_2 J; \lambda_1 0}|^2 , \quad (15b)$$

$$\Sigma_T = \frac{1}{8^4 M_J^2} \sum_{\lambda_2, J, \lambda_1} H_{\lambda_2 J; \lambda_1 +} H^*_{\lambda_2 J; \lambda_1 -} , \quad (15c)$$

$$\Sigma_I = \frac{1}{8^4 M_J^2} \sum_{\lambda_2, J, \lambda_1} \mathrm{Re}(H_{\lambda_2 J; \lambda_1 +} H^*_{\lambda_2 J; \lambda_1 0}) , \quad (15d)$$

where the s-channel helicity amplitudes $H_{\lambda_2 J; \lambda_1 l}$ are given in eq. (29) of ref. [9] and their expressions are given in table 1. λ_1 and λ_2 refer to the helicities of gluons g_1 and g_2, respectively, while J and l refer to the helicities of the J/ψ and the virtual photon, respectively. The final results for the azimuthal dependence are identical to those of the previous approach. The advantage of the helicity approach is that the dependence of certain kinematical factors is now more apparent: by inspection of table 1 one immediately sees that

$$\Sigma_L \sim Q^2 ,$$
$$\Sigma_T \sim \sin^2\theta_J \sim p_\perp^2 ,$$
$$\Sigma_I \sim (-Q^2)^{1/2} \sin\theta_J \sim (-Q^2)^{1/2} p_\perp . \quad (16)$$

Thus the coefficient of cos 2Φ goes to zero as the transverse momentum of the J/ψ goes to zero, while the coefficient of cos Φ goes to zero both when Q^2 or $p_\perp^2$ goes to zero.

The numerical results on B and C are given in fig. 3 (real lines) where the integrals over ν, $p_\perp^2$ and z are done under the kinematical constraints $p_\perp^2 \geqslant 0.5$ $(\mathrm{GeV}/c)^2$ and $z \leqslant 0.95$. From fig. 3 one sees that the coefficient of the cos Φ term is negative. For $Q^2 = 0$

Table 1
s-channel helicity amplitudes $H_{\lambda_2 J;\lambda_1 l}$ for the constituent scattering process, $\gamma^* g_1 \to J g_2$, J being a heavy reconance (J = J/ψ, ϒ, ---). Entries are to be multiplied by the common denominator $[(s - M_J^2)(M_J^2 - u)(t - M_J^2 + Q^2)]^{-1}$.

$$
\begin{aligned}
H_{++;++} &= H_{--;--} = 8M_J[-(s - Q^2 - M_J^2)\, t(1 + \cos\theta_J) \\
&\qquad + 2(s - M_J^2)(s - Q^2 + t) + 2t(Q^2 + t)] \\
H_{++;+-} &= H_{--;-+} = 4M_J^3(s - M_J^2)(s - Q^2)\sin^2\theta_J/s \\
H_{++;-+} &= H_{--;+-} = 4M_J^3(s - M_J^2)(s - Q^2)\sin^2\theta_J/s \\
H_{++;--} &= H_{--;++} = -8M_J^3 t Q^2(1 - \cos\theta_J)/s \\
H_{+-;++} &= H_{-+;--} = 4M_J Q^2(s - Q^2)(s - M_J^2)\sin^2\theta_J/s \\
H_{+-;+-} &= H_{-+;-+} = 4M_J(s - Q^2)(s - M_J^2)(1 + \cos\theta_J)^2 \\
H_{+-;-+} &= H_{-+;+-} = 4M_J(s - Q^2)(s - M_J^2)(1 - \cos\theta_J)^2 \\
H_{+-;--} &= H_{-+;++} = 4M_J Q^2(s - Q^2)(s - M_J^2)\sin^2\theta/s \\
H_{+0;++} &= -H_{-0;--} = -8\sqrt{2}\, Q^2[(s - Q^2)(s - M_J^2) + M_J^2 t]\sin\theta_J/\sqrt{s} \\
H_{+0;+-} &= -H_{-0;-+} = 4\sqrt{2}\, M_J^2(s - Q^2)(s - M_J^2)\sin\theta_J(1 + \cos\theta_J)/\sqrt{s} \\
H_{+0;-+} &= -H_{-0;+-} = -4\sqrt{2}\, M_J^2(s - Q^2)(s - M_J^2)\sin\theta_J(1 - \cos\theta_J)/\sqrt{s} \\
H_{+0;--} &= -H_{-0;++} = -8\sqrt{2}\, Q^2 M_J^2 t\sin\theta_J/\sqrt{s} \\
H_{++;+0} &= -H_{--;-0} = -8\sqrt{2}\, M_J(-Q^2/s)^{1/2}[(s - Q^2)(s - M_J^2) + M_J^2 t]\sin\theta_J \\
H_{+-;+0} &= -H_{-+;-0} = 4\sqrt{2}\, M_J(-Q^2/s)^{1/2}(s - Q^2)(s - M_J^2)\sin\theta_J(1 + \cos\theta_J) \\
H_{-+;+0} &= -H_{+-;-0} = -4\sqrt{2}\, M_J(-Q^2/s)^{1/2}(s - Q^2)(s - M_J^2)\sin\theta_J(1 - \cos\theta_J) \\
H_{--;+0} &= -H_{++;-0} = -8\sqrt{2}\, M_J(-Q^2/s)^{1/2} M_J^2 t\sin\theta_J \\
H_{+0;+0} &= H_{-0;-0} = -8(-Q^2)^{1/2}[(s - Q^2)(s - M_J^2) + 2M_J^2 t](1 + \cos\theta_J) \\
H_{+0;-0} &= H_{-0;+0} = -8(-Q^2)^{1/2}[(s - Q^2)(s - M_J^2) + 2st](1 - \cos\theta_J)
\end{aligned}
$$

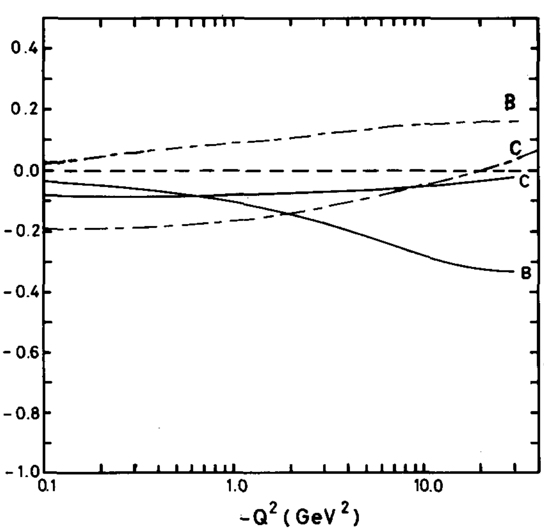

Fig. 3. The coefficient of cos Φ (B) and the coefficient of cos 2Φ (C) as a function of Q^2, all other variables are integrated over. The kinematical constraints $z \leqslant 0.95$ and $p_\perp^2 \geqslant 0.5$ (GeV/c)^2 were applied on the integration region. The solid lines are for the vector gluons (QCD) and the dashed–dotted lines are for the scalar gluons.

it must be zero as was discussed above, for increasing values of Q^2 it decreases continuously to reach −0.35 at $Q^2 = -30$ GeV^2. The coefficient of the cos 2Φ term does not have a kinematic zero for $Q^2 = 0$ but is always very small even for the largest values of Q^2 considered. On this point our results disagree with the photon–gluon fusion model considered in ref. [1] where a very large value for this term was found. Experimental absence of a cos 2Φ term would clearly favour the model considered here. Concerning the cos Φ term, we see from fig. 3 that it will dominate at large values of Q^2 over the cos 2Φ term. In this point our results are similar to those of Leveille and Weiler for the diffractive J/ψ production process. Those authors found a dominating cos Φ term for the elastic cross-section. We checked that our results are not sensitive to the precise value of the mass of the charmed quark. In our calculations we took $m_c \simeq M_J^2/2$, changing M_J from 3.09 to 2.7 GeV/c^2, thus effectively changing the charmed quark mass, modifies our results for B by about 5% while C becomes smaller by about 10%.

For comparison we have also calculated the azimuthal distribution in the scalar gluon model. In this case the corresponding $\mathcal{H}_i$ (i = 1–4) are given as follows:

$$\mathcal{H}_1 = -(1/16D)M_J^4 \times [(M_J^2 - Q^2)(M_J^2 + Q^2 - t) + (s - Q^2)(Q^2 - u)] , \quad (17a)$$

$$\mathcal{H}_2 = (1/16D)(M_J^2 - Q^2) \times [(u - Q^2)^2(3M_J^2 + Q^2)/M_J^2 + 4Q^2(M_J^2 - Q^2)] , \quad (17b)$$

$$\mathcal{H}_3 = (1/16D)(M_J^2 - Q^2) \times [(s - Q^2)^2(3M_J^2 + Q^2)/M_J^2 + 4Q^2(M_J^2 - Q^2)] , \quad (17c)$$

$$\mathcal{H}_4 = (1/16D)(M_J^2 - Q^2) \times [-(s - Q^2)(Q^2 - u)(3M_J^2 - Q^2)/M_J^2 + 4Q^2 t] , \quad (17d)$$

with

$$D \equiv (s - M_J^2)^2(M_J^2 - u)^2(t - M_J^2 + Q^2)^2 ,$$

as before. The numerical results for B and C are included in fig. 3 (dashed–dotted lines). B is positive while C is negative and rather small in the scalar gluon case.

The dependence on azimuthal angle would thus provide a phenomenological way of distinguishing between the two versions of the gluon-fusion model.

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