Charmed-particle production in neutrino scattering

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Inclusive production of charmed vector mesons has been studied in high-energy neutrino scattering. The cross sections calculated in the triple-Regge limit are of the order of 10^{-41} cm^2/GeV^3. The production mechanism in the central region is also discussed.

I. INTRODUCTION

Considerable attention has been focused on two extremely narrow resonances at 3.095 and 3.684 GeV[1,2] (hereafter referred to as ψ and ψ' respectively). These resonances have been interpreted as bound states of a charmed and an anticharmed quark.[3] It has been concluded from the narrow widths of ψ and ψ' that the mass[4] of the charmed quark should be greater than 1.850 GeV; otherwise ψ' will decay into charmed hadrons. A further experimental support for this scheme has come from the experimental observation of radiative decay of ψ' to a low-lying state of charmonium. The experimental results are in agreement with the qualitative predictions of such models. A natural consequence of the SU(4) scheme[3] for hadrons is to incorporate charmed particles as members of various multiplets of this group. Several mass formulas have been derived for charmed hadrons, and it is expected that their masses will be around 2 GeV[5,6]. It should be noted that both strong and electromagnetic interactions conserve the new quantum number charm. Therefore, charmed hadrons must be produced in pairs in these interactions. Consequently, the cross section for production of charmed hadrons in strong and electromagnetic interactions is lower than that for the production of ordinary hadrons. On the other hand, it is possible to produce a single charmed hadron through weak interactions. Therefore, high-energy neutrino scattering will provide opportunities to study production of charmed particles.

Our aim here is to investigate diffractive production of charmed vector mesons in high-energy neutrino scattering. To be specific, we consider the model of Glashow, Iliopoulos, and Maiani.[7] In this model the weak current is written as

$$J_\mu = \bar{\mathcal{P}}\gamma_\mu(1+\gamma_5)(\mathfrak{N}\cos\theta_C + \lambda\sin\theta_C) + \bar{\mathcal{P}}'\gamma_\mu(1+\gamma_5)(\lambda\cos\theta_C - \mathfrak{N}\sin\theta_C) + \bar{\nu}\gamma_\mu(1+\gamma_5)e + \bar{\nu}'\gamma_\mu(1+\gamma_5)\mu\,. \quad (1)$$

Here $\mathcal{P}$, $\mathfrak{N}$, and λ refer to the usual quark and $\mathcal{P}'$ is the charmed quark; θ_C is the Cabibbo angle. We are interested in the reaction $\nu + p \to \mu^- + F^{*+} + X$. The presence of the factor $\cos\theta_C$ with $\theta_C \simeq 15°$ in the amplitude will make the cross section for F^* production dominate over those for other vector mesons.

It is well known that the Regge-pole model provides useful guidance to the study of hadronic interactions at high energies. This model has also been applied to the study of the scattering of weak and electromagnetic currents off hadrons. The model has also been applied with considerable success to deep-inelastic electron-proton scattering.[8,9] It has been shown that the Mueller-Regge technique is applicable to diffractive dissociation in deep-inelastic electroproduction processes.[10] We take the viewpoint that the strong interaction governs the essential dynamics in the production of charmed particles through weak currents. Therefore, we are tempted to apply Mueller-Regge phenomenology to the reaction of our interest to obtain a qualitative description of the inclusive process.

The rest of the paper is arranged as follows: The kinematics is given in Sec. II, and the model is described in Sec. III. We make a few comments in Sec. IV.

II. KINEMATICS

We first define kinematical variables for the process $\nu + p \to \mu^- +$ anything as follows:

p (p') = initial (final) momentum of lepton,

P = initial momentum of hadron,

$q = p - p'$,

$\nu = q \cdot P$,

θ = scattering angle of lepton in laboratory,

$Q^2 = -q^2 = 4EE'\sin^2(\frac{1}{2}\theta)$,

scaling variable $\omega = \dfrac{2\nu}{Q^2}$.

The cross section for the above process in the zero-lepton-mass limit is given by

$$\frac{d\sigma}{d\Omega' dE'} = \frac{G^2 E'^2}{2\pi^2}\left[2W_1(q^2, \nu)\sin^2(\tfrac{1}{2}\theta) + W_2(q^2, \nu)\cos^2(\tfrac{1}{2}\theta) - \frac{E+E'}{M} W_3(q^2, \nu)\sin^2(\tfrac{1}{2}\theta)\right]. \quad (2)$$

The structure functions W_1, W_2, and W_3 are related to the total scattering cross section of the spin-averaged proton and the longitudinal (transverse) component of the current in the following manner:

$$\sigma_S(q^2, \nu) = \left(1 + \frac{\nu}{Q^2 M^2}\right) W_2 - W_1, \quad (3a)$$

$$\sigma_R(q^2, \nu) = W_1 + \frac{1}{2M}\left(\frac{\nu^2}{M^2} - q^2\right)^{1/2} W_3, \quad (3b)$$

$$\sigma_L(q^2, \nu) = W_1 - \frac{1}{2M}\left(\frac{\nu^2}{M^2} - q^2\right)^{1/2} W_3. \quad (3c)$$

Thus

$$W_2 = \frac{\sigma_S + \frac{1}{2}(\sigma_L + \sigma_R)}{1 + \nu/Q^2 M^2}. \quad (3d)$$

Now the inclusive cross section for $\nu + p \to \mu^- + F^{*+} + X$ is given by

$$E_F{}^* \frac{d\sigma}{d\Omega' dE' d^3k} = \frac{d\sigma}{d\Omega' dE'} \frac{E_F{}^*}{\sigma} \frac{d\sigma}{d^3k}. \quad (4)$$

Here $E_F{}^* d\sigma/d^3k$ is the inclusive cross section for $J^{+i} + p \to F^{*+} + X$, and σ is the total cross section for $J^{+i} + p$. Since $d\sigma/d\Omega' dE'$ is determined from deep-inelastic neutrino scattering, we have to concentrate on $E_F{}^* d\sigma/d^3k$. This cross section is governed purely by strong-interaction dynamics. The appropriate kinematical variables for this inclusive process are

$$s = (P+q)^2 = M^2 + q^2 + 2\nu, \quad (5a)$$

$$t = (q-k)^2, \quad (5b)$$

$$u = (P-k)^2, \quad (5c)$$

$$M_X{}^2 = (P+q-k)^2 = \text{square of the missing mass}, \quad (5d)$$

k = momentum of the charmed hadron F^{*+}.

In terms of the Feynman variable x we have

$$\frac{M_X{}^2}{s} = (1 - \bar{x}) + O(1/s), \quad (6a)$$

$$\bar{x} = \left(x^2 + \frac{k_\perp^2 + m_F{}^{*2}}{\frac{1}{4}s}\right)^{1/2}. \quad (6b)$$

The cross section $E_F{}^* d\sigma/d^3k$ is related to $d\sigma/dt\, dM_X{}^2$ as

$$E_F{}^* \frac{d\sigma}{d^3k} = \frac{s}{\pi} \frac{d\sigma}{dt\, dM_X{}^2}.$$

Thus, finally,

$$E_F{}^* \frac{d\sigma}{d\Omega' dE' d^3k} = \frac{G^2 E'^2}{2\pi^2} \frac{W_2}{\sigma} \frac{s}{\pi} \frac{d\sigma}{dt\, dM_X{}^2} \quad (7)$$

in the small-θ limit. Since we are interested only in the Pomeron exchange, $\sigma_L = \sigma_R$.

III. THE MODEL

In this section we consider the inelastic cross section in various kinematical regions. Since F^{*+} is produced diffractively, we assume that it couples to the Pomeron. Furthermore, the contribution to the inclusive cross section will be negligible from the small-u region (target fragmentation). We consider the kinematical region where $s \to$ large, $M_X{}^2 \to$ large, and $s/M_X{}^2 \to$ large. This corresponds to the triple-Regge region, and the cross section is

$$\frac{d\sigma^i}{dt\, dM_X{}^2} = \frac{1}{8\pi} |\beta_{F^*P}(t, Q^2)|^2 \left(\frac{sQ^2}{M_X{}^2 m_F{}^{*2}}\right)^{2\alpha_P(t)} \times g_P(t, t, 0)\beta_{pP}^{(0)} M_X{}^{2\alpha_P(0)} \cos^2\theta_C. \quad (8)$$

Equation (8) is obtained in the limit $Q^2/m_F{}^{*2} < 1$. Here $\beta_{F^*P}(t, Q^2)$ is the residue function for the coupling of the current, F^*, and the Pomeron. $g_P(t,t,0)$ is the triple-Pomeron coupling constant, and $\beta_{pP}(0)$ is the Pomeron-proton-proton residue at $t=0$. $g_P(t,t,0)$ and $\beta_{pP}(0)$ are known from purely hadronic interactions. To determine $\beta_{F^*P}(t, Q^2)$ we adopt the following procedure. One knows that, in the case of electroproduction, scaling of W_2 requires that[8,9]

$$\beta_{\gamma p}^i(\alpha_0, q^2) \to (1/q^2)^{\alpha_0} \tilde{\beta}^i(0),$$

with the normalization

$$s\sigma_t^{pp} = |\beta_{pP}^{(0)}|^2 s^{\alpha_P(0)},$$

$$s\sigma_{\gamma p}^i(s, Q^2) = \beta_{pP}^{(0)} \tilde{\beta}_{\gamma p}^I(0)\omega^{\alpha_P(0)}.$$

We introduce the ansatz

$$\beta_{F^*P}(Q^2 \alpha_P(t)) \sim \tilde{\beta}(t)(1/Q^2)^{\alpha_P(t)} \quad (9)$$

so that

$$s^2\left(\frac{d\sigma}{dt}\right)_{\nu p \to \nu p} = \frac{1}{8\pi} |\beta_{pP}(t)|^2 |\tilde{\beta}(t)|^2 s^{2\alpha_P(t)}. \quad (10)$$

Since the differential cross section for pp elastic scattering is given by

$$s^2\left(\frac{d\sigma}{dt}\right)_{pp \to pp} = \frac{1}{16\pi} |\beta_{pP}(t)|^4 s^{2\alpha_P(t)}, \quad (11)$$

by straightforward substitution of Eqs. (10) and

(11) in (8) we get

$$s^2\frac{d\sigma}{dt\,dM_X{}^2}=\frac{1}{16\pi}\left(\frac{s}{m_{F^*}{}^2M_X{}^2}\right)^{2\alpha_P(t)}M_X{}^{2\alpha_P(0)}\times|\beta_{pP}^{(t)}|^2g_P(t,t,0)\beta_{pP}(0)\cos^2\theta_C\times\frac{(d\sigma/dt)_{Vp\to Vp}}{(d\sigma/dt)_{pp\to pp}},\qquad(12)$$

where

$$G(t)=|\beta_{pP}(t)|^2g_P(t,t,0)\beta_{pP}(0)$$

is determined from purely hadronic interactions. It can be parametrized as[10]

$$G(t)=Ge^{Bt}$$

with

$$G=4\text{ mb GeV}^{-2},\quad B=4.5\text{ GeV}^{-2}.$$

$(d\sigma/dt)_{Vp\to Vp}$ is parametrized as Ae^{bt}. Assuming that $\sigma(Vp\to Vp)$ is the same as $\sigma(\phi p\to\phi p)$ and $A=\sigma_{\phi p\to\phi p}{}^2/16\pi$, we choose $b\simeq 8$ GeV^{-2} and $(d\sigma/dt)_{pp\to pp}=40e^{10t}$ in the s value of our interest (see Fig. 2). Therefore, Eq. (12) reduces to

$$s^2\frac{d\sigma}{dt\,dM_X{}^2}=\frac{\sigma_{\phi p\to\phi p}{}^2}{16\pi}\left(\frac{\omega}{M_X{}^2}\right)^{2\alpha_P(t)}\left(\frac{Q^2}{m_{F^*}{}^2}\right)^{2\alpha_P(t)}\times M_X{}^{2\alpha_P(0)}\frac{Ge^{(B+b)t}\cos^2\theta_C}{(d\sigma/dt)_{pp\to pp}}.\qquad(13)$$

Since all the parameters in Eq. (13) are fixed, the inclusive cross section could be evaluated in the triple-Regge region.

We next turn to evaluate the cross section in the central region. The diagram relevant in this region is shown in Fig. 1(c). The inclusive cross section near $x=0$ (double-Regge region) is given by

$$E_{F^*}\frac{d\sigma}{d^3k}\bigg|_{Jp\to F^*X}=\beta_{F^*P}(\alpha_0,Q^2)\beta_{pP}(0)g_{PF^*}(k_\perp{}^2)=\tilde\beta_{F^*P}(0)\beta_{pP}(0)g_{PF^*}(k_\perp{}^2)\left(\frac{1}{Q^2}\right)^{\alpha_P(0)}.\qquad(14)$$

In order to determine this cross section we consider the model of Farrar and Field.[11] In their model, near $x\simeq 0$,

$$E_\psi\frac{d\sigma}{d^3k}\bigg|_{pp\to\psi X}=f(0)g_{\psi P}(k_\perp{}^2),\qquad(15)$$

and $g_{\psi P}(k_\perp{}^2)$ has the following parametrization:

$$g_{\psi P}(k_\perp{}^2)=\frac{m_\psi{}^{11}e^{-\kappa_\perp}}{(m_\psi{}^2+k_\perp{}^2)^{5.5}}.\qquad(16)$$

Note that in the double-Regge limit

$$E_\psi\frac{d\sigma}{d^3k}\bigg|_{pp\to\psi X}=\beta_{pP}{}^2(0)g_{\psi P}(k_\perp{}^2).\qquad(17)$$

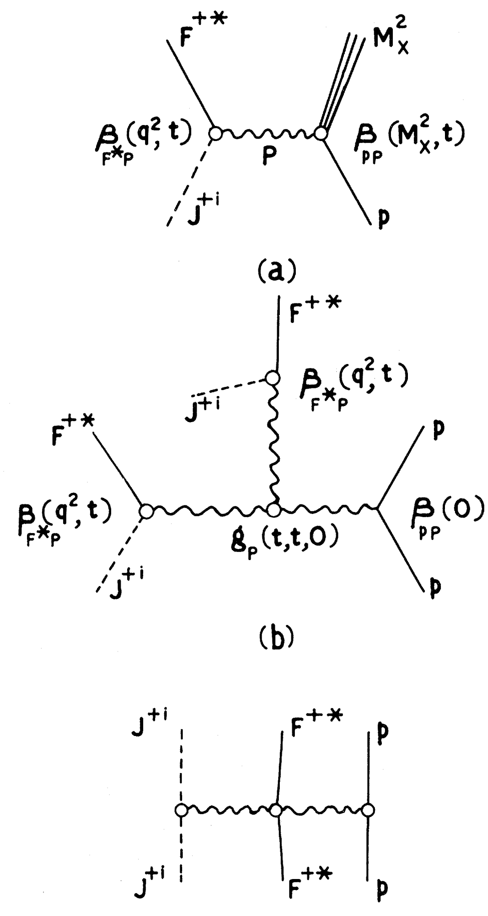

FIG. 1. (a) is the diagram for the inclusive reaction $J^i+p\to F^{*+}$ + anything. Here i stands for components S, L, and R of the current. (b) corresponds to the above diagram in the triple-Regge limit, and (c) is the diagram in the central (double-Regge) region.

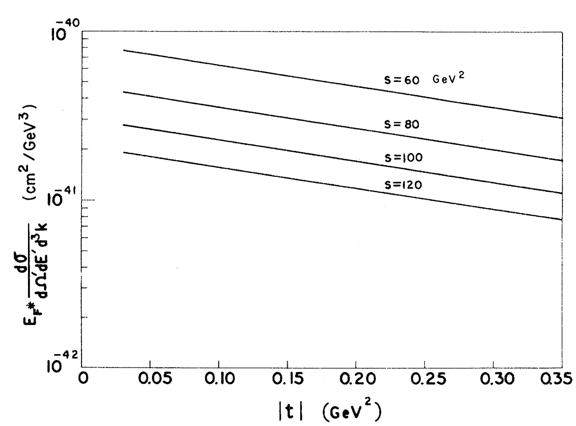

FIG. 2. The inclusive cross section in the triple-Regge limit. The curve is plotted for $s=60$, 80, 100, 120 GeV^2, $Q^2=2$ GeV^2, and $m_{F^*}=2.2$ GeV.

Assuming that g_{F^*P} has the same form as $g_{\psi P}$, we get

$$\frac{E_F\, d\sigma/d^3k|_{Jp\to F^*X}}{E_\psi\, d\sigma/d^3k|_{pp\to\psi X}} = \frac{\tilde\beta_{F^*P}(0)}{\beta_{pP}(0)}\left(\frac{1}{Q^2}\right)^{\alpha_P(0)} \frac{m_{F^*}^{11}(m_\psi^2+k_\perp^2)^{5.5}}{m_\psi^{11}(m_{F^*}^2+k_\perp^2)^{5.5}}$$

$$= \frac{\sigma_{\gamma p}}{\sigma_{pp}}\left(\frac{m_{F^*}}{m_\psi}\right)^{11}\left(\frac{m_\psi^2+k_\perp^2}{m_{F^*}^2+k_\perp^2}\right)^{5.5}. \qquad (18)$$

In deriving the last equation we have replaced the ratio $\tilde\beta_{F^*P}(0)/\beta_{pP}(0)$ by the ratio of the total cross section for photoabsorption to that for pp. Since ψ production in pp collisions is known, it will be possible to estimate the F^*-production cross section in neutrino scattering. We comment here that a similar argument also holds for F^* production in pp scattering.

IV. DISCUSSIONS

We intend to summarize our main results in this section and discuss various features of our model.

(1) In this paper we have calculated the inclusive cross section for production of F^{*+}. In order to estimate cross sections we have used factorization and the fact that the current-hadron amplitude is of the same order of magnitude as the photon-hadron amplitude. In fact, one can formally define current-hadron coupling as follows:

$$\langle 0|J_\mu^{+i}|V^i\rangle = \epsilon_\mu \frac{m_{V^i}^{\,2}}{g_i}, \qquad (19)$$

analogous to

$$\langle 0|J_\mu^{em}|V^i\rangle = \epsilon_\mu \frac{m_{V^i}^{\,2}}{g_i}. \qquad (20)$$

(2) Throughout this paper we used $\alpha_P(0)=1$. There is always the problem of the triple-Pomeron coupling and the Pomeron intercept. It is worth mentioning here that one can choose the intercept of the Pomeron to be less than unity if $g_P(t,t,0)\neq 0$. However, we feel that such technical details could be taken care of by more refined arguments.

(3) It might be interesting to investigate the cross section for $J^i+p\to F^{*+}+p$. We note that if we assume Pomeron exchange we get

$$s^2\frac{d\sigma}{dt} = \frac{1}{8\pi}|\beta_{F^*P}(t,Q^2)|^2|\beta_{pP}(t)|^2\left(\frac{s}{Q^2}\right)^{2\alpha_P(t)}, \qquad (21)$$

and if we use the property of the residue function for large Q^2 [Eq. (7)] we get

$$s^2\frac{d\sigma}{dt} = \frac{1}{8\pi}|\tilde\beta(t,0)|^2|\beta_{pP}(t)|^2\left(\frac{s}{Q^2}\right)^{2\alpha_P(t)}\frac{1}{Q^4}. \qquad (22)$$

Thus one can relate this cross section to the photoproduction cross section of ψ on the proton target.

(4) The cross section for production F^{*+} in the triple-Regge limit is shown in Fig. 2. The inclusive cross sections for $\nu+p\to\mu^-+F^{*+}+X$ are of the order of 10^{-41} cm^2/GeV^3, when $Q^2=2$ GeV^2, $m_{F^*}=2.2$ GeV, and $M_X^{\,2}=0.1s$. We evaluate cross sections for $s=60$, 80, 100, 120 GeV^2.

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[1] J. J. Aubert *et al.*, Phys. Rev. Lett. 33, 1404 (1974); J.-E. Augustin *et al.*, *ibid.* 33, 1406 (1974); C. Bacci *et al.*, *ibid.* 33, 1408 (1974).
[2] G. S. Abrams *et al.*, Phys. Rev. Lett. 33, 1453 (1974).
[3] For detailed references see M. K. Gaillard, B. W. Lee, and J. L. Rosner, Rev. Mod. Phys. 47, 277 (1975).
[4] See addendum to Ref. 3.
[5] S. Okubo *et al.*, Phys. Rev. Lett. 34, 38 (1975).
[6] A. Kazi, G. Kramer, and D. H. Schiller DESY Report No. 75/11 (unpublished).
[7] S. L. Glashow, J. Iliopoulos, and L. Maiani Phys. Rev. D 2, 1258 (1970).
[8] H. D. I. Abarbanel, M. L. Goldberger, and S. B. Treiman, Phys. Rev. Lett. 22, 500 (1969).
[9] H. Pagels, Phys. Lett. 34B, 229 (1971).
[10] Z. F. Ezawa and J. Maharana, Phys. Rev. D 9, 2164 (1974).
[11] G. Farrar and R. D. Field, Phys. Lett. 58B, 180 (1975).