

Hadroproduction of charmonium at large x_F

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(Received 6 March 1995)

We present a higher-twist production mechanism for η_c and χ_J charmonium at large momentum fraction x_F in pion-nucleon collisions. The higher-twist contribution is essentially independent of x_F and therefore becomes dominant at large x_F , where the leading-twist contribution falls off as $(1 - x_F)^3$. We show that the higher-twist mechanism produces longitudinally polarized χ_1 and χ_2 . For the χ_2 , this is clearly different from the leading-twist prediction of transverse polarization. For the χ_1 , the polarization of the leading- and higher-twist contributions is qualitatively similar.

PACS number(s): 13.85.Ni, 12.38.Bx, 13.75.Gx, 14.40.Gx

I. INTRODUCTION

Heavy-quark-antiquark bound states (quarkonium) were among the first systems to be successfully studied using perturbative QCD. Much of the quarkonium decay data were adequately explained by the color-singlet model [1–3]. In this model, the quarkonium states are taken to be nonrelativistic color-singlet $Q\bar{Q}$ states, and decay amplitudes are written as convolutions of the quarkonium wave function with the lowest-order perturbative amplitudes for $Q\bar{Q}$ annihilation. It could then be expected that quarkonium states are dominantly produced by the leading-twist mechanisms where the subprocess (e.g., $gg \rightarrow J/\psi + g$) is related to a decay mechanism (e.g., $J/\psi \rightarrow ggg$) by crossing [4,5]. However, the resulting predictions are contradicted by several experiments, from fixed-target [6–8] to the highest collider energies [9,10].

In this article, we shall concentrate on the production of charmonium at low transverse momentum $P_\perp$ and large longitudinal momentum fraction x_F in fixed-target pion-nucleon collisions. We have earlier pointed out that the anomalous dependence of the production cross section on the mass number A of a nuclear target [11] suggests the importance of higher-twist, nonfactorizable¹ production mechanisms of the J/ψ in this region [13,14]. In a recent work [15], we calculated the polarization of J/ψ 's produced in pion-nucleon reactions in the color-singlet model at leading twist, including contributions from direct parton fusion into J/ψ and from radiative decays of the $\chi_{1,2}$ charmonium states. We found a discrepancy between these leading-twist predictions and experimental data [6,7]. The purpose of this paper is to present a specific higher-twist mechanism for the production of $C = +1$ charmonia (the η_c and χ_J) at large x_F .

We shall show that higher-twist mechanisms are likely to be dominant at large x_F and that their characteristic features may explain some of the charmonium polarization data.

The polarization of J/ψ charmonium produced in pion-nucleon collisions has been measured by the Chicago-Iowa-Princeton Collaboration [6] for $0.3 < x_F < 1.0$. They observe an abrupt change from unpolarized production to longitudinal polarization at $x_F \approx 0.95$. This effect is reminiscent of the longitudinal polarization effect [16,17] observed in the production of dileptons in pion-nucleon collisions, which may be explained by a higher-twist production mechanism first suggested by Berger and Brodsky [18–22].

The higher-twist charmonium production mechanism which we present here is closely related to the higher-twist Drell-Yan mechanism. We have chosen to study the $C = +1$ states because they couple to two gluons, which makes this the simplest case to calculate. We shall show that the $\chi_{1,2}$ are produced longitudinally polarized at large x_F . For the χ_2 , this is strikingly different from the usual leading-twist gluon-fusion mechanism $gg \rightarrow \chi_2$, which only produces transversely polarized χ_2 [15,23]. For the χ_1 , on the other hand, the leading- and higher-twist contributions have similar polarization properties, so that the signature of the higher-twist mechanisms is not as clear as for the χ_2 .

Thus a measurement of the polarization of the χ_2 will provide important information about the production dynamics of charmonium. An observation of the predicted longitudinal polarization of the $\chi_{1,2}$ at large x_F would support the helicity conservation rule between the pion and the large x_F system [18], which is already supported by data for the J/ψ and the Drell-Yan virtual photons.

In Sec. II of this article we outline the calculation of the charmonium production cross sections. We illustrate their x_F dependence in η_c production, in which case compact analytical expressions can be derived. In Sec. III, we present numerical results for the polarization of the $\chi_{1,2}$. In Sec. IV, we summarize our results and discuss the relation of this model to other models of charmonium production.

¹By “nonfactorizable,” we mean “nonfactorizable in terms of single parton distributions.” Note that certain higher-twist processes have been shown to be factorizable in terms of multi-parton correlation functions [12].

II. OUTLINE OF THE CALCULATION

The production of systems that carry a large momentum fraction x_F requires a perturbative momentum exchange involving all the valence partons of the projectile [18,24]. Thus a correlation scale is present in the reaction. The cross sections become suppressed by powers of the ratio of this correlation scale and the hard momentum scale, e.g., f_π^2/m_c^2 in charm production from pions. These mechanisms are therefore of higher twist and are usually neglected. At large x_F , however, inverse powers of $(1-x_F)$ may compensate the suppression, and higher-twist mechanisms may become dominant [25].

Here, we present a higher-twist mechanism for $C = +1$ charmonium production. Our mechanism is described by the Feynman diagram in Fig. 1. Note that whereas the mechanism of Ref. [18] produces both the higher-twist component of the cross section and the perturbative tail of the leading-twist component, our mechanism is purely higher twist. This is evident from the Feynman diagrams. As in the Drell-Yan case, we treat the interaction with the nuclear target simply as the annihilation of one of the pion's valence quarks with an on-shell quark from the target. This is the simplest higher-twist mechanism for charmonium production, but not necessarily the most important one. There may be a significant contribution from higher-order processes where a gluon from the target interacts with the pion, because at high energies and large x_F the parton distribution of the target is sampled at small x_{target} . Moreover, the simplifying assumption that the target just contributes one parton implies that

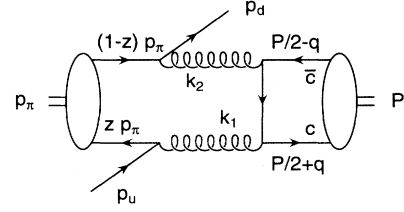


FIG. 1. One of the two Feynman diagrams that describe the higher-twist charmonium production mechanism considered in this article. The four-momenta flow from left to right. The other diagram is obtained by crossing the gluon lines.

the cross section is factorizable in terms of the target parton distribution. An analysis [13] of J/ψ production data [11] suggests that there are important nonfactorizable production mechanisms.

We treat the pion bound-state effects as in Ref. [26] and the charmonium binding as in the color-singlet model [27]. In other words, the amplitude for the reaction $\pi^- + u \rightarrow d + {}^{2S+1}L_J$ is written as a convolution of three factors: the pion distribution amplitude $\phi(z)$, the perturbative-QCD amplitude for $\bar{u}d + u \rightarrow d + c\bar{c}$, and the nonrelativistic wave function $\Phi(P, q)$ of the charmonium state. Here, z is the light-cone momentum fraction carried by the $\bar{u}$ quark in the pion, P is the total four-momentum of the charmonium state, and $2q$ is the relative four-momentum of the charm quarks. The amplitude is

$$A(\pi^- + u \rightarrow d + {}^{2S+1}L_J) = \frac{4\pi\alpha_s\sqrt{2}f_\pi}{9} \bar{u}_C(p_d)\gamma^{\mu_2}\gamma_5\phi_\pi\gamma^{\mu_1}u_C(p_u) \int dz \frac{\phi(z)}{zs(1-z)u} A_{\mu_1\mu_2}(gg \rightarrow {}^{2S+1}L_J), \quad (2.1)$$

where $s = (p_\pi + p_u)^2$, $u = (p_\pi - p_d)^2$, C is a color index, and $A_{\mu_1\mu_2}$ is the truncated amplitude for $gg \rightarrow {}^{2S+1}L_J$, which is the convolution of the quarkonium wave function $\Phi(P, q)$ with the hard amplitude $\mathcal{O}_{\mu_1\mu_2}$ for $gg \rightarrow c\bar{c}$,

$$\mathcal{O}_{\mu_1\mu_2} = \frac{2\pi\alpha_s}{\sqrt{3}} \left[\gamma_{\mu_1} \frac{1}{2} \not{P} - \not{k}_1 + \not{q} + m_c \right] \gamma_{\mu_2} + (1 \leftrightarrow 2), \quad (2.2)$$

as explained in Ref. [27]. For the $C = +1$ states, the relations between $A_{\mu_1\mu_2}$ and $\mathcal{O}_{\mu_1\mu_2}$ can be written as

$$A_{\mu_1\mu_2}(gg \rightarrow {}^1S_0) = \frac{R_S(0)}{\sqrt{16\pi M}} \text{Tr} [\mathcal{O}_{\mu_1\mu_2}(\not{P} - M) \gamma_5]_{q=0}, \quad (2.3)$$

$$A_{\mu_1\mu_2}(gg \rightarrow {}^3P_0) = \frac{iR'_P(0)}{\sqrt{16\pi M}} \text{Tr} \left[\frac{\partial \mathcal{O}_{\mu_1\mu_2}}{\partial q^\alpha} (\gamma^\alpha - P^\alpha/M)(\not{P} + M) + 6\mathcal{O}_{\mu_1\mu_2} \right]_{q=0}, \quad (2.4)$$

$$A_{\mu_1\mu_2}(gg \rightarrow {}^3P_1) = -R'_P(0)\sqrt{\frac{3}{32\pi M^3}} \epsilon^{\alpha\beta\mu\nu} P_\mu \epsilon_\nu^*(\lambda) \text{Tr} \left[\frac{\partial \mathcal{O}_{\mu_1\mu_2}}{\partial q^\alpha} \gamma_\beta(\not{P} + M) - 2\mathcal{O}_{\mu_1\mu_2} \gamma_\alpha \not{P} \gamma_\beta/M \right]_{q=0}, \quad (2.5)$$

$$A_{\mu_1\mu_2}(gg \rightarrow {}^3P_2) = iR'_P(0)\sqrt{\frac{3}{16\pi M}} \epsilon^{\alpha\beta*}(\lambda) \text{Tr} \left[\frac{\partial \mathcal{O}_{\mu_1\mu_2}}{\partial q^\alpha} \gamma_\beta(\not{P} + M) \right]_{q=0}, \quad (2.6)$$

where $M = 2m_c$ is the mass of the charmonium state, $R_S(0)$ is the value of the S -wave wave function at the spatial origin, $R'_P(0)$ is the derivative of the P -wave wave function at the spatial origin, and $\epsilon_\nu(\lambda)$ and $\epsilon^{\alpha\beta}(\lambda)$ are the polarization vector and tensor of the χ_1 and χ_2 . The convolutions of the pion distribution amplitude with $\mathcal{O}_{\mu_1\mu_2}$ and its derivative with respect to q , evaluated at the nonrelativistic limit $q = 0$, are

$$\left[\int dz \frac{\phi(z)}{z(1-z)} \mathcal{O}_{\mu_1\mu_2} \right]_{q=0} = \frac{4\pi\alpha_s}{\sqrt{3}(s-u)} \left[I_0(z_0) G_{\mu_1\mu_2}^{(0+)} + I_1(z_0) G_{\mu_1\mu_2}^{(1-)} \right], \quad (2.7)$$

$$\left[\frac{\partial}{\partial q^\alpha} \int dz \frac{\phi(z)}{z(1-z)} \mathcal{O}_{\mu_1\mu_2} \right]_{q=0} = \frac{8\pi\alpha_s}{\sqrt{3}(s-u)^2} \left\{ 2p_{\pi\alpha} \left[I_0(z_0) G_{\mu_1\mu_2}^{(0-)} - I_1(z_0) G_{\mu_1\mu_2}^{(1+)} \right] + (s-u) I_0(z_0) [g_{\alpha\mu_1}\gamma_{\mu_2} + g_{\alpha\mu_2}\gamma_{\mu_1}] \right. \\ \left. + 2(p_{u\alpha} + z_0 p_{\pi\alpha}) \left[I'_0(z_0) G_{\mu_1\mu_2}^{(0-)} - I'_1(z_0) G_{\mu_1\mu_2}^{(1+)} \right] \right\}, \quad (2.8)$$

where

$$I_n(z_0) = \int dz \frac{\phi(z)}{z(1-z)} \frac{z^n}{z - z_0 + i\epsilon}, \quad (2.9)$$

$$z_0 = \frac{M^2 - u}{s - u}, \quad (2.10)$$

$$G_{\mu_1\mu_2}^{(0\pm)} = P_{\mu_1}\gamma_{\mu_2} - \gamma_{\mu_1}\not{p}_u\gamma_{\mu_2} \\ \pm [P_{\mu_2}\gamma_{\mu_1} - \gamma_{\mu_2}(\not{p}_\pi - \not{p}_d)\gamma_{\mu_1}], \quad (2.11)$$

$$G_{\mu_1\mu_2}^{(1\pm)} = \gamma_{\mu_2}\not{p}_\pi\gamma_{\mu_1} \pm \gamma_{\mu_1}\not{p}_\pi\gamma_{\mu_2}. \quad (2.12)$$

The singularities at $z = 0, 1$ from the gluon propagators are expected to be canceled by the distribution amplitude $\phi(z)$, whereas the singularity at $z = z_0$ from the heavy quark propagator is regularized by the $+i\epsilon$ prescription, which gives the amplitude a nontrivial phase [28].

The presence of the singularity at $z = z_0$ raises the possibility that long-distance contributions to the amplitude could invalidate the use of QCD perturbation theory. However, the transverse momentum of an on-shell intermediate heavy quark is opposite to that of the outgoing heavy quarks. Therefore the probability of forming a quarkonium bound state in the final state becomes small if the intermediate heavy quark propagates over a long distance. This is the same argument that has been applied to hadron Compton scattering [28]. Note that in Compton scattering several singularities are present and one has to demonstrate that the singularities do not pinch.

Below, we shall present results for the symmetric (asymptotic) [26] and two-humped [29] distribution amplitudes²

$$\phi(z) = 6z(1-z), \quad (2.13)$$

$$\phi(z) = z(1-z)[26 - 100z(1-z)]. \quad (2.14)$$

We shall first study η_c production with a symmetric pion distribution amplitude, in which case the analytic

expressions are quite compact. The square of the $\pi u \rightarrow d\eta_c$ amplitude, averaged over the spin and color of the u quark, is

$$|\overline{A}|^2 = \frac{32(4\pi\alpha_s)^4 |R_S|^2 f_\pi^2}{27\pi M} \frac{1}{su(s-u)^2} \left\{ -s^2(L^2 + \pi^2) \right. \\ \left. + 2s(s+u) [L(1+z_0L) + \pi^2 z_0] \right. \\ \left. - (s-u)^2 [(1+z_0L)^2 + \pi^2 z_0^2] \right\}, \quad (2.15)$$

where $L \equiv \ln[(s-M^2)/(M^2-u)]$. The differential cross section is

$$\frac{d\sigma}{dx dP_\perp^2} = \frac{1}{16\pi s_{\text{tot}}} \left[f_{u/N}(s/s_{\text{tot}}) + f_{\bar{d}/N}(s/s_{\text{tot}}) \right] \\ \times \frac{|\overline{A}|^2}{x(1-x)s}, \quad (2.16)$$

where $s_{\text{tot}} = (p_\pi + p_{\text{nucleon}})^2$. In this and the following formulas, we use the light-cone momentum fraction $x = (P^0 + P^3)/(p_\pi^0 + p_\pi^3)$ instead of the longitudinal momentum fraction x_F . There is a simple relation between these two: $x_F = x - (M^2 + P_\perp^2)/(xs_{\text{tot}})$. The two contributions are from $\pi + u \rightarrow d + {}^{2S+1}L_J$ and $\pi + \bar{d} \rightarrow \bar{u} + {}^{2S+1}L_J$, respectively.

In the limit $M^2 \gg P_\perp^2$ with x finite, the cross section simplifies to

$$\frac{d\sigma}{dx dP_\perp^2} = \frac{512\pi^2 \alpha_s^4 |R_S|^2 f_\pi^2}{27M^5 s_{\text{tot}}} \\ \times \left[f_{u/N} \left(\frac{M^2}{xs_{\text{tot}}} \right) + f_{\bar{d}/N} \left(\frac{M^2}{xs_{\text{tot}}} \right) \right] \frac{1}{P_\perp^2} \\ \times x \left\{ \left[1 - (1-x) \ln \frac{1-x}{x} \right]^2 + \pi^2 (1-x)^2 \right\}. \quad (2.17)$$

Now, taking the limit $x \rightarrow 1$ in (2.17) means taking the double limit

$$M^2 \gg P_\perp^2, \quad x \rightarrow 1 \quad \text{with} \quad M^2(1-x) \gg P_\perp^2, \quad (2.18)$$

i.e., the same limit as for the Drell-Yan process in Ref. [18]. Although the limit (2.18) is not quite reached in present-day pion-nucleon reactions [22], it illustrates the behavior of the $P_\perp$ -integrated cross section at large x since it takes into account the fact that as x is increased, the dominant region of the integration shifts to smaller $P_\perp$. The cross section (2.17) simplifies to

²The simplest model for the pion's valence state, $\phi(z) = \delta(z - 1/2)$, leads to rather nonrealistic shapes of the cross sections since the $+i\epsilon$ prescription fails to regularize the singularity of the quark propagator if the distribution amplitude does not vary smoothly over the pole. The same is actually true of the higher-twist contribution [18] to the Drell-Yan cross section. However, the x_F dependence of the angular distribution coefficients λ, μ, ν in the Drell-Yan cross section is smooth even in the δ -function case.

$$\frac{d\sigma}{dx dP_\perp^2} = \frac{512\pi^2 \alpha_s^4 |R_S|^2 f_\pi^2}{27M^5 s_{\text{tot}}} \times \left[f_{u/N} \left(\frac{M^2}{s_{\text{tot}}} \right) + f_{\bar{d}/N} \left(\frac{M^2}{s_{\text{tot}}} \right) \right] \frac{1}{P_\perp^2}. \quad (2.19)$$

The integral $\int dP_\perp^2$ over the perturbative region where $P_\perp^2/(1-x) \gg \Lambda_{\text{QCD}}^2$ brings a logarithmic factor. Apart from this logarithm, the higher-twist contribution is suppressed by f_π^2/M^2 with respect to the leading-twist, leading-order contribution from the subprocess $gg \rightarrow \eta_c$ [5], which is

$$\int_0^\infty dP_\perp^2 \frac{d\sigma_{\text{LT}}}{dx dP_\perp^2} = \frac{\pi^2 \alpha_s^2 |R_S|^2}{3M^3 s_{\text{tot}}} \frac{1}{x} f_{g/\pi}(x) f_{g/N} \left(\frac{M^2}{x s_{\text{tot}}} \right). \quad (2.20)$$

From (2.19) we also see that $d\sigma/dx$ remains constant as $x \rightarrow 1$, up to the logarithmic factor from the $P_\perp$ integration. The leading-twist contribution, on the other hand, falls off according to the power behavior of the pion's gluon distribution, i.e., $(1-x)^3$ from the spectator counting rules³ [30]. At very large x , this will compensate the suppression by f_π^2/M^2 , and the higher-twist component will become dominant.

III. POLARIZATION OF CHARMONIUM

We now consider the production of the $\chi_{1,2}$ states within our higher-twist model. The cross sections for absolute values of the helicity $|\lambda| = 0, 1, 2$ can be extracted by using the covariant polarization sums given in, e.g., Ref. [34]. However, the analytical expressions that are obtained by using a symbolic manipulation program such as REDUCE [35] are too lengthy to be given here (some of them contain over 1000 terms) or even to be used in numerical integration routines. Instead, we have evaluated the double-differential polarized cross sections $d\sigma/dx dP_\perp^2$ ($|\lambda| = 0, 1, 2$) numerically at a characteristic value of the transverse momentum, $P_\perp^2 = M^2(1-x)$.

In Fig. 2 we plot the ratio

$$R(1) = \frac{\sum_{\lambda=\pm 1} d\sigma/dx dP_\perp^2(\lambda)}{d\sigma/dx dP_\perp^2(\lambda=0)} \quad (3.1)$$

of the polarized cross sections for the χ_1 as a function of x_F , using the symmetric and two-humped distribution amplitudes. The ratio of the polarized leading-twist cross sections $d\sigma/dx$ [15] is also shown. In Fig. 3 we plot the ratios $R(1)$ and $R(2)$ for the χ_2 . The leading-

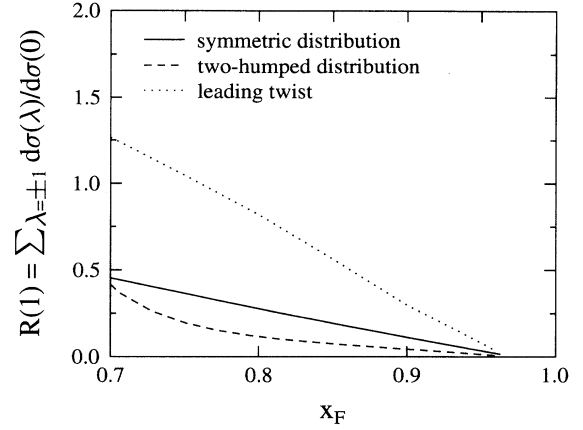


FIG. 2. The ratio $R(1) = \sum_{\lambda=\pm 1} d\sigma(\lambda)/d\sigma(0)$ of the polarized χ_1 production cross sections plotted as a function of x_F . The higher-twist contributions have been evaluated at a characteristic value of the transverse momentum, $P_\perp^2 = M^2(1-x)$, using the symmetric and two-humped pion distribution amplitudes. The leading-twist contribution [15] has been integrated over $P_\perp$.

twist, leading-order gluon-fusion mechanism $gg \rightarrow \chi_2$ gives $d\sigma/dx(0, \pm 1) = 0$, up to corrections of the order of 15% from transverse momentum smearing of the pion [15].

As $x_F \rightarrow 1$, the contribution from our mechanism to both the χ_1 and χ_2 cross sections becomes longitudinally polarized ($\lambda = 0$). For the χ_2 , this is in striking contrast with the leading-twist prediction. For the χ_1 , however, the polarization of the higher-twist component is qualita-

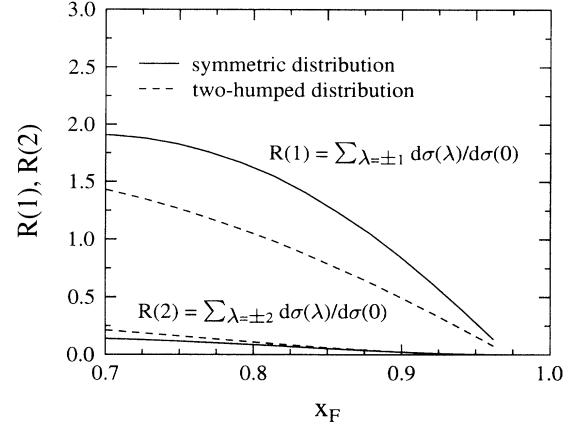


FIG. 3. The ratios $R(1) = \sum_{\lambda=\pm 1} d\sigma(\lambda)/d\sigma(0)$ and $R(2) = \sum_{\lambda=\pm 2} d\sigma(\lambda)/d\sigma(0)$ of the polarized χ_2 production cross sections, plotted as a function of x_F . The higher-twist contributions, which are shown here, have been evaluated at a characteristic value of the transverse momentum, $P_\perp^2 = M^2(1-x)$, using the symmetric and two-humped pion distribution amplitudes. The leading-twist prediction, with transverse momentum smearing and higher-order corrections neglected, is $d\sigma(0, \pm 1) = 0$.

³Experimental determinations of the large x behavior of the gluon distribution [31–33] are based on fitting a purely leading-twist model to Drell-Yan, prompt-photon, and/or charmonium data. The resulting powers of $(1-x)$ therefore partly reflect a hardening of the cross section due to the higher-twist component.

tively similar to that of the leading-twist component, so that polarization cannot be used to discriminate between leading- and higher-twist mechanisms as for the χ_2 .

IV. DISCUSSION

When a charmonium state is produced at a very large momentum fraction x_F , all of the projectile hadron's momentum must be transferred to the charm quarks. This requires a higher-twist mechanism, where all the valence quarks of the projectile couple perturbatively to the charm quarks. The higher-twist nature of these processes is due to the soft correlation scale, e.g., f_π , from the integration over the transverse momentum distribution of the valence state. In this article, we have presented a higher-twist mechanism for the production of the η_c and χ_J charmonia in pion-nucleon collisions. The contribution from this mechanism dominates over the leading-twist contribution at large x_F . As seen from Eq. (2.19) for the η_c , the higher-twist cross section is almost flat as $x \rightarrow 1$, whereas the leading-twist cross section falls off as $(1-x)^3$.

In our model, both the χ_1 and χ_2 are produced longitudinally polarized at large x_F . For the χ_2 , this result is in striking contrast with the leading-twist prediction of purely transverse polarization; for the χ_1 , the polarization of the leading- and higher-twist contributions is qualitatively similar. A direct measurement of the polarization of the χ_2 would give important information about the production mechanisms of charmonium.⁴

The mechanism considered in this paper is the simplest but not necessarily the most important of many possible higher-twist production mechanisms. In mechanisms such as this and the Drell-Yan mechanism of Ref. [18], where the valence state of the pion interacts with a single parton from the target to produce a large x_F system in the final state, the parton distribution of the target is at high energies sampled at small values of x_{target} . Because of the dominance of gluon distributions over quark distributions at small x_{target} , there may be significant contributions from higher-order (in α_s) mechanisms where the target contributes a gluon instead of a quark.

Mechanisms where the target contributes a single parton are factorizable in terms of the parton distribution of the target. Thus they would not bring about the violation of factorization seen in the reaction $\pi + A \rightarrow J/\psi + X$ [13]. Instead, the nuclear number dependence of the cross section would be consistent with the nuclear number dependence of the target parton distribution, as measured,

e.g., in deep inelastic scattering. Experimentally, this is the case in the Drell-Yan reaction. More general higher-twist mechanisms outlined in Ref. [25], where a spectator valence quark interacts softly with the target, would give factorization-violating cross sections.

In this article, we have compared the leading- and higher-twist components of the charmonium production cross sections within the color-singlet model. There is also a potentially important contribution to the production of the χ_J states from mechanisms which involve nonperturbative transitions between color-singlet and color-octet $Q\bar{Q}$ states [37]. These mechanisms are necessary in a rigorous calculation of the χ_J decay widths [38], which in the color-singlet model are infrared divergent at $O(\alpha_s^3)$. In the rigorous analysis, there is a contribution proportional to $\alpha_s^2 \langle \chi_J | \mathcal{O}_8(^3S_1) | \chi_J \rangle$, where the infrared singularity has been absorbed into the nonperturbative matrix element $\langle \chi_J | \mathcal{O}_8(^3S_1) | \chi_J \rangle$. In χ_J production, there is a corresponding contribution from the nonperturbative production matrix element $\langle 0 | \mathcal{O}_8^{XJ}(^3S_1) | 0 \rangle$. The magnitude of this matrix element can be measured in B -meson decays into χ_J charmonium [39,40]. In our mechanism, where the χ_J states couple to two gluons, there is no infrared singularity that would signal the breakdown of the color-singlet model. Nevertheless, the octet contribution is of the same order in the perturbative and nonrelativistic expansions as the singlet contribution. To estimate the importance of color-octet mechanisms in χ_J production at higher twist, one should calculate the cross section for the process $\pi^- + u \rightarrow ^3S_1 + d$ where the $Q\bar{Q}(^3S_1)$ heavy-quark-antiquark pair is in a color-octet state. We postpone the analysis of these contributions to future work.

In spite of these reservations, the present calculation illustrates the perturbative dynamics behind the hadron helicity conservation at large x_F , which has been observed in the Drell-Yan process [16,17] and in J/ψ production [6].

Finally, we note that the present mechanism and other similar mechanisms where the charmonium state couples to virtual gluons allow the production of χ_1 at the same order in α_s as χ_2 . In the leading-twist case, the production of χ_1 in the leading-order subprocess of on-shell gluon fusion, $gg \rightarrow \chi_1$, is forbidden by the Landau-Yang theorem [41], and consequently the ratio of the total χ_1 and χ_2 production cross sections is predicted to be small. The experimental data show equal cross sections [8]. If higher-twist production mechanisms are important down to low x_F , as the failure of the leading-twist predictions of charmonium polarization [15] seems to indicate, they could explain the total cross section anomaly.

ACKNOWLEDGMENTS

This work was supported by the Academy of Finland under Project No. 8579. The author wishes to thank NORDITA and the theory group of SLAC for their warm hospitality during his visits when much of the work reported in this article was done. He is grateful to P. Hoyer, S. J. Brodsky, and W.-K. Tang for important discussions.

⁴The decay angular distributions of P -wave charmonia have actually been measured recently [36]. However, no separation was made between the χ_1 and χ_2 in the angular distribution analysis, and due to limited statistics, the errors are still quite large. With the present statistics, it is also impossible to determine the polarization as a function of the momentum fraction x_F .

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