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PHYSICS LETTERS B

Physics Letters B 621 (2005) 139–144

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$D_s^+ - D_s^-$ asymmetry in photoproduction

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Received 15 April 2005; received in revised form 24 May 2005; accepted 17 June 2005

Available online 28 June 2005

Editor: T. Yanagida

Abstract

Considering of the possible difference in strange and antistrange quark distributions inside nucleon, we investigate the $D_s^+ - D_s^-$ asymmetry in photoproduction in the framework of heavy-quark recombination mechanism. We adopt two distribution models of strange sea, those are the light-cone meson–baryon fluctuation model and the effective chiral quark model. Our results show that the asymmetry induced by the strange quark distributions is distinct, which is measurable in experiments. And, there are evident differences between the predictions of our calculation and previous estimation. Therefore, the experimental measurements on the $D_s^+ - D_s^-$ asymmetry may impose a unique restriction on the strange–antistrange distribution asymmetry models.

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PACS: 12.38.Lg; 12.39.Hg; 13.60.Le

High energy charmed hadron production plays an important role in studying strong interactions. Due to the large charm quark mass, the charm quark involving processes can often be factorized, i.e., into a hard process, which describes the interaction scale much larger than Λ_{QCD} , the typical QCD renormalization scale; and a soft part, which demonstrates the hadronization effects. In recent years, a large production asymmetries between the charmed and

anticharmed mesons have been measured in fixed-target hadro- and photoproduction experiments [1–7]. Among these, photoproduction is thought to give more clean signature than the hadroproduction ones, because there is only one hadron in the initial state. However, it is intriguing to notice that the experimental observation on the asymmetries of the charmed hadron production are greatly in excess of the predictions of perturbative QCD (pQCD) [8–11].

For large transverse momentum processes, the QCD factorization theorem [12] enables us to calculate the cross sections of heavy hadron production by the pQCD. To be specific, for the D_s photoproduc-

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tion, the differential cross section may be expressed as the convolution of the parton distribution function, the partonic cross section, and the fragmentation function, like

$$d\sigma[\gamma + N \rightarrow D_s + X] = \sum f_{i/N} \otimes d\hat{\sigma}[\gamma + i \rightarrow c + \bar{c} + X] \otimes D_{c \rightarrow D_s}, \quad (1)$$

where $f_{i/N}$ is the distribution function of the parton i in a nucleon N ; $d\hat{\sigma}(\gamma + i \rightarrow c + \bar{c} + X)$ is the pQCD calculable cross section of partonic subprocess; and $D_{c \rightarrow D_s}$ is the nonperturbative fragmentation function. In this picture, c and \bar{c} are produced symmetrically at leading order in α_s . The asymmetry appears only in the next-to-leading order (NLO), or higher, corrections. However, the charm–anticharm asymmetries predicted by NLO correction are an order of magnitude smaller than the asymmetries observed in photoproduction experiments [13]. That is, the charmed and anticharmed photoproduction asymmetry may not be fully explained by the charm–anticharm quark production asymmetry. Moreover, the theoretical predictions through above mechanism (1) cannot account for the differences among the asymmetries of D mesons with different light quark flavors.

There have already been some early attempts to explain the observed asymmetries [14,15]. In these approaches, the asymmetry is supposed to appear due to the nonperturbative hadronization effects. Hence, they are all sensitive to unknown distribution functions of partons in the remnant of the target nucleon or photon after the hard scattering. In comparison, the heavy-quark recombination mechanism proposed by Braaten et al. [16] can give a more reasonable explanation to the D meson and Λ_c production asymmetries. In the heavy-quark recombination mechanism, a light parton (q) that participates in the hard-scattering process recombines with a heavy quark (c) or an antiquark (\bar{c}) and subsequently hadronize into the final-state heavy-light meson. The recombination happens only when the light-quark in the final state has momentum of $O(\Lambda_{\text{QCD}})$ in the heavy-quark, or antiquark, rest frame. Namely, the light-quark and the heavy-quark recombine in a small phase space. By virtue of the heavy quark symmetry, $SU(3)$ flavor symmetry and the large N_c limit of QCD, the heavy-quark recombination mechanism gives a simple and predicative

explanation for the asymmetries with two nonperturbative parameters. Since the light-quark q in the recombination model can be either u , d or s quark, it can account for the difference of asymmetries among different light flavors.

In Ref. [17], Braaten et al. calculated the charm/anticharm production ratio and asymmetry of D mesons using the heavy-quark recombination mechanism and confronted their results to the experimental data. In their consideration, the asymmetry of D_s meson comes from the process in which the \bar{c} (c) and light valence-quark of nucleon recombine into a D^- (D^+) meson, while the recoiling c (\bar{c}) quark fragments to D_s^+ (D_s^-) meson. That is, the $D_s^+ - D_s^-$ asymmetry stems from the excess of u and d over \bar{u} and \bar{d} in the nucleon. Because the s and \bar{s} content of the nucleon are identical in their assumption, the asymmetry of D_s meson has the opposite sign as that of D^+ and D^- meson and is relatively small. The experimental data on D_s exist [5,7], but with very large errors, and hence do not tell whether there is an asymmetry or not. Different from their consideration, recent years, many studies show that there is striking strange/antistrange sea asymmetry in the momentum distribution inside the nucleon [18,19]. Stimulated by this idea, in this work, we calculate the D_s meson production asymmetry induced by the strange–antistrange quark distribution asymmetry within nucleon by employing the heavy-quark recombination mechanism.

According to the heavy-quark recombination mechanism, the cross section of D_s meson photoproduction may schematically formulated as:

$$d\sigma[\gamma + N \rightarrow D_s + X] = f_{q/N} \otimes \sum d\hat{\sigma}[\gamma + \bar{s} \rightarrow (c\bar{s})^n + \bar{c}] \times \rho[(c\bar{s})^n \rightarrow D_s], \quad (2)$$

where $(c\bar{s})^n$ represents that the s in the final state has small relative momentum in the \bar{c} rest frame, and n is the color and angular momentum quantum numbers of $(c\bar{s})$ intermediate state. $d\hat{\sigma}[\gamma + \bar{s} \rightarrow (c\bar{s})^n + \bar{c}]$ is the perturbative QCD calculable partonic subprocess. The factor $\rho[(c\bar{s})^n \rightarrow D_s]$ is the probability of the $(c\bar{s})^n$ state to evolve into a final state, here, the D_s .

In our consideration the D_s meson may be produced via two different schemes, i.e.,

$$(a) \quad d\hat{\sigma}[\gamma + \bar{s} \rightarrow (c\bar{s})^n + \bar{c}] \rho[(c\bar{s})^n \rightarrow D_s], \quad (3)$$

$$(b) \quad d\hat{\sigma}[\gamma + q \rightarrow (\bar{c}q)^n + c] \\ \times \sum_{\bar{D}} \rho[(\bar{c}q)^n \rightarrow \bar{D}] \otimes D_{c \rightarrow D_s}, \quad (4)$$

while in Ref. [17] only the second one, the recombination–fragmentation mechanism, was taken into account. In process (a), the $(c\bar{s})^n$ recombines into the D_s meson directly; in process (b), $(\bar{c}q)^n$ recombines into the D_s^- , D^- or \bar{D}^0 meson, and the recoiling c quark fragments to the D_s^+ meson.

The calculation of the partonic cross section $d\hat{\sigma}[\gamma + q \rightarrow (\bar{c}q)^n + c]$ in pQCD is straightforward, and our calculation confirms Ref. [17] results. That is,

$$\frac{d\hat{\sigma}}{dt}[\bar{c}q(^1S_0^{(1)})] \\ = \frac{256\pi^2 e_c^2 \alpha_s^2 m_c^2}{81 S^3} \left[-\frac{S}{U} \left(1 + \frac{\kappa T}{S}\right)^2 \right. \\ \left. + \frac{m_c^2 S}{U^2} \left(-\frac{S^3}{T^3} + \frac{2(1+\kappa)S}{T} + 4\kappa + \frac{\kappa^2 T}{S} \right) \right. \\ \left. + \frac{2m_c^4 S^3}{T^3 U^2} \left(1 + \frac{\kappa T}{S}\right) \right], \quad (5)$$

$$\frac{d\hat{\sigma}}{dt}[\bar{c}q(^3S_1^{(1)})] \\ = \frac{256\pi^2 e_c^2 \alpha_s^2 m_c^2}{81 S^3} \left[-\frac{S}{U} \left(1 + \frac{2U^2}{T^2}\right) \left(1 + \frac{\kappa T}{S}\right)^2 \right. \\ \left. + \frac{m_c^2 S}{U^2} \left(\frac{S^3}{T^3} + \frac{4(2+\kappa)S^2}{T^2} + \frac{2(3+7\kappa)S}{T} \right. \right. \\ \left. \left. + 4\kappa(3+\kappa) + \frac{3\kappa^2 T}{S} \right) + \frac{6m_c^4 S^3}{T^3 U^2} \left(1 + \frac{\kappa T}{S}\right) \right]. \quad (6)$$

Here, $\kappa = e_q/e_c$ is the ratio of the electric charge fractions of light and charm quarks. The Lorentz invariants are defined as $S = (p_q + p_\gamma)^2$, $T = (p_\gamma - p_c)^2 - m_c^2$ and $U = (p_\gamma - p_{\bar{c}})^2 - m_c^2$. p_q , p_γ and p_c are the momenta of the light quark, photon, and c quark, respectively. It is noted that because the relatively small momentum of the light quark in $(c\bar{s})^n$ system, the higher angular momentum excited states are suppressed by power of Λ_{QCD}/p_T or Λ_{QCD}/m_c . For a leading order estimation, we take only the contributions from 1S_0 and 3S_1 states.

To calculate the total cross section ($d\sigma[\gamma + N \rightarrow D_s + X]$) of the production of D_s meson, we need

to know the distribution of strange sea quark. According to the light-cone meson–baryon fluctuation model [18], the strange quark–antiquark asymmetry of the nucleon sea is generated by the intrinsic strangeness fluctuations in the proton wavefunction. In this model, the asymmetry stems from the intermediate $K^+ \Lambda$ configuration of the incident nucleon, which has the lowest off-shell light-cone energy and invariant mass [20]. The momentum distributions of the intrinsic strange and antistrange quarks in the $K^+ \Lambda$ system can be parameterized as

$$s(x) = \int_x^1 \frac{dy}{y} f_{\Lambda/K^+ \Lambda}(y) q_{s/\Lambda} \left(\frac{x}{y}\right), \quad (7)$$

$$\bar{s}(x) = \int_x^1 \frac{dy}{y} f_{K^+/K^+ \Lambda}(y) q_{\bar{s}/K^+} \left(\frac{x}{y}\right), \quad (8)$$

where $f_{\Lambda/K^+ \Lambda}(y)$, $f_{K^+/K^+ \Lambda}(y)$ are the probabilities of finding Λ , K^+ in the $(K^+ \Lambda)$ state; $q_{s/\Lambda}(\frac{x}{y})$, $q_{\bar{s}/K^+}(\frac{x}{y})$ are probabilities of finding strange and antistrange quark in Λ or K^+ states. To estimate these quantities, two simple functions of the invariant mass $M^2 = \sum_{i=1}^2 \frac{K_{\perp i}^2 + m_i^2}{x_i}$ for the two-body wavefunction are given [21],

$$\psi_{\text{Gaussian}}(M^2) = A_{\text{Gaussian}} \exp(-M^2/2\alpha^2), \quad (9)$$

$$\psi_{\text{power}}(M^2) = A_{\text{power}} (1 + M^2/\alpha^2)^{-p}, \quad (10)$$

where the α represents the typical internal momentum scale. In our analysis in this work, we simply adopt the Gaussian type wavefunction for use.

In recently, another model, which based on the effective chiral quark theory, for the distribution of strange sea quark is proposed by Ma et al. [19]. In this model, the strange quark distribution is determined by both the constituent quark distribution and the quark splitting function. For instance,

$$s(x) = P_{sK^+/u} \otimes u_0 + P_{sK^0/d} \otimes d_0, \quad (11)$$

$$\bar{s}(x) = V_{\bar{s}/K^+} \otimes P_{K^+/s/u} \otimes u_0 \\ + V_{\bar{s}/K^0} \otimes P_{K^0/s/d} \otimes d_0. \quad (12)$$

Here, $P_{j\alpha s/i}$ is the splitting function representing the probability of finding a constituent quark j together with a spectator Goldstone (GS) boson ($\alpha = \pi, K, \eta$)

in a parent constituent quark. u_0 and d_0 denote the constituent quark distributions; $V_{i/\alpha}$ is the quark i distribution within the GS boson α .

Using the resultant production cross section, it is straightforward to calculate the asymmetry of the $D_s^+ - D_s^-$ photoproduction, which are defined as

$$\alpha[D_s] = \frac{\sigma_{D_s^+} - \sigma_{D_s^-}}{\sigma_{D_s^+} + \sigma_{D_s^-}}. \quad (13)$$

In our calculation, the two nonperturbative input parameters, ρ_{sm} and ρ_{sf} , are extracted from experiments by fitting to the E687 and E691 data [5,7].

$$\begin{aligned} \rho_{sm} &= \rho_{\text{eff}}[c\bar{u}(^1S_0) \rightarrow D^0] = \rho_{\text{eff}}[c\bar{d}(^1S_0) \rightarrow D^+] \\ &= \rho_{\text{eff}}[c\bar{s}(^1S_0) \rightarrow D_s^+] = \rho_{\text{eff}}[c\bar{u}(^3S_1) \rightarrow D^{*0}] \\ &= \rho_{\text{eff}}[c\bar{d}(^3S_1) \rightarrow D^{*+}] = \rho_{\text{eff}}[c\bar{s}(^3S_1) \rightarrow D_s^{*+}] \\ &= 0.15, \end{aligned} \quad (14)$$

and

$$\begin{aligned} \rho_{sf} &= \rho_{\text{eff}}[c\bar{u}(^3S_1) \rightarrow D^0] = \rho_{\text{eff}}[c\bar{d}(^3S_1) \rightarrow D^+] \\ &= \rho_{\text{eff}}[c\bar{s}(^3S_1) \rightarrow D_s^+] = \rho_{\text{eff}}[c\bar{u}(^1S_0) \rightarrow D^{*0}] \\ &= \rho_{\text{eff}}[c\bar{d}(^1S_0) \rightarrow D^{*+}] = \rho_{\text{eff}}[c\bar{s}(^1S_0) \rightarrow D_s^{*+}] \\ &= 0. \end{aligned} \quad (15)$$

Here, the subscripts sm and sf stand for spin-matched and spin-flipped situations, respectively. Considering the possible uncertainties, the above values are in agreement with those extracted from the hadroproduction [22].

As for the fragmentation function of charm quark to the D_s meson, we exploit the well-known Peterson fragmentation function,

$$D_{c \rightarrow D_s}(z) = P(z; \epsilon) f_{c \rightarrow D_s}. \quad (16)$$

Here, $P(z; \epsilon)$ [23] is the Peterson function and $f_{c \rightarrow D_s}$ [24] is the fragmentation probability. In practice, the parameter ϵ is set to be 0.06, as in Ref. [17], which falls in the region of experimental fitting allowance. All through this Letter, the charm quark mass are set to be 1.5 GeV, and the factorization and renormalization scales are set to be at the meson transverse mass m_T , the $\sqrt{p_\perp^2 + m_{D_s}^2}$. The parton distribution function (PDF) of CTEQ6L [25] is used. For simplicity, we use the average photon beam energy of the E687 experiment, $\langle E_\gamma \rangle = 200$ GeV, in the calculation. To be more accurate, the D_s^* feed down effect is included in our

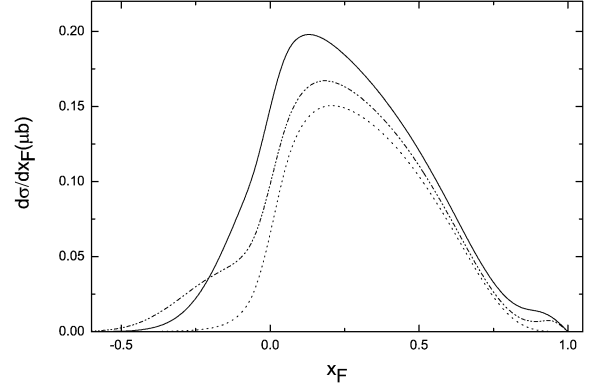


Fig. 1. The inclusive cross sections $d\sigma/dx_F$, calculated by using the light-cone meson–baryon fluctuation model, for D_s production with $\rho_{sm} = 0.15$ and $\rho_{sf} = 0$. The solid, dash-dotted, and dotted lines correspond to the production of D_s^+ , D_s^- in recombination mechanism, and D_s^\pm at the leading order photon–gluon interaction with the fragmentation mechanism, respectively.

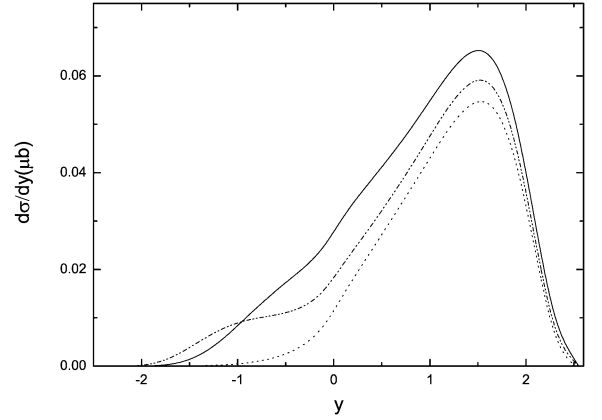


Fig. 2. Inclusive cross sections $d\sigma/dy$ calculated with the light-cone meson–baryon fluctuation model. The curves are described as in Fig. 1.

analysis, but those from other higher excited states are thought to be small and simply neglected.

In Figs. 1 and 2, based on the light-cone meson–baryon fluctuation model, the dependence of the inclusive cross sections of D_s^+ and D_s^- production on the Feynman variable x_F and the rapidity y , respectively, are shown. The corresponding effective chiral quark model results are presented in Figs. 3 and 4. For comparison, the cross section of the leading order photon–gluon fusion process, which will give a symmetric D_s^\pm production through fragmentation, is also

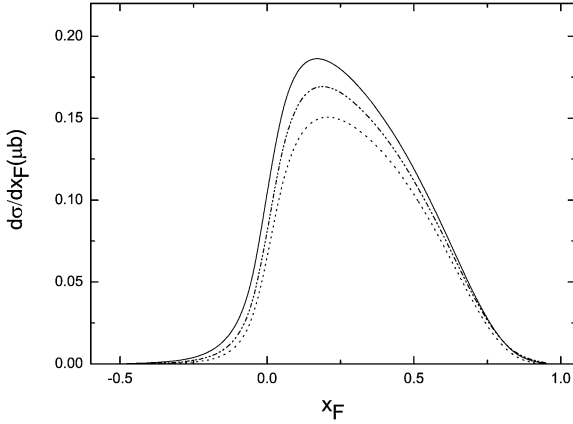


Fig. 3. Inclusive cross sections $d\sigma/dx_F$ calculated with the effective chiral quark model. The curves are described as in Fig. 1.

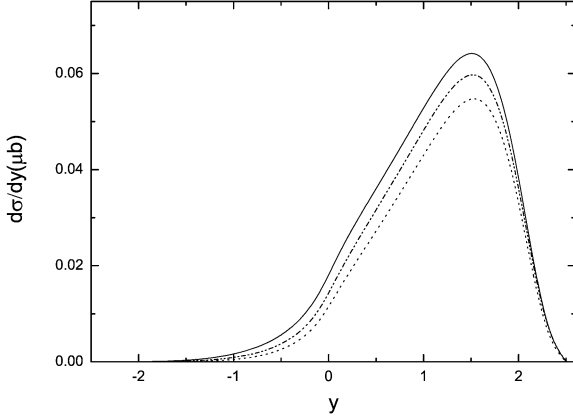


Fig. 4. Inclusive cross sections $d\sigma/dy$ calculated with the effective chiral quark model. The curves are described as in Fig. 1.

presented. We find that the results appear to be insensitive to the variations of the factorization scale and the Peterson parameter ϵ . From Figs. 1–4, we see that cross sections peak in the region of small x_F and large rapidity y . Although the contributions from the recombination mechanism are not the dominant ones, they may give out the asymmetry. Both for the light-cone meson–baryon fluctuation model and the effective chiral quark model, in the region of the experimental cut $x_F > 0$, the corresponding cross sections for D_s^+ are bigger than those for the D_s^- . Thus, the D_s photoproduction asymmetry appears with the effect of strange–antistrange quark asymmetric distributions.

Our predicted D_s production asymmetries, in comparison with Ref. [17], are shown in Fig. 5. We find

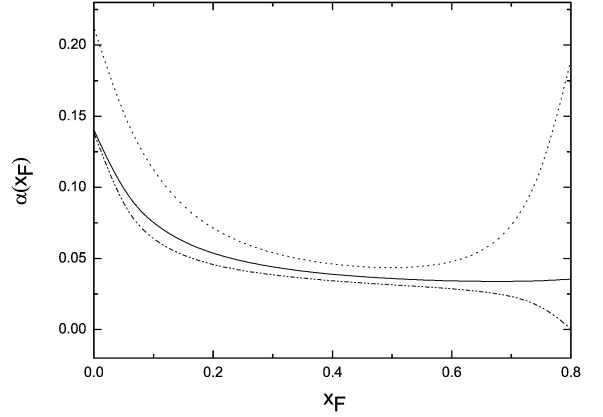


Fig. 5. The asymmetry $\alpha[D_s]$ versus x_F . The dotted and dash-dotted lines correspond to the results from the light-cone meson–baryon fluctuation model and the effective chiral quark model, respectively. The solid line from the Ref. [17] result.

that, by adopting the light-cone meson–baryon fluctuation model, the production asymmetry of D_s meson is about 1.2 times larger than that in Ref. [17]; and by using the effective chiral quark model, it is about 80% bigger only. If we extend to the x_F negative region, the asymmetries coming from the intrinsic strange sea momentum distribution are very big and diverge very much from the prediction of Ref. [17]. In this region, the asymmetry from the light-cone meson–baryon fluctuation model is even flipped relative to the positive x_F region, while the effective chiral quark model and the fragmentation–recombination scheme give the results with the same sign as in the $x_F > 0$ region. Since the obtained data are in the $x_F > 0$ region, here we will not show the results in the $x_F < 0$ region. From Fig. 5, it is obvious that the three different asymmetry producing schemes (models) give results diverging with each other with the x_F increase. This scaling difference leaves the experiments with an opportunity to decide the physical reality.

In summary, in this Letter, the heavy-quark recombination mechanism was employed, which gets a first success in explaining the D meson and Λ_c asymmetry production. We have studied the $D_s^+ - D_s^-$ asymmetry in the photoproduction. Our point is that this asymmetry can be induced not only by the recombination–fragmentation mechanism, Eq. (4), but also by the strange and antistrange distribution asymmetry inside the nucleon, Eq. (3). And, we find the latter effect is even bigger than the former one depending on the em-

ployed model and phase space region. Two QCD relevant strange quark distribution models were used, that is the light-cone meson–baryon fluctuation model and the effective chiral quark model. After including the process (3), the predicted asymmetry increases anyhow. However, what we discussed should be the dominant contributions to the D_s photoproduction asymmetry. Although there remains some uncertainties, such as from the breakdown of $SU(3)$ symmetry, the use of large N_c limit and the NLO corrections to the leading order photon–gluon fusion, etc., our results are adequate to make qualitative predictions, and more detailed discussion on the relevant uncertainties can be found in Refs. [16,17]. It is noticed that the nowadays D_s asymmetry experimental data are very limited and preliminary. With enough data collection in the future, it is expected that the experiment can inversely impose a strong restriction on the strange and antistrange quark distributions by measuring the D_s production asymmetry.

Acknowledgements

We thank Y. Jia and B.Q. Ma for helpful discussions on this topic. This work was supported in part by the Natural Science Foundation of China (NSFC).

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