Λ_c^+/Λ_c^- asymmetry in hadroproduction from heavy-quark recombination

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(Received 17 June 2003; revised manuscript received 22 June 2004; published 21 September 2004)

Asymmetries in the hadroproduction of charm particles directly probe power corrections to the QCD factorization theorems. In this paper, the heavy-quark recombination mechanism, a power correction that explains charm meson asymmetries, is extended to charm baryons. In this mechanism, a light quark participates in the hard-scattering that creates a charm quark and they hadronize together into a charm baryon. This provides a natural and economical explanation for the Λ_c^+/Λ_c^- asymmetries measured in πN and pN collisions.

DOI: 10.1103/PhysRevD.70.054021 PACS numbers: 12.39.Hg, 12.38.Bx, 13.85.Ni, 14.20.Lq

The production of charm particles in fixed-target hadroproduction experiments exhibits large asymmetries that are commonly referred to as the "leading particle effect" [1–4]. Charm hadrons that have a valence parton in common with the beam hadron are produced in greater numbers than other charm hadrons in the forward region of the beam. Asymmetries have also been observed in the production of light particles, such as pions and kaons. Asymmetries in charm particles are particularly interesting, because one can exploit the fact that the charm quark mass m_c is much larger than Λ_{QCD} to make closer contact with fundamental aspects of quantum chromodynamics (QCD). The large mass guarantees that even at small transverse momentum the production process involves short-distance effects that can be treated using perturbative QCD. Furthermore, the nonperturbative long-distance effects of QCD can be organized as an expansion in $\Lambda_{\mathrm{QCD}}/m_c$.

There have been many measurements of the asymmetries for charm mesons [1]. Several proposed charm production mechanisms are able to explain these asymmetries by tuning nonperturbative parameters [5,6]. Recent experiments have also measured the asymmetry for the charm baryon Λ_c^+ [2–4], defined by

$$\alpha[\Lambda_c] = \frac{\sigma[\Lambda_c^+] - \sigma[\Lambda_c^-]}{\sigma[\Lambda_c^+] + \sigma[\Lambda_c^-]}. \tag{1}$$

The WA89 [2] and SELEX [4] experiments observe a large positive asymmetry for Λ_c produced in the forward direction of p and Σ^- beams. These asymmetries are consistent with the leading particle effect, but much larger than those observed for charm mesons. For π^- beams, the leading particle effect predicts no Λ_c asymmetry, but a small positive asymmetry has been observed by the E791 [3] and SELEX [4] experiments. Explaining the Λ_c asymmetries is a severe challenge for any of the proposed mechanisms for generating charm asymmetries [6,7].

The factorization theorems of perturbative QCD [8] imply that the cross section for Λ_c^+ in a collision between two hadrons h, h' is given by

$$d\sigma[hh' \to \Lambda_c^+ + X] = \sum_{i,j} f_{i/h} \otimes f_{j/h'} \otimes d\hat{\sigma}[ij \to c\bar{c} + X] \otimes D_{c\to\Lambda_c^+} + \cdots. \tag{2}$$

Here $f_{i/h}$ is a parton distribution, $d\hat{\sigma}[ij \to c\bar{c} + X]$ is the parton cross section, and $D_{c\to\Lambda_c^+}$ is the fragmentation function for a c quark hadronizing into a Λ_c^+. The ellipsis represents corrections that are suppressed by powers of $\Lambda_{\mathrm{QCD}}/m_c$ or $\Lambda_{\mathrm{QCD}}/p_T$. The leading order processes $gg \to c\bar{c}$ and $q\bar{q} \to c\bar{c}$ produce c and $\bar{c}$ symmetrically. Charge conjugation invariance requires that $D_{c\to\Lambda_c^+} = D_{\bar{c}\to\Lambda_c^-}$, so $\alpha[\Lambda_c] = 0$ at leading order in perturbation theory. Next-to-leading order perturbative corrections [9,10] generate asymmetries that are an order of magnitude too small to explain the data. Therefore the observed asymmetries in charm production must come from the power corrections to Eq. (2).

Recent work has shown that D meson asymmetries in hadroproduction and photoproduction can be explained by an $O(\Lambda_{\mathrm{QCD}}/m_c)$ power correction called heavy-quark recombination [11–13]. In the $c\bar{q}$ recombination process, a light antiquark $\bar{q}$ that participates in the hard scattering emerges with momentum of $O(\Lambda_{\mathrm{QCD}})$ in the rest frame of the charm quark c, and the $c\bar{q}$ pair then hadronizes into a D meson. In this paper, we extend the heavy-quark recombination mechanism to charm baryons. The most important process is cq recombination, which is like $c\bar{q}$ recombination except the $\bar{q}$ is replaced by a light quark q and the cq diquark hadronizes into a charm baryon. Note that the additional light quark needed to form a baryon is created in the hadronization process. The cq produced in the short-distance process picks up a light quark to form a baryon in the same manner that a single heavy-quark picks up a light antiquark to form a meson in the conventional fragmentation process. We will show that this mechanism can explain the observed Λ_c asymmetries.

Other proposed power corrections to heavy-quark production include heavy-quark coalescence [14] and mechanisms based on intrinsic charm in the proton [15,16]. Contributions from intrinsic charm are suppressed by $\Lambda_{\rm QCD}^2/m_c^2$ [17]. Heavy-quark coalescence, in which the heavy quark combines with a spectator from one of the incident hadrons, is suppressed by $\Lambda_{\rm QCD}/m_c$ [14] and could be as important as heavy-quark recombination at low p_T. Calculation of this contribution requires knowledge of the momentum fractions of spectators and, therefore, predictions for differential cross sections are sensitive to poorly constrained nonperturbative functions. On the other hand, the only nonperturbative inputs that enter heavy-quark recombination cross sections are multiplicative factors. The shapes of all distributions are therefore determined by perturbatively calculable hard-scattering cross sections and by the parton distributions of the colliding hadrons. Therefore, heavy-quark recombination is a more economical and predictive mechanism for generating charm asymmetries.

The $c\bar{q}$ recombination cross section for a D meson is

$$d\hat{\sigma}[\bar{q}g \to D] = \sum_n d\hat{\sigma}[\bar{q}g \to c\bar{q}(n) + \bar{c}]\rho[c\bar{q}(n) \to D]. \tag{3}$$

The factor $\rho[c\bar{q}(n) \to D]$ takes into account the nonperturbative hadronization of a $c\bar{q}$ with color and spin quantum numbers n into a final state that includes the D meson. Since the process is inclusive, the quantum numbers of the $c\bar{q}$ produced in the short-distance process can be different from that of the D. The color and spin quantum numbers can both be changed by the emission of soft gluons in the hadronization process. The flavor of the recombining $\bar{q}$ can also be different from that of the valence antiquark in the D, but this requires creating a light quark-antiquark pair which is suppressed in the large N_c limit. Neglecting such contributions, the heavy-quark recombination cross section for D^+ depends on four independent parameters:

$$\rho_1 = \rho[c\bar{d}(^1S_0^{(1)}) \to D^+], \quad \tilde{\rho}_1 = \rho[c\bar{d}(^3S_1^{(1)}) \to D^+],$$
$$\rho_8 = \rho[c\bar{d}(^1S_0^{(8)}) \to D^+], \quad \tilde{\rho}_8 = \rho[c\bar{d}(^3S_1^{(8)}) \to D^+]. \tag{4}$$

Explicit expressions for these parameters in terms of nonperturbative QCD matrix elements can be found in Ref. [18]. They scale with the heavy-quark mass as $\Lambda_{\rm QCD}/m_c$. Analogous parameters for D^0 and D^- mesons are obtained by using isospin symmetry and charge conjugation invariance, while parameters for D^{*+} states are related to those in Eq. (4) by heavy-quark spin symmetry. One might have expected the heavy-quark recombination cross sections to involve a convolution with a nonperturbative function that depends on the fraction of the light-cone momentum of the charmed hadron carried by the light parton. However, to lowest order in $\Lambda_{\rm QCD}/m_c$, only the leading moment of such a distribution is relevant. Therefore, $c\bar{q}$ recombination cross sections are calculable using perturbative QCD up to the four multiplicative factors ρ_1, $\tilde{\rho}_1$, ρ_8, and $\tilde{\rho}_8$.

The direct $c\bar{q}$ recombination process is not expected to be a significant source of charm baryons, since baryon production requires creating at least two light quark-antiquark pairs and is therefore suppressed by $1/N_c^2$ relative to Eq. (4). The leading recombination mechanism for charm baryon production is cq recombination. A leading order Feynman diagram for this process is shown in Fig. 1. Creation of a light quark-antiquark pair is required for the cq to hadronize into either a charm meson or a charm baryon, so there is a $1/N_c$ suppression in either case. This factor makes cq recombination a subleading mechanism for charm mesons, but the leading mechanism for charm baryons. The cq recombination cross section for Λ_c^+ has the form

$$d\hat{\sigma}[qg \to \Lambda_c^+] = \sum_n d\hat{\sigma}[qg \to cq(n) + \bar{c}] \times \eta[cq(n) \to \Lambda_c^+]. \tag{5}$$

The factor $\eta[cq(n) \to \Lambda_c^+]$ takes into account the nonperturbative hadronization of a cq with color and spin quantum numbers n into a final state that includes the Λ_c^+.

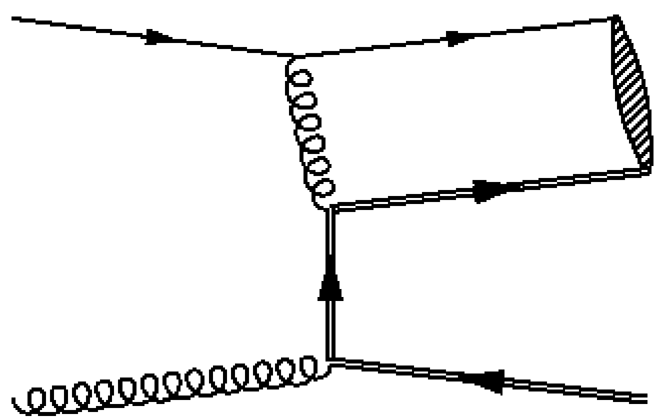

FIG. 1. A Feynman diagram for the heavy-quark recombination process $g + q \to cq(n) + \bar{c}$. The double solid lines and single solid lines represent charm quarks and light quarks, respectively. The shaded blob represents the hadronization of the diquark $cq(n)$.

Isospin symmetry implies that $\eta[cq(n) \to \Lambda_c^+]$ is the same for $q = u, d$, while it is suppressed by $1/N_c$ for $q = s$. There are four possible color and spin states of the cq and, therefore, four independent η parameters:

$$\eta_3 = \eta[cu(^1S_0^{(\bar{3})}) \to \Lambda_c^+], \;\; \tilde{\eta}_3 = \eta[cu(^3S_1^{(\bar{3})}) \to \Lambda_c^+],$$
$$\eta_6 = \eta[cu(^1S_0^{(6)}) \to \Lambda_c^+], \;\; \tilde{\eta}_6 = \eta[cu(^3S_1^{(6)}) \to \Lambda_c^+]. \quad (6)$$

These parameters scale as $\Lambda_{\rm QCD}/m_c$, so cq recombination gives a power-suppressed contribution to the cross section.

The parton cross sections for cq recombination can be calculated using a straightforward generalization of the method described in Ref. [12] for $c\bar{q}$ recombination. Charge conjugation is applied to the line of spinors and Dirac matrices associated with the recombining light quark in Fig. 1. Angular momentum states can then be projected out using the operators of Ref. [12]. The amplitude is projected onto the appropriate color representation, which is either the $\bar{3}$ or 6 of SU(3). A simple prescription for projecting onto the leading moment of the light-cone momentum fraction of the q can be found in Ref. [12]. The parton cross sections for cq recombination are

$$\frac{d\hat{\sigma}}{d\hat{t}}[qg \to cq(n) + \bar{c}] = \frac{2\pi^2\alpha_s^3}{27}\frac{m_c^2}{\hat{s}^3}F(n|\hat{s}, \hat{t}), \quad (7)$$

where $\hat{s}$, $\hat{t}$, and $\hat{u}$ are the standard parton Mandelstam variables for $g + q \to cq(n) + \bar{c}$. The functions $F(n|\hat{s}, \hat{t})$ for the four color and spin channels are

$$F(^1S_0^{(\bar{3})}|\hat{s}, \hat{t}) = -\frac{16U}{S}\left(1 - \frac{ST}{U^2}\right) - \frac{m_c^2}{T}\left(3 + \frac{28U}{T} + \frac{16U^2}{T^2} - \frac{16T^2}{U^2}\right) + \frac{4m_c^4S}{UT^2}\left(3 + \frac{4T}{U} + \frac{8U}{T}\right), \quad (8)$$

$$F(^3S_1^{(\bar{3})}|\hat{s}, \hat{t}) = 3F(^1S_0^{(\bar{3})}|\hat{s}, \hat{t}) - 32\left(\frac{T}{U} - \frac{U^2}{T^2}\right) - \frac{4m_c^2}{T}\left(8 - \frac{6U}{T} - \frac{16U^2}{T^2} + \frac{13T}{U} + \frac{15T^2}{U^2}\right), \quad (9)$$

$$F(^1S_0^{(6)}|\hat{s}, \hat{t}) = -\frac{4U}{S}\left(2 - \frac{5ST}{U^2}\right) - \frac{m_c^2}{T}\left(27 + \frac{14U}{T} + \frac{8U^2}{T^2} - \frac{20T^2}{U^2}\right) + \frac{2m_c^4S}{UT^2}\left(9 + \frac{10T}{U} + \frac{8U}{T}\right), \quad (10)$$

$$F(^3S_1^{(6)}|\hat{s}, \hat{t}) = 3F(^1S_0^{(6)}|\hat{s}, \hat{t}) - \frac{8U}{S}\left(3 + \frac{5ST}{U^2} + \frac{5U}{T} + \frac{2U^2}{T^2}\right) + \frac{4m_c^2S}{U^2}\left(27 - \frac{U}{T} - \frac{U^2}{T^2} - \frac{8U^3}{T^3}\right), \quad (11)$$

where $S = \hat{s}$, $T = \hat{t} - m_c^2$, and $U = \hat{u} - m_c^2$.

The parton cross sections for both cq and $c\bar{q}$ recombination are strongly peaked in the forward direction of the incoming q or $\bar{q}$. For example, consider the ratio of the parton cross sections for cq recombination to that for $gg \to c\bar{c}$, which dominates the fragmentation term in the cross section. We define θ to be the angle between the incoming q and outgoing cq in the parton center-of-momentum frame, so that $T = -S(1 - \beta\cos\theta)/2$, where $\beta = \sqrt{1 - 4m_c^2/S}$. In the backward direction $\theta = \pi$, the ratios of the parton cross sections for recombination and fragmentation scale as

$$\frac{d\hat{\sigma}[qg \to cq(^1S_0^{(\bar{3},6)}) + \bar{c}]}{d\hat{\sigma}[gg \to \bar{c}c]}\Bigg|_{\theta=\pi} \sim \alpha_s\frac{m_c^6}{S^3}, \quad (12)$$

$$\frac{d\hat{\sigma}[qg \to cq(^3S_1^{(\bar{3},6)}) + \bar{c}]}{d\hat{\sigma}[gg \to \bar{c}c]}\Bigg|_{\theta=\pi} \sim \alpha_s\frac{m_c^2}{S}. \quad (13)$$

Thus the heavy-quark recombination contribution is suppressed. For $\theta = \pi/2$, the recombination contribution is suppressed by the same factor in all four channels n:

$$\frac{d\hat{\sigma}[qg \to cq(n) + \bar{c}]}{d\hat{\sigma}[gg \to \bar{c}c]}\Bigg|_{\theta=\pi/2} \sim \alpha_s\frac{m_c^2}{S}. \quad (14)$$

Finally, in the forward direction $\theta = 0$, the ratios scale as

$$\frac{d\hat{\sigma}[qg \to cq(n) + \bar{c}]}{d\hat{\sigma}[gg \to \bar{c}c]}\Bigg|_{\theta=0} \sim \alpha_s, \quad (15)$$

so there is no kinematic suppression of the recombination contribution. The forward enhancement of the $c\bar{q}$ cross section gives charm meson asymmetries which are largest near $x_F \approx 1$. For Λ_c produced in pN collisions, the fragmentation cross section is smaller relative to cq recombination, so the asymmetry is large even for $x_F = 0.2$.

In addition to direct recombination of cq into Λ_c^+, we need to include two additional effects: cq recombination into a heavier charm baryon that subsequently decays into Λ_c^+ and "opposite-side recombination," in which a c produced in a $\bar{c}\,\bar{q}$ or $\bar{c}q$ recombination process fragments into Λ_c^+. The cross sections for the latter process are

$$d\hat{\sigma}[qg \to \Lambda_c^+ + X] = \sum_n d\hat{\sigma}[qg \to \bar{c}q(n) + c] \times \sum_{\overline{D}}\rho[\bar{c}q(n) \to \overline{D}] \otimes D_{c\to\Lambda_c^+}, \quad (16)$$

$$d\hat{\sigma}[\bar{q}g \to \Lambda_c^+ + X] = \sum_n d\hat{\sigma}[\bar{q}g \to \bar{c}\,\bar{q}(n) + c] \times \sum_{\overline{B}}\eta[\bar{c}\,\bar{q}(n) \to \overline{B}] \otimes D_{c\to\Lambda_c^+}. \quad (17)$$

The recombination factors in Eq. (16) and (17) are summed over $\overline{D}$ mesons whose valence partons are $\bar{c}q$

and over antibaryons $\overline{B}$ whose valence partons include $\bar{c}\,\bar{q}$. The process in Eq. (16) gives rise to a small Λ_c asymmetry even in the absence of recombination into charm baryons.

The feed-down from heavier charm baryons is taken into account by defining inclusive η parameters:

$$\eta_{\rm inc}[cq(n)\to\Lambda_c^+] = \eta[cq(n)\to\Lambda_c^+] + \sum_B \eta[cq(n)\to B]Br[B\to\Lambda_c^+ + X]. \tag{18}$$

The sum over B includes all charm baryons that decay into Λ_c^+. They include Σ_c^+, Σ_c^{*+}, and the negative-parity excitations $\Lambda_c^+(2593)$ and $\Lambda_c^+(2625)$ states, all of which have branching fractions into Λ_c^+ of nearly 100%. Charm baryons with strangeness do not contribute to $\eta_{\rm inc}$.

In our analysis, we choose $m_c = 1.5$ GeV, use the one-loop running α_s with $\Lambda_{\rm QCD} = 200$ MeV, and set the renormalization and factorization scales equal to $\sqrt{p_T^2 + m_c^2}$. We use the parton distributions GRV 98 LO [19] for the proton and GRS [20] for the pion. For the fragmentation function for $c\to\Lambda_c^+$, we use

$$D_{c\to\Lambda_c^+}(z) = f_{\Lambda_c^+}\delta(1-z), \tag{19}$$

where $f_{\Lambda_c^+} = 0.076$ is the inclusive fragmentation probability [21]. We also use delta-function fragmentation functions for the charm mesons, since this reproduces the shapes of their momentum distributions more accurately than Peterson fragmentation functions [10,15]. In the opposite-side $\bar{c}q$ recombination cross section, Eq. (16), we include the D and D^* multiplets, but neglect the excited charm mesons. The best one-parameter fit to all the D meson asymmetries measured by E791 gives $\rho_1 = 0.15$ with $\tilde{\rho}_1 = \rho_8 = \tilde{\rho}_8 = 0$. This value of ρ_1 is larger than the value $\rho_1 = 0.06$ obtained in Ref. [13] using Peterson fragmentation functions. The fit to the asymmetries can be improved only slightly by using multiple ρ parameters. The one-parameter fit also gives reasonable results for the shapes of the x_F and p_T^2 distributions. The fits to the x_F and p_T^2 distributions can be improved significantly by using multiple ρ parameters. Note that the sum of recombination parameters appearing in the opposite-side $\bar{c}\,\bar{q}$ recombination cross section, Eq. (17), differs from the inclusive parameter in Eq. (18). The two are related if the sum in Eq. (17) is dominated by the lowest mass $J^P = \frac{1}{2}^+$ and $\frac{3}{2}^+$ SU(3) multiplets. Then, using charge conjugation and a simple quark counting argument, we estimate $\sum_{\overline{B}}\eta[\bar{c}\,\bar{q}(n)\to\overline{B}] \approx 3/2\eta_{\rm inc}[cu(n)\to\Lambda_c^+]$ for $q = u, d, s$.

Because of the large uncertainties associated with the parton distributions of the Σ^-, we focus on the Λ_c asymmetry data from π^- and p beams measured by the E791 and SELEX experiments. The opposite-side recombination mechanism in Eq. (16) generates a positive asymmetry even if all the η parameters vanish. The predictions for $\alpha[\Lambda_c]$ as a function of x_F are shown as dotted lines in Fig. 2 for the pion beam and in Fig. 3 for the proton beam. We have used the value $\rho_1 = 0.15$ that gives the best single-parameter fit to the charm meson asymmetry data. This gives a reasonable fit to the pion beam data, but it severely underpredicts the asymmetry for the proton beam. The χ^2 per degree of freedom is 10.6.

We now fit the Λ_c asymmetry data from the E791 and SELEX experiments by allowing one of the four cq recombination parameters in Eqs. (8)–(11) to be nonzero and constraining the other three to be zero. Since the tree level definitions of the η parameters are positive definite, we consider only non-negative values. The best single-parameter fit and the value of the χ^2 per degree of freedom is

$$\eta_3 = 0.80, \quad \tilde{\eta}_3 = \eta_6 = \tilde{\eta}_6 = 0: \quad \chi^2/\text{dof} = 1.12. \tag{20}$$

If we take $\tilde{\eta}_3$ ($\tilde{\eta}_6$) to be the nonzero parameter, the χ^2 per degree of freedom decreases monotonically from 10.6 to 1.11 (1.61), respectively, as $\tilde{\eta}_3$ ($\tilde{\eta}_6$) increases from 0 to ∞. If we take η_6 to be the nonzero parameter, the χ^2 per degree of freedom has a minimum value of 1.11 at $\eta_6 = 2.78$. We consider values of the η parameters greater than one to be unphysical, because the parameters are supposed to scale as $\Lambda_{\rm QCD}/m_c$. Significant improvements in the fit can be obtained with physically reasonable values of these parameters. For instance, $\tilde{\eta}_3 = 0.1$ gives $\chi^2/\text{dof} = 1.21$. However the best one-parameter fit with a physically reasonable value of the parameter is the one in Eq. (20). We show $\alpha[\Lambda_c]$ as a function of x_F as a solid line in Figs. 2 and 3 for the pion beam and the proton beam, respectively. The one-parameter fit agrees well with both the pion beam and proton beam data. The small

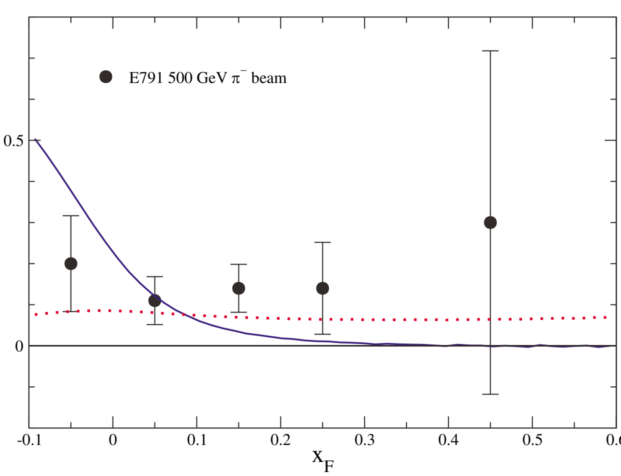

FIG. 2 (color online). $\alpha[\Lambda_c]$ vs x_F for a 500 GeV π^- beam [3]. The solid curve is the best single-parameter fit with $\eta_{3,\rm inc} = 0.80$, while the dotted curve is in the absence of cq recombination.

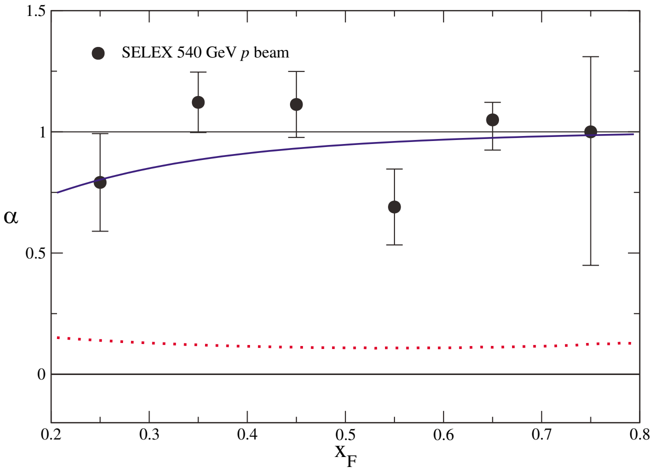

FIG. 3 (color online). $\alpha[\Lambda_c]$ vs x_F for a 540 GeV p beam [4]. Fit parameters for the solid and dotted curves are the same as Fig. 2. The horizontal line at $\alpha = 1$ is the physical upper bound.

value of the χ^2 in Eq. (20) indicates that one could not improve the fits significantly by allowing multiple η parameters to be nonzero. The one-parameter fit also yields good agreement with the observed p_T^2 dependence of the asymmetries as well as the shape of the x_F distributions of Λ_c^+ in both pion and proton beams. The agreement can of course be improved by fitting multiple η parameters. We conclude that the large Λ_c^+/Λ_c^- asymmetry from the proton beam is naturally explained by the cq recombination mechanism.

We have shown that heavy-quark recombination provides a natural and economical explanation of the production asymmetries for charm baryons as well as charm mesons. Further work includes a more systematic analysis of all the charm asymmetry data from hadroproduction experiments and the prediction of charm and bottom asymmetries in present and future experiments. Previous analyses of D meson asymmetries [12,13] do not include contributions from opposite-side cq recombination into charm baryons. This is particularly important for D_s mesons since any asymmetry is generated by opposite-side recombination.

E. B. and M. K. are supported in part by DOE Grant No. DE-FG02-91-ER4069. Y. J. is supported by NSF Grant No. PHY-0100677. T. M. is supported in part by DOE Grants No. DE-FG02-96ER40945 and No. DE-AC05-84ER40150.

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