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Intrinsic heavy-quark states

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The postulate that ordinary hadrons contain intrinsic charm-quark states (such as $|uudc\bar{c}\rangle$ in the proton) at the 1% level is shown to explain two sets of unexpected experimental results: (1) the copious diffractive production of charmed hadrons at large longitudinal momentum in high-energy proton-nucleon and pion-nucleon collisions, and (2) the anomalously large number of same-sign dimuon events observed in deep-inelastic neutrino reactions. We also predict cross sections for open b and t production for high-energy hadron-hadron collisions.

I. INTRODUCTION

Although the nucleon is usually regarded as a three-quark bound state, its actual Fock-state structure in quantum chromodynamics (QCD) must be much more complicated. If we define the state at equal time on the light cone (equivalent to the infinite-momentum frame) then the proton has a general decomposition in terms of color-singlet eigenstates of the free Hamiltonian[1,2]

$$|uud\rangle,\quad |uudg\rangle,\quad |uudq\bar{q}\rangle,\ldots. \tag{1}$$

Since hadrons are color singlets, all infrared divergences cancel and each of the amplitudes $\langle p|uud\rangle$, etc., have a well-defined probability.[2,3]

In this paper we shall explore the consequences of heavy-quark pairs $Q\bar{Q}$ in the Fock-state decomposition of the bound-state wave function of ordinary mesons and baryons. Although proton states such as $|uudc\bar{c}\rangle$ and $|uudb\bar{b}\rangle$ are surely rare, the existence of hidden charm and other heavy quarks within the proton bound state will lead to a number of striking phenomenological consequences.

It will be important to distinguish two types of contributions to the hadron quark and gluon distributions: extrinsic and intrinsic. Extrinsic quarks and gluons are generated on a short time scale in association with a large-transverse-momentum reaction; their distributions can be derived from QCD bremsstrahlung and pair-production processes and lead to standard QCD evolution. The intrinsic quarks and gluons exist over a time scale independent of any probe momentum, and are associated with the bound-state hadron dynamics. In particular, we expect the presence of intrinsic heavy quarks $c\bar{c}$, $b\bar{b}$, etc., within the proton state by virtue of gluon-exchange and vacuum-polarization graphs as illustrated in Fig. 1, where all the colored particles are confined by the effective QCD potential. In fact, by using such a mechanism in the MIT bag model, Donoghue and Golowich[5] have estimated that the probability of finding a five-quark $|uudc\bar{c}\rangle$ configuration bound within the nucleon bag is at the order of 1 to 2%; i.e., the mean number $n_{c/p}$ of intrinsic charmed quarks within the proton is 0.01 to 0.02.

For heavier $Q\bar{Q}$ configurations the vacuum-polarization mechanism of Fig. 1 evidently leads to the scaling

$$\frac{n_Q}{n_c}\sim\left(\frac{m_c}{m_Q}\right)^2, \tag{2}$$

i.e., $n_b/n_c\sim 0.1$ and $n_t/n_c\sim 0.005$ for $m_t=20$ GeV.

The most striking property of an intrinsic heavy-quark state such as $|uudQ\bar{Q}\rangle$ is that the heavy constituents tend to carry the largest fraction of the momentum of the hadron: $\langle x_Q\rangle > \langle x_q\rangle$.

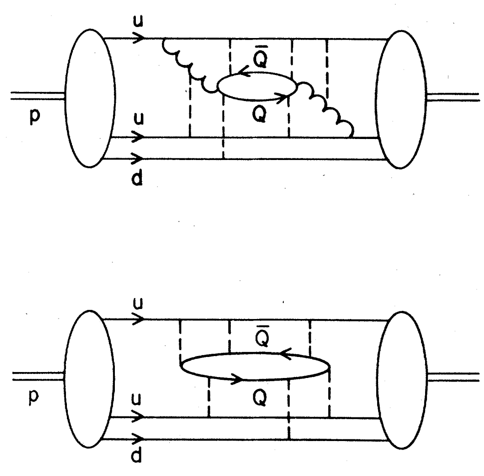

FIG. 1. Diagrams which give rise to the intrinsic heavy quarks ($Q\bar{Q}$) within the proton. Curly and dashed lines represent transverse and longitudinal-scalar (instantaneous) gluons, respectively.

This general feature follows simply from the fact that the constituents in a moving bound state tend to have the same velocity; thus to first approximation the quark momentum $p_Q = E_Q v$ and $x_Q = (p_Q + E_Q)/(p+E)$ scale with its mass. In contrast, the perturbative production of extrinsic quark pairs occurs dominantly at wee $x \sim 0$; the extrinsic quark distribution falls even faster than the gluon distributions at large x.

The existence of the intrinsic charm component in the proton wave function then leads to the following consequences.

(1) The deep-inelastic structure functions of nucleons (and other light hadrons) contain charm quarks at large x_{Bj} (Bjorken x).

(2) Intrinsic charm states will be diffractively dissociated in high-energy hadron collisions, producing open charm ($pN \to \Lambda_c X$, $pN \to DX$, etc.) in the large-x_L beam and target fragmentation regions. The recent observation of such processes at CERN ISR and Fermilab[6-12] gives a strong confirmation of this picture. We also predict that hidden-charm states ψ, χ, η_c, etc., will also be diffractively produced in the fragmentation region, but at much lower rates compared to open charm. As we discuss in Sec. IX, the standard "fusion" subprocesses $gg \to \eta_c$, $q\bar{q} \to \psi$, etc., undoubtedly dominate hidden-charm production in hadron collisions.

(3) Entirely new types of charge-current reactions become possible using the intrinsic charm components of nucleons. For example, the reaction $\nu\bar{c} \to \mu^-\bar{b}$, $\bar{b} \to \mu^- X$ illustrated in Fig. 2 yields a source of *same-sign* dimuon pairs $\nu N \to \mu^-\mu^- X$ (with the second μ^- in the current fragmentation region). Our predictions for these reactions are roughly consistent with recent data (see Sec. VII). These events can occur on $\bar{c}$ quarks with large x_{Bj} and are accompanied by charm particles produced in the target fragmentation region.

The existence of intrinsic charm quarks in the proton wave function gives a new perspective on ordinary hadron structure. It must then also follow that intrinsic gluons, $u\bar{u}$ and $d\bar{d}$ sea quarks, and strange quarks exist in the hadron Fock

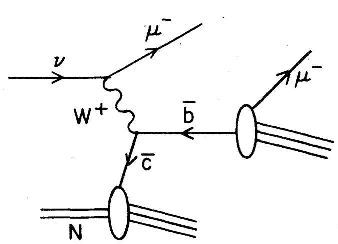

FIG. 2. Same-sign dimuon pair production from the intrinsic charm component of nucleons.

state,[13] apart from distributions generated by QCD evolution. The assumption that the nucleon wave function can be regarded as a simple three-quark amplitude at some fixed scale Q_0 thus becomes untenable.

The most important phenomenological consequences of the intrinsic heavy-quark components are the implications for b- and t-flavored hadron production in $p\bar{p}$ collisions at CERN SPS and Fermilab collider energies. If our prediction, Eq. (2), for heavy-quark probabilities is correct then the diffractive cross sections for open b and t at large energy ($\sqrt{s} \gtrsim 1$ TeV) will be surprisingly large,

$$\sigma(p\bar{p} \to Q\bar{Q}X) \sim \frac{m_c^2}{m_Q^2}\,\sigma(p\bar{p} \to c\bar{c}X)$$

$$\sim \begin{cases} 70\ \mu\text{b (b) with } m_b = 5\text{ GeV}, \\ 3\ \mu\text{b } (t) \text{ with } m_t = 20\text{ GeV}, \end{cases} \tag{3}$$

using an extrapolated value 700 μb for $\sigma(p\bar{p} \to c\bar{c}X)$ (see Sec. V). Furthermore, since the intrinsic transverse momentum associated with the $|uudQ\bar{Q}\rangle$ states tends to increase monotonically with the quark mass (see Sec. III) the Q and $\bar{Q}$ hadron jets are expected to be produced at relatively large transverse momentum.

This paper is organized as follows: In Sec. II we briefly review the experimental results on charm production and discuss how conventional models fail to explain the data. Section III contains a general discussion on nonperturbative properties of hadronic wave functions, and model distributions for $|qqqQ\bar{Q}\rangle$ and $|q\bar{q}Q\bar{Q}\rangle$ states are derived. In Sec. IV we calculate inclusive hadron spectra for open-charm production and the energy dependence is discussed in Sec. V. Section VI contains predictions for t and b cross sections at ISR and Tevatron energies. Consequences of the intrinsic charm pictures for neutrino production of same-sign dimuons and "wrong-sign" muons are derived in Secs. VII and VIII, respectively. Several important factors of ψ suppression relative to open-charm productions are given in Sec. IX and finally Sec. X contains the conclusions.

II. CHARM PRODUCTION AT HIGH ENERGIES

One of the most important results obtained at ISR has been the observation of charm production in pp collisions with remarkably large cross sections and quite unexpected momentum distributions.[6-12]

(1) The cross sections observed for $pp \to \Lambda_c^+ X$ and $pp \to D^+ X$ at $\sqrt{s} = 53$ and 63 GeV are of the order of 0.1 to 0.5 mb.

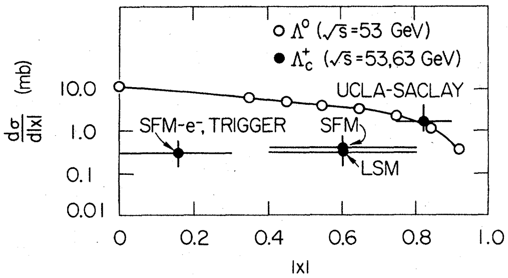

FIG. 3. $d\sigma/d|x|$ for Λ^0 at 53 GeV and Λ_c^+ at 53 and 63 GeV (Refs. 7–9). The smooth curve is a fit to the Λ^0 data points.

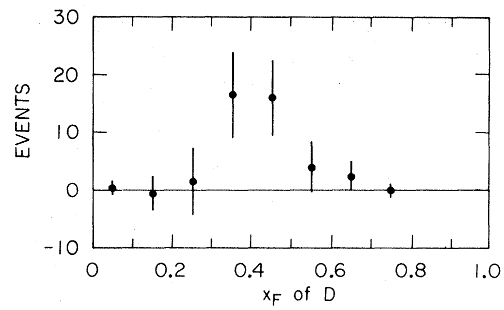

FIG. 5. x_F distribution of events in D peak found in π-p interactions at $\sqrt{s}=20$ GeV (Ref. 12).

(2) The charmed hadrons are produced abundantly in the forward regions of phase space (see Figs. 3–5), contrary to standard expectations. In particular, the D^+, which shares no valence quarks with the proton, would have been expected to be suppressed in the photon fragmentation region. In fact, its Feynman-x distribution seems to be roughly flat in the measured region ($0 \leq x_F \leq 0.4$) and relatively flat in p_T. Moreover, at least one experiment,[7] which triggers a single proton in the other c.m. hemisphere, strongly suggests a diffractive dissociation mechanism for the production.[14] More recently, the Illinois-Fermilab-Harvard-Oxford-Tufts collaboration at Fermilab[13] has observed the reaction $\pi^- p \rightarrow D\bar{D}pX$ at $E_{\rm lab}=217$ GeV ($\sqrt{s}=20.2$ GeV) with a slow-recoil proton trigger. The diffractive $\pi^- p \rightarrow D^0\bar{D}^0 pX$ cross sections are of order 10–40 μb. The charged $D^\pm$ production again shows abundant production of charm at large x_F, peaking at $x_F \sim 0.4$ (see Fig. 5).

It is difficult to understand the physics of these charm-production cross sections from the standpoint of conventional dynamical mechanisms.[15] Although it is expected that perturbative QCD should be applicable to heavy-quark production, the predicted QCD cross sections are of order 10–50 μb at ISR energies[16] and fall steeply with x_F, certainly faster than the gluon distributions. In triple-Regge models, charm production at large x_F is strongly suppressed by the low intercept of the charmed particle trajectories. In models based on conventional soft hadronizations mechanisms, charm pair production is strongly suppressed by factors such as $\exp(-\pi\alpha_R' M_Q^2)$. Most important, it is impossible to explain the observed forward produced D^+ in pp collisions from any valence-quark-recombination or quark-fragmentation model, since none of the proton valence quarks are contained in a D^+.

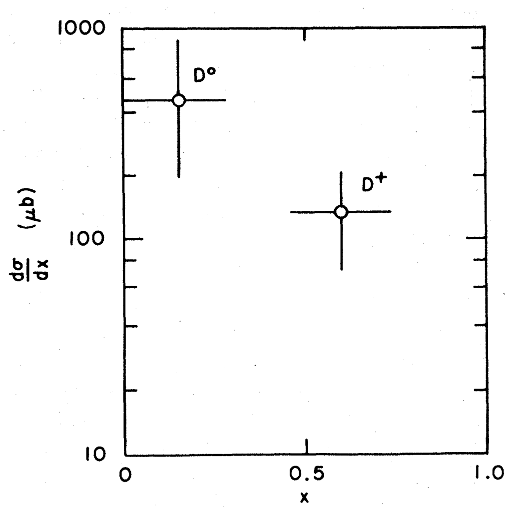

FIG. 4. $d\sigma/dx$ for D^+ (Ref. 10) and D^0 (Ref. 11) at $\sqrt{s}=53$ GeV.

It is evident from the accumulating data that a new dynamical mechanism is involved in open-charm production. Since the experiments indicate that a short time-scale perturbative picture of charm production is not adequate, we shall explore the consequences of the existence of "intrinsic" (long time-scale) charm components in the proton bound state.[17] As discussed in the Introduction, we assume that the number of charmed quarks per nucleon is at the 1 to 2% level.

It is clear that intrinsic charm states are easily dissociated into open-charm hadrons at high-energy collisions (see Fig. 6). The production can be diffractive at high energies since only small

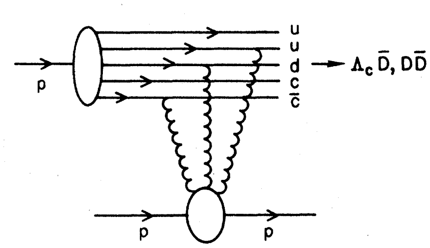

FIG. 6. Diffractive production of open charm from the intrinsic charm component of nucleons.

momentum transfers are needed. In fact, a 300-μb charm-production cross section can be easily understood if ~1% of all pp inelastic events involve dissociation of the intrinsic charm state.

Since the mass of the diffractively excited state M should be much larger than the heavy-quark mass one does not expect cross sections at this magnitude until high energies. A detailed discussion of the energy dependence is given in Sec. V.

III. HADRONIC WAVE FUNCTIONS

As we have discussed in the Introduction, the hadronic wave function in QCD has a well-defined Fock-state decomposition at equal time on the light cone in terms of quark and gluon momentum states.[4] The wave-function amplitude for each Fock component has the form

$$\psi_{(n)}(\vec{k}_{\perp i}, x_i), \quad 0 \le x_i \le 1, \quad i = 1, \ldots, n, \tag{4}$$

where by momentum conservation $\sum_{i=1}^{n} x_i = 1$ and $\sum_{i=1}^{n} k_{\perp i} = 0$. The x_i are the light-cone (infinite-momentum-frame) momentum fractions $(k^0 + k^3)/(p^0 + p^3)$ for each constituent. (Spin labels are suppressed.) The state is off the $p^- = p^0 - p^3$ mass shell

$$p^- - \sum_{i=1}^{n} k^- = \frac{1}{p^+}\left[M^2 - \sum_{i=1}^{n}\left(\frac{k_\perp^2 + m^2}{x}\right)_i\right]. \tag{5}$$

The standard quark and gluon distributions are obtained from integrating the square of the wave functions up to the momentum scale Q of the probe and summing over all Fock states:

$$G_{q/p}(x, Q^2) \propto \sum_n \int [d^2k_\perp][dx]\delta(x - x_q)|\psi_{(n)}(k_{\perp i}, x_i)|^2 \times \theta(k_{\perp i}^2 < Q^2). \tag{6}$$

The "extrinsic" quarks and gluons correspond to the standard bremsstrahlung and $q\bar{q}$ pair-production processes of perturbative QCD. These perturbative contributions yield wave functions with minimal power-law falloff,

$$|\psi(k_{\perp i})|^2 \sim \frac{1}{k_{\perp i}^2}, \tag{7}$$

and lead to the logarithmic evolution of the structure functions. In contrast, the intrinsic contributions to the quark distribution are associated with the bound-state dynamics and necessarily have a faster falloff in $k_{\perp i}$ ($\psi \sim 1/k_\perp^2$ or faster[2]). The intrinsic states thus contribute to the initial quark and gluon distributions. A simple illustration of extrinsic and intrinsic $|uudq\bar{q}\rangle$ contributions to the deep-inelastic structure functions is shown in Figs. 7(a) and 7(b). We see that the existence of gluon-exchange graphs, plus vacuum-polarization insertions, automatically yields an intrinsic $|uudq\bar{q}\rangle$ Fock state.[13] Even if we imagined that in the nucleon rest frame there are effectively only three quarks in the nucleon, the gluon-exchange diagrams automatically generate $|uudg\rangle$ and $|uudq\bar{q}\rangle$ components when boosted to infinite momentum or the light cone. In fact, the magnitude of the p-Δ hyperfine splitting due to transversely polarized gluon exchange must yield a lower bound on the intrinsic gluon and $q\bar{q}$ components.[18] A complete calculation must take into account the binding of the gluon and $q\bar{q}$ constituents inside the hadron (see Fig. 1) so that the analysis is necessarily nonperturbative.

We also note that the normalization of the $|uudq\bar{q}\rangle$ state is not necessarily tied to the normalization of the $|uudg\rangle$ components since the latter only refers to transversely polarized gluons. Figure 1(b) shows that $q\bar{q}$ pairs arise from the longitudinal-scalar (instantaneous) parts of the vector potential.

The general form of a Fock-state wave function is

$$\psi(k_{\perp i}, x_i) = \frac{\Gamma(k_{\perp i}, x_i)}{M^2 - \sum_{i=1}^{n}\left(\dfrac{m^2 + k_\perp^2}{x}\right)_i}, \tag{8a}$$

where Γ is the truncated wave function or vertex function. The actual form of Γ must be obtained from the nonperturbative theory, but following Ref. 4 it is reasonable to take Γ as a decreasing function of the off-energy-shell variable $\mathcal{E} = M^2 - \sum_{i=1}^{n}[(m^2 + k_\perp^2)/x]i$. Independent of the form $\Gamma(\mathcal{E})$, we can read off some general features of the quark distributions.

(1) In the limit of zero binding energy ψ becomes singular and the fractional momentum distributions peak at the values $x_i = m_i/M$. More generally, $\mathcal{E}$ is minimal and the longitudinal momentum distributions are maximal when the constituents with the largest transverse mass $m_\perp = (m^2 + k_\perp^2)^{1/2}$ have the largest light-cone fraction x_i. This is equivalent to the statement that constituents in a moving bound state tend to have the same rapidity.

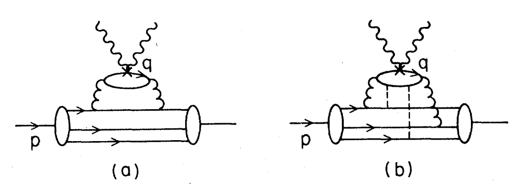

FIG. 7. (a) Example with contribution to the deep-inelastic structure functions from an extrinsic quark q; (b) from an intrinsic quark q.

(2) The intrinsic transverse momentum of each quark in a Fock state generally increases with the quark mass. In the case of a power-law wave function $\psi \sim (\mathcal{E})^{-\beta}$ we have $\langle k_\perp^2 \rangle \propto \langle m_Q^2 \rangle$; for an exponential wave function $\psi \sim \exp(-\beta \mathcal{E}^{1/2})$ the dependence is $\langle k_\perp^2 \rangle \propto m_Q$.

In the limit of large $k_\perp$ one can use the operator product expansion near the light cone (or equivalently gluon-exchange diagrams) to prove that, modulo logarithms, the Fock-state wave functions fall off as inverse powers of $k_\perp^2$.[2] For our purpose, which is to illustrate the characteristic shape of the Fock states containing heavy quarks, we will choose a simple power-law form for the Fock-state longitudinal-momentum distributions

$$P_{(n)}(x_1, \ldots, x_n) = N_{(n)} \frac{\delta\left(1 - \sum_{i=1}^{n} x_i\right)}{\left(M^2 - \sum_{i=1}^{n} \frac{\hat{m}_i^2}{x_i}\right)^2}, \tag{8b}$$

where the $\hat{m}_i^2$ are identified now as effective transverse masses $\hat{m}_i^2 = m_i^2 + \langle k_\perp^2 \rangle_i$ and the $\langle k_\perp^2 \rangle$ are average transverse momentum. With this choice, single-quark distributions have power-law falloffs $(1-x)^2$ and $(1-x)^3$ for mesons and baryons, respectively.

For example, consider a $|\bar{q}Q\rangle$ state, e.g., a D meson. Here the momentum distributions of the two quarks are according to Eq. (8b) given by

$$P(x_1, x_2) = N \frac{\delta(1 - x_1 - x_2)}{\left(m_D^2 - \frac{\hat{m}_c^2}{x_1} - \frac{\hat{m}_u^2}{x_2}\right)^2}. \tag{9}$$

From this expression we obtain the charmed-quark distribution

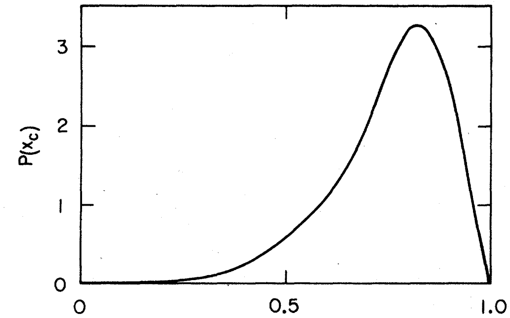

FIG. 8. The x distribution of the charmed quark in a D meson.

$$P(x_1) = \int_0^{1-x_1} P(x_1, x_2)\, dx_2 = N' \frac{1}{\left(1 - \frac{1}{x_1} - \frac{\epsilon}{1 - x_1}\right)^2}, \tag{10}$$

where $N' = N/m_D^4$, $\epsilon = \hat{m}_u^2/m_D^2$, and we take $m_D^2 \approx \hat{m}_c^2$. We see from Fig. 8 that the c quark tends to carry most of the D-meson momentum ($\langle x_1 \rangle = 0.73$). This leading feature of the c quark is due to the fact that the quarks should have roughly the same velocity in order for the hadron to "stay together." This can be seen more explicitly by minimizing the off-shell denominator in Eq. (9),

$$\frac{\hat{m}_c^2}{x_1^2} = \frac{\hat{m}_u^2}{x_2^2}, \tag{11}$$

keeping the transverse masses fixed. (A related idea has previously been considered by Bjorken and Suzuki[19] in the context of charm fragmentation into hadrons.)

We now turn to the discussion of $|uudQ\bar{Q}\rangle$ and $|u\bar{d}Q\bar{Q}\rangle$ states. For a $|uudc\bar{c}\rangle$ proton Fock state the momentum distribution is given by

$$P(x_1, \ldots, x_5) = N \frac{\delta\left(1 - \sum_{i=1}^{5} x_i\right)}{\left(m_p^2 - \sum_{i=1}^{5} \frac{m_i^2}{x_i}\right)^2}. \tag{12}$$

In the limit of heavy quarks $\hat{m}_4^2 = \hat{m}_5^2 = \hat{m}_6^2 \gg m_p^2$, $\hat{m}_i^2$ $(i = 1, 2, 3)$ we get

$$P(x_1, \ldots, x_5) = N_5 \frac{x_4^2 x_5^2}{(x_4 + x_5)^2} \delta\left(1 - \sum_{i=1}^{5} x_i\right), \tag{13}$$

where $N_5 = 3600\, P_5$ is determined from $\int dx_1 \ldots \times dx_5 P(x_1, \ldots, x_5) = P_5$, where P_5 is the $|uudc\bar{c}\rangle$ Fock-state probability. Integrating over the light quarks (x_1, x_2, and x_3) we get the charmed-quark distributions

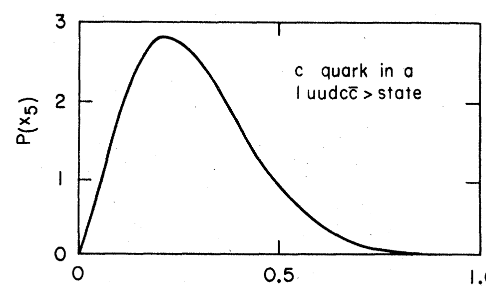

FIG. 9. The x distribution of the charmed quark in a $|uu\, dc\bar{c}\rangle$ state.

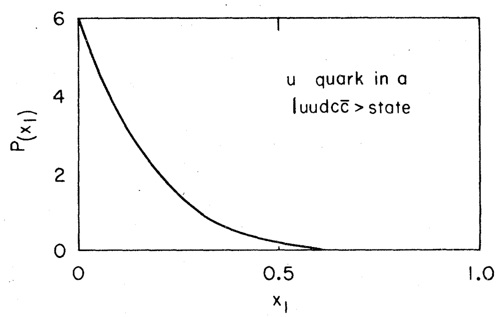

FIG. 10. The x distribution of a light quark in a $|uudc\bar{c}\rangle$ state.

$$P(x_4, x_5) = \tfrac{1}{2}N_5 \frac{x_4^2 x_5^2}{(x_4+x_5)^2}(1-x_4-x_5)^2 . \qquad (14)$$

By performing one more integration we obtain the charmed-quark distribution

$$P(x_5) = \tfrac{1}{2}N_5 x_5^2[\tfrac{1}{3}(1-x_5)(1+10x_5+x_5^2) - 2x_5(1-x_5)\ln(1/x_5)] , \qquad (15)$$

which has average $\langle x_5\rangle = \frac{2}{7}$ and is shown in Fig. 9. This is to be contrasted with the corresponding light-quark distribution derived from Eq. (13) and shown in Fig. 10:

$$P(x_1) = 6(1-x_1)^5 P_5 . \qquad (16)$$

We estimate P_5 from the magnitude of "diffractive" production of $\Lambda_c (pp \to p\Lambda_c X)$,[6]

$$P_5 = P(|uudc\bar{c}\rangle) = \frac{\sigma_{\Lambda_c}}{\sigma_{\text{tot}}} \simeq \frac{300\ \mu\text{b}}{40\ \text{mb}} \approx 0.01 , \qquad (17)$$

in agreement with the bag-model estimate.[5] The charmed-quark distribution $c(x) = P(x_5)$ should be measurable in leptoproduction for high enough Q^2 and $W^2 > W_{\text{th}}^2 \approx 20\ \text{GeV}^2$. Hence to measure $c(x)$ at, e.g., $x = 0.5$, requires $Q^2 > 20\ \text{GeV}^2[x = Q^2/(Q^2 + W^2)]$. In Fig. 11 we show the same quantity, both $c(x)$ from the intrinsic sea and the one calculated from perturbative QCD (Ref. 20) at $Q^2 = 50\ \text{GeV}^2$.

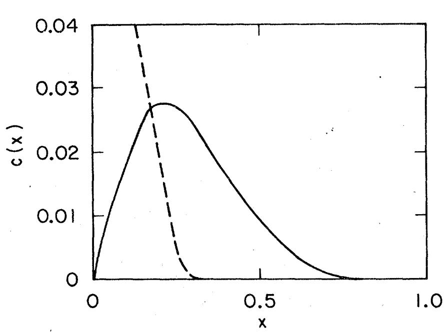

FIG. 11. Comparison of the x distribution of the charm quark in the intrinsic sea (solid line) and the sea calculated from perturbative QCD (Ref. 20) at $Q^2 = 50$ GeV^2 (dashed line).

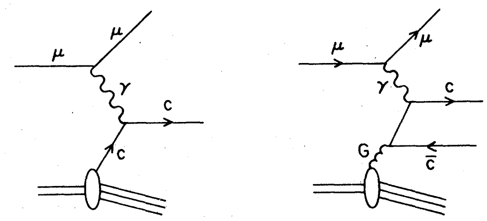

FIG. 12. Leptoproduction of charm from the intrinsic charm sea and via the photon-gluon-fusion model, respectively.

We have neglected the small Q^2 evolution of the intrinsic charmed sea.

From Fig. 11 we see that intrinsic charmed quarks in the proton are "rare" but not "wee." Of course, high-x_{Bj} leptoproduction data would provide a most crucial test of the intrinsic idea. At present, data only exist for small $x < 0.13$ and seem to be adequately described by the proton-gluon-fusion model.[21] Note that the intrinsic charm component gives single charm-quark jets in the current fragmentation region, whereas the photon-gluon-fusion model[22] gives two jets (charm and anticharm) (see Fig. 12).

We next compute the corresponding x distributions for meson Fock states, e.g., $|u\bar{d}c\bar{c}\rangle$. Again neglecting meson and light-quark masses compared to the heavy-quark mass, we get

$$P(x_1, \ldots, x_4) = N_4\left(\frac{x_3 x_4}{x_3+x_4}\right)^2 \delta\left(1 - \sum_{i=1}^{4} x_i\right) , \qquad (18)$$

where $x_3 = x_c$, $x_4 = x_{\bar{c}}$, and $N_4 = 600 P_4$. By integrating over x_1, x_2, x_3 we obtain the single-quark distribution

$$P(x_4) = N_4 x_4^2[\tfrac{1}{2}(1+4x_4-5x_4^2) + x_4(2+x_4)\ln x_4] , \qquad (19)$$

which is shown in Fig. 13 and has $\langle x_4\rangle = \frac{1}{3}$. The corresponding light-quark spectrum from Eq. (18)

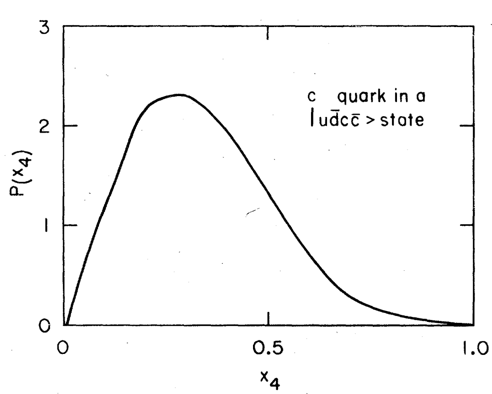

FIG. 13. The x distribution of the charm quark in the $|u\bar{d}c\bar{c}\rangle$ state.

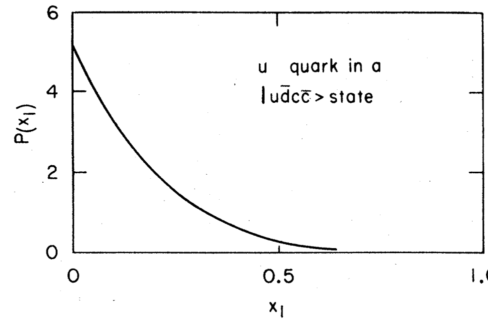

FIG. 14. The x distribution of the light quark in the $|u\bar{d}c\bar{c}\rangle$ state.

is given by

$$P(x_1)=5(1-x_1)^4P_4\,, \tag{20}$$

with $\langle x_1\rangle=\frac{1}{6}$ and is shown in Fig. 14.

It should be remarked that the $\hat{m}_Q\to\infty$ approximations used above are quite accurate. It turns out that keeping the masses implies that the denominator (x_3+x_4) in Eq. (13) is just replaced by $(x_3+x_4+\epsilon x_3x_4)$, where $\epsilon=m_p^2/\hat{m}_Q^2$ leading to only a few percent corrections. The light-quark distributions [Eqs. (16) and (20)] will be hidden in the perturbative extrinsic sea of the hadron in large-momentum-transfer experiments.

IV. INCLUSIVE CHARM PRODUCTION AND HADRON-HADRON COLLISIONS

As we discussed in Sec. II it is reasonable to assume that the intrinsic heavy-quark states, $|qqqc\bar{c}\rangle$ and $|q\bar{q}c\bar{c}\rangle$, should be "fragile" and can be materialized into open-charm hadrons in high-energy, low-momentum-transfer reactions. In principle, the Λ_c and D spectra can be calculated from the strong overlap between the five-quark and the charmed-hadron state wave functions, allowing for decays of excited state, etc. For the purpose of obtaining the x_F distributions we shall use a simple recombination mechanism for the quarks involved in the states. Neglecting its binding energy, the Λ_c spectrum is given by combining the u, d, and c quark in $|uudc\bar{c}\rangle$ to obtain

$$P(x_{\Lambda_c})=N_5\int_0^1\prod_{i=1}^{5}dx_i\,\delta(x_{\Lambda_c}-x_2-x_3-x_4)\times\left(\frac{x_4x_5}{x_4+x_5}\right)^2\delta\left(1-\sum_{i=1}^{5}x_i\right) \tag{21}$$

(see Fig. 15) with $\langle x_{\Lambda_c}\rangle=\frac{1}{7}+\frac{1}{7}+\frac{2}{7}=\frac{4}{7}$. The ISR data for $(d\sigma/dx)(pp\to\Lambda_cX)$ [see Fig. 3(a)] are consistent with the prediction of Eq. (21) that charmed baryons are produced in the forward fragmentation region, although the existing data are too scarce for a detailed comparison. We expect that the low-x region for charm production will be filled in by both perturbative and higher-Fock-state intrinsic contributions. The corresponding distribution for $D^-(\bar{c}d)$ is given by

$$P(x_{D^-})=N_5\int_0^1\prod_{i=1}^{5}dx_i\,\delta(x_{D^-}-x_3-x_5)\times\left(\frac{x_4x_5}{x_4+x_5}\right)^2\delta\left(1-\sum_{i=1}^{5}x_i\right), \tag{22}$$

with $\langle x_{D^-}\rangle=\frac{1}{7}+\frac{2}{7}=\frac{3}{7}$, and is shown in Fig. 16. The $D^+(c\bar{d})$ distribution would, in principle, be obtained from the $|uudc\bar{c}d\bar{d}\rangle$ Fock state of the proton, where the $d\bar{d}$ could be extrinsic or intrinsic. Assuming that the $\bar{d}$ momentum is small, the D^+ distribution should be close to that of the c quark shown in Fig. 12. These predictions apply for forward production ($x_F\gtrsim 0.1$), where perturbative contributions and higher-Fock-state contributions can be neglected.

For the π fragmentation region in the reaction $\pi p\to D\bar{D}X$ the x_F distribution of the D's is obtained in the same way:

$$P(x_D)=N_4\int\prod_{i=1}^{4}dx_i\,\delta(x_D-x_1-x_3)\times\left(\frac{x_3x_4}{x_3+x_4}\right)^2\delta\left(1-\sum_{i=1}^{4}x_4\right) \tag{23}$$

with $\langle x_D\rangle=\frac{1}{2}$. $P(x_D)$ is shown in Fig. 17 and again we observe that the intrinsic model gives a consistent picture of the forward nature of the charm-production data (compare Fig. 5).

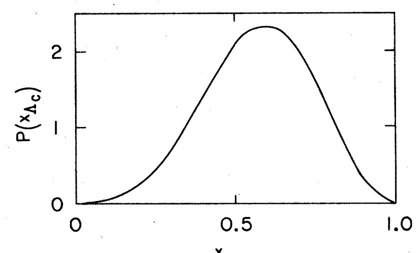

FIG. 15. The x distribution of the Λ_c from the intrinsic charm $|uudc\bar{c}\rangle$ component of the proton.

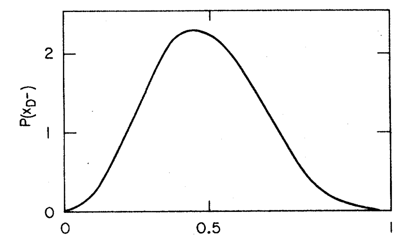

FIG. 16. The x distribution of the D^- from the intrinsic charm $|uudc\bar{c}\rangle$ component of the proton.

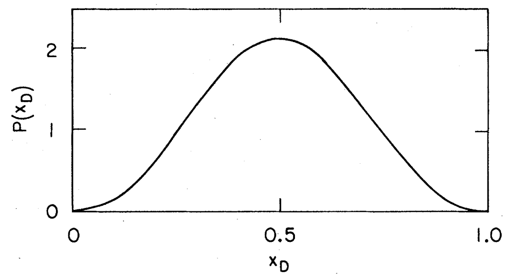

FIG. 17. The x distribution of the D meson from the intrinsic charm component of the pion.

V. ENERGY DEPENDENCE OF HADROPRODUCTION OF HEAVY QUARKS

For perturbative heavy-quark production mechanisms such as the gluon-gluon-fusion model[16] the energy dependence of the cross section essentially comes from the lower limit $m_Q/(2\sqrt{s})$ of convolution integrals, and gives rise to a logarithmic energy dependence. To study the energy dependence of the "diffraction" mechanism with intrinsic heavy quarks we will use the empirical formula for high-mass diffraction[23]:

$$\frac{d\sigma}{dM^2}=\sigma_0\frac{1}{M^2} \tag{24}$$

valid for $M^2 \gtrsim 2$ GeV^2. The integrated charm cross section is given by

$$\sigma=\sigma_0^c\int_{M_0{}^2}^{M_1{}^2}\frac{dM^2}{M^2}=\sigma_0^c\ln\frac{(1-x_1)s}{(M_{\Lambda_c}+M_D)^2}, \tag{25}$$

where in this case $M_0{}^2$ is the threshold value for associated production of a pair of hadrons containing charmed quarks. The upper limit $M_1{}^2$ is determined from the kinematical relation $M_1{}^2=s(1-x_1)$, where x_1 is the lower fractional momentum cut on the recoiling proton. In the ISR $pp\to p_1\Lambda_c X$ experiment[7] one triggers on events with $x_1\geq 0.8$. If we assume that essentially all the charm cross section $\sigma_c\sim 300$ μb is due to diffractive production, then we can determine $\sigma_0=77$ μb. We thus predict that at SPS and Fermilab energies ($s\cong 400$–600 GeV^2), the total $pp\to$ charm cross section should be of the order of 150 μb. The beam dump experiment (model dependent),[24] which can be interpreted in terms of charm production, gives cross sections of up to 80 ± 40 μb. The pion-induced diffractive charm cross section measured at Fermilab[13] at $s=416$ GeV^2 is of order 10–40 μb. This could imply that the probability for intrinsic charm is smaller for the pion than for the proton.

VI. PRODUCTION OF b AND t QUARKS

We now estimate the associated production of b- and t-flavored hadrons from the intrinsic sea from Eq. (24). As discussed in the Introduction, we expect a $\sim 1/m_Q{}^2$ dependence for the suppression of intrinsic heavy quarks. The prediction for the diffractive b cross section at the ISR ($\sqrt{s}=63$ GeV^2) is

$$\sigma^b=\left(\frac{m_c}{m_b}\right)^2\sigma_0^c\int_{(2m_b)^2}^{0.2s}\frac{dM^2}{M^2}\approx 15\ \mu\text{b} \tag{26}$$

using the value for σ_0^c determined from charm production. The Λ_b spectrum should have the same shape as the Λ_c. This cross section should be compared with those predicted from perturbative mechanisms[16] for the central region: $\sigma^b\cong 0.1$ μb. If $m_t\gtrsim 20$ GeV, then t-quark-hadron production is kinematically suppressed at ISR energies. At Tevatron collider energies ($\sqrt{s}=2.10^3$ GeV) we expect for b production, with $M_1{}^2=0.25s$,

$$\sigma^b=73\ \mu\text{b},$$

and for t production (assuming $m_t=20$ GeV)

$$\sigma^t=\left(\frac{m_c}{m_t}\right)^2\sigma_0^c\int_{(2m_t)^2}^{0.2s}\frac{dM^2}{M^2}\approx 3\ \mu\text{b}. \tag{27}$$

The corresponding values from hard mechanisms[16] at these energies are

$$\sigma^b\approx 2\ \mu\text{b and } \sigma^t\approx 0.1\ \mu\text{b}.$$

Not only are the intrinsic cross sections larger than perturbative contributions, but also the combinatorial background is reduced in diffractive configurations. In addition, since the transverse momentum of heavy quarks in the intrinsic state generally increases with the quark mass (see Sec. II), the heavy-quark hadrons are expected to be produced at relatively large transverse momentum. We also note that concerning diffractive production of heavy quarks on nuclear targets one expects a $A^{2/3}$ dependence from the intrinsic charm model. This is in contrast to the perturbative hard-scattering cross section, which should be proportional to A.

VII. SAME-SIGN DIMUON PRODUCTION

The presence of intrinsic charm in the proton at the 1 and 2% level makes possible a number of new charged- and neutral-weak-current interactions. We will consider neutral-current phenomena in Sec. VIII. Perhaps the most interesting implication of intrinsic charm for νN and $\bar\nu N$ charged-current reactions is the production of b quarks ($\bar\nu c\to\mu^+b$ and $\nu\bar c\to\mu^-\bar b$). The subsequent leptonic decay of the b and $\bar b$ then leads to same-sign muon pairs (see Fig. 2). We will also briefly discuss t production via $b\to t$ charged currents.

In principle, charged-current b production can proceed via the process $\bar\nu u\to\mu^+b$, $\bar\nu c\to\mu^+b$, and

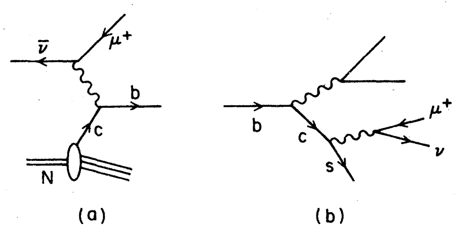

FIG. 18. (a) Production of b quark from the intrinsic charm component of the proton; (b) decay of b quark into μ^+.

$\bar{\nu}t \to \mu^+ b$. The t-quark intrinsic sea is negligible in this connection $[P(uudt\bar{t}) \sim (m_c^2/m_t^2)P(uudc\bar{c})$; see Eq. (2)]. For the first two reactions we need the u-b and c-b couplings. From a recent analysis[25] of the Kobayashi-Maskawa matrix elements[26] we learn that the $u \to b$ coupling is small (0.004 $\pm$ 0.004). For the $c \to b$ charged-current coupling there exist two solutions: $V_{cb} \cong 0.48\, G_F$ and 0.22 G_F, where G_F is the standard weak coupling constant. The differential cross section for $\bar{\nu}p \to \mu^+ b$ [see Fig. 18(a)] or $\nu p \to \mu^- \bar{b}$ is given by

$$\frac{d\sigma}{dx\,dy} = \frac{m_N E_\nu x}{4\pi}(2\sqrt{2}G_F)^2 |V_{cb}|^2 c(x)(1-y)^2 . \quad (28)$$

The integrated cross section $\sigma_{c\to b}(E_\nu)$ as a function of beam energy is obtained by using the kinetical restriction $2m_N\nu - Q^2 \geq 2m_N m_b + m_b^2$.[27] In Fig. 19 we show $\sigma_{c\to b}(E_\nu)$ for the two V_{cb} solutions. The subsequent decay of the produced b quark as in $b \to c \to s\mu^+\nu$ gives rise to same-sign muons [see Fig. 18(b)]. The resulting same-sign muon-

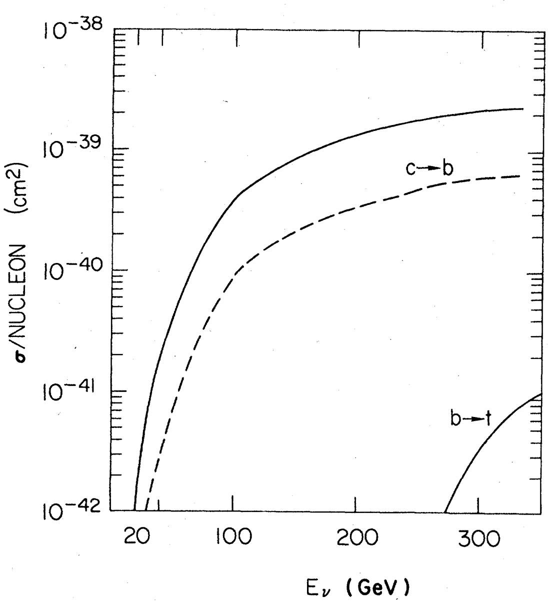

FIG. 19. Energy dependence of the production cross section of b quarks (solid and dashed lines corresponding to two V_{cb} solutions) and t quarks.

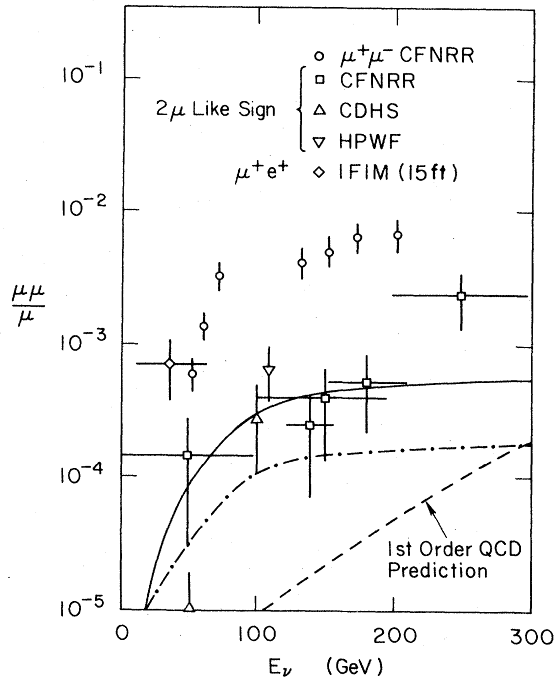

FIG. 20. The ratio of the same-sign dimuon events to the normal charged-current events. The solid and dashed-dotted lines are a theoretical prediction (ignoring experimental cuts) from the intrinsic charm component for $\mu^+\mu^+$ and $\mu^-\mu^-$, respectively, using $V_{cb} = 0.48$. Data from Ref. 30. Also shown (dashed line) is the QCD first-order prediction (Ref. 29).

production cross section is given by $\sigma_{c\to b}(E_\nu) \times B(c \to \mu)$, where the branching ratio $B(c \to \mu)$ is $\simeq 0.15$.[28] In Fig. 20 the ratio $\mu^\pm\mu^\pm/\mu$ is plotted for isoscalar targets and compared with the experimental results. In fact, the measured[21] ratio $(\mu^\pm\mu^\pm)/\mu^+\mu^-$) is as large as 0.1, an order of magnitude larger than what is predicted from first-order perturbative QCD associated charm production[29] (see Fig. 20). Intrinsic charm, however, is seen to give a reasonable explanation of the large same-sign dimuon rate. In addition, in the BEBC experiment[31] a correlation of p_{μ^+} and total hadronic energy is observed; this could perhaps be inter-

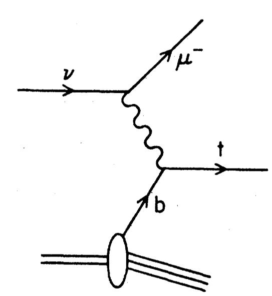

FIG. 21. Production of t quark from the intrinsic b component of hadron.

preted as the production of a bottom baryon. It should be remarked, however, that the approximate agreement in Fig. 20 for same-sign dimuons is not very sensitive to the x_{Bj} distribution of the intrinsic charm sea, but only depends on its magnitude. Further experiments which can discriminate the x distribution of $c(x)$ are necessary. The b- ($\bar{b}$-) quark production in the charged-current fragmentation region should be accompanied by associated charm $\bar{c}$ (c) production in the target fragmentation region.

Finally we consider the production of t quarks through the $b \to t$ mechanism of Fig. 21. In Fig. 19 the result for $\sigma_{b\to t}(E)$ is displayed assuming the mass suppression m_b^{-2}.

The occurrence of intrinsic states such as $|uuds\bar{s}\rangle$ and $|uudc\bar{c}\rangle$ in the proton also allows the possibility of producing rather exotic baryons in exclusive charged-current reactions. Two examples are shown in Fig. 22, in which s-c baryons and c-b baryons are produced. The process illustrated in Fig. 22(b) would not be expected for the standard perturbative sea since the formation of the final baryon requires b quarks with large x.

VIII. WRONG-SIGN MUON PRODUCTION

The intrinsic charm component in the proton can give abundant wrong-sign single-muon ν ($\bar{\nu}$) reactions, $\nu N \to \mu^+ X$ and $\bar{\nu} N \to \mu^- X$. There is a rather stringent experimental bound on this ratio from the CERN-Dortmund-Heidelberg-Saclay (CDHS) experiment[32]

$$R_{\text{expt}} = \frac{\sigma(\nu N \to \mu^+ X)}{\sigma(\nu N \to \mu^- X)} < 1.6 \times 10^{-4}. \tag{29}$$

We will show that the intrinsic-charm model gives a cross section consistent with this bound after taking into account an estimate of the experimental cuts.

Intrinsic charm contributes to the reaction $\nu N \to \mu^+ X$ through the neutral-current interaction depicted in Fig. 23. The cross section for $\nu N \to \nu c X$ is given by[33]

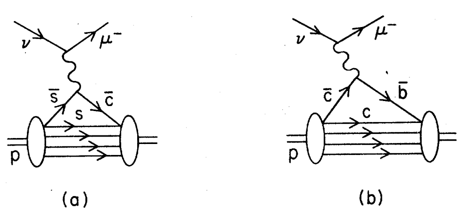

FIG. 22. Exclusive production of s-c baryons and c-b baryons.

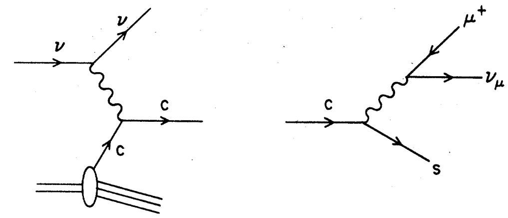

FIG. 23. Wrong-sign single-muon production from the intrinsic charm component.

$$\frac{d\sigma}{dx\,dy} = \frac{m_N E_\nu x}{4\pi}(2\sqrt{2}G_F)^2[(\tfrac{1}{2} - \tfrac{2}{3}\sin^2\theta_W)^2 + (\tfrac{2}{3}\sin^2\theta_W)^2(1-y)^2]c(x). \tag{30}$$

If there were no experimental bias, we could obtain its ratio to the normal charged-current cross section in Eq. (30) by simply integrating over y,

$$\left.\frac{(d\sigma/dx)(\nu N \to \nu c X)}{(d\sigma/dx)(\nu N \to \mu^- X)}\right|_{\text{uncorrected}} \simeq 0.135\,\frac{c(x)}{d(x)}, \tag{31}$$

where we used $\sin^2\theta_W \simeq 0.229$.[34]

An important constraint of the CDHS experiment comes from the cut in the total visible energy: $E_{\text{vis}} \equiv E_{\text{hadron}} + E_{\mu^+} > E_{\text{cut}}$. The probability for the neutral-current event to survive the experimental cut is less than

$$\int_{y_c}^{1} dy\,\frac{d\sigma}{dx\,dy}\Big/\int_0^1 dy\,\frac{d\sigma}{dx\,dy} < 1 - y_c, \tag{32}$$

where $y_c \equiv E_{\text{cut}}/E_\nu$. (We have overestimated this probability by neglecting the energy carried by the neutrino from the decay of the c quark.) The narrow band beam of the experiment has an approximately flat energy spectrum[35] extending up to E_0. Therefore, the correction factor due to the cut is on the average

$$\langle 1 - y_c\rangle = 1 - \frac{E_{\text{cut}}}{E_0 - E_{\text{cut}}}\ln\frac{E_0}{E_{\text{cut}}} \simeq 0.29, \tag{33}$$

where we used $E_0 \simeq 190$ GeV, $E_{\text{cut}} \simeq 100$ GeV. If we combine the factors (31), (32), and (33) with our estimate $\langle c(x)/d(x)\rangle \simeq 10^{-2}$ and the branching ratio $B(c \to \mu^+ X) \cong 0.15$, we obtain

$$R_{\text{expt}} = \frac{\sigma(\nu N \to \nu\mu^+ X)}{\sigma(\nu N \to \mu^- X)} \simeq 0.135\,\frac{c(x)}{d(x)}\,0.15 \times 0.29 \simeq 5.8 \times 10^{-5} \tag{34}$$

consistent with the experimental bound, Eq. (29).[36]

IX. HADRONIC J/ψ PRODUCTION

In addition to charmed mesons and baryons, the J/ψ may also be produced diffractively from the intrinsic charm component of the proton. Compared to the charm-production cross section at Fermilab energies[12]

$$\sigma(\pi N \to DX) \simeq 10\text{–}40\ \mu\text{b}\,, \tag{35}$$

J/ψ production data around 200 GeV give[37]

$$\sigma(\pi N \to \psi X) \simeq 100\ \text{nb}\,.$$

Further, the observed x_F distribution appears to be more strongly peaked near $x \simeq 0$ compared to what would be expected from the intrinsic charm distribution. Evidently most of the ψ production comes from other central production mechanisms such as gluon and $q\bar{q}$ fusion.[38] In order for the intrinsic charm model to be consistent, there must be a large suppression factor for the ψ production from the intrinsic charm compared to the D production:

$$\left.\frac{\sigma(\pi N \to \psi X)}{\sigma(\pi N \to DX)}\right|_{\text{intrinsic charm}} \lesssim 5\times 10^{-5}\,. \tag{36}$$

Our general picture is as follows: After the incoming proton is diffractively dissociated, the $|uudc\bar{c}\rangle$ system can form either open- or hidden-charm states with roughly the same energy dependence. The probability of ψ production then depends on color and flavor combinatorics and the "semilocal duality" probability of forming the color-singlet state in the correct mass range. In fact, there are a number of factors which act to suppress forward ψ production compared to open charm.

(1) In the decay of the $|uudc\bar{c}\rangle$ state, the probability that the $\bar{c}$ quark combines to form a $c\bar{c}$ system is about $\frac{1}{4}$ (*flavor suppression*). Similarly, the flavor-suppression factor for the $|u\bar{d}c\bar{c}\rangle$ state is about $\frac{1}{2}$.

(2) A $c\bar{c}$ system can be formed in either a color octet $c\bar{c}$ or singlet $c\bar{c}$ state. The color octet $c\bar{c}$ state should interact with other colored particles and is most likely to decay into open-charm particles such as D's. Therefore, we can take only the color-singlet combination of $c\bar{c}$ for ψ production. This occurs only $\frac{1}{9}$ of the time (*color suppression*).

(3) If the color-singlet $c\bar{c}$ system has a mass larger than the $D\bar{D}$ threshold, it will decay strongly into charmed particles rather than ψ production. Therefore, we have to require that the invariant mass $M_{c\bar{c}}$ is below the $D\bar{D}$ threshold (*mass suppression*):

$$2m_c < M_{c\bar{c}} < 2m_D\,. \tag{37}$$

If we neglect transverse momenta, the simplified model of the intrinsic charm state $|uudc\bar{c}\rangle$ gives a mass spectrum [see Eq. (14)]

$$\begin{aligned} \frac{dP}{dM_{c\bar{c}}^2} &= \frac{N}{m_c^2}\left(\frac{m_c^2}{M_{c\bar{c}}^2}\right)^4 \times \int_{4m_c^2/M_{c\bar{c}}^2}^{1} dx\,(1-x)^2 \frac{\sqrt{x}}{[x-(4m_c^2/M_{c\bar{c}}^2)]^{1/2}}\,, \\ M_{c\bar{c}}^2 &= \frac{m_c^2}{x_c} + \frac{m_c^2}{x_{\bar{c}}}\,, \\ x &= x_c + x_{\bar{c}}\,, \\ N &= 3600\ P_5\,, \end{aligned} \tag{38}$$

which is shown in Fig. 24. The approximate mass spectrum near threshold is given by

$$\frac{dP}{dM_{c\bar{c}}^2} \simeq \frac{15}{32}\frac{(M_{c\bar{c}}^2 - 4m_c^2)^{5/2}}{m_c^7}\,. \tag{39}$$

In addition, the transverse momenta of c quarks suppress the mass spectrum near threshold even more strongly. We shall simulate this effect by effectively increasing m_c ($m_c^2 \to \hat{m}_c^2 = m_c^2 + \langle k_\perp^2\rangle$).[39] If we take $\hat{m}_c \simeq 1.7$ GeV for instance, we obtain

$$\int_{4\hat{m}_c^2}^{4m_D^2} dM_{c\bar{c}}^2 \frac{dP}{dM_{c\bar{c}}^2} \simeq 10^{-2}\,. \tag{40}$$

(4) Even if the $c\bar{c}$ system is below $D\bar{D}$ threshold, it may be realized as χ, η_c and ψ' states which do not decay into ψ's. We estimate this suppression factor as $\frac{1}{3}$ (*channel suppression*). If we combine the factors in (1)–(4) we obtain the very rough theoretical estimate

$$\left.\frac{\sigma(\pi N \to \psi X)}{\sigma(\pi N \to DX)}\right|_{\text{intrinsic charm}} \simeq 5\times 10^{-5}\,. \tag{41}$$

Despite these uncertainties, it is clear that although the intrinsic charm model does predict ψ production in the forward fragmentation region, the rate is at a very suppressed level.

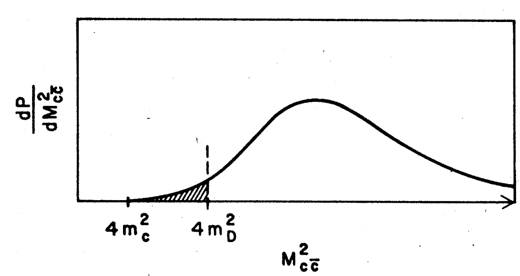

FIG. 24. The $c\bar{c}$ mass spectrum in the intrinsic charm state $|uudc\bar{c}\rangle$. The shaded area corresponds to the χ, η_c, ψ, and ψ' production.

X. CONCLUSION

The postulate that ordinary hadrons contain intrinsic charm states[17] at the 1% level can explain two sets of unexpected experimental results: (1) the copious diffractive production of charmed hadrons at large longitudinal momentum in high-energy proton-nucleon and pion-nucleon collisions, and (2) the anomalously large number of same-sign dimuon events observed in deep-inelastic neutrino reactions. Clearly, much more experimental work is necessary to verify and test the model and its predictions. Most important, it will be crucial to verify that the charm-quark distribution $c(x)$ in nucleons is appreciable at large $x_{\rm Bj}$ in deep-inelastic lepton scattering.

The concept that the wave function of an ordinary hadron has finite mixing with virtual bound states containing heavy-quark pairs is a new dynamical concept and is somewhat at variance with standard parton-model ideas. Nevertheless, as we have discussed in Sec. I, there is no reason in the context of QCD why this probability should be zero. In fact, bag models indicate that the mean number of charm quarks in the proton is of order 1%.[5] More generally it may be possible that in lattice calculations or other nonperturbative approaches one may be able to calculate such probabilities (e.g., the mixing of $c\bar{c}$ states in a glueball) because of the heavy-quark mass parameters.

We emphasize that although the virtual heavy-quark states in a pion or proton are far off-shell, all of the constituents are bound within the hadron because of color forces. This leads inevitably to the prediction that the heavy quarks carry the largest fraction x of the hadronic momentum in the infinite-momentum frame for the heavy-quark Fock states. This is in striking contrast to the x distribution of heavy quarks created in a high-momentum-transfer collision by perturbative mechanisms.

The suppression mechanism (described in Sec. IX) for hadronic ψ production at large x_F presumably also has consequences for the decay of the B meson $(b\bar{q})$ into ψX where the b quark decays into $c\bar{c}s$. Taking into account all possible decay modes and kinematical suppression one concludes in the free-quark model $\Gamma(b\to c\bar{c}X)/\Gamma(b\to \text{all})\sim 14-20\%$.[40] According to the discussion in Sec. IX we expect a strong suppression $(\sim 10^{-3})$ when instead considering $\Gamma(B\to\psi X)$. Generally speaking, the decay of hadrons into hidden-heavy-quark states is expected to provide a powerful probe into nonperturbative hadron dynamics.

We have argued on general grounds that the probability of a hadron to contain an intrinsic heavy-quark pair should fall as

$$P_{Q\bar{Q}}\propto 1/R^2 m_Q^{\,2}\,, \tag{42}$$

where m_Q is the heavy-quark mass and R is a hadron size parameter. If this very slow falloff in m_Q is correct, there will be copious diffractive large-x_L production of heavy-quark (b,t) systems in very-high-energy hadron collisions. Because of the large cross section and the large luminosity possible in hadron-hadron collisions, diffractive production of heavy quarks could become an important method for the discovery and measurement of heavy-quark systems.

The presence of intrinsic charm in nucleons makes possible a new array of charged- and neutral-weak-current interactions, the most dramatic one being the charged-current interaction which changes the intrinsic charm quarks into bottom quarks at large x_{Bj}. In particular we predict large neutrino-induced cross sections for b-quark production in the current fragmentation region with $\bar{c}$-quark production in the target fragmentation region.

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[1] The states are defined at equal $\tau=t+z$ in the light-cone gauge $A^+=A^0+A^3=0$.
[2] G. P. Lepage and S. J. Brodsky, Phys. Rev. D 22, 2157 (1980); S. J. Brodsky, Y. Frishman, G. P. Lepage, and C. Sachrajda, Phys. Lett. 91B, 239 (1980), and references therein.
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