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Production of particles containing charmed quarks in hadronic collisions*

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The ability to obtain the correct cross section for production of the new 3.1-GeV resonance at Brookhaven National Laboratory using the natural analog of the Drell-Yan annihilation mechanism leads us to explore further implications of this approach. We present in detail the longitudinal- and transverse-momentum dependence of the cross section, as well as its dependence upon beam energy. We point out correlations among final-state multiplicity, longitudinal momentum of the resonance, and resonance mass peculiar to the parton-model annihilation mechanism. We also examine the alterations which are expected upon employing a $\bar{p}$ beam. Finally, we obtain cross sections for charmed meson and baryon production. $\bar{D}^{0*}$ and D^{-*} mesons should have particularly large production cross sections.

INTRODUCTION

Since the introduction of the parton model as an explanation of the scaling observed in deep-inelastic scattering, all confrontations of the model and its extensions with experiment have met with a remarkable degree of success. Thus, parton-model predictions for and interpretations of the new resonances observed at SLAC, Brookhaven, and elsewhere[1] are of particular interest. Especially interesting is the possibility that the production of the J(3.1 GeV) at Brookhaven in pN collisions proceeds via a Drell-Yan annihilation type of process[2] (Fig. 1). Because of the many constraints upon the distribution of quarks within nucleons obtained[3-5] from earlier theoretical and experimental work, quite specific predictions for this annihilation mechanism are possible.[6]

In an earlier paper[7] (paper I) we demonstrated that the magnitude of the cross section for J(3.1) production at Brookhaven is quite consistent with its interpretation[8] as a $\mathcal{P}'$-$\bar{\mathcal{P}}'$ resonance ($\mathcal{P}'$ is the fourth "charmed" quark), produced via $\mathcal{P}'\bar{\mathcal{P}}'$ annihilation, provided (a) that the portion of the quark distribution functions associated with Pomeron-like behavior is approximately SU(4)-symmetric (see footnote 10 of paper I for a bit of theoretical justification), so that there are as many $\mathcal{P}'$ ($\bar{\mathcal{P}}'$) quarks carrying a given fraction x of a nucleon's momentum as there are $\bar{\mathcal{P}}$, $\bar{\mathfrak{N}}$, λ, and $\bar{\lambda}$, and (b) that it is not necessary to produce charmed hadrons in association with J in the final state. In terms of the annihilation process this is equivalent to the assumption that the "cores" left behind by the annihilating quarks in Fig. 1 are capable of communicating in such a way that both their quark-like quantum numbers and their "charm" are neutralized. This latter assumption is important at Brookhaven energies (e.g., $p_{\rm lab} = 28.5$ GeV/c for the initial experiment) at which the threshold suppression associated with the introduction of large final-state masses would be dramatic. It was also shown in paper I that the annihilation mechanism predicts a purely kinematic reduction of the production cross section by a factor of 20 in going from the 3.1-GeV resonance to the 3.7-GeV resonance. This coupled with the smaller branching ratio of J(3.7) to the e^+-e^- channel [perhaps $\frac{1}{5}$ that of J(3.1)] is consistent with the failure[9] of Ting's group to see the higher-massed J at the 1% level.

Given this "success" it is clearly of importance to pursue this approach. In this paper we will give more details on the mass and energy dependence of the annihilation cross section and on the distribution of the J(3.1) in longitudinal and transverse momentum. These latter distributions turn out to provide distinctive signatures capable of testing the specific quark distribution function forms derived in Refs. 3 and 4 using scaling laws[10] developed for high-transverse-momenta phenomenology. In addition, interesting correlations between final-state multiplicity and the longitudinal momentum of the resonance are discussed.

Thirdly, we discuss expectations for J production in $p\bar{p}$ collisions via the annihilation mechanism. Even at $p_{\rm lab} = 9$ GeV/c the cross sections are substantial due to the ability to use the small $\mathcal{P}\bar{\mathcal{P}}$ and $\mathfrak{N}\bar{\mathfrak{N}}$ admixture in the J wave function, which probes valence quark components of the p and $\bar{p}$

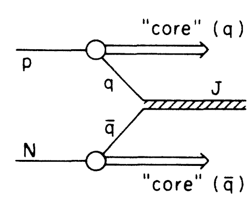

FIG. 1. Drell-Yan annihilation-type diagram for meson production.

wave functions simultaneously. (Also important is the smaller threshold final-state mass possible in $p\bar{p}$ collisions.) This contrasts with the case of $J\,(3.1)$ production in pN collisions where the contribution from the normal quark admixture in J is small compared to that of the $\mathcal{P}'\bar{\mathcal{P}}'$ channel.

We will also give expectations for charmed (F, D, etc.)[11] meson production in pN collisions. The expected cross sections can be quite large as in the case of two of the D-type mesons containing a $\mathcal{P}$ or $\mathcal{N}$ quark which can come from the valence component of the nucleonic quark wave function. An estimate of charmed baryon production is presented.

REVIEW AND PRELIMINARIES

First let us briefly review the results of paper I and Ref. 6. The integrated cross section for J production in pN collisions is (asymptotically)

$$\int \sigma^{\text{resonance}}_{p+N\to J+X}(Q^2)\,dQ^2 = 4\pi^2 \sum g_{Ji\bar{i}}^{\,2}/4\pi \int [f_i^p(x_1)f_{\bar{i}}^N(x_2) + f_{\bar{i}}^p(x_1)f_i^N(x_2)]\,\delta(x_1x_2s - m_J^2)\,dx_1dx_2\,. \tag{1}$$

[The index i runs over all quark types ($\mathcal{P}, \mathcal{N}, \lambda, \mathcal{P}'$), and the f_i represent the probabilities of finding a quark with a fraction x of the proton (or neutron) linear momentum P as P becomes large.] We ignore off-shell dependence of $g_{Ji\bar{i}}^{\,2}/4\pi$. Later we will return to justify this assumption. For the $J\,(3.1)$ particle (for all our other estimates involving direct coupling we make analogous assumptions) we take

$$g_{J\mathcal{P}'\bar{\mathcal{P}}'}^{\,2}/4\pi \approx 1, \tag{2}$$

corresponding to the fact that it is primarily composed of $\mathcal{P}'$ and $\bar{\mathcal{P}}'$, and

$$\frac{g_{J\mathcal{P}\bar{\mathcal{P}}}^{\,2}}{4\pi} = \frac{g_{J\mathcal{N}\bar{\mathcal{N}}}^{\,2}}{4\pi} = \tfrac{1}{2}R_J\,\frac{g_{Je^+e^-}^{\,2}}{4\pi}\,, \tag{3}$$

where R_J is the rate of J to hadrons divided by that of J to e^+e^-. $g_{Je^+e^-}$ is the coupling of J to the e^+e^- channel. To facilitate comparison with paper I we use

$$g_{Je^+e^-}^{\,2}/4\pi = 2.4\times10^{-6}\,,\qquad R_J = 25 \tag{4}$$

corresponding to a total width Γ_J for J of 65 keV. It is now clear that Γ_J is larger than this so that those few results sensitive to the normal quark ($\mathcal{P}\bar{\mathcal{P}}$ and $\mathcal{N}\bar{\mathcal{N}}$) admixture in J should be scaled accordingly once Γ_J is well determined.

Actually there is also evidence that the estimate (2) may be too small. In particular let us estimate $g_{J\mathcal{P}'\bar{\mathcal{P}}'}$ for the situation when both the quarks are on shell. To do so we examine the J's decay into e^+e^- using the ideas of vector-meson dominance. Briefly, if we write

$$|J\rangle = g_{J\mathcal{P}'\bar{\mathcal{P}}'}|\mathcal{P}'\bar{\mathcal{P}}'\rangle \tag{5}$$

and assume that all other vector mesons couple weakly to $\mathcal{P}'\bar{\mathcal{P}}'$, then the $\mathcal{P}'$ charge, $2/3e$, must be given in the usual notation [see Fig. 2(a)] by

$$\frac{em_J^{\,2}}{g_J}\,\frac{1}{m_J^{\,2}}\,g_{J\mathcal{P}\bar{\mathcal{P}}} = \frac{2}{3}e. \tag{6}$$

As mentioned, we have an approximate value for $g_{Je^+e^-}$ which in this model is given by [Fig. 2(b)]

$$\frac{1}{g_J}e^2 = g_{Je^+e^-}\,. \tag{7}$$

Using (4) we obtain

$$g_J = 16.7, \tag{8}$$

which implies

$$g_{J\mathcal{P}'\bar{\mathcal{P}}'}^{\,2}/4\pi = 9.8. \tag{9}$$

A larger e^+e^- width for J will reduce this value. However, there is one reason to be skeptical of this approach. Vector-meson dominance, at least in this naive form, does not appear to work well for J photoproduction unless the Jp total cross section is substantially smaller than the ϕp.[12] On the other hand, a similar approach in the case of the ρ works very well in the following sense. Experimentally from ρ decay into e^+e^- we know

$$g_\rho^{\,2}/4\pi \approx 2. \tag{10}$$

Obtaining the correct charge for the $\mathcal{P}$ and $\mathcal{N}$ quarks requires

$$\frac{g_\rho}{\sqrt{2}} = \sqrt{2}\,g_{\rho\mathcal{P}\bar{\mathcal{P}}} = \sqrt{2}\,g_{\rho\mathcal{N}\bar{\mathcal{N}}}, \tag{11}$$

where we have written

$$|\rho\rangle = g_{\rho\mathcal{P}\bar{\mathcal{P}}}(|\mathcal{P}\bar{\mathcal{P}}\rangle - |\mathcal{N}\bar{\mathcal{N}}\rangle). \tag{12}$$

This yields

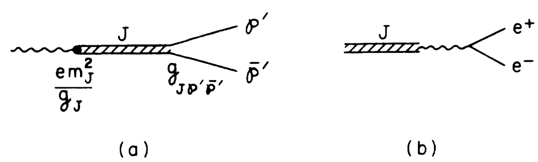

FIG. 2. (a) Vector-meson dominance calculation of the $\mathcal{P}'$ quark charge. (b) Vector-meson dominance calculation of the $J\to e^+e^-$ width.

$$\frac{g_{\rho\mathcal{P}\bar{\mathcal{P}}}^2}{4\pi} = \frac{g_{\rho\mathcal{N}\bar{\mathcal{N}}}^2}{4\pi} \approx \frac{g_\rho^2}{16\pi} \approx \tfrac{1}{2} \tag{13}$$

The approximate total width of the ρ should then, in a parton-model–like approximation, be

$$\Gamma_\rho = \frac{m_\rho}{3}\left(\frac{g_{\rho\mathcal{P}\bar{\mathcal{P}}}^2}{4\pi} + \frac{g_{\rho\mathcal{N}\bar{\mathcal{N}}}^2}{4\pi}\right) \approx 200 \text{ MeV}, \tag{14}$$

a bit large but not bad.

Thus the cross-section predictions should perhaps be somewhat larger than the values we quote on the basis of (2). Of course, we have ignored the $J(3.7)$ in this treatment. For a reasonable e^+e^- partial width, $1/g_{J(3.7)}$ must be substantial, say, of the size of $1/g_{J(3.1)}$. Since the total $\mathcal{P}'$ charge is fixed, if the estimate (9) is to apply to the $J(3.1)$, $g_{J(3.7)\mathcal{P}'\bar{\mathcal{P}}'}$ must be much smaller. Currently popular dynamical models would not predict so great a difference between $g_{J(3.1)\mathcal{P}'\bar{\mathcal{P}}'}$ and $g_{J(3.7)\mathcal{P}'\bar{\mathcal{P}}'}$, so in all likelihood these couplings will be comparable, with the result that (9) should be reduced by a factor of roughly 4 (for, say, e^+e^- widths in the ratio of the resonance masses) when applied to either the 3.1- or the 3.7-GeV resonance separately.

The quark distribution functions f we employ are those of Ref. 3 with the exception that the "sea" or Pomeron components, s, used are taken to be SU(4)-symmetric. Thus

$$s_{\mathcal{P}} = f_{\bar{\mathcal{P}}} = s_{\mathcal{N}} = f_{\bar{\mathcal{N}}} = f_\lambda = f_{\bar{\lambda}} = f_{\mathcal{P}'} = f_{\bar{\mathcal{P}}'} = \tfrac{12}{20}\, 2\, \frac{(1-x)^7}{x}. \tag{15}$$

Note that the sea component is the only contribution to all except the $\mathcal{P}$ and $\mathcal{N}$ quark distribution functions. The factor of $\frac{12}{20}$ relative to the SU(3)-symmetric sea case guarantees that the total sea contribution to deep-inelastic scattering (to which the $\mathcal{P}'$ and $\bar{\mathcal{P}}'$ quarks are assumed to contribute) is consistent with the sum rules and theoretical constraints of Ref. 3 which fix its over-all normalization. This reduction in each individual quark's sea component leads to some reduction in the e^+e^- background (due to the original off-shell photon Drell-Yan process) and in the "normal" quark contribution to J production, relative to that quoted in paper I.

As discussed in paper I and Ref. 6 the asymptotic formula (1) cannot be used at Brookhaven energies. It is necessary to incorporate threshold effects. At $p_{\text{lab}} = 28.5$ GeV the maximum allowed resonance mass is

$$\sqrt{s} - 2m_{\text{proton}} \approx 5.64 \text{ GeV}$$

since two baryons must appear in the final state. As the J mass is not too much smaller than this the threshold effects will be substantial. The technique for incorporating thresholds was developed in Ref. 6 and discussed in paper I. However, because of its importance to many of the calculations to be presented, we briefly review it in an appendix. We only note here that all corrections of order $m_{\text{proton}}/\sqrt{s}$ (or more generally $m_{\text{core}}/\sqrt{s}$, the core being the residue left behind after extraction of a quark from a proton) are treated exactly, while some approximations are made at the m^2/s level. (Terms in $M_{\text{resonance}}^2/s$ are, of course, treated exactly.) The maximum value of τ, where

$$\tau = M_{\text{resonance}}^2/s, \tag{16}$$

is thus

$$\tau_{\max} \approx \frac{s - 4m_{\text{core}}\sqrt{s}}{s} \approx \frac{(\sqrt{s} - 2m_{\text{core}})^2}{s}. \tag{17}$$

RESULTS AND DISCUSSIONS

We now turn to the calculations of interest. Let us begin by discussing very briefly the results of Fig. 3. Figure 3 gives the energy dependence of the production cross sections for the 3.1- and 3.7-

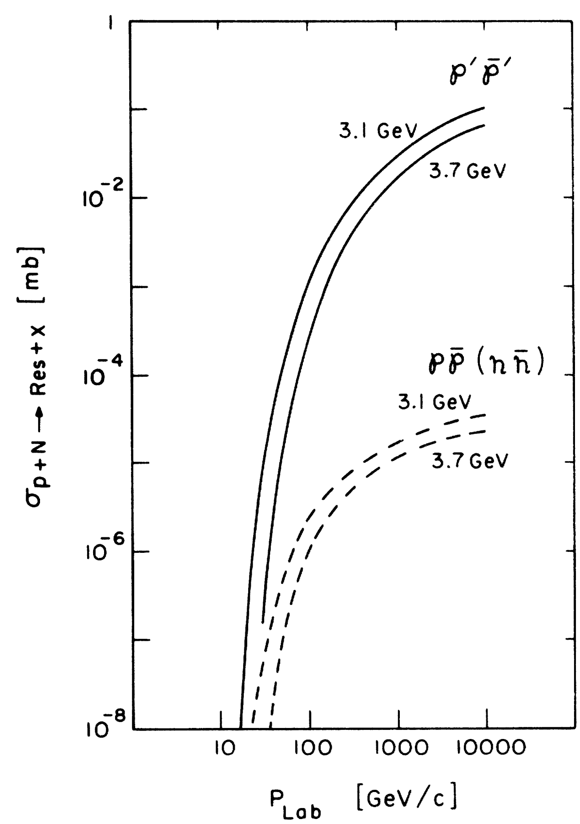

FIG. 3. Energy dependence of the $\mathcal{P}'\bar{\mathcal{P}}'$ and $\mathcal{P}\bar{\mathcal{P}}$ $(\mathcal{N}\bar{\mathcal{N}})$ annihilation mechanisms for the production of $J(3.1)$ and $J(3.7)$. Cross sections do not include any final-state branching ratio. They assume $g_{J\mathcal{P}'\bar{\mathcal{P}}'}^2/4\pi = 1$ and $g_{J\mathcal{P}\bar{\mathcal{P}}}^2/4\pi$ and $g_{J\mathcal{N}\bar{\mathcal{N}}}^2/4\pi$ as given in the text.

GeV resonances, using the coupling constants for the various quarks given earlier. We ignore Fermi motion effects in the nucleus which are possibly significant at Brookhaven energies. From Fig. 3 it is immediately apparent that the $\mathcal{P}'\bar{\mathcal{P}}'$ formation mechanism has far stronger energy dependence than that due to normal quark annihilation; this is expected from the general threshold rules originally developed in their most complete form for application to high-p_T phenomenology.[13] According to these rules the threshold damping is roughly of the form

$$\int \sigma^{\text{resonance}}\, dQ^2 \propto (1-\tau/\tau_{\max})^{2n_s-1}, \tag{18}$$

where n_s is the number of spectator quarks not participating directly in the quark-antiquark resonance formation vertex. Thus n_s is the total number of quarks in the cores of Fig. 1. For the $\mathcal{P}'\bar{\mathcal{P}}'$ mode both cores contain a minimum of 4 quarks, while for the $\mathcal{P}\bar{\mathcal{P}}$ ($\mathcal{N}\bar{\mathcal{N}}$) mode one core can contain as few as 2 quarks (corresponding to the valence quark state of the proton.) Thus

$$\begin{aligned} 2n_s-1 &= 15 \quad \mathcal{P}'\bar{\mathcal{P}}', \\ 2n_s-1 &= 11 \quad \mathcal{P}\bar{\mathcal{P}}\ (\mathcal{N}\bar{\mathcal{N}}). \end{aligned} \tag{19}$$

The rough form (18) gives approximately the correct shape in the regions of Fig. 3 near threshold. This is so despite the fact that we found it necessary to include subasymptotic corrections, related to the core masses, in the calculations relevant to Fig. 3. The proton's wave-function dependence, discussed in the Appendix, upon the transverse momentum of one of the quarks is precisely that which guarantees a smooth asymptotic-subasymptotic connection.

At this point it should be noted that, analogously to the situation in high-p_T phenomenology, there exist production mechanisms other than the annihilation process. One of many possibilities is that sketched in Fig. 4. This process obeys the interchange theory "requirement" that only wave-function vertices appear. (Processes with explicit vector gluons, for instance, always seem to be absent in high-p_T phenomenology.) Such processes are characterized by fewer spectators (5 in Fig. 4) and hence less threshold damping. However, they involve more "active" quarks participating in the formation of heavy particles and, of course, a heavier final state mass—a charmed baryon and meson in addition to the J are required in the final state of Fig. 4. Thus one generally requires energies above those at Brookhaven before such processes can even take place, in addition to which the increased number of active quarks is likely to lead to some off-shell damping effects in the process by which the quark and $\mathcal{P}\mathcal{N}$ pair produce a J plus a charmed baryon. As the energy increases at fixed resonance mass, the direct annihilation mechanism rises very rapidly and will, if dominant at low energies, continue to dominate away from threshold. Verification of this rapid energy dependence is a crucial test for the present approach. A more detailed investigation of such alternative mechanisms is being performed by Blankenbecler and collaborators.

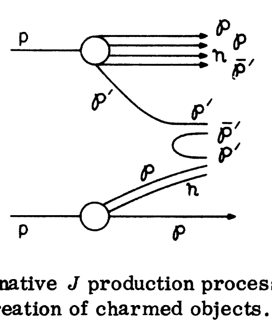

FIG. 4. Alternative J production process involving simultaneous creation of charmed objects.

We now turn to a comparison of expectations for J-like resonance production in pN collisions with those for $p\bar{p}$ collisions. The quark distribution functions for the $\bar{p}$ are, of course, simply obtained by charge conjugation from those of the p. For pN collisions the core masses are both taken equal to the proton mass. In contrast, because the cores need no longer have baryonlike quantum numbers in the $p\bar{p}$ case (the communication mechanism between cores presumably allows annihilation to pions and other light particles), we take $m_{\text{core}} = 0.1 m_{\text{proton}}$. Final-state branching ratios are not included.

In Fig. 5(a) we present total resonance production cross sections for pN collisions at 28.5 GeV/c as a function of resonance mass. The two types of curves correspond to (a) the $\mathcal{P}'\bar{\mathcal{P}}'$ annihilation channel and (b) the $\mathcal{P}\bar{\mathcal{P}}$ and $\mathcal{N}\bar{\mathcal{N}}$ "normal" quark annihilation mode. For (b), the normal quark coupling constants are held fixed at the values given earlier, Eq. (3), appropriate to J (3.1 GeV).

Note that the cross section for both cases is zero or negligible over the mass range studied for $p_{\text{lab}} = 9$ GeV/c, and note that for $p_{\text{lab}} = 28.5$ GeV/c the $\mathcal{P}'\bar{\mathcal{P}}'$ mode dominates the $\mathcal{P}\bar{\mathcal{P}}$ ($\mathcal{N}\bar{\mathcal{N}}$) mode. This suppression, due to the small admixture of normal quarks in J, occurs in spite of the ability to take the $\mathcal{P}$ (or $\mathcal{N}$) from the (weakly, $x \to 1$, threshold damped) valence component of a nucleon wave function, in comparison to the $\mathcal{P}'\bar{\mathcal{P}}'$ mode in which both quarks come from (strongly threshold damped) sea components.

J-like resonance production in $p\bar{p}$ collisions [Fig. 5(b)] is different in this respect. From Fig.

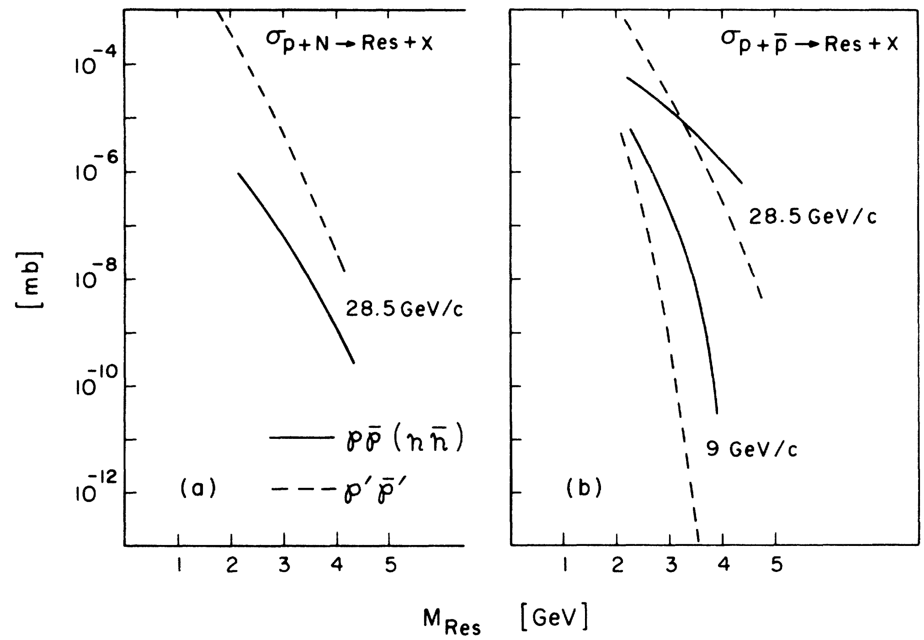

FIG. 5. Comparison of (a) J production in pp collisions and (b) J production in $p\bar{p}$ collisions via the two annihilation mechanisms. For core masses see text.

5(b) we see that at $p_{lab} = 28.5$ the normal quark mode is comparable to the $\mathcal{P}'\bar{\mathcal{P}}'$ mode. This merely reflects the ability to take the $\mathcal{P}$($\mathcal{N}$) quark from the proton valence component and the $\bar{\mathcal{P}}$($\bar{\mathcal{N}}$) quark from the antiproton valence component. The general magnitude of even the $\mathcal{P}'\bar{\mathcal{P}}'$ production mode is larger (by approximately a factor of 4 at 3.1-GeV resonance mass) than for pN collisions because of the smaller core masses.

Unfortunately a more realistic energy for a $\bar{p}$ beam at Brookhaven is $p_{lab} = 9$ GeV/c. Throughout most of the mass range the $\mathcal{P}'\bar{\mathcal{P}}'$ component is negligible at this energy. However, the total production cross section due to the normal quark mode is not insubstantial. For instance, at a resonance mass of 3.1 GeV the cross section is about $\frac{1}{30}$ that at $p_{lab} = 28.5$ GeV/c in pp collisions. Should it be possible to observe the annihilation cross section at the lower energies in $p\bar{p}$ collisions, for which the normal quark annihilation mode dominates, its dependence upon center-of-mass longitudinal momentum should be decidedly weaker than that of the J cross section in pN collisions (arising from the $\mathcal{P}'\bar{\mathcal{P}}'$ mode). Further discussion of this dependence will appear shortly.

Before turning to this topic, note that the above results indicate that there may be little cross section to gain from using either pion or $\bar{p}$ beams at Fermilab energies. The reason is that the $\mathcal{P}'\bar{\mathcal{P}}'$ mode dominates the $\mathcal{P}\bar{\mathcal{P}}$ ($\mathcal{N}\bar{\mathcal{N}}$) mode above 28.5 GeV/$c$ even for the $p\bar{p}$ case most favorable to the latter. (At 28.5 GeV/c the two are comparable but the $\mathcal{P}'\bar{\mathcal{P}}'$ rises far more rapidly with energy.) Since charmed quarks are not contained in p, $\bar{p}$ or π valence wave-function components, one may as well use the high-intensity p beams.

Let us now turn to the dependence of the resonance cross section upon longitudinal and transverse momentum in the center of mass. We define a variable y (essentially the usual Feynman longitudinal fraction) such that

$$y = p_L^{\text{resonance}} / \tfrac{1}{2}\sqrt{s} \,. \tag{20}$$

Here $p_L^{\text{resonance}}$ is the momentum of the resonance in the initial beam direction; p_T will denote the magnitude of the resonance's momentum transverse to the beam direction. In order to make a comparison with the background in the e^+e^- decay channel we include the e^+e^- branching ratio ($\frac{1}{27}$) in our calculation of the cross sections we now present. We give curves for the J (3.1) at $p_{lab} = 28.5$ GeV/c.

Figure 6(a) gives the p_T distributions

$$\frac{d\sigma^{\text{resonance}}_{p+N\to J+X\to e^+e^-+X}}{dp_T^{\,2}} \quad\text{and}\quad \frac{d^2\sigma^{\text{background}}_{p+N\to ``\gamma"+X\to e^+e^-+X}}{dp_T^{\,2}\,dQ^2}\,,$$

where the resonance cross section is, as before, integrated over its width and the background cross section refers to that from the original off-shell photon Drell-Yan process. (Ref. 6 discusses these calculations of the background.) Because higher powers of transverse momentum are associated with the sea component of a nucleon's wave function

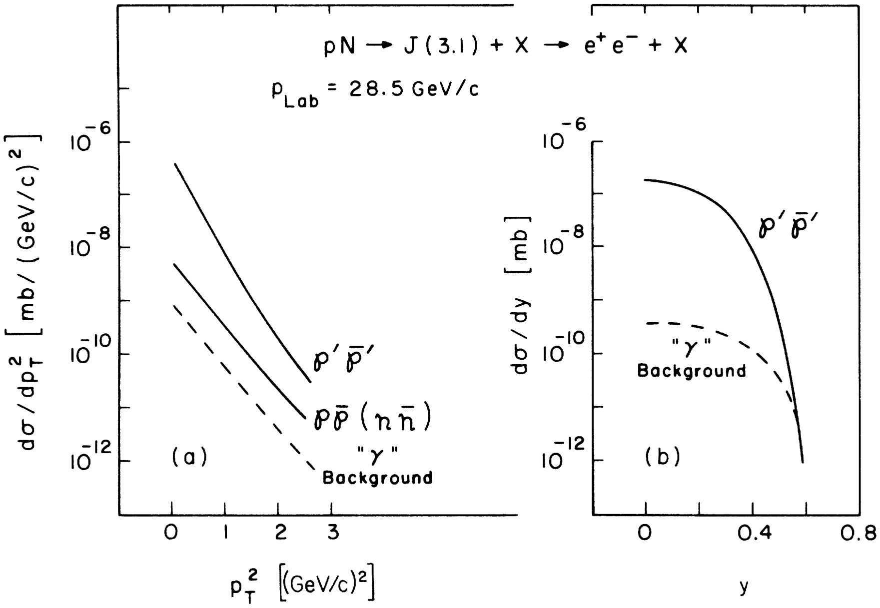

FIG. 6. p_T and y dependence of the resonance production cross section, computed for the $\mathcal{P}'\bar{\mathcal{P}}'$ and $\mathcal{P}\bar{\mathcal{P}}(\mathcal{N}\bar{\mathcal{N}})$ modes, and of the off-shell photon background. The final-state branching ratio to the e^+e^- channel is included. The differential cross sections integrate to 1.2×10^{-7} mb, 2.1×10^{-9} mb, and 0.32×10^{-9} mb/GeV^2 respectively for the $\mathcal{P}'\bar{\mathcal{P}}'$, $\mathcal{P}\bar{\mathcal{P}}(\mathcal{N}\bar{\mathcal{N}})$ and "γ" background modes.

than with the valence component (see the Appendix), the $\mathcal{P}'\bar{\mathcal{P}}'$ mode cross section has steeper p_T^2 dependence than either the $\mathcal{P}\bar{\mathcal{P}}$ $(\mathcal{N}\bar{\mathcal{N}})$ mode or the off-shell photon background. The latter two have very similar dependence upon p_T^2 since the quark annihilations involved are quite similar. The $\mathcal{P}'\bar{\mathcal{P}}'$ cross section behaves roughly as $e^{-4.6p_T^2}$ while the latter two are approximately described by $e^{-2.7p_T^2}$.

It is convenient at this point to ask ourselves what happens if we include the effects of the J wave function. Since the annihilating quarks are off-shell by varying degrees depending upon their longitudinal and transverse momenta, some variation of the $g_{J\mathcal{P}'\bar{\mathcal{P}}'}$ coupling is possible. We argue, however, that this variation is small and does not, in any case, necessarily lead to any suppression of the coupling when the quarks are further off-shell than when they are not. The reasoning is as follows. First note that the off-shell quark propagators in Fig. 1 have, by definition, already been incorporated into the quark distribution function forms. [Part of the $(1-x)$ and $k_\perp$ damping arises from these propagators.] Secondly we appeal to the ideas[10,13] developed for high-p_T phenomenology in examining the amputated $\mathcal{P}'\bar{\mathcal{P}}'$ J vertex. According to these ideas one should employ a model of the J in which the $\mathcal{P}'$ and $\bar{\mathcal{P}}'$ interact (assuming they are spin $\frac{1}{2}$) via vector-gluon exchange, so that the characteristic coupling constant is dimensionless. When examining the form factor of the J (a slightly different kinematical situation than the present one), if we assume that no anomalous dimensions arise as a result of the interaction and that the J wave function is finite at the origin, then the J will have a monopole form factor; the single power falloff arises entirely from one off-shell quark propagator while the associated amputated wave function exhibits no falloff as this quark becomes far off-shell. This type of assumption leads to a very consistent phenomenology of high-p_T interactions. Thus the amputed wave function of the J does not exhibit off-shell damping asymptotically and is thus unlikely to exhibit substantial off-shell dependence even in the nonasymptotic region.

Turning now to the longitudinal momentum or y dependence of the $J(3.1)$ cross section, we present in Fig. 6(b) the results for the $\mathcal{P}'\bar{\mathcal{P}}'$ mode and for the off-shell photon background. [As for the p_T^2 distributions, the $\mathcal{P}\bar{\mathcal{P}}$ $(\mathcal{N}\bar{\mathcal{N}})$ resonance production mode has very similar y dependence to that of the background.] As expected from the stronger threshold damping of the sea quark distributions, the $\mathcal{P}'\bar{\mathcal{P}}'$ mode is much more sharply cut-off as y approaches its kinematical limit, given to order $m_{\text{core}}/\sqrt{s}$ by

$$y_{\max} \approx 1 - \tau - \frac{2m_{\text{core}}}{\sqrt{s}}(1+\tau)^{1/2}. \qquad (21)$$

This limit is somewhat more restrictive than that of the most general production process

$$y_{max} = [(s - M_{resonance}^2 - 4m_{proton}^2)^2 - 16M_{resonance}^2 m_{proton}^2]^{1/2}/s. \tag{22}$$

If the $\mathcal{P}'\bar{\mathcal{P}}'$ mode does in fact dominate, it is clear from the above results that the MIT group[1] may have slightly overestimated their total cross section by assuming a uniform distribution in y out to the full kinematic limit.

At this point it is perhaps worth mentioning a few interesting correlation effects, involving y, that are predicted by the parton-model quark-antiquark annihilation mechanism. In particular, we examine the total residual multiplicity as a function of the resonance's longitudinal momentum, $\frac{1}{2}y\sqrt{s}$. We employ a fully asymptotic approximation. First recall[14] that the intuition associated with the concepts of short-range correlations in combination with parton-model concepts can be used to motivate a rise in total multiplicity in two-particle collisions proportional to $\ln(s)$. In the parton picture, each incoming particle fragments into $\ln E$ final-state particles so that for a collision between A and B the total multiplicity behaves roughly as

$$n \propto \ln E_A + \ln E_B = \ln E_A E_B \approx \ln \tfrac{1}{4} s . \tag{23}$$

Generalization[15] of these concepts to deep-inelastic scattering requires considering a slightly different situation. One must ask what happens to a final state consisting of a parton moving at large momentum relative to a "core" of partons not struck by the massive photon. Given one major assumption beyond those mentioned above, one again discovers that the multiplicity of the final state is a function of the invariant energy of the parton and core (which in the deep-inelastic case is also the invariant s of the initial massive photon plus proton). This extra assumption is that the multiplicity plateau associated in the parton model with e^+e^- annihilation has the same height as that associated with, say, pp scattering. Deep-inelastic scattering data[16] seems to support this conjecture.

In the present situation we have a unique opportunity to test these ideas. Resonance production at a given mass and given y fixes the energies of the two residual cores:

$$E_{core\,1} \approx \tfrac{1}{2}\sqrt{s}\,(1 - x_1), \quad E_{core\,2} \approx \tfrac{1}{2}\sqrt{s}\,(1 - x_2), \tag{24}$$

where asymptotically

$$x_1 = \frac{y + (y^2 + 4\tau)^{1/2}}{2}, \quad x_2 = \frac{-y + (y^2 + 4\tau)^{1/2}}{2} . \tag{25}$$

Thus the multiplicity associated with production of a massive resonance should be most naturally a function of the total invariant mass of the two cores

$$E_{core\,1} E_{core\,2} = \tfrac{1}{4} s\,[1 + \tau - (y^2 + 4\tau)^{1/2}]. \tag{26}$$

Clearly this is a very unique prediction and experimental verification would be of great significance.

The final portion of this paper deals with expectations for production in pN collisions of other particles containing a $\mathcal{P}'$ quark. In particular we present the results for production of all the various types of charmed mesons, at several energies, as a function of resonance mass; we also estimate the cross section for charmed baryon production. All estimates employ the direct coupling mode analogous to $\mathcal{P}'\bar{\mathcal{P}}'$ annihilation for the J. If anything this is a better assumption than that of ignoring the normal quark mode in J production simply because all decay modes to uncharmed states are weak in nature. As before we also ignore any off-shell effects that might be present in the coupling. The only real difficulty arises in estimating the core masses.

As an example consider $\bar{D}^{0*}$ production. $\bar{D}^{0*}$ consists of a $\mathcal{P}$ quark and $\bar{\mathcal{P}}'$ quark. The least massive final state associated with $\bar{D}^{0*}$ production in pN collisions consists of one normal baryon and one charmed baryon. If we label the $\bar{\mathcal{P}}'$ quark as 1, we should associate a charmed baryonlike mass with core 1 while associating a normal protonlike mass with core 2. The smallest[11,17] mass estimated for a charmed baryon is of the order of 3 GeV. Though larger values are also possible we choose to employ the value 3×0.939 GeV for the $\mathcal{P}'\mathcal{P}\mathfrak{N}$, $\mathcal{P}'\mathcal{P}\mathcal{P}$, etc. baryons. Thus our low-energy estimates where these threshold mass effects are substantial may be too high.

A slightly different type of case is that of, for instance, the D^{0*} meson consisting of a $\bar{\mathcal{P}}$ and a $\mathcal{P}'$. Here the least massive final state is almost certainly that of 2 normal baryons plus a charmed meson of the $\bar{\mathcal{P}}'\mathcal{P}$ or $\bar{\mathcal{P}}'\mathfrak{N}$ type. Labeling the $\mathcal{P}'$ quark as 1, core 1 will then consist of, at least, a charmed meson plus a nucleon while core 2 will have the quantum numbers of a nucleon. Thus, again, the smallest mass one can associate with core 1 is of the order 3 GeV, the lightest charmed mesons presumably having a mass of the order of 2 GeV. Thus in either case we take

$$m_{core\,1} = 3 \times 0.939 \text{ GeV}, \quad m_{core\,2} = 0.939 \text{ GeV}. \tag{27}$$

Needless to say, the above considerations must be symmetrized with respect to the target and beam.

In Figs. 7(a)–7(d) we summarize the various possible types of meson production cross sections, presenting each distinct type as a function of mass at the energies $p_{lab} = 20$, 28.5, and 300 GeV/c. Figure 7(a), which we give for reference, applies to J-like particles composed of $\mathcal{P}'\bar{\mathcal{P}}'$. Figure 7(b) applies to $\bar{D}^{0*}$ production, Fig. 7(c) to D^{-*} production, and Fig. 7(d), for annihilation of a quark and antiquark both belonging to the sea component, is appropriate for all other charmed meson types.

We wish to point out just a few features of these results. Note that production of a $\bar{D}^{0*}$ or D^{-*} resonance of mass 2 GeV at $p_{\text{lab}} = 28.5$ GeV/c is of the order of 400 times the production cross section for $J(3.1)$ at the same energy. In addition certain decay modes such as the $K\pi$ channel should have quite substantial branching ratios,[11] so that proper choice of a final state would result in only a minor loss from the total cross section. Even for the other charmed mesons the production cross section at a resonance mass of 2 GeV is comparable (at $p_{\text{lab}} = 28.5$ GeV/c) to that for $J(3.1)$, the lower resonance mass compensating for the increased core mass.

Finally, in Fig. 7(e) are the estimates for production of a charmed baryon of the $\mathcal{P}'\mathcal{P}\mathcal{N}$ variety. The core, associated with extracting a $\mathcal{P}\mathcal{N}$ pair from a nucleon, presumably communicates with the $\bar{\mathcal{P}}'$ left in the core, associated with $\mathcal{P}'$ extraction from the other nucleon, to form a charmed meson whose mass we take as 2.2 GeV. The remaining core on the $\mathcal{P}'$ side has nucleon quantum

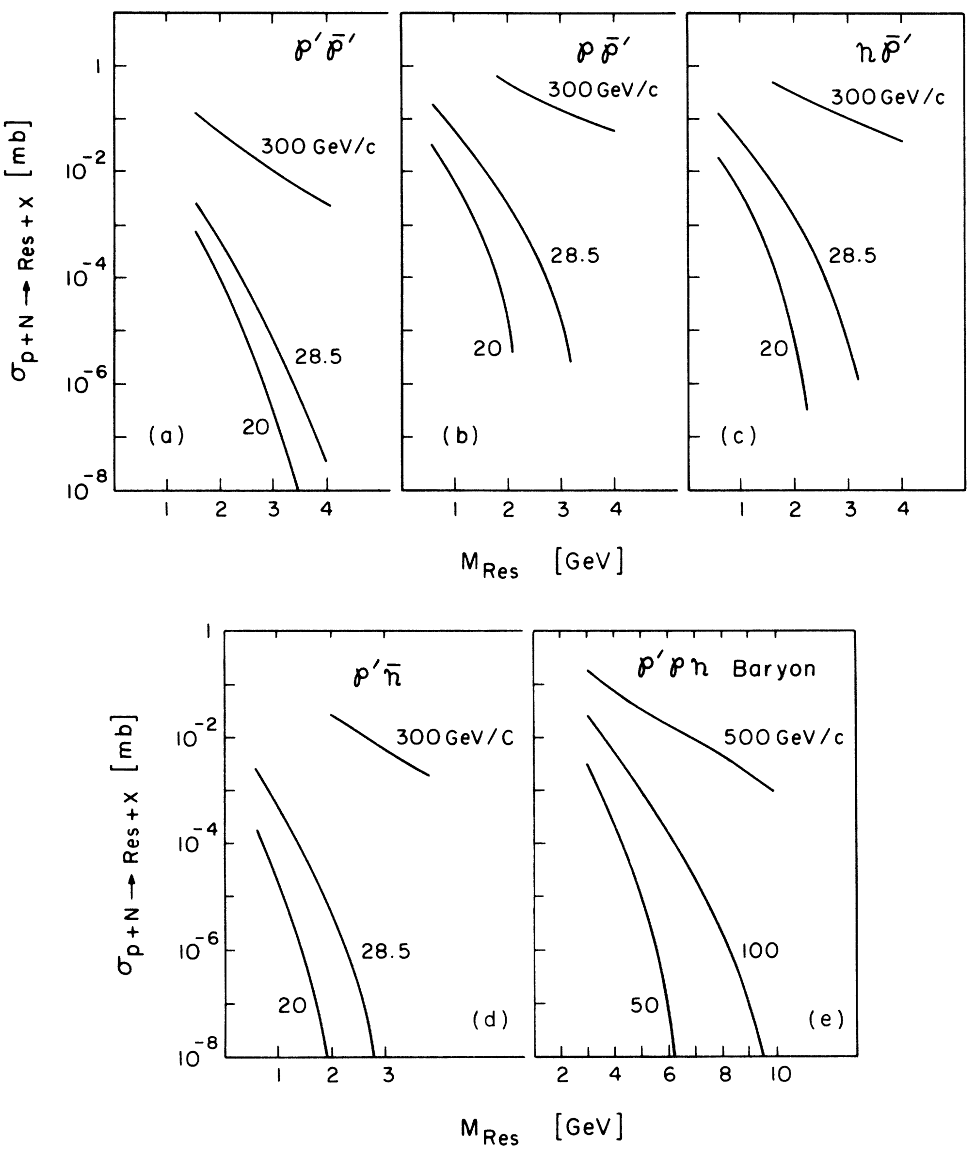

FIG. 7. Production cross sections for: (a) J-like mesons; (b) $\bar{D}^{0*}$; (c) D^{-*}; (d) D^{+*}, D^{0*}, F^{+*}, and F^{-*} mesons; and (e) the most favorable charmed baryon case, C_0^+. For all but (a), one core mass is $3\times m_{\text{proton}}$ while the other is $1\times m_{\text{proton}}$. For the reasoning see the text. No final branching ratios are incorporated. $g_{q\bar{q}}^2/4\pi = 1$ for all cases.

numbers with an associated mass of the order of 0.939 GeV. We employ a simplified distribution function describing the probability of extracting a $\mathcal{P}\mathcal{N}$ pair from a nucleon. We take

$$f_{\mathcal{P}'\mathcal{N}\text{pair}}(x_{\text{pair}}) = 0.5 \times \frac{1 - x_{\text{pair}}}{x_{\text{pair}}}, \quad (28)$$

which has the threshold damping expected theoretically for a residual core consisting of only one quark (before communication occurs). The normalization is chosen so that the $\mathcal{P}\mathcal{N}$ pair carries $\frac{1}{4}$ of the nucleon's momentum.

Because of the uncertainty in the mass of a charmed baryon we have given results (at three energies $p_{\text{lab}} = 50$, 100, and 500 GeV/c) over a large mass range. Clearly production cross sections are substantial even at moderate energies, such as 50 GeV/c, for charmed baryon masses near the lower end of the possible range. Unfortunately the standard Brookhaven energy, $p_{\text{lab}} = 28.5$ GeV/c, corresponding to a maximum final-state mass of about 7.4 GeV is too near the minimum final-state mass of about 6 GeV, appropriate to the present case, for the cross section to be measureable. We should also again mention the likelihood that off-shell effects will, in this case of production of a dipole form-factored baryon, lead to some reduction in the present estimates.

CONCLUSION

In conclusion we should attempt to examine the above predictions in light of currently available data.[18] To begin with we concentrate on the $\mathcal{P}'\overline{\mathcal{P}}'$ annihilation mode. The main uncertainty in our predictions arises from the possibility that there are substantially fewer charmed quarks than λ, $\bar{\lambda}$ quarks in a nucleon, i.e., that the charmed distribution functions should have a substantially smaller normalization than that appropriate to the SU(4)-symmetric case. (The charm coupling constant $g_{J\mathcal{P}'\overline{\mathcal{P}}'}{}^2/4\pi$ could also be slightly different from the value of 1 used here but not by more than a factor of 2 or so.)

This possibility may be tested by examining a related process—namely inclusive photoproduction of $J(3.1)$. The (pointlike or bare) photon contains approximately as many charmed $\mathcal{P}'$ quarks as normal $\mathcal{P}$ quarks. The distribution of $\mathcal{P}'$ quarks in the bare photon is known—both in shape and normalization. Thus one may calculate the Drell-Yan $\mathcal{P}'\overline{\mathcal{P}}'$ annihilation contribution to inclusive photoproduction.[19]

Both $\gamma + N \to J(3.1) + X(\langle p_{\text{lab}}\rangle \simeq 150$ GeV/c) and $n + N \to J(3.1) + X$ ($\langle p_{\text{lab}}\rangle \simeq 250$ GeV/c) have been measured at Fermilab. The cross sections[18] (excluding the $\simeq 50\%$ contribution from coherent diffractive production in the photoproduction case are approximately

$$1.5 \times 10^{-32}\ \text{cm}^2/\text{nucleon}$$

and

$$3 \times 10^{-32}\ \text{cm}^2/\text{nucleon}, \quad |y| \geq 0.22,$$

respectively (y is the Feynman momentum fraction variable). At the above $\langle p_{\text{lab}}\rangle$ values we obtain theoretically, for an SU(4)-symmetric $q\bar{q}$ sea, about

$$2 \times 10^{-31}\ \text{cm}^2/\text{nucleon}$$

and

$$4 \times 10^{-30}\ \text{cm}^2/\text{nucleon}, \quad |y| \geq 0.24,$$

both substantially above the experimental values, by factors of $\simeq 10$ and $\simeq 100$, respectively. (For the latter case, integration over all y values would approximately double the theoretical cross section.) An obvious possibility presents itself; since the $\mathcal{P}'$ nucleonic distribution enters linearly in the photoproduction cross section but quadratically in the purely hadronic cross section reducing its normalization by a factor of 10 puts both theoretical estimates in reasonable agreement with experiment. Thus the theory is viable provided the quark-antiquark sea is dominated by normal quarks.

We should note that the situation is less appealing if the inclusion of color is appropriate. We take the photon to be a color singlet for the purpose of illustration. Color inclusion reduces both theoretical cross sections by a factor of 9 ($\frac{1}{3}$ from color matching requirements and $\frac{1}{3}$ because of the reduced average size of $g_{J\mathcal{P}'\overline{\mathcal{P}}'}{}^2/4\pi$ necessary to obtain the $\mathcal{P}'$ charge of $\frac{2}{3}$). The SU(4)-symmetric sea photoproduction prediction then roughly agrees with experiment while the hadronic cross section is still a factor of 10 too high.

It is possible that one should remove the two-body (J, p) final state from the experimental data before comparing to the annihilation theory. This seems arbitrary to us but would help in the present case. The photoproduction cross section would be most affected and is probably reduced to less than $\frac{1}{3}$ of the full cross section. Thus reducing the $\mathcal{P}'\overline{\mathcal{P}}'$ sea normalization by a factor of 3 while including color would restore agreement.

It is because of this type of uncertainty that we reemphasize the importance of the p_T and y distributions. Those measured in $n + \text{Be} \to J(3.1) + X$ ($|y| \geq 0.24$) at Fermilab are remarkably similar to the theoretical predictions given earlier. In addition, the photoproduction cross section is theoretically predicted to exhibit a slightly weaker p_T dependence than the nucleon-induced production cross section (consistent with observation) while displaying a markedly different y distribution which is relatively flat in the positive y (photon fragmentation) region but rapidly decreasing for negative y. We also mention that the $pp\to J+X$ cross section at CERN ISR[18] ($\sqrt{s}=52.7$), which, as expected, is somewhat higher than the Fermilab cross section (our model predicts a factor of 2.5, consistent with preliminary results), also confirms that the y dependence is flat, near $y\simeq 0$, as predicted by the annihilation model, even though at larger y theory predicts and experiment confirms a relatively steep y dependence. Thus current indications are promising in this respect.

The factor of 10 reduction, ignoring color, in the $\mathcal{P}'$ distribution function in nucleons implies that the D and F cross sections of the figures should be reduced by a factor of 10. Nonetheless, they are large so that D's and F's should be observed at Fermilab energies (i.e., $p_{\rm lab}\gtrsim 200$ GeV/c). Their detection may be tricky because of their many possible decay modes. Failure to observe them at the expected total cross-section level would constitute an important failure of the $\mathcal{P}'\bar{\mathcal{P}}'$ mode approach. D's and F's (the pseudoscalar charmed mesons) should, of course, have quite similar cross sections (perhaps smaller by the spin statistical weight factor of 3). Inclusion of color in the manner mentioned earlier would imply a reduction by a factor of 30 in the figure cross sections.

The above adjustments do, however, worsen agreement with the Brookhaven result. The normal quark mode would be dominant at $p_{\rm lab}=28.5$ GeV/c and yields (for a sea dominated by uncharmed quarks)

$$\sigma(pN\to J+X)\simeq 0.1\times 10^{-34}\ {\rm cm}^2/{\rm nucleon}$$
$$\searrow e^+e^-$$

(this includes the latest branching ratio information $R\simeq 15$ and $\Gamma_J\simeq 70$ KeV). This is factor of 10 below the experimental estimate of 10^{-34} cm^2/nucleon. This latter estimate assumes (probably incorrectly) a completely flat y distribution for J and is probably too high by a factor of 2 to 3.

It is clear that the y distribution must be measured and compared to the theoretical prediction, as well as used to obtain an accurate experimental cross section, before a final decision can be made. At this lower energy it is certainly not impossible that non-$q\bar{q}$ annihilation contributions could also be important. Such modes would clearly be essential if color, which reduces the theoretical annihilation cross section due to normal quarks by a further factor of 9, is included.

The above comments are obviously speculative in nature at the present time but should serve to illustrate the possibilities. The most important experimental unknowns are the D and F meson cross sections at Fermilab energies. We anxiously await definite results.

APPENDIX[6]

In order to include subasymptotic corrections we must first write the δ function of Eq. (1) including nonasymptotic corrections. The exact form of the δ function is

$$s\delta((k_1+k_2)^2-M_J^2), \tag{A1}$$

where k_1 and k_2 are the momenta of the off-shell annihilating quarks in Fig. 1. In terms of x_1 and x_2, the momentum fractions mentioned earlier, and $k_1^\perp$ and $k_2^\perp$, the transverse-momentum fluctuations of the quarks about the direction of the initial beam particles, we have ($\tau=M_J^2/s$) from Ref. 6

$$\begin{aligned}\frac{(k_1+k_2)^2}{s}-\tau=&\,x_1x_2(1-2m_{\rm proton}^2/s)\\&+(x_1+x_2)m_{\rm proton}^2/s-\tau-\frac{2k_1^\perp\cdot k_2^\perp}{s}\\&-\frac{x_1\sigma_1+k_1^{\perp 2}}{1-x_1}-\frac{x_2\sigma_2+k_2^{\perp 2}}{1-x_2}.\end{aligned} \tag{A2}$$

σ_1 and σ_2 are the squared masses of the "cores" left behind by the quarks of Fig. 1. For J production we take $\sigma_1=\sigma_2=m_{\rm proton}^2$ corresponding to a communication mechanism capable of removing their charm and quark quantum numbers.

In addition we must expose the $k_1^\perp$ and $k_2^\perp$ integrations in Eq. (1) by writing

$$f_i(x_1)=\int d^2k_1^\perp f_i(x_1,k_1^\perp),\quad \text{etc.} \tag{A3}$$

The final ingredient is the form of $f(x_1,k_1^\perp)$. The distribution functions depend most directly upon the off-shell quark momenta, e.g.,

$$k_i^2=x_i m_{\rm proton}^2-\frac{x_i\sigma_i+k_i^{\perp 2}}{1-x_i}. \tag{A4}$$

We shall see that threshold effects probe x_i near 1 and hence we keep only those terms in k_i^2 singular in $1/(1-x_i)$. As $x_i\to 1$ the threshold power associated with a given component (valence, sea, ...) of f_i is taken to obey the theoretically motivated power laws of Refs. 3–5,

$$f_i(x_i)\underset{x_i\to 1}{\sim}\frac{(1-x_i)^{2N_i-1}}{1-x_i}, \tag{A5}$$

where N_i is the number of quarks in the core associated with quark i. Since the power suppression (of order $2N_i$) in the limit $x_i\to 1$ arises from the off-shell damping in $k_i^{\perp 2}$ power. Specifically, we take, for each component of f_i,

$$f_i(x_i,k_i^\perp)=P_if_i(x_i)/(k_i^{\perp 2}+\sigma_i)^{2N_i}, \tag{A6}$$

where P_i is chosen so that $\int f_i(x_i, k_i^\perp) d^2k_i^\perp = f_i(x_i)$. This approximation to the $k_\perp$ dependence is valid in the region $x_i \to 1$ which is that probed as threshold is approached. (Away from threshold the $k_i^\perp$ integration is essentially unrestricted and the specific form employed immaterial.)

The maximum possible τ which satisfies the function is obtained with

$$k_1^\perp = k_2^\perp = 0,$$
$$\sigma_1 = \sigma_2 = m_{\rm proton}^2,$$
and
$$x_1 = x_2 \simeq 1 - \frac{m_{\rm proton}}{\sqrt{s}} \tag{A7}$$

and is

$$\tau_{\max} \approx s - 4m_{\rm proton}/\sqrt{s}\ . \tag{A8}$$

For this value the $k_i^\perp$ integrations are totally restricted so that the cross section goes to zero. This procedure is exact to order $m_{\rm proton}/\sqrt{s}$; terms of order m^2/s (other than $m_{\rm resonance}{}^2/s$) are not treated with total precision. For Brookhaven energies, and above, the approximation is adequate.

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[1]J. J. Aubert *et al*., Phys. Rev. Lett. 33, 1404 (1974); J. E. Augustin *et al*., *ibid*. 33, 1406 (1974); C. Bacci *et al*., *ibid*. 33, 1408 (1974); G. S. Abrams *et al*., *ibid*. 33, 1453 (1974).

[2]S. D. Drell and T.-M. Yan, Ann. Phys. (N.Y.) 66, 576 (1971). The worry that diffractive effects modify the simple prediction now appears to be unfounded. See J. L. Cardy and G. A. Winbow, Phys. Lett. 52B, 95 (1974).

[3]J. F. Gunion, Phys. Rev. D 10, 242 (1974).

[4]G. Farrar, Nucl. Phys. B77, 429 (1975).

[5]R. Blankenbecler and S. J. Brodsky, Phys. Rev. D 10, 2973 (1974) also discuss the theoretical considerations.

[6]G. Chu and J. F. Gunion, Phys. Rev. D 11, 3672 (1974). We note that we do not include the colored parton degrees of freedom in the discussion of the present paper. These decrease our estimates by a factor of 9 [assuming $g_{J\mathcal{P}'\bar{\mathcal{P}}'}{}^2/4\pi$ for each of the 3 $\mathcal{P}'$ quarks is $\frac{1}{3}$ the noncolored value quoted in (3); the analogous reduction in the case of $g_{\rho\mathcal{P}\bar{\mathcal{P}}}$ $(g_{\rho\mathfrak{N}\bar{\mathfrak{N}}})$ is required in order to preserve the total ρ width in the parton model approximation—see text].

[7]J. F. Gunion, Phys. Rev. D 11, 1796 (1975).

[8]T. Appelquist and H. D. Politzer, Phys. Rev. Lett. 34, 43 (1975); A. De Rújula and S. L. Glashow, *ibid*. 34, 46 (1975); C. G. Callan *et al*., *ibid*. 34, 52 (1975).

[9]J. J. Aubert *et al*., Phys. Rev. Lett. 33, 1624 (1974).

[10]S. J. Brodsky and G. Farrar, Phys. Rev. Lett. 31, 1153 (1973); F. Muradyn and A. Tavkhelidze, Nuovo Cimento Lett. 7, 719 (1973).

[11]We use the notation of M. K. Gaillard *et al*., Rev. Mod. Phys. 47, 277 (1975).

[12]This assumes that the photoproduction cross section is in fact not greatly above the 20 nb value rumored from the experiment of H. Lee *et al*. at Fermilab

[13]For a review see J. F. Gunion, in *Particles and Fields—1974*, proceedings of the 1974 Williamsburg meeting of the Division of Particles and Fields of the American Physical Society, edited by Carl E. Carlson (A.I.P., New York, 1975), p. 551. See also Ref. 3, R. Blankenbecler and S. J. Brodsky, Ref. 5, and R. Blankenbecler, S. J. Brodsky, and J. F. Gunion, Phys. Rev. D (to be published).

[14]J. Kogut and L. Susskind, Phys. Rep. 8C, 175 (1973); R. P. Feynman, *Photon-Hadron Interactions* (Benjamin, N. Y., 1972).

[15]J. Bjorken, Phys. Rev. D 7, 282 (1973) and R. Cahn, J. W. Cleymans, and E. W. Colglazier, Phys. Lett. 43B, 323 (1973).

[16]J. Ballam *et al*., submitted to the XVI International Conference on High Energy Physics, Chicago-Batavia, 1972 (unpublished), and SLAC Report No. SLAC-PUB-1195 (unpublished).

[17]S. Borchardt, V. S. Mathur, and S. Okubo, Phys. Rev. Lett. 34, 38 (1975).

[18]We quote results presented at the 1975 Meeting of the American Physical Society, Washington, by the various groups.

[19]The photoproduction calculation has been performed independently by T. Goldman, SLAC Report No. SLAC-PUB-1538 (unpublished). Our SU(4)-symmetric-sea photoproduction cross sections are smaller than his, *for a given value of* $g_{J\mathcal{P}'\bar{\mathcal{P}}'}$, by roughly a factor of 2. His normalization for the $\mathcal{P}'$ ($\bar{\mathcal{P}}'$) distribution does not agree with the limitations on the total sea's contribution to deep-inelastic scattering.