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LARGE TRANSVERSE MOMENTUM PROCESSES AND THE CONSTITUENT INTERCHANGE MODEL*

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Abstract: The predictions of the constituent interchange model are found to be consistent with the normalization, as well as the scaling laws and angular dependence, of measured large p_T meson and baryon cross sections. The normalization of the hadronic couplings to valence quarks is computed. Predictions for quantum number correlations between the trigger particles and away side jets are discussed. We also contrast the predictions of the CIM and quark-quark scattering models.

Note: A final version of this work, taking into account more realistic parameterizations of the structure functions, color factors, etc., will be published elsewhere.

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I. Introduction

Hadronic collisions involving the production of particles at large transverse momentum have the exciting potential of being able to resolve the underlying structure of hadrons and the interactions of their constituents at very short distances. The phenomenological features which have emerged from the recent ISR and Fermilab experiments—particularly the jet structure and the scaling laws of the inclusive cross sections—appear to be consistent with the properties expected from underlying two-body hard scattering subprocesses.[1-4] The data[4] for single particle cross sections, charge, momentum, and angular correlations are now so extensive that the constraints on models are overwhelmingly restrictive.

In this paper we will present a comparison of this data with the predictions of the constituent interchange model[2] (CIM). The central postulate of the CIM is that the dominant short distance subprocesses are quark-hadron interactions (e.g., $qM \to qM$, $qB \to qB$, and the reactions related by crossing, $q\bar{q} \to M\bar{M}$, etc.) which may be computed from an underlying scale-invariant field theory. We emphasize that such diagrams contribute in any quark model since their amplitude normalization is already fixed from the hadronic Bethe-Salpeter wavefunctions, elastic form factors, momentum sum rules for structure functions, etc. In fact, as we show in this paper, the CIM predictions are consistent not only with the scaling laws and angular dependence of the measured exclusive and inclusive large p_T cross sections, but also with their normalization. The new preliminary data from the British-French-Scandinavian group (BFS) presented at this meeting by Møller[5] on charge and momentum correlations also appear to support the basic features of the CIM subprocesses, in particular, the prediction of strong quantum number correlations between the trigger particles and the away side jet.

It should also be emphasized that dominance of the CIM diagrams at present energies is not incompatible with the assumption of a fundamental quark-gluon field theory such as quantum chromodynamics. In particular, the single gluon exchange term for quark-quark scattering,

$$\frac{d\sigma}{dt} = \frac{2}{9} \frac{4\pi\alpha_s^2}{t^2} \frac{s^2+u^2}{2s^2} \quad , \tag{1.1}$$

has been shown in Reference 6 to give a contribution below present data for $\frac{d\sigma}{d^3p/E}(pp \to \pi X)$ for all $p_T \lesssim 8$ GeV, assuming $\alpha_s(p_T^2) \lesssim .4$ (a conservative value). The CIM contributions will then dominate at lower p_T simply because of the relatively large effective hadron-quark coupling strengths. We note though that the $qq \to qq$ cross section could still be an important contribution to jet-trigger experiments in which the effect of trigger bias is removed.

II. CIM Predictions

In the constituent intercharge model and other hard scattering models, the inclusive cross section for $A+B \to C+X$ at large p_T can be written as a convolution over structure functions $G_{a/A}(x_a, \vec{k}_{Ta})$, $G_{b/B}(x_b, \vec{k}_{Tb})$, and $G_{C/c}(x_C, \vec{k}_T^C)$ times the square of the matrix element for the subprocesses $a+b \to c+d$ (see Fig. 1). In a scale-invariant theory, dimensional counting[7] predicts at large p_T

$$\frac{d\sigma}{dt}(a+b \to c+d) \Rightarrow \frac{1}{\left(p_T^2\right)^{n_{active}-2}} f(\theta_{c.m.}) , \quad (2.1)$$

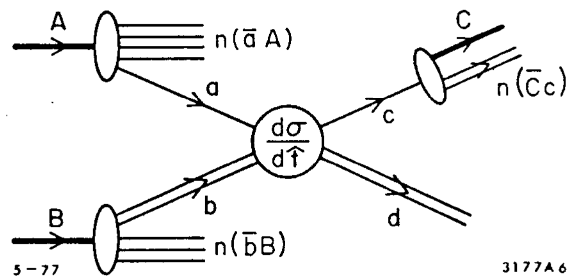

Fig. 1. Hard scattering subprocess contribution $ab \to c+d$ to the inclusive cross section $A+B \to C+X$.

where $n_{active} = n_a+n_b+n_c+n_d$ is the number of elementary fields in the subprocess, and[8] $G_{a/A}(x_a) \sim (1-x_a)^{2n(\bar{a}A)-1}$ at $x_a \to 1$, where $n(\bar{a}A)$ is the number of elementary particles left behind in the fragmentation of $A \to a$. These predictions are based on the short distance behavior of lowest order terms in renormalizable perturbation theories assuming a finite Bethe-Salpeter hadronic wave-function. Detailed discussions and comparisons with exclusive processes, form factors, large angle scattering, and structure functions are given in Refs. 4, 7, and 8.

The result of the convolution then gives the counting rules[7,8]

$$E\frac{d\sigma}{d^3p}(A+B \to C+X) = \sum_{abcd} \frac{1}{\left(p_T^2+m^2\right)^{n_{active}-2}} f(\epsilon, \theta_{c.m.})$$

$$\underset{\epsilon \to 0}{\sim} \sum_{abcd} \frac{1}{\left(p_T^2+m^2\right)^{n_{active}-2}} \epsilon^F f(\theta_{c.m.}) \qquad (2.2)$$

where $\epsilon = \mathcal{M}^2/s = (1-x_T)$ at $\theta_{c.m.} = \pi/2$. Here n_{active} is the number of active fields in the high p_T subprocess (e.g., $n_{active} = 4$ for $qq \to qq$, 6 for $qM \to qM$) and $F = 2n_{spect} - 1$ where $n_{spect} = n(\bar{a}A) + n(\bar{b}B) + n(\bar{C}c)$ is the minimum number of elementary constituents that "waste" the momentum in the fragmentations $A \to a$, $B \to b$, $c \to C$ (e.g., $n_{spec} = 5$ and $F=9$ for $qq \to qq$ or $qM \to qM$ in $pp \to MX$). In general, one predicts that aside from normalization effects, the subprocesses with the minimum n_{active} (minimum p_T^{-1} power) and minimum n_{spect} (minimum F power) will dominate the cross section at large p_T, and small ϵ. Thus, given the fact that the $qq \to qq$ term has a small predicted normalization, the dominant terms for $pp \to \pi^\pm, K^+X$ will come from the $qM \to qM$ subprocess[2] (Fig. 2a):

$$\frac{d\sigma}{d^3p/E}(pp \to \pi^\pm, K^+, X) \sim \frac{\epsilon^9}{\left(p_T^2 + m^2\right)^4} f(\theta_{c.m.}). \quad (2.3)$$

Here m^2 represents terms of order $\langle \vec{k}_T^2 \rangle$, m_q^2, etc. All other quark-hadron subprocesses lead to a higher power of $1/p_T$ or ϵ. In the case of K^- production, the dominant contribution at high p_T small ϵ will come from the "fusion" subprocess[3,2] $q\bar{q} \to K^-M$ (Fig. 2b)

$$\frac{d\sigma}{d^3p/E}(pp \to K^-X) \sim \frac{\epsilon^{11}}{\left(p_T^2 + m^2\right)^4} f(\theta_{c.m.}) . \quad (2.4)$$

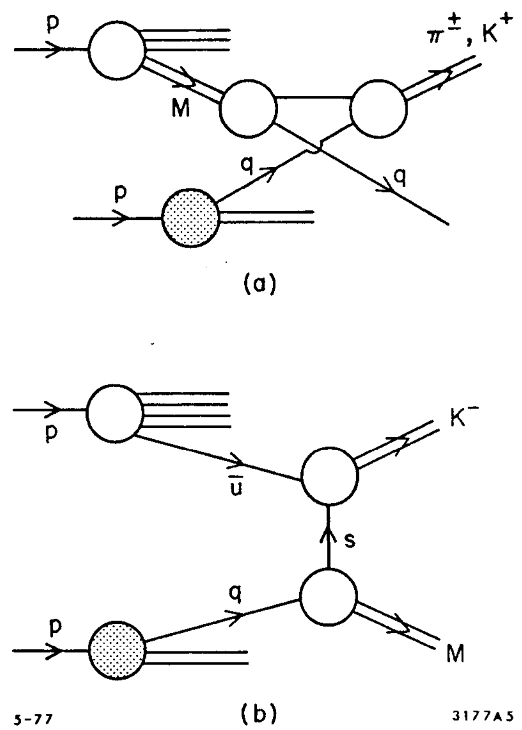

Fig. 2. Dominant CIM contribution to (a) $pp \to \pi^\pm$, K^+X and (b) $pp \to K^-X$.

A comparison of the CIM predictions with the experimentalists' fits to the Chicago-Princeton-Fermilab[9] data for $pp \to \pi^\pm, K^\pm, p^\pm X$ is shown in Table I. The agreement seems remarkable. For example, as shown in Fig. 3, the best fit for the Chicago-Princeton $\theta_{c.m.} = 90^\circ$ data for $pp \to \pi^+X$ is $p_T^{-8.2}(1-x_T)^{9.0}$ (with uncertainties in n and F order ±0.5). The relative suppression of $Ed\sigma/d^3p\,(pp \to \pi^-X)/Ed\sigma/d^3p\,(pp \to \pi^+X) \sim (1-x_T)$ evidently reflects the relative suppression of the d/u quark ratio in the proton structure function at large x.

Table I. Scaling predictions for $E d\sigma/d^3p = C p_T^{-n} (1-x_T)^F$.

Large p_T Process	Leading CIM Subprocess	Predicted	Observed (CP)[9]
		n//F	n//F
$pp \to \pi^+ X$	$qM \to q\pi^+$	8//9	8.5//9.0
π^-	$qM \to q\pi^-$	8//9	8.5//9.9
K^+	$qM \to qK^+$	8//9	8.4//8.8
K^-	$q\bar{q} \to K^+K^-$	8//11	8.9//11.7
	$qM \to qK^-$	8//13	
$pp \to pX$	$q(qq) \to Mp$	12//5	11.7//6.8
	$qB \to qp$	12//7	
$pp \to \bar{p}X$	$q\bar{q} \to B\bar{p}$	12//11	(8.8//14.2)[a]
	$qM \to qM$	8//15	
	$M\bar{M} \to q\bar{q}$	8//15	
$\pi p \to \pi X$	$q\bar{q} \to M\pi$	8//5	
	$qM \to q\pi$	8//7	
	$q(qq) \to B\pi$	12//3	
	$\pi q \to \pi q$	8//3	

[a] The $\bar{p}$ fit has large uncertainties and is compatible with n=12, F=11.

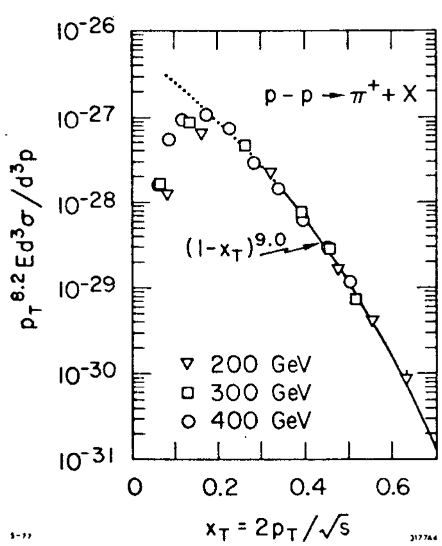

Fig. 3. Scaling law fit to the cross section $pp \to \pi^+ X$, $\theta_{c.m.} \cong 90^\circ$, $x_T = 2p_T/\sqrt{s} > 0.3$. From Ref. 9.

A crucial check on the identification of the underlying subprocess is the angular dependence of its cross section. The leading CIM contributions at high p_T to $pp \to \pi^+ X$ arise from the basic process

$$\frac{d\sigma}{d\hat{t}}(u\pi^+ \to u\pi^+) = \frac{C}{\hat{s}\hat{u}^3} \tag{2.5}$$

and by $\hat{u} \to \hat{s}$ crossing

$$\frac{d\sigma}{d\hat{t}}(d\pi^+ \to d\pi^+) = \frac{C\hat{u}}{\hat{s}^5} , \tag{2.6}$$

These predictions can be obtained by explicit calculation, or by using quark counting and the fact that the $u\pi^+ \to u\pi^+$ amplitude corresponds to spin 1/2 exchange in the u channel. It is easy to see that $d\pi^+ \to d\pi^+$ term gives a small contribution compared to the leading $1/\hat{s}\hat{u}^3$ term.

The angular dependence of the subprocess can be directly determined from experiment either from the correlated angular dependence of the away side jet[10] or the angular dependence of the $pp \to \pi X$ inclusive cross section.[11] In both cases the experimental p_T data are best fit with the form

$$\frac{d\sigma}{d\hat{t}} \propto \frac{1}{\hat{s}\hat{t}^3} \quad \text{or} \quad \frac{1}{\hat{s}\hat{u}^3} \tag{2.7}$$

(equivalent because of the pp symmetry). It should be emphasized that this angular dependence implies elementary spin 1/2 exchange in the t or u channel and is evidently difficult to reconcile with a subprocess based on quark-quark scattering.

The convolution of the distributions $G_{M/p} \sim (1-x)^5/x$, $G_{u/p} \sim (1-x)^3/x$ and the cross section for $uM \to u\pi^+$ gives the CIM prediction $E d\sigma/d^3p\,(pp \to \pi^+) \propto p_T^{-8}\,\epsilon^9$, with the angular dependence given in Eq. (2.7). An immediate and important question is whether we can understand and predict the normalization of the cross sections as well. This will be discussed in detail in the next section. The CIM subprocesses also make detailed predictions for the quantum number flow of the valence quarks in large p_T reactions. We discuss this and the general question of jets and correlations in Section IV.

III. Normalization of CIM Subprocesses

A. The Meson-Quark-Antiquark Coupling

The magnitude of the amplitude $\mathcal{M}(u\pi^+ \to u\pi^+)$ required for the CIM predictions (see Fig. 2a) is directly related to the normalization of the Bethe-Salpeter vertex function for $\pi^+ \to u\bar{d}$ which in turn can be fixed by the normalization of the pion form factor or equivalently, the momentum sum rule for its structure function. The connection is clear from Fig. 4a-c. For simplicity we shall at first ignore the minor effects of spin and parametrize the large angle amplitude in Fig. 4a as $\mathcal{M}(u\pi^+ \to u\pi^+) = g^2/u$ where g represents the $\pi^+ \to \bar{u}d$ vertex function (i.e., coupling constant); g has dimensions of mass. Note that g refers to the effective coupling of the pion to its valence $q\bar{q}$ component, the wavefunction which dominates both the large angle elastic scattering amplitude and the meson structure function for x near 1.

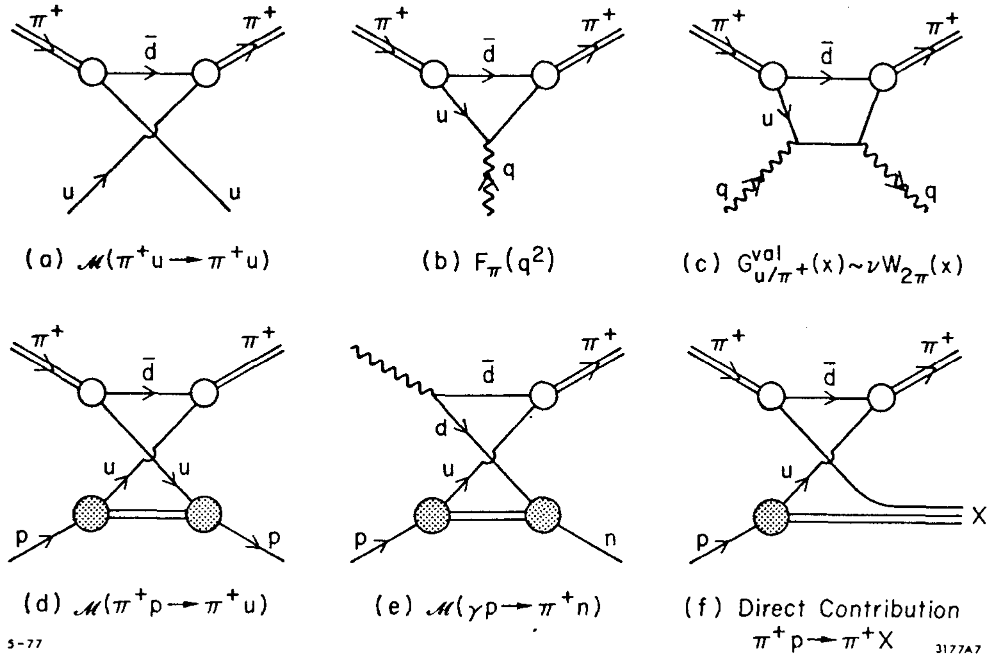

Fig. 4. Contribution of the $\pi^+u \to \pi^+u$ valence scattering amplitude (a), to the pion form factor (b), valence structure function (c), large angle $\pi^+p \to \pi^+p$ scattering (d), and inclusive scattering (f) (direct contribution). The relationship of photoproduction (e) to elastic scattering at large angles (c) is also shown.

The contribution of the valence state to the pion structure function is then

$$G^{val}_{u/\pi^+}(x) = \int d^2k_T \frac{g^2}{2(2\pi)^3} \frac{x(1-x)}{\left[\vec{k}_T^2 + M^2(x)\right]^2} \tag{3.1}$$

where $M^2(x) = m_u^2(1-x) + m_d^2(x) - x(1-x)m_\pi^2$, which we shall treat as a phenomenological constant. The fraction of the pion momentum carried by the valence quark in the pion is then

$$f^{val}_{u/\pi^+} \equiv \int_0^1 dx\, x\, G^{val}_{u/\pi}(x) = \frac{g^2}{4\pi}\frac{1}{4\pi}\frac{1}{\langle M^2(x)\rangle}\int_0^1 dx\, x^2(1-x) \tag{3.2}$$

A reasonable estimate is $\langle M^2(x)\rangle \sim .25$ GeV^2 (to set the mass scale of the pion form factor correctly) and $f^{val}_{u/\pi} \sim 0.05$ (from the empirical behavior of the fragmentation functions $D_{\pi^+/u}(x)$ at $x \gtrsim 0.8$ where $D_{\pi^+/u}(x) \sim G^{val}_{\pi^+/u}(x)$. This gives the rough estimate $g^2/4\pi \sim 1$–2 GeV^2. We note that more accurate information on $G_{\pi^+/u}(x)$ in the valence region could be obtained from forward pair production in the Drell-Yan process $\pi^+p \to \ell^+\ell^-X$.

An important cross check to determine the coupling of the meson to its valence component is the magnitude of large angle meson-nucleon scattering and photoproduction.

Quark exchange diagrams such as those shown in Fig. 4d for π^+p scattering give an excellent parametrization of the fixed angle scaling behavior and angular dependence of the cross section $d\sigma/dt \propto s^{-1}t^{-4}u^{-3}$. A simple calculation, apparent from the impulse approximation structure of the diagrams, gives

$$\frac{d\sigma}{dt}(\pi^+p \to \pi^+p) \cong 4\,\frac{d\sigma}{dt}(\pi^+u \to \pi^+u)\; F_p^2(t) \left\langle \frac{1}{x^2} \right\rangle \tag{3.3}$$

(The factor of 4 comes from the two coherent diagrams. The $\pi^+d \to \pi^+d$ term is relatively small. The factors of x^{-1} occur here because the $\pi q \to \pi q$ amplitude is proportional to $(xu)^{-1} = \hat{u}^{-1}$ compared to the $eq \to eq$ coupling in $F_p(t)$ which is proportional to $\hat{x}s/t$.) Empirically, $d\sigma/dt \sim 0.4\ \mathrm{nb/GeV}^4$ at $t=u=-10\ \mathrm{GeV}^2$, giving $g^2/4\pi \sim 1.1\ \mathrm{GeV}^2$, taking $\langle x \rangle = 1/3$ (as expected from the proton valence wavefunction).

Alternatively we can consider the ratio of pion photoproduction $\gamma p \to \pi p$ and $\pi p \to \pi p$ scattering at fixed angle. The measured cross sections[13] are consistent with the dimensional counting predictions $d\sigma/dt \cong s^{-7} f(\theta_{c.m.})$ and $d\sigma/dt \cong s^{-8} f(\theta_{c.m.})$, respectively. In the CIM, the amplitudes only differ by the replacement of the direct photon coupling by the composite meson coupling, in Fig. 4e. Thus we have

$$\frac{\frac{d\sigma}{dt}(\gamma p \to \pi^+ n)}{\frac{d\sigma}{dt}(\pi^+p \to \pi^+p)} \cong \frac{2}{3}\,\frac{\bar{\lambda}^2 \alpha}{g^2/4\pi}\,\langle x \rangle s \tag{3.4}$$

where $\bar{\lambda}^2$ is the average quark charge. Using the measured ratio[13] at $s = 10\ \mathrm{GeV}^2$, $\bar{\lambda}^2 = 5/9$, and $\langle x \rangle = 1/3$, we find $g^2/4\pi \sim 1.2\ \mathrm{GeV}^2$. Of all determinations of g^2 this invokes the least number of assumptions for parameter values, and thus should be the most reliable. We also note that the near equality of $d\sigma/dt\,(\pi^+p \to \pi^+p)$ and $d\sigma/dt\,(K^+p \to K^+p)$ at $\theta_{c.m.} = 90^\circ$, $s = 10\ \mathrm{GeV}^2$ implies that $g^2/4\pi$ is to first approximation SU(3) symmetric. We will discuss the implications of Eq. (3.4) for the inclusive γ/π ratio at high p_T in the next section. We can also predict the ratio of $d\sigma/dt\,(\gamma p \to \gamma p)$ to $d\sigma/dt\,(\gamma p \to \pi p)$ from a form similar to (3.4).

B. Normalization of Inclusive Reactions

Let us now try to predict the magnitude of inclusive large p_T reactions using the above coupling constant. The simplest contribution to $\pi p \to \pi X$ comes from the "direct" scattering graph, Fig. 4f. One only expects this "quasi-exclusive" diagram to be important at quite large $x_R = 1-\epsilon$ in analogy to the dominance of triple Regge contribution at large x_L. A simple estimate gives

$$\frac{d\sigma}{d^3p/E}(\pi p \to \pi X) \cong \frac{s}{\pi\left(m_X^2 - t\right)} \nu W_2(x) \frac{d\sigma}{d\hat{t}}(\pi q \to \pi q) \quad ,$$

$$x = x_{bj} = -t/\left(m_X^2 - t\right) \quad , \tag{3.5}$$

where $d\sigma/d\hat{t}\,(\pi q \to \pi q)$ is evaluated at $\hat{s} = xs$, $\hat{u} = xu$, $\hat{t} = t$. The derived cross section behaves as $x_T(1-x_T)^3/p_T^8$. Using $g^2/4\pi = 1\ \mathrm{GeV}^2$, this direct contribution is in fact smaller than the observed cross section (but it should become dominant in $d\sigma/d^3p\,(\pi p \to \pi X)$ at $x_T \gtrsim 0.6$).

Let us now try to predict the cross section for $pp \to \pi X$ for the various contributing CIM subprocesses. For completeness, we give the general formula for the contribution of subprocesses each parametrized as

$$\frac{d\sigma}{d\hat{t}}(ab \to Cd) = \frac{\pi D}{s^{N-T-U}(-t)^T(-u)^U} \tag{3.6}$$

to the inclusive cross section for $A+B \to C+X$: $(\epsilon = 1-x_R)$

$$\frac{d\sigma}{d^3p/E}(A+B \to C+X) = \sum_{ab \to Cd} \frac{(1-x_R)^F}{\left(p_T^2\right)^N}(1+x_R z)^{-F^+}(1-x_R z)^{-F^-} I \ . \tag{3.7}$$

The coefficient is

$$I = D f_{a/A} f_{b/B}\, 2^{F^+ + F^-} \frac{\Gamma(a+2)\,\Gamma(b+2)}{\Gamma(a+b+2)} J \tag{3.8}$$

where $J(z,\epsilon)$ is a slow function of $z = \cos\theta_{c.m.}$ and ϵ, and $J(z=0) = 1$. Here $xG_{a/A} \sim (1-x)^a$, $xG_{b/B} \sim (1-x)^b$, $F^+ = T+1+b-N$, and $F^- = U+1+a-N$. Typically, processes involving a fragmentation or decay process $a+b \to c+d$ with $c \to C+X$ are relatively suppressed because of the higher p_T of the subprocess, and these will be neglected for the simple and rough estimates given here.

Thus let us consider the contribution of the subprocesses $Mq \to K^+q$ to $pp \to K^+X$, summing over the possible contributing meson states (see Fig. 2a). Here $d\sigma/dt = (g^4/16\pi)\, s^{-1}u^{-3}$, so $D = (g^2/4\pi)^2$, N=4, T=0, U=3. We take $G_{M/p}(x) \sim (1-x)^5$ and estimate $f_{M/p} \sim f_{\bar{q}/p}/f^{val}_{\bar{q}/\pi} \sim .03/.10 \sim .3$, $f_{q/p} \sim 0.3$ (only q=u makes a sizable contribution). Note that the f's are the fraction of momentum carried by both valence and non-valence states. The sum over mesons includes K^+, K^{+*}, K^o, K^{o*}, etc.; hence $\Sigma f_{M/p} \sim 4\, f_{M/p} \sim 1.2$.

If we take $g^2/4\pi = 1.2\ \mathrm{GeV}^2$, as determined from exclusive processes, then Eq. (3.7) gives

$$\frac{d\sigma}{d^3p/E}(pp \to K^+X) \cong 1.9\, \frac{(1-x_R)^9}{p_T^8} \left[\frac{(1-x_R z)^{-5} + (1+x_R z)^{-5}}{2}\right] \tag{3.9}$$

in GeV units. This is the prediction for the "prompt" K^+, those which are created directly in the subprocess. We estimate that the contribution from decays, etc., would multiply (3.9) by ~2 or 3. The Chicago-Princeton data[9] at $z = \cos\theta_{c.m.} = 0$ fits

$$\frac{d\sigma}{d^3p/E}(pp \to K^+X) \sim 5.1\, \frac{(1-x_R)^9}{p_T^8} . \tag{3.10}$$

After accounting for other subprocesses (e.g., $q\bar{q} \to K^+\bar{M}$), this seems satisfactory agreement. The fact that $Ed\sigma/d^3p\,(pp \to \pi^+X)/Ed\sigma/d^3p\,(pp \to K^+X) \sim 2.2$ in the data can be accounted for from extra resonance decay contributions for the pion and extra diagrams such as $\sim d\pi^+ \to d\pi^+$.

In the case of K^- production, the counting rules predict that the dominant contribution at large x_R should be the $q\bar{q} \to K^-M$ subprocess (Fig. 2b). By crossing we obtain (ignoring spin factors)

$$\frac{d\sigma}{dt}(q\bar{q} \to M\bar{M}) = (g^4/16\pi)\, \hat{t}/\hat{s}^2\hat{u}^3 \tag{3.11}$$

and a=3, b=7, F=11. Using (3.7) we obtain

$$E\frac{d\sigma}{d^3p}(pp \to K^-X) = (0.1\, \langle d\rangle)\, \frac{(1-x_R)^{11}}{p_T^8} \left[\frac{(1-x_R z)^{-3}(1+x_R z)^{-2} + (z \to -z)}{2}\right] \tag{3.12}$$

where <d> is the number of contributing recoil meson states. The data are consistent with

$$\frac{Ed\sigma/d^3p\,(pp \to K^-X)}{Ed\sigma/d^3p\,(pp \to K^+X)} \sim 0.9\,(1-x_R)^2 \tag{3.13}$$

so we require $\langle d \rangle \sim 5$ to 10 to completely account for this ratio (this is not an unreasonable estimate for the total number of contributing meson states). The subprocess $K^- q \rightarrow K^- q$ gives a $(1-x_R)^{13}/p_T^8$ contribution but its normalization is hard to estimate.

It should be emphasized that these calculations are only approximate due to uncertainties in the effects of spin, color, the small variation of J and the transverse momentum integrations. The main point here is that to within factors of 2 or 3 we find that the CIM diagrams immediately and simply account for the normalization of the inclusive cross section given the known non-zero coupling of the hadronic state to its valence quark components.

We can also proceed to calculate the normalization of the baryon subprocesses. From the magnitude of elastic pp scattering and the proton structure function sum rules, we find a coupling strength $h^2/4\pi \sim 30\ \mathrm{GeV}^4$ for the effective proton $\rightarrow q + (qq)$ coupling (where the (qq) system is at relatively low mass):

$$\frac{d\sigma}{d\hat{t}}(B+q \rightarrow B+q) \sim \frac{h^4}{16\pi^2} \frac{1}{\hat{s}^2 \hat{t}^2 \hat{u}^2} \tag{3.14}$$

The subprocess $B+q \rightarrow p+q$ then gives the contribution ($z=0$)

$$\frac{d\sigma}{d^3p/E}(pp \rightarrow pX) \sim 120 \frac{(1-x_T)^7}{p_T^{12}} \quad , \tag{3.15}$$

where we have used the estimates $\Sigma f_{B/p} \sim 1.2$ and $\Sigma f_{q/p} \sim 0.5$. The CP data are consistent with this scaling behavior; the experimental coefficient is ~ 170.

We must also consider the direct $pq \rightarrow pq$ contribution. In fact, using (3.7) we find this gives $E d\sigma/d^3p\,(pp \rightarrow pX) \sim 200\ (1-x_T)^3\ x_T^4/p_T^{12}$, and thus exceeds the contribution of (3.15) for $x_T \gtrsim 0.45$. It will be interesting to see if a change in the $(1-x_T)$ power from $F=7$ to $F=3$ is observed at the higher x_T values. There is also the possibility of an additional $p_T^{-8}\ (1-x_T)^7$ contribution from the subprocess $q+q \rightarrow B+\bar{q}$ but presumably the coupling constant for such large p_T processes is of order $(g^2/4\pi)^2$ and thus gives negligible contributions until considerably larger p_T values.

We have also computed the contribution of the fusion process $q\bar{q} \to B\bar{B}$ for $pp \to \bar{p}X$ production, using the crossed ($s \longleftrightarrow t$) version of (3.14). This gives

$$\frac{Ed\sigma/d^3p\,(pp \to \bar{p}X)}{Ed\sigma/d^3p\,(pp \to pX)} \cong \frac{\langle d \rangle}{10}\,(1-x_T)^4 \tag{3.16}$$

where $\langle d \rangle$ is the number of opposite side baryon states. The value $\langle d \rangle \sim 3\text{-}5$ gives a consistent fit to experiment.

One can also work out in a similar way the various contributions to meson-induced processes. The subprocesses based on $Mq \to Mq$ are again predicted to dominate in the present x_T range and reasonable agreement is obtained with experiment. We also find that the formulae and normalizations are consistent with the exclusive-inclusive connection.

Finally, we note that we can readily predict the cross section for direct photon production simply by replacing the valence meson contribution in $Mq \to Mq$ by a photon to obtain $Mq \to \gamma q$. We predict

$$\frac{Ed\sigma/d^3p\,(pp \to \gamma X)}{Ed\sigma/d^3p\,(pp \to K^+X)} \sim \frac{2\alpha\bar{\lambda}_q^2}{g^2/4\pi}\,p_T^2 \tag{3.17}$$

at fixed x_T and $\theta_{c.m.}$. This gives $\gamma/\pi^0 \sim 0.005\,p_T^2$, or about 1/4 the value reported by Darriulat et al.[14] A similar estimate follows directly from the ratio of exclusive cross sections and crossing.

Finally, we note that our normalization estimate for the production of real photons can be extended to virtual photons, and it has been shown to agree with the data for massive lepton pair production in both its predicted magnitude and p_T dependence.[15]

IV. Correlations and High p_T Processes

One of the most important discriminants between models for high p_T production is the nature of the flow of the valence quantum numbers, momentum, and multiplicity produced in association with the high p_T trigger. The new preliminary ISR data from the British-French-Scandinavian group[5] gives a first look at the detailed effects associated with the quantum number of the trigger particle. The experiment utilizes the split field magnet facility combined with a wide angle spectrometer.

Figure 5 shows that the total momentum of charged particles on the same side and within one unit of rapidity y of the trigger particle (at $\theta_{c.m.} = 90^o$) increases very slowly with p_T for π^- and π^+ triggers, and not at all for K^-. In the CIM such behavior is expected since the trigger particle can be produced directly and alone in the subprocess or by low mass resonance decay. In models based on simple $qq \rightarrow qq$ scattering, extra momentum which scales with the trigger momentum is expected in the same side jet (although this effect could be reduced somewhat by transverse momentum fluctuations[11]). Furthermore, if the meson is produced as a fragment of a scattered valence u or d quark, then the greatest amount of same side momentum would have been expected in association with K^- than with K^+, π^+, or π^- triggers, just the opposite to what is seen!

A very dramatic feature of the preliminary BFS data is shown in Fig. 6. This shows the number of fast (p_T > 1.5 GeV/c) positive or negative particles per event in the away side jet ($|y| < 1$) opposite a $\pi^\pm$, $K^\pm$, or $p^\pm$ trigger at 90^o with $3 < p_T^{trig} < 4.5$ GeV/c. One sees that there is significantly more fast positives than negatives opposite a K^- trigger, an effect not seen for π^-, π^+ and K^+ triggers. This is a direct indication that

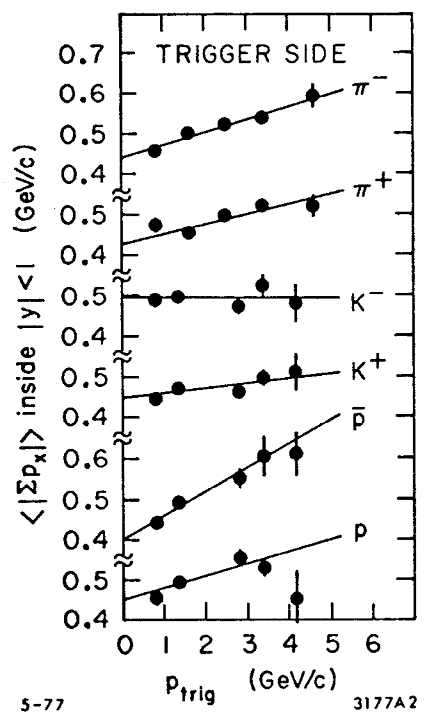

Fig. 5. The total momentum of particles along the 90^o trigger particle for various charged particles. From Ref. 5.

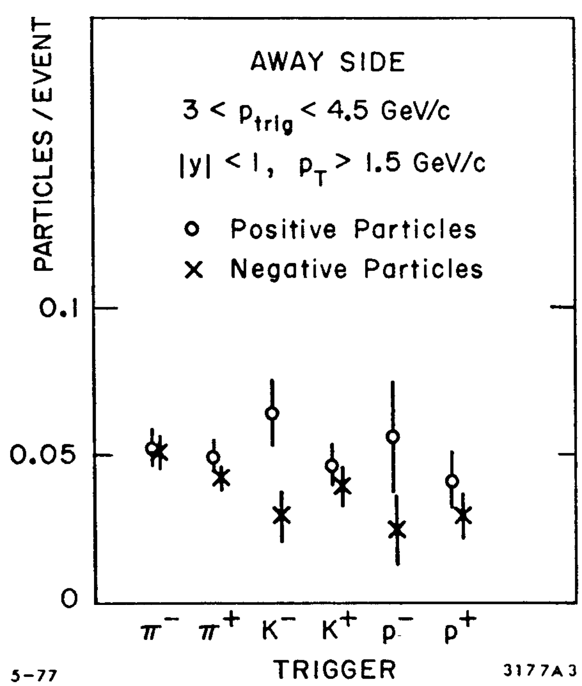

Fig. 6. Number of fast positive and negative particles on the side away from a 90^o trigger for various trigger types. From Ref. 5.

there is a strong quantum number correlation between the trigger and away side jet. Such a correlation is not expected in a $qq \rightarrow qq$ model since the away side quark is not correlated in any obvious way to the trigger side quark: the away side jet should have quantum numbers completely independent of the trigger.

In the CIM, this charge correlation for the K^- trigger is a natural prediction of the model. In the case of $\pi^\pm$ or K^+ triggers the leading subprocess contribution is $qM \rightarrow qM$ scattering which produces an away side jet corresponding to u or d quark fragmentation. The average away charge is then[16]:

$$\frac{1}{3}\left[2\left(\frac{2}{3} - n_Q\right) + \left(-\frac{1}{3} - n_Q\right)\right] = \frac{1}{3} - n_Q$$

where n_Q is the average charge of quarks in the sea ($\sim$0.07).[11,16] The $q\bar{q} \rightarrow M\bar{M}$ terms give additional contributions opposite in sign to the trigger charge. In the case of the K^- trigger, the dominant CIM subprocess is $q\bar{q} \rightarrow K^-M$ where M is either a positive or neutral strange mesonic system. The away side jet is thus predicted to have charge +2/3 on the average. (In both cases this average charge estimate would increase slightly if we assume that $G_{u/p} > 2G_{d/p}$ at large x.) These predictions for the mean charge of the jet can be compared with the BFS data of Fig. 7 which shows the presence of a strong positively charged system in the jet recoiling against the K^- trigger with $p_{trig} > 2.5$ GeV.

One possible modification of the quark-quark scattering description would be to introduce a strong $q\bar{q} \rightarrow s\bar{s}$ quark-antiquark annihilation contribution specifically for K^- production. Although this would yield a quantum number correlation between the away and same side jets, the mean charge of the $\bar{s}$ system would not yield a sufficiently strong positive charge on the away side. In addition, the magnitude of the

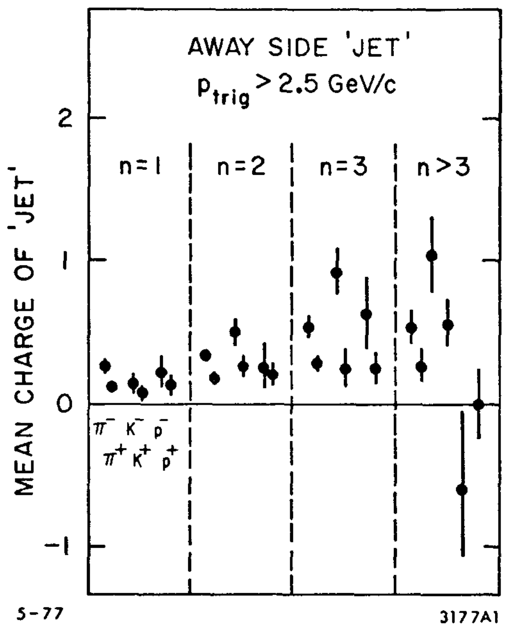

Fig. 7. Net mean charge of jet recoiling on the side away from the 90° trigger for various trigger types. See Ref. 5 for details of the definition of jet used here. (The data for $\bar{p}$ production is probably not statistically significant.) n is the number of charged particles seen in the jet.

$q\bar{q} \to s\bar{s}$ cross section is small if one crosses a form like $d\sigma/dt = s^{-1}t^{-3}$ for $qq \to qq$ to $q\bar{q} \to q\bar{q}$.

In general the distinctive quantum number flow of the CIM subprocesses can be used to predict a whole range of charge correlations associated with a high p_T trigger, corresponding to quark and multiquark jets in the fragmentation regions of the beam, target, or recoil jet. The quantum number retention[16] of these jets can also be tested in deep inelastic lepton scattering and the jets produced in the Drell-Yan process $A+B \to \ell^+\ell^-X$.

V. Jet Triggers and the CIM

Although the CIM appears to predict single particle data at large p_T very well, it is not clear that it can successfully account for the entire large jet trigger rate seen in the FNAL calorimeter experiment reported at this meeting.[17] The dominant jet-trigger contribution in the CIM comes from $Mq \to M'q$ subprocesses giving $d\sigma/d^3p_J/E_J \propto p_{TJ}^{-8}(1-x_{TJ})^9$. Since each meson in the pseudoscalar and vector SU(3) nonets contribute, and either the q or M system can provide the trigger, this gives a contribution at least 20 to 40 times the single meson rate at the same p_T. In addition there are contributions from other subprocesses $q\bar{q} \to M\bar{M}$, $M\bar{M} \to q\bar{q}$, $qB \to qB^*$, $q+qq \to M^*+B^*$, etc. which also provide jet triggers at high p_T. It may thus not be impossible to understand jet trigger cross sections which are 100 or more times larger than the single rate. However, one should also not rule out the possibility that because of the absence of the single-particle trigger bias, some jet trigger events could be due to QCD scale-invariant $qq \to qq$ scattering or processes involving gluon jets such as $gg \to gg$, $Mq \to gq$, etc. It will be crucial to have knowledge of the scaling behavior in p_T^J and x_T^J in order to begin to unravel these various contributions.

VI. Conclusions

As a summary it may be useful to contrast the basic assumptions and predictions of the CIM and quark-quark scattering models. The scaling laws of the CIM assume a basic scale-free theory, modulo logarithmic corrections characteristic of renormalizable perturbation theories.[18] Given that α_s is numerically small, the leading subprocess for $pp \to \pi^\pm$, $K^+ + X$ in the FNAL and ISR s, p_T range is then $qM \to qM$. The <u>calculated</u> subprocess cross section is $d\sigma/dt\,(qM \to qM) = \pi D/\hat{s}\,\hat{u}^3$ where the constant $D = (g^2/4\pi^2)$ is determined by the valence meson wavefunction normalization (see Section III). This form then correctly predicts the p_T, $\theta_{c.m.}$, x_T, and yields the normalization of the inclusive cross section.

In the approach of Feynman and Field,[17] Hwa et al.,[19] and others[20] sufficient scale-breaking is assumed so that literal quark-quark scattering can be taken to represent the large p_T subprocess. The form $d\sigma/dt = C/s\,t^3$ or $C/s\,u^3$ is then found to be a best simple fit to the data. (It should be remarked, though, that such a form, which corresponds to elementary spin 1/2 exchange in the t or u channel, is not natural for elastic qq scattering.) Both the $qM \to qM$ and $qq \to qq$ subprocesses correctly predict the $\sim(1-x_T)^9$ behavior of the inclusive cross section at fixed x_T and $\theta_{c.m.}$. Also, each model can account for the π^+/π^- and K^-/K^+ x_T dependence. Such ratios tend to be model-independent because one must pick up the same number of non-valence quarks somewhere in the inclusive process independent of the subprocess.

In the case of $pp \to pX$, the CP data[9] show a dramatic change in the p_T power to p_T^{-12} at fixed x_T and $\theta_{c.m.}$ (see Table I). In the CIM this is a natural consequence of the dominance of the $Bq \to Bq$ subprocesses, whose normalization is determined from $pp \to pp$ elastic scattering. (The calculated normalization of $qq \to B\bar{q}$ and $q + qq \to M + B$ turns out to be small in the present kinematic regime.) The CIM also predicts the observed $(1-x_T)^7$ behavior. In contrast, the $qq \to qq$ models, as interpreted by Feynman and Field, would lead to a p_T^{-8} $(1-x_T)^{11}$ behavior. One must then invoke new contributions such as the direct $pq \to pq$ subprocess (which gives an incorrect $(1-x_T)^3$ behavior) or perhaps $q + (qq) \to q + (qq)$ scattering.[11] New assumptions must then be introduced in the quark scattering model in order to calculate such additional processes. It then becomes doubly mysterious why processes such as $qM \to qM$ should not be considered for meson production.

In the case of the $\pi p \to \pi X$ cross sections, the $qq \to qq$ Feynman-Field model gives an excellent fit to the cross section provided $\nu W_2 \propto xG_{q/\pi}(x)$ goes to a finite constant ~ 0.15 at $x=1$. This assumption can be directly tested by checking for a flat non-vanishing Drell-Yan massive pair production cross section $d\sigma/dm^2\,dx_L(\pi p \to \ell^+\ell^- X)$ in the forward region, $x_L \sim 1$. In the CIM, the $\pi p \to \pi X$ cross section is computed from the subprocesses

$Mq \to Mq$, $q\bar{q} \to M\bar{M}$, as well as direct $\pi q \to \pi q$ scattering and is consistent with the data.

The CIM has the advantage of simultaneously predicting large p_T exclusive processes as well as inclusive cross sections in form as well as normalization. In the CIM one makes a natural progression from the proton form factor to the Compton amplitude to meson photoproduction to meson-proton scattering to inclusive cross sections, i.e. each case utilizing the same basic quark-exchange mechanism (see Fig. 4). In the case of the $qq \to qq$ model, there is no corresponding theory of exclusive reactions. For example, if $d\sigma/dt\,(qq \to qq) \sim C/st^3$ as determined by Feynman and Field[11] with $C = 2.3\,b \cdot GeV^6$ then one might expect a contribution $d\sigma/dt\,(pp \to pp) \sim C/st^3\, F_p^4(t)$. However, the predicted normalization is then four orders of magnitude smaller than experiment at $s = 20\ GeV^2$, $\theta_{c.m.} = \pi/2$. The angular dependence is also incompatible with the data, and the amplitude does not cross properly to $p\bar{p} \to p\bar{p}$.

Both the CIM and $qq \to qq$ models share the general features of hard scattering models for jet production angular correlations, etc. The predictions are in fact often indistinguishable since the same subprocess form is used. However, as we have emphasized here, the new preliminary charge correlation measurements of the BFS group,[5] particularly the K^- trigger data, implies quantum number correlations between the trigger and away side systems. Although such correlations are natural features of the CIM approach, it is not natural in a $qq \to qq$ model.

Finally, we again note that the CIM approach is not incompatible with the eventual dominance of a $\alpha_s^2 p_T^{-4}\,(1-x_T)$ scaling term from QCD in the single particle production cross section at very high p_T, probably well beyond $p_T = 8$ GeV. This $qq \to qq$ scattering contribution could, however, still make a significant $p_T^{-4}\,(1-x_T^J)^7$ contribution to the jet trigger cross section as presently measured.

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