MODELS FOR LEFT-RIGHT MULTIPLICITIES*

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ABSTRACT

In the context of various models for production processes, we examine what can be learned from data on the number of charged particles per event in the left and right c.m. hemispheres, leftright multiplicity distributions. The appropriate generating functional formalism is developed. We make explicit calculations of the generating function in both the multiperipheral and Mueller Regge approaches. We explain how left-right multiplicity distributions are used to separate the components of a hybrid, diffraction plus short-range-order, model. Questions concerning the factorization of the leading Regge singularity are also discussed.

(Submitted to Phys. Rev.)

^{*}Work performed under the auspices of the U. S. Atomic Energy Commission.

I. INTRODUCTION

Recent data from the National Accelerator Laboratory (NAL), from the CERN intersecting storage ring (ISR), and from the Serpukhov machine have stimulated an active interest in building models for charged particle multiplicity distributions as well as for one and two particle inclusive distributions.¹⁻³ The inclusive distributions are of interest theoretically because, as shown by Mueller,⁴ they give direct information about Regge singularities through a generalized optical theorem. Moreover, Mueller's constraints

$$= \int \rho_1(\mathbf{p}, \mathbf{s}) d^3\mathbf{p}/E$$
 (I. 1a)

$$\langle n(n-1) \rangle - \langle n \rangle^{2} = \int C_{2}(\underline{p}_{1}, \underline{p}_{2}, s) \frac{d^{3}\underline{p}_{1}}{\underline{E}_{1}} \frac{d^{3}\underline{p}_{2}}{\underline{E}_{2}} , \qquad (I. 1b)$$

etc., relate the moments of the multiplicity distribution to the inclusive distributions and correlation functions.

Multiplicity distributions have proved useful in testing general theoretical pictures of hadron scattering but since several different shapes of correlation function can have the same integral it is important to move toward more finely grained distributions. At the same time, for data of limited statistics there is an obvious advantage to studying integrated quantities. An interesting compromise between completely integrated and completely differential measurements consists of determining the number of particles (which in practice are charged) going left ("backward") and going right ("forward") in the center-of-mass system. Data of this type will be referred to as left-right multiplicity (LRM) distributions, $\sigma(n_L, n_R)$. From this data, more refined information about inclusive distributions is obtained than in Eq. (I. 1):

$$< n_{L} > = \int_{p_{L} < 0} \frac{d^{3} \underline{p}}{\underline{E}} \rho_{1}(\underline{p}, s)$$
 (I. 2a)

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$$<\mathbf{n}_{\mathrm{L}}\mathbf{n}_{\mathrm{R}}> - <\mathbf{n}_{\mathrm{L}}> <\mathbf{n}_{\mathrm{R}}> = \int_{\mathbf{p}_{\mathrm{L}1}<0} \frac{\mathrm{d}^{3}\underline{\mathbf{p}}_{1}}{\underline{\mathbf{E}}_{1}} \int_{\mathbf{p}_{\mathrm{L}2}>0} \frac{\mathrm{d}^{3}\underline{\mathbf{p}}_{2}}{\underline{\mathbf{E}}_{2}} C_{2}(\underline{\mathbf{p}}_{1},\underline{\mathbf{p}}_{2},\mathbf{s})$$
 (I. 2b)

etc.

Left-right multiplicity distributions are extremely easy to obtain in bubble chamber experiments and many types of counter experiments. What we would like to show here is that LRM already contain enough information to answer some crucial questions about the dynamics of production processes.^{5,6} In particular, they provide a series of quantitative tests for the factorization of the leading Regge singularity.⁷ These tests will be reviewed in Section II of this paper along with a discussion of the meaning and content of factorization. For completeness, an introduction to the generating function formalism is included.⁸ This formalism is convenient for the discussion of these multiplicity distributions, as well as to the closely related question of charge transfer across c.m. hemispheres. Differences between charged multiplicities and total multiplicities are discussed in this section as well.

To supply concrete examples of what can be learned from LRM, we turn in Section III to model calculations. Two main classes of models are distinguished: those constructed directly for inclusive distributions and those constructed directly for exclusive distributions. Inclusive models use the Mueller-Regge ideas; exclusive models are constructed, for the most part, using the multiperipheral picture, although adaptations of the multiperipheral picture to include diffraction are discussed as well. In addition, the Bjorken-Feynman-Wilson fluid analogy is here considered as a simple reformulation and extension of the multiperipheral picture. In discussions we shall loosely include all three of these exclusive classes of models in the "multiperipheral picture". Separate

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from this multiperipheral picture is the fragmentation picture. We discuss some simple predictions of this approach, but since recent data seems to discredit it as an explanation for data in the central region, ⁹ we will not present detailed calculations.

In Section IV, we apply what we have learned to the available data on leftright multiplicities and to the related quantity charge transfer. The claim⁷ that data on $\sigma(n_L, tot)$ at 205 GeV/c provide direct evidence for a diffractive, or fragmentation mechanism is examined in more detail, and the possibility of separating out the fraction of events which correspond to diffractive fragmentation is discussed. Data on $\langle n_L \rangle_R$ vs n_R at current energies is shown to favor a situation in which the left-right cross sections do not factorize, although the situation is not clear cut. Finally, we show data on charge transfer in pp collisions can give some information on the size of the corrections to factorization.

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II. GENERAL CONSIDERATIONS AND GENERATING FUNCTION TECHNIQUES

The generating function formalism has proved to be an exceptionally valuable tool for dealing with multiplicity distributions in hadronic production processes.⁸ In this section we introduce the notation that we will use for the generating function, and define in terms of this function what we mean by inclusive, exclusive, or "mixed" quantities. Many of the results in this section are contained in a previous paper.⁷ The treatment here is made more general so that in addition to left-right multiplicities, the related question of charge transfer can be discussed. The various experimental quantities we discuss are indicated in Table I.

A. Notation

It may make our later definitions clearer if we give first a simple example of a generating function in terms of the cross section $\sigma(n, s)$ for producing n particles. Given the complete set $\{\sigma(n, s)\}$ at a given c.m. energy \sqrt{s} , the function of the parameter z,

$$Q(z,s) \equiv \sum_{n} \sigma(n,s) z^{n}$$
, (II. 1a)

is formed. Clearly all the information contained in $\{\sigma(n, s)\}$ is contained in the function Q(z, s). In fact, $\sigma(n, s)$ can be recovered using

$$\sigma(n, s) = \frac{1}{n!} Q^{(n)}(0, s)$$
 (II. 1b)

where $Q^{(n)}(z,s) = (\partial^n / \partial z^n) Q(z,s)$. The advantage of the generating function is that inclusive moments can also be simply obtained

$$f_{m}(s) = \frac{1}{m!} \left(\frac{\partial^{m}}{\partial z^{m}} \right) \ln Q(z, s) \Big|_{z=1}$$
(II. 1c)

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where it is easy to check that

$$f_1(s) = \langle n \rangle$$

 $f_2(s) = \langle n(n-1) \rangle - \langle n \rangle^2$, (II. 1d)

etc. The definition we will use for <u>exclusive</u> quantities will correspond to the situation (II. 1b), where Q and its derivatives are evaluated at z=0; <u>inclusive</u> quantities arise from (II. 1c), where Q and its derivatives are evaluated at z=1. If we distinguish between particles either by quantum number or by binning events in different kinematic regions so that more than one type of z will be present in Q, then the possibility of "mixed" quantities exists, where some z's are unity and the remainder are zero.

A typical example of a mixed quantity is the cross section $\sigma(n_{ch}, s)$ where n_{ch} represents the number of charged particles and a sum is taken over the unseen neutrals. Such mixed quantities are also frequently called "semi-inclusive".

These expressions for the integrated quantities $\sigma(n, s)$ and $f_m(s)$, correspond to the usual definitions. The $\sigma(n, s)$ are integrals over the exclusive differential cross section $d\sigma_n/(d^3\underline{p}_1/\underline{E}_1 \dots d^3\underline{p}_n/\underline{E}_n)$ for producing exactly n particles (treated here as identical). The $f_m(s)$ are integrals over the inclusive correlation functions $C_m(\underline{p}_1, \dots, \underline{p}_m)$, where

$$\sigma(s) = \sum_{n} \sigma(n, s)$$

$$\rho_1(\underline{\mathbf{p}}_1, \mathbf{s}) = \frac{1}{\sigma(\mathbf{s})} \sum_{\mathbf{n}} \frac{\mathrm{d}\sigma_{\mathbf{n}}}{\mathrm{d}^3 \underline{\mathbf{p}}_1 / \mathbf{E}_1}$$
(II. 1e)

$$\rho_2(\underline{\mathbf{p}}_1, \underline{\mathbf{p}}_2, \mathbf{s}) = \frac{1}{\sigma(\mathbf{s})} \sum_{\mathbf{n}} \frac{\frac{d\sigma_{\mathbf{n}}}{\mathbf{p}_1}}{\frac{d^3 \underline{\mathbf{p}}_1}{\mathbf{E}_1} \frac{d^3 \underline{\mathbf{p}}_2}{\mathbf{E}_2}}$$

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$$C_2(\underline{p}_1, \underline{p}_2, s) = \rho_2(\underline{p}_1, \underline{p}_2, s) - \rho_1(\underline{p}_1, s) \rho_1(\underline{p}_2, s)$$
 (II. 1f)

etc.

It is an interesting exercise to take one step from completely integrated quantities toward differential cross sections by imagining an experiment with exactly two momentum bins. In this case we will assume that each bin consists of an entire hemisphere in the c.m. system centered around the beam or target directions. To discuss the production of charged particles in the left and right c.m. hemispheres, we form the generating function

$$Q(z_{+L}, z_{-L}, z_{+R}, z_{-L}, s) =$$

$$= \sum \sigma(n_{+L}, n_{-L}, n_{+R}, n_{-L}, s) z_{+L}^{n_{+L}} z_{-L}^{n_{-L}} z_{+R}^{n_{+R}} z_{-R}^{n_{-R}}$$
(II. 2)

where the summation here, and in what follows, goes over all values of those n's appearing in the summand. The notation is self-evident; $n_{\pm L}$ is the number of "±" particles in the left ("backward") and $n_{\pm R}$ are in the right ("forward") hemisphere. We assume a sum over the unseen neutral particles: They can be explicitly included, if needed, by introducing the extra factors $z_{0L}^{n_{0L}} z_{0R}^{n_{0R}}$ with the completely exclusive cross sections $\sigma(n_{\pm L}, n_{-L}, n_{0L}; n_{\pm R}, n_{-R}, n_{0L})$ inside the summation (II. 2).

Because of charge conservation, the four n's in (II.2) are not independent. A smaller set is

$$n_{L} = n_{+L} + n_{-L}$$
(II. 3a)

$$n_{\rm R} = n_{+\rm R} + n_{-\rm R}$$
(II. 3b)
$$u = \frac{1}{2} \left((n_{+\rm R} - n_{-\rm R}) - (n_{+\rm L} - n_{-\rm L}) \right)$$
$$- \frac{1}{2} (q_{\rm a} - q_{\rm b})$$
(II. 3c)

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where n_L is the number of charged particles in the left hemisphere, n_R the number in the right hemisphere, and u is the net charge transfer from the right to the left. Here q_a and q_b are the charges of the initial particles. Charge conservation eliminates one n_i from the original set through the constraint

$$q_a + q_b = (n_{+L} + n_{+R}) - (n_{-L} + n_{-R})$$
 . (II. 3d)

Thus (II. 3c) can be written equivalently as

$$u = n_{+R} - n_{-R} - q_a$$

= $n_{-L} - n_{+L} + q_b$. (II. 3c')

The set (II. 3) suggests introducing a new set of variables into the generating function

$$z_{L} = (z_{+L} z_{-L})^{1/2}$$

$$z_{R} = (z_{+R} z_{-R})^{1/2}$$

$$x = \left(\frac{z_{+R} z_{-R}}{z_{+L} z_{-R}}\right)^{1/2}$$

$$x' = \left(\frac{z_{+L} z_{+R}}{z_{-L} z_{-R}}\right)^{1/2}$$
(II. 4)

Note that $xx' = z_{+R}/z_{-R}$ depends only on "right" quantities and $x'/x = z_{+L}/z_{-L}$ depends only on "left" quantities. In terms of Eq. (II.4), (II.2) becomes

$$Q(z_{+L}, z_{-L}, z_{+R}, z_{-L}, s) \equiv (xx')^{1/2} q_{a} (x'/x)^{1/2} q_{b} Q(z_{L}, z_{R}, x) =$$

$$= (xx')^{1/2} q_{a} (x'/x)^{1/2} q_{b} \sum \sigma(n_{L}, n_{R}, u) z_{L}^{n_{L}} z_{R}^{n_{R}} x^{u}$$
(II. 5)

where u ranges over both positive and negative values. The charge constraint implicit here has been discussed by B. Webber.⁸ Certain simple symmetry

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properties of Q follow immediately. If the incident charge $q_a + q_b$ is even (odd) then $n_L + n_R$ is even (odd), and hence

$$Q(-z_{L}, -z_{R}, x) = (-)^{q_{a}+q_{b}} Q(z_{L}, z_{R}, x)$$
 (II. 6)

We will return later to this symmetry and its implications for the factorization of the generating function. Another simple property is obtained if

$$\sigma(\mathbf{n}_{\mathrm{L}},\mathbf{n}_{\mathrm{R}},-\mathbf{u}) = \sigma(\mathbf{n}_{\mathrm{L}},\mathbf{n}_{\mathrm{R}},\mathbf{u}) \quad , \tag{II.7}$$

as would be the case, for example, in pp reactions. Consequently the generating function would then satisfy

$$Q(z_{L}, z_{R}, 1/x) = Q(z_{L}, z_{R}, x)$$
 (II.8)

In what follows we will often specialize to the case where a summation has been made over charge transfer

$$Q(z_{L}, z_{R}) = Q(z_{L}, z_{R}, 1)$$
 (II. 9)

We want to now examine the consequences for the generating function (II. 5) of certain general theoretical notions. In particular, we will assume that the cross sections are power bounded and that there exists exclusive cluster decomposition in the rapidities of the produced particles. ¹⁴ When we adopt these assumptions, it becomes convenient to consider explicitly the power of s with which the generating function (II. 2) grows for z_i 's not unity:

$$Q(z_{+L}, z_{-L}, z_{+R}, z_{-R}, s) \equiv \exp\{p(z_{+L}, z_{-L}, z_{+R}, z_{-R}, s)Y\}$$
(II. 10a)
$$Y \equiv \ln s/s_0$$
(II. 10b)

where p will in general be a function of s. Asymptotically in s, we expect p to have a limiting value for each choice of its arguments z_i . The expansion

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coefficients of p for $z_{\pm L}^{}, z_{\pm R}^{}$ around unity, $p_{ijkl}^{}$ defined in

$$p = \sum p_{ijkl} \frac{(z_{+L}-1)^{i}}{i!} \frac{(z_{-L}-1)^{j}}{j!} \frac{(z_{+R}-1)^{k}}{k!} \frac{(z_{-R}-1)^{l}}{l!}$$
(II. 11)

are related directly to integrals of inclusive correlation functions (recall Eqs. (II. 1) and the associated discussion). A few simple examples are as follows:

$$Y p_{1000} = \langle n_{+L}(s) \rangle = \int_{-Y/2}^{0} dy \rho_{1}(Y)$$
 (II. 12)

$$Y p_{1010} = \langle n_{+L} n_{+R} \rangle - \langle n_{+L} \rangle \langle n_{+R} \rangle$$
$$= \int_{-Y/2}^{0} dy_1 \int_{0}^{Y/2} dy_2 C_2^{++}(y_1, y_2)$$
(II. 13)

etc. Within the context of the <u>short range order</u> correlation picture, in which for example, $C_2(y_1, y_2) \rightarrow e^{-|y_1 - y_2|/L}$ as $|y_1 - y_2| \rightarrow \infty$, a sharp distinction exists between the coefficients p_{ijkl} which involve particles all in one hemisphere, and those involving particles in both hemispheres. If the correlation functions become integrable in the rapidity differences at high energy, then it is true that for $n_L \neq 0$, $n_R \neq 0$,⁷

$$\int_{-Y/2}^{0} dy_1 \dots dy_{n_L} \int_{0}^{Y/2} dy'_1 \dots dy'_{n_R} C_{n_L + n_R} \begin{pmatrix} y_1 \dots y_{n_L}, y'_1 \dots y'_{n_R} \end{pmatrix}$$

 \sim constant

(II. 14)

as s (or equivalently Y) $\rightarrow \infty$. The implication is that whenever i+j > 1 and k+l > 1, then

$$p_{iikl} \rightarrow 0$$

as $s \rightarrow \infty$. In contrast

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in general are nonzero, and finite as $s \rightarrow \infty$. Therefore, if the short range order hypothesis is correct, we can obtain a meaningful separation of p,

$$p(z_{+L}, z_{-L}, z_{+R}, z_{-R}, s) = a + p_L(z_{+L}, z_{-L}, s) + p_R(z_{+R}, z_{-R}, s) + p_c(z_{+L}, z_{-L}, z_{+R}, z_{-R}, s) ,$$
(II. 15)

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and we are guaranteed that $p_c \rightarrow 0$ asymptotically. Charge conservation implies further that as $s \rightarrow \infty$,

$$\begin{split} &\exp\left(Y\,p_{L}(z_{+L},z_{-L})\right) \sim (x'/x)^{1/2}\,q_{b}\,f^{n}(z_{L}) \\ &\exp\left(Y\,p_{R}(z_{+R},z_{-R})\right) \sim (xx')^{1/2}\,q_{a}\,f^{n}(z_{R}) \end{split}$$

(read " f^n " as "function of"), and so asymptotically

$$p_{L}(z_{+L}, z_{-L}, s) \rightarrow p_{L}(z_{L})$$

$$p_{R}(z_{+R}, z_{-R}, s) \rightarrow p_{R}(z_{R})$$
(II. 16)

The separation (II. 15) can also be written, therefore, as

$$p(z_{+L}, z_{-L}, z_{+R}, z_{-R}, s) = a + p_L(z_L) + p_R(z_R) + p_c(z_L, z_R, x, x')$$
 (II. 15')

(suppressing the s dependence), where again $p_c = O(1/Y)$ asymptotically.

To see that (II. 15') is indeed what we would have expected for short range models, but is not expected in general, we remind the reader of the energy behavior, for pp reactions, of the mean square fluctuation of the net charge transfer,

$$\ll u^2 \gg = Y \frac{\partial^2}{\partial x^2} p \Big|_{x=1}$$
 (II. 17)

(This notation and the results below come from Ref. 6.) In models with short range order $\ll u^2 >>$ is expected to be constant, as implied by (II. 15'), whereas

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in fragmentation models $\ll u^2 \gg \sim \sqrt{s}$. This can be interpreted as implying that fragmentation models have nonfactorizable singularities which will show up as strong energy dependence of moments of left going particles as a function of right going particles. We will discuss this point further in Section IV. D where data on mean squared charge transfer is presented.

To help understand the way in which (II. 15') depends on the short range order hypothesis, a useful exercise is to compare the predictions for the two particle correlation function $C_2(y_1, y_2)$ in short range order, fragmentation and hybrid models. In the c.m. system there are four regions of interest to us: (1) Quadrant I, where both particles are in the right hemisphere $(y_1 > 0, y_2 > 0)$ the integral over this region is p_{02} Y (we are here suppressing information concerning charges); (2) Quadrant II $(y_1 > 0, y_2 < 0)$ whose integral is p_{11} Y; (3) Quadrant III $(y_1 < 0, y_2 < 0)$ whose integral is p_{20} Y; and (4) Quadrant IV $(y_1 < 0, y_2 > 0)$ whose integral is p_{11} Y (Y = $\ln s/s_0$).

Short range order. In short range order models, $C_2(y_1, y_2)$ depends only on y_1-y_2 . The maximum height of C_2 becomes constant asymptotically, with the growing part of the integral coming mainly from regions I and II. Ignoring energy momentum conservation, as these effects don't change the logarithmic term, we obtain for Quadrant II (or IV)

$$p_{11}Y = \frac{1}{2} \int_{-Y/2}^{0} dy_1 \int_{0}^{Y/2} dy_2 C_2(y_1, y_2)$$

= constant

(II. 18)

and for Quadrant I (or III)

$$p_{02}Y = \int_{-Y/2}^{0} \int_{-Y/2}^{0} dy_1 dy_2 C_2(y_1, y_2) \propto \frac{Y}{2}$$
(II. 19)

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These results coincide with (II. 14) which lead to factorization of the generating function.

<u>Fragmentation</u>. Contrast this with the fragmentation model¹⁵ predictions for C_2 : the only nonzero contributions to C_2 are in Quadrants I and III; and the maximum height (which is at $y_1 = y_2 = 0$) increases like \sqrt{s} as $\rightarrow \infty$. The behavior in regions II and IV is less dramatic, and results primarily from fragmentation events created in one hemisphere yielding products which end up in the opposite hemisphere. As noted by Chou and Yang,⁶ these events can account for a conzero amount of charge transfer which grows like \sqrt{s} asymptotically. If the size (area) of this region is denoted by ϵ^2 , we expect

$$p_{11} Y \propto \epsilon^2 \sqrt{s}$$
 (II. 20)

and so we predict $p_c(z_L, z_R, x)$ (Eq. (II. 15') to be extremely large asymptotically, contrary to the short range order prediction. In fragmentation models we definitely do not expect factorization of the generating function.

<u>The hybrid model</u>. A hybridization of the above two models is that there remains a small fraction of events in σ_{inel} which are due to diffractive fragmentation, the exact fraction depending upon the energy.¹⁶ If these events could be removed by subtraction, then the remainder would be of short range order. The correlation function, as usually defined

$$C_{2}(y_{1}, y_{2}) = \frac{1}{\sigma_{in}} \frac{d^{2}\sigma}{dy_{1}dy_{2}} - \frac{1}{\sigma_{in}^{2}} \frac{d\sigma}{dy_{1}} \frac{d\sigma}{dy_{2}}$$
(II. 21)

will behave differently then in short range order models due to $\sigma_{\rm in}$ having two components

$$\sigma_{\rm in} = \sigma_{\rm D} + \sigma_{\rm M}, \qquad \sigma_{\rm D} \ll \sigma_{\rm M}$$
 (II. 22)

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In the central region, we assume $d\sigma_D/dy_1$, and $d^2\sigma_D/dy_1dy_2 = 0$ so that

$$C_{2}(y_{1}, y_{2}) \approx \left(1 - \frac{\sigma_{D}}{\sigma_{M}}\right) \left[\frac{1}{\sigma_{M}} \frac{d^{2}\sigma_{M}}{dy_{1} dy_{2}} - \frac{1}{\sigma_{M}^{2}} \frac{d\sigma_{M}}{dy_{1}} \frac{d\sigma_{M}}{dy_{2}}\right] + \frac{\sigma_{D}}{\sigma_{M}} \left[\frac{1}{\sigma_{M}^{2}} \frac{d\sigma_{M}}{dy_{1}} \frac{d\sigma_{M}}{dy_{2}}\right]$$
(II. 23)

to first order in σ_D / σ_M . For large values of $|y_1 - y_2|$, the first bracket tends to zero whereas the second bracket is constant (in the plateau region). The integrals over this plateau region then give asymptotically the results

$$p_{20} Y \propto \frac{\sigma_{\rm D}}{\sigma_{\rm M}} \left(\frac{Y}{2}\right)^2$$
(II. 24)
$$p_{11} Y \propto \frac{\sigma_{\rm D}}{\sigma_{\rm M}} \left(\frac{Y}{2}\right)^2 .$$

The violation of factorization is determined by the (expected) small value of $\sigma_{\rm D}/\sigma_{\rm M}$. This can be made more reasonable by noting that this model could result from two terms, each of which factorizes

$$\sigma(n_{L}, n_{R}) = a^{D}(n_{L}) a^{D}(n_{R}) + a^{M}(n_{L}) a^{M}(n_{R}) . \qquad (II. 25)$$

The factorization property (II. 15') is therefore seen to test the form of correlation functions. What is needed to obtain factorization is the integrability of correlation functions (both exclusive and inclusive) in rapidity differences. This type of integrability is not found in all models. The resolution of whether it exists or not will ultimately come from experiment so we now turn to the problem of experimental tests for factorization in left-right multiplicities.

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B. Factorization and Factorization Tests

From the above results we recall that a short range order picture implies certain asymptotic properties (II.15') for the generating function. These can be written in the form

$$Q(z_{L}, z_{R}, x) \sim g(z_{L}, z_{R}, x) \exp\{p_{L}(z_{L})Y\} \exp\{p_{R}(z_{R})Y\}$$
. (II. 26)

Specifically, this means that

$$\sigma(\mathbf{n}_{\mathrm{L}},\mathbf{n}_{\mathrm{R}}) = \frac{1}{\mathbf{n}_{\mathrm{L}}!} \left(\frac{\partial}{\partial_{\mathrm{L}}}^{\mathbf{n}_{\mathrm{L}}} \right) \frac{1}{\mathbf{n}_{\mathrm{R}}!} \left(\frac{\partial}{\partial_{\mathrm{R}}}^{\mathbf{n}_{\mathrm{R}}} \right) Q(\mathbf{z}_{\mathrm{L}},\mathbf{z}_{\mathrm{R}},\mathbf{1}) \Big|_{\mathbf{z}_{\mathrm{L}}=\mathbf{z}_{\mathrm{R}}=0}$$
$$\sim \mathbf{a}_{\mathrm{L}}(\mathbf{n}_{\mathrm{L}}) \mathbf{a}_{\mathrm{R}}(\mathbf{n}_{\mathrm{R}}) [\mathbf{1} + O(\mathbf{1}/\mathbf{Y})] , \qquad (II.27)$$

the <u>left-right</u> cross section factorizes to leading order in $Y = \ln s/s_0$. This factorization property can be tested experimentally by studying the energy dependence of ratios of left-right multiplicity cross sections and comparing with the SRO prediction

$$\frac{\sigma(N, M) \sigma(K, L)}{\sigma(N, L) \sigma(K, M)} = 1 + O(1/Y) \quad . \tag{II.28}$$

Another set of tests of $(\Pi, 26)$ involves the set of averages

$$<^{n}_{R}>_{L} vs_{L} n_{L}$$
 (II. 29)
 $<^{n}_{R}(^{n}_{R}-1)>_{L} vs_{L} n_{L}$,

etc. Some discussion of the tests (II. 29) and recent NAL and ISR data has been presented recently.⁷ The lack of correlation expected in (II. 27) implies

$$< n_{R} > L \sim cY \quad as \quad Y \to \infty$$
 (II.30)

where c is independent of n_L . As we will discuss in more detail in Section IV the data on this test are not conclusive. In particular, the important question of whether (II. 30) is true for $n_L=1$ has not yet been answered with current data.

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The energy behavior of the cross sections provides yet another test of the factorization (II. 26). In particular, the following cross sections, in pp collisions are related if $p_c \rightarrow 0$ in (II. 15'):

$$\sigma(\mathbf{n}_{\mathrm{L}}, \mathrm{tot}) = \sum_{\mathbf{n}_{\mathrm{R}}} \sigma(\mathbf{n}_{\mathrm{L}}, \mathbf{n}_{\mathrm{R}}) \simeq (\mathrm{s/s}_{0})^{\mathbf{a}+\mathbf{p}_{\mathrm{L}}(0)+\mathbf{p}_{\mathrm{R}}(1)}$$

$$\sigma(\mathbf{n}_{\mathrm{L}}, \mathbf{n}_{\mathrm{R}}) \simeq (\mathrm{s/s}_{0})^{\mathbf{a}+\mathbf{p}_{\mathrm{L}}(0)+\mathbf{p}_{\mathrm{R}}(0)} \qquad (\mathrm{II.31})$$

$$\sigma(\mathrm{tot}) \simeq (\mathrm{s/s}_{0})^{\mathbf{a}+\mathbf{p}_{\mathrm{L}}(1)+\mathbf{p}_{\mathrm{R}}(1)}$$

ignoring logarithmic factors. As an example, if the exclusive and total cross sections behave asymptotically like $(s/s_0)^{-1}$ and $(s/s_0)^0$ respectively, then $\sigma(n_L, tot) \sim (s/s_0)^{-1/2}$: i.e., the mixed quantity has an energy exponent half way between the exclusive and inclusive quantity exponents. The vanishing of $p_c(z_L = 0, z_R = 0)$ in (II. 15') asymptotically can be tested by looking at the energy behavior of the quantities (for a general reaction now)

$$R(N, M) = \frac{\ln \sigma(N, \text{tot}) + \ln \sigma(\text{tot}, M)}{\ln \sigma(N, M) + \ln \sigma(\text{tot})}$$
(II. 32)

which are expected to approach 1 + O(1/Y) as $Y \to \infty$.

It is important to note that these tests (II. 27), (II. 30) and (II. 32) are to be made at several energies so we can study the Y dependence of the correction terms.

At finite energy, the effect of charge conservation is to destroy the factorization expected asymptotically. In terms of the generating function for pp scattering as an example, charge conservation imposes that the number of charges in the final state is even and greater than two. This means, for example, that the generating function has the symmetry (II.6)

$$Q(-z_{L}, -z_{R}) = Q(z_{L}, z_{R})$$
 (II. 33)

From this it is easy to illustrate that even if the cross sections factorize,

$$\sigma(\mathbf{n}_{\mathrm{L}}, \mathbf{n}_{\mathrm{R}}) = a(\mathbf{n}_{\mathrm{L}}) a(\mathbf{n}_{\mathrm{L}}) , \qquad (\mathrm{II.34})$$

the generating function $Q(z_L, z_R)$ does not, except asymptotically. A particularly instructive case is

$$a(n_{L}) = \frac{m^{n_{L}}}{n_{L}!}$$
(II. 35)

which leads to the generating function

$$Q(z_{L}, z_{R}) = \sum_{n_{R}+n_{L}} \sum_{\text{even, } \geq 2} a(n_{L}) a(n_{R}) z_{L}^{n_{L}} z_{R}^{n_{R}} \qquad n_{L}, n_{R} \neq 0$$

= sinh (m $\rm z^{}_L)$ sinh (m $\rm z^{}_R)$ + cosh (m $\rm z^{}_L)$ cosh (m $\rm z^{}_R)$ - 1

$$= \cosh \left[m(z_{\rm L} + z_{\rm R}) \right] - 1 \tag{II.36}$$

which obviously satisfies the symmetry property (II.33), but does not factorize. At high energies with $m \sim c \ln s/s_0$, the behavior of $\ln Q(z_L, z_R)$ has the expected (factorized) form q.v. (II.15')

$$\ln Q(z_L, z_R) \sim c(z_L + z_R) Y + const. + O(1/Y)$$
. (II. 37)

C. What has to Factorize

In a short range order picture it appears well known that the leading J-plane singularity is an isolated pole and that the above factorization constraints involving left-right multiplicities are among the consequences of the existence of this leading pole. It is important to keep in mind that several currently popular models for production processes do not possess short range order. In particular

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we have seen that the hybrid diffraction plus multiperipheral model which has been the subject of much current theoretical speculation¹⁶ has the feature that the correlation functions $C_n(y_1, \ldots, y_n)$ approach nonzero constants in the limit that $|y_i - y_i| \rightarrow \infty$ for any pair of rapidities y_i , y_i in the central, or plateau region. It is then possible for the integrated correlations in (II. 14) to grow like $(Y/2)^{n_L+n_R}$, both in the case that either n_L or $n_R = 0$, and for n_L , $n_R \neq 0$. There is no a priori reason to expect factorization from these models. It is also clear that fragmentation models need not factorize, as discussed in Section II.A, due to the long range correlations which can occur in them. Thus, the questions discussed here are not trivial and deserve experimental investigation. Indeed there are important theoretical reasons to believe that the leading J-plane singularity, even if it is not a simple pole, may factorize anyway. If this is true, we then have (for example) powerful restrictions on the type of hybrid models that are possible. The vanishing of $p_c(z_L, z_R, x)$ which is automatic in the short range order picture may occur in complicated ways in other pictures.

One of the early results of Regge pole theory was the observation that Regge pole residues must factorize.¹⁷ The impression has been gained apparently by a sizeable segment of the physics community that factorization is a property of poles only. In fact it has been shown by Kawai¹⁸ that partial wave unitarity implies that any infinity in the partial wave amplitude must reflect certain factorization properties. Since this result bears on our investigations here, we briefly review Kawai's argument.

Suppose we have m coupled channels of scalar particles with the signatured, reduced partial wave amplitude $B_{ii}(J,t)$ in the "t" channel. Time reversal

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invariance makes B_{ij} symmetric

$$B_{ij}(J,t) = B_{ij}(J,t)$$
 (II. 38)

If the amplitude is analytic and properly bounded so the unitarity relation can be continued to complex J, we can write t channel elastic unitarity as

$$B_{ij}(J,t) - \widetilde{B}_{ij}(J,t) = 2 i\rho B_{ia}(J,t) \widetilde{B}_{ja}(J,t)$$
(II. 39)

where channel "a" is the elastic channel and \tilde{B}_{ij} is obtained from B_{ij} by a continuation around the elastic cut. Using (II. 38) and (II. 39) it is easy to verify that $B_{aa}\tilde{B}_{ja} = B_{ja}\tilde{B}_{aa}$; substituting $B_{ja}\tilde{B}_{aa}/B_{aa}$ for \tilde{B}_{ja} in (II. 39) we obtain

$$B_{aa}(B_{ij} - \tilde{B}_{ij}) = 2 i\rho B_{ia} B_{ja} \tilde{B}_{aa} . \qquad (II.40)$$

If the partial wave amplitude B_{aa} is singular,

$$\lim_{J \to \alpha} |B_{aa}(J,t)| = \infty , \qquad (II.41)$$

we see that (II. 39) implies

$$\lim_{J \to \alpha} \widetilde{B}_{aa}(J,t) = (2 i_{\rho})^{-1} \quad . \tag{II.42}$$

Thus, the partial wave amplitude B_{ij} has from (II. 40) the approximate behavior as $J \rightarrow \alpha$

$$B_{ij} \simeq \frac{B_{ia}B_{ja}}{B_{aa}}$$

since \widetilde{B}_{ij} is not expected to be singular, whereas B_{ia} and B_{ja} are. It is then easy to see that the singular behavior (II.41) as $J \rightarrow \alpha$, leads to

$$B_{ij}(J,t) B_{kl}(J,t) \sim B_{ik}(J,t) B_{jl}(J,t)$$
 (II.43)

We thus have factorization, and are allowed to write

$$B_{ij}(J, t) = \gamma_i(t) \gamma_j(t) f(J, t) + \hat{b}_{ij}(J, t)$$
(II.44)

where

$$\lim_{J \to \alpha} |f(J,t)| = \infty$$
 (II. 45)

$$\lim_{\mathbf{J}\to\alpha} \frac{\mathbf{b}_{\mathbf{ij}}(\mathbf{J},\mathbf{t})}{\mathbf{f}(\mathbf{J},\mathbf{t})} = 0 \quad . \tag{II.46}$$

This factorization property depends upon unitarity and on the fact that the partial wave amplitude is infinite at $J = \alpha$. It does not depend on this singularity being a simple pole. The most singular part of a <u>hard</u> branch point must factorize as well; a <u>soft</u> branch point such as $(J-\alpha)^{1/2}$ where the partial wave amplitude vanishes need not factorize.

The factorization of reduced partial wave amplitude in coupled two-body channels does not, of course, lead directly to the factorization of left-right multiplicity distributions. It is possible to construct a chain of inference which would lead to this result, however. If (II.43) could be extended to a situation where the indices i, j ranged over quasi-two-body as well as two-body channels we could make contact with production processes. Such an extension is plausible but has not been proved. Such a proof would require considerable theoretical effort because of technical complications involving many-particle-partial wave amplitudes such as complex helicity.¹⁹

If we <u>assume</u> the extension can be made and allow the indices i, j to represent the internal quantum numbers of a many-body-final state, then we can obtain a constraint on production cross sections when the leading J-plane singularity obeys (II. 44) - (II. 46). The constraint is that the production cross section factorizes whenever any rapidity gap becomes large. That is, if we order the rapidities $y_i < y_{i+1} < y_{i+2}$

$$d\sigma(\mathbf{y}_{1}\cdots\mathbf{y}_{n}) |_{\mathbf{y}_{i}} - \mathbf{y}_{i+1}| \rightarrow \infty \qquad \mathbf{a}(\mathbf{y}_{1}\cdots\mathbf{y}_{i}) |_{\mathbf{x}(\mathbf{y}_{i+1}\cdots\mathbf{y}_{n})} + O((\mathbf{y}_{i} - \mathbf{y}_{i+1})^{-\epsilon}) \qquad (II.47)$$

This is still not sufficient to guarantee the factorization of integrated cross sections (II. 28) unless the correction term is integrable in $(y_i - y_{i+1})$, that is $\epsilon > 1$. (II. 48)

It is therefore interesting to look at differential cross sections directly and test this exclusive cluster decomposition¹³ (II. 47) as well as the simpler constraints on the integrated cross section (II. 28).

Using the same assumption of a factorizable hard singularity in the discontinuity of a Mueller graph we can conclude that inclusive cross sections satisfy (II. 47) as well as exclusive ones.

III. MODEL CALCULATIONS

Some of the underlying dynamical features of production processes which can be learned by studying left-right multiplicities will be illustrated here by means of some explicit calculations in simple models. Two pictures of hadron dynamics will be stressed; an inclusive one as best represented by Mueller-Regge models (MRM), and an exclusive picture, as best represented by multiperipheral models (MPM). It appears well known to some that the multiperipheral picture always leads to, or is expected to always lead to an equivalent Mueller-Regge inclusive description. Recent work has taken the first step in establishing an explicit connection between these two pictures.²⁰ Although we believe the physics of MRM and MPM to be the same, it is nevertheless instructive to treat them separately in that certain aspects of data can be more -conveniently described from one picture than the other.

A. Multiperipheral Picture

The essence of the MPM can be obtained by making some crude simplifying approximations. The approximations consist of neglecting the transverse momentum dependence of the matrix element and assuming all invariant momentum transfers along the multiperipheral chain to be small so that we can reformulate the model in one (longitudinal) dimension. ¹⁰ In this reformulation, the strong ordering limit is taken so that the only energy-momentum constraints are those fixing the total allowed interval of rapidity. If the matrix element depends only on momentum transfers, we shall call such a model the nearest neighbor multiperipheral model (NNMM). ²¹ For n secondary particles produced, the differential cross section in the NNMM has the form

$$\frac{d^{n}\sigma}{dy_{1}\cdots dy_{n}} = g_{a}g_{b}g^{n}e^{(2\alpha-2)Y}f(y_{1}-y_{a})f(y_{2}-y_{1})\cdots f(y_{n}-y_{n-1})f(y_{b}-y_{n}) \quad (III. 1)$$

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where g is the Reggeon-Reggeon-particle coupling squared; y_a and y_b are the initial rapidities of the beam and target $(Y = y_b - y_a)$; $\left[e^{2\alpha y} f(y)\right]$ is the multiperipheral kernel; and α is the leading trajectory in the kernel. In the Chew-Pignotti model, ¹⁶ the kernel is $(s_{ii+1}/s_0)^{2\alpha} \sim e^{2\alpha (y_{i+1}-y_i)}$ so that f(y) = 1. The function f(y) will in general not be unity due to the existence of daughter trajectories, and may be replaced by a matrix if more than one input trajectory is considered, or if isospin or charge is to be included explicitly.²²

With the particles and rapidities labelled as in Fig. 1, the cross section for n_R particles in the region $[y_a, Y_R + y_a]$ and n_L particles in the region $[y_b - Y_L, y_b]$ (where $Y_L + Y_R = Y$) is in this approximation

$$\sigma(n_{\rm L}, n_{\rm R}) = g_{\rm a}g_{\rm b}g^{\rm n} e^{(2\alpha-2)Y} \int_{\substack{II \\ i=1}}^{n_{\rm L}+1} dx_{i}^{\dagger} \delta(\Sigma x_{i}^{\dagger} - Y_{\rm L}) \int_{\substack{II \\ j=1}}^{n_{\rm R}+1} dx_{j} \delta(\Sigma x_{j}^{\dagger} - Y_{\rm R}) \prod_{i=1}^{n_{\rm L}} f(x_{i}^{\dagger}) \prod_{j=1}^{n_{\rm R}} f(x_{j}) f(x_{n_{\rm L}+1}^{\dagger} + x_{n_{\rm R}+1}) \quad (III. 2)$$

Here $x_i^{!}$ and $x_j^{}$ are rapidity differences, and $Y_L^{}$, $Y_R^{}$ will be treated as arbitrary. The generating function which determines the left-right multiplicities,

$$Q(z_{L}, z_{R}) = \sum_{n_{L}n_{R}} \sigma(n_{L}, n_{R}) z^{n_{L}} z^{n_{R}},$$
 (III. 3)

will here refer to total rather than charged multiplicities for convenience. It is straightforward to include charge in the NNMM but the equations become cumbersome. To evaluate (III. 3), we first take a double Laplace transform of

$$\begin{aligned} \mathbf{Q}_{\mathbf{n_L}\mathbf{n_R}}(\mathbf{Y}_{\mathbf{L}},\mathbf{Y}_{\mathbf{R}}) &\equiv \sigma(\mathbf{n_L},\mathbf{n_R}) / \left[\sigma_0(\mathbf{Y}) \mathbf{g}^{\mathbf{n_L}+\mathbf{n_R}} \right] \\ & \left(\sigma_0(\mathbf{Y}) \equiv \mathbf{g}_{\mathbf{a}} \mathbf{g}_{\mathbf{b}} \mathbf{e}^{(2\alpha-2)\mathbf{Y}} \right) , \end{aligned} \tag{III.4}$$

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which is defined as

$$Q_{n_{L}n_{R}}(\alpha_{L},\alpha_{R}) \equiv \int_{0}^{\infty} dY_{L} e^{-\alpha_{L}Y_{L}} \int_{0}^{\infty} dY_{R} e^{-\alpha_{R}Y_{R}} Q_{n_{L}n_{R}}(Y_{L},Y_{R})$$
$$= \overline{f}(\alpha_{L})^{n_{L}} \overline{f}(\alpha_{R})^{n_{R}} \frac{\overline{f}(\alpha_{L}) - \overline{f}(\alpha_{R})}{\alpha_{R} - \alpha_{L}}$$
(III.5)

where $f(\alpha)$ is the Laplace transform of f(y). Summing $Q_{n_L n_R}(z_L g)^{n_L}(z_R g)^{n_R}$ we obtain

$$Q(\alpha_{\rm L}, \alpha_{\rm R}) \equiv \Sigma \ Q_{\rm n_{\rm L}n_{\rm R}} (z_{\rm L}g)^{\rm n_{\rm L}} (z_{\rm R}g)^{\rm n_{\rm R}}$$
$$= \frac{1}{1 - z_{\rm L}g \ \bar{f}(\alpha_{\rm L})} \frac{\bar{f}(\alpha_{\rm L}) - \bar{f}(\alpha_{\rm R})}{\alpha_{\rm R}^{-\alpha_{\rm L}}} \frac{1}{1 - z_{\rm R}g \ \bar{f}(\alpha_{\rm R})} (III.6)$$

In the usual case where $f(\alpha)$ is analytic for Re $\alpha > 0$ and has the behavior $f(\alpha) \sim 1/\alpha$ as $\alpha \to 0$, the vanishing of $1 - zg \bar{f}(\alpha)$ determines the leading asymptotic behavior of the inverse transform of (III.6). Denoting this solution by p(z)

1 = zg f[p(z)], (III. 7)

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we obtain for (III.3)

$$Q(z_{L}, z_{R}) \sim g_{a}g_{b} e^{(2\alpha-2)Y} \frac{z_{L}g - z_{R}g}{p(z_{L}) - p(z_{R})} \frac{\rho(z_{L})}{z_{L}g} e^{(z_{L})Y_{L}} \frac{\rho(z_{R})}{z_{R}g} e^{p(z_{R})Y_{R}}$$
(III.8)

where

$$\rho(z) \equiv zp'(z) = -\frac{1}{zg f'(p)}$$
 (III.9)

We note that the asymptotic behavior of σ_{inel} is

$$\sigma_{\text{inel}} = g_{a}g_{b} e^{(2\alpha-2)Y} \frac{\rho(z)}{zg} e^{p(z)Y}$$
(III. 10)

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which can be obtained directly from (III. 1) or from (III.8) with $z_L = z_R = z$. The correlation between left and right depends upon the factor

 $(z_L-z_R)/[p(z_L)-p(z_R)]$, and we can compute explicitly the function $p_c(z_L, z_R)$ which appears in (II.15). With the constraint $p_c(z_L, 1) = 0$ used in defining p_c , we obtain

$$p_{c}(z_{L}, z_{R}) \approx \frac{1}{Y} \left\{ \ln \left(\frac{gz_{L} - gz_{R}}{p(z_{L}) - p(z_{R})} \right) - \ln \left(\frac{g - gz_{R}}{p(1) - p(z_{R})} \right) - \ln \left(\frac{g - gz_{R}}{p(1) - p(z_{R})} \right) - \ln \left(\frac{g - gz_{L}}{p(1) - p(z_{L})} \right) + \ln \frac{g}{\rho(1)} \right\}$$
(III. 11)

To proceed further we would need a specific model for the kernel f(y). Two simple examples may serve to illustrate what can result.

The first example is the well known Chew Pignotti model whose kernel is f(y) = 1 and $\alpha = 1/2$. For this model $\overline{f}(\alpha) = 1/\alpha$, p(z) = zg and $(z_Lg - z_Rg)/[p(z_L) - p(z_R)] = 1$. Therefore there are no left-right correlations, and $p_c(z_L, z_R) = 0$.²⁴

The second example is the "hard core" model of Chew and Snider.²³ In this model $f(y) = \theta(y-b)$, and $\bar{f}(\alpha) = e^{-\alpha b}/\alpha$. The correlation can be noted directly by computing the Laplace inverse of (III.5)

$$Q_{n_{L}n_{R}}(\alpha_{L},\alpha_{R}) = \frac{e^{-\alpha_{L}n_{L}b}}{\alpha_{L}} \frac{e^{-\alpha_{R}n_{R}b}}{\alpha_{R}} \frac{\frac{e^{-\alpha_{R}b}}{\alpha_{R}} - \frac{e^{-\alpha_{L}b}}{\alpha_{L}}}{\alpha_{L} - \alpha_{R}} \quad .$$
(III. 12)

By standard techniques, we invert to obtain

$$Q_{n_{L}n_{R}}(Y_{L}, Y_{R}) = \frac{(Y_{R} - (n_{R}^{+1})b)^{n_{R}}}{n_{R}!} \frac{(Y_{L} - n_{L}b)^{n_{L}}}{n_{L}!} + \frac{\sum_{k=0}^{n_{L}} \left\{ \frac{(Y_{R} - n_{R}b)^{n_{L}^{+n_{R}^{-k}}}}{(n_{L}^{+}n_{R}^{-k})!} \frac{(Y_{L} - (n_{L}^{+1})b)^{k}}{k!} - \frac{\left[\frac{(Y_{R} - (n_{R}^{+1})b)^{n_{L}^{+n_{R}^{-k}}}}{(n_{L}^{+}n_{R}^{-k})!} \frac{(Y_{L} - n_{L}b)^{k}}{k!} \right\}.$$
(III. 13)

It can be checked that for $Y_L = Y_R = 1/2 Y$, $Q_{n_L n_R}$ is symmetric in n_L and n_R , and further that

$$\sum_{n_{L}+n_{R}} Q_{n_{L}n_{R}} = \frac{(Y - (n+1)b)^{n}}{n!} , \qquad (III.14)$$

the result for the total multiplicity distribution in the "hard core" NNMM. There are clearly left-right correlations in (III. 13) which result from the sum of components. To compute $p_c(z_L, z_R)$, we return to the asymptotic expression (III. 11). The power exponent p(z) is obtained by solving (III. 7)

$$1 = zg \frac{e^{-p(z)b}}{p(z)}$$
 (III. 15)

which will clearly result in a nonzero expression for $p_c(z_L, z_R)$. We leave this hard core NNMM for the time being since without additional input, it is incapable of describing existing multiplicity data. The basic problem is that it predicts a distribution narrower than Poisson whereas data favor a distribution (ignoring diffraction effects) slightly broader than Poisson.

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B. The KUH Model

The fact that there exists a mathematical isomorphism between the multiperipheral model as formulated by DeTar¹⁰ and the nearest neighbor dimensional fluid has led to a great deal of speculation that the fluid analogy might be a way of extending our understanding of production processes.¹³

Kac, Uhlenbeck and Hemmer¹² (KUH) have discussed in detail a one dimension fluid where the pair interaction consists of a hard core repulsion (as in the Chew-Snider model) plus an exponential attraction

$$V_{\text{pair}}(y) = \gamma a e^{-\gamma y} \theta(y-b) \quad . \tag{III. 16}$$

The KUH fluid is of particular interest for hadronic processes because the corresponding generating function Q can be evaluated explicitly, and because the potential (III. 16) contains a weak interaction which is not of the nearest neighbor, or multiperipheral type but might represent long range interactions. There is some evidence that this type of physical mechanism does play a role in data and so the KUH model will be taken as an example of a model with complicated final state interactions.

KUH show that in the limit of $\gamma \to 0$ and 'a' fixed (i.e., $\int_0^\infty dy \ V_{\text{pair}}(y)$ fixed) the resultant energy exponent of $\sigma_{\text{tot}}(z) = s^{p(z)}$ satisfies

$$p(z) = \frac{\rho(z)}{1 - \rho(z)b} - a_{\rho}(z)^{2}$$
(III. 17)

where $\rho(z) = zp'(z)$ as before. This equation, known as the "Van der Waals equation of state" has been applied to NAL and Serpukhov multiplicity data with some success. ^{3e} Even if the KUH model were to have no physical connection to high energy data, the formal isomorphism still makes calculations in this model instructive.

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The technique which makes the interaction (III. 16) soluble is an identity which reduces the calculation of the cross section involving non-nearest neighbor interactions to that of a nearest neighbor cross section. The price one pays is the introduction of n auxiliary integrations (for calculating σ_n). The essential point (the reader should refer to Ref. 12 for details) is that instead of (III. 1), the KUH model has the form

$$\frac{d^{n}Q_{n}}{dy_{1}\cdots dy_{n}} = \int_{-\infty}^{\infty} dz_{1} \int_{-\infty}^{\infty} dz_{n} \Psi_{e}(z_{1}) \prod_{i=1}^{n-1} \left[K(y_{i+1} - y_{i} | z_{i}, z_{i+1}) \right] \theta(y_{i+1} - y_{i} - b) \Psi_{e}(z_{n}) \quad .$$
(III. 18)

The explicit form of the kernel, K, in (III. 18) is known, but is not important to us here. Once (III. 18) is Laplace transformed, the important quantities are the eigenvalues and eigenfunctions of the integral equation

$$\int_{-\infty}^{\infty} dz' K(\alpha | z, z') \Psi(z') = \lambda \Psi(z) \quad . \tag{III.19}$$

It is known¹² that K_{α} is a symmetric, positive definite Hilbert Schmidt kernel which has a discrete set of <u>positive</u> eigenvalues $\lambda_i(\alpha)$ whose maximum is $\lambda_0(\alpha)$ and which go to zero as $i \rightarrow \infty$. The kernel therefore has a uniformly and absolutely convergent expansion in terms of its complete orthonormal set of eigenfunctions.

$$K_{\alpha}(z, z') = \sum_{\lambda} \lambda_{i}(\alpha) \Psi_{i}(z, \alpha) \Psi_{i}(z', \alpha)$$
(III. 20)

These facts allow the evaluation of $Q_n(\alpha)$:

$$Q_{n}(\alpha) = \sum_{i=0}^{\infty} A_{i}(\alpha)^{2} \lambda_{i}^{n+1}(\alpha)$$
(III. 21)

$$A_{i}(\alpha) = \int_{-\infty}^{\infty} dz \Psi_{i}(z) \Psi_{e}(z) \quad . \tag{III. 22}$$

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The calculation now proceeds as with the NNMM. The total cross section is

$$Q(\alpha) = \sum_{i=0}^{\infty} \frac{A_i(\alpha)^2 \lambda_i(\alpha)}{1 - zg \lambda_i(\alpha)}$$
(III. 23)

and the behavior of Q(Y) is determined by the leading singularities of (III. 23). If the interaction (III. 16) has a nonzero γ , the asymptotic behavior is determined by the vanishing of

$$1 - zg \lambda_0(\alpha) = 0$$
 (III. 24)

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where $\lambda_0(\alpha)$ is analytic for Re $\alpha > 0$ and has the behavior $e^{-\alpha b}/\alpha$ as $\alpha \to \infty$. In the limit as $\gamma \to 0$, however, we have a feature not present in the simple NNMM in that an infinite number of terms in (III.23) contribute because the eigenvalues begin bunching around the function $\omega(\alpha)$, the maximum eigenvalue. It is shown in Ref. 12 that in this case $\lambda_i(\alpha) \simeq \omega(\alpha)$ and from the completeness of the eigenfunctions

$$Q(\alpha) = \frac{\omega(\alpha)}{1 - zg \ \omega(\alpha)}$$
(III. 25)

The vanishing of this denominator leads to the behavior

$$Q(Y) = \frac{\rho(z)}{zg} e^{p(z)Y}$$
(III. 26)

for the total cross section, where p(z) and $\rho(z)$ satisfy (III. 17).

The steps leading to $Q(\alpha_L, \alpha_R)$ for the left-right multiplicities are analogous to those leading to (III.6), but for the subtleties involving the eigenvalues and eigenfunctions. No new complications arise which did not occur in obtaining (III.23), and so we quote the result

$$Q(\alpha_{\rm L}, \alpha_{\rm R}) = \sum_{i, j=0}^{\infty} \frac{A_i(\alpha_{\rm L})}{1 - z_{\rm L}g \lambda_i(\alpha_{\rm L})} \frac{A_j(\alpha_{\rm R})}{1 - z_{\rm R}g \lambda_i(\alpha_{\rm R})}$$
$$< i, \alpha_{\rm L} | j, \alpha_{\rm R} > \frac{\lambda_j(\alpha_{\rm R}) - \lambda_i(\alpha_{\rm L})}{(\alpha_{\rm L} - \alpha_{\rm R})}$$
(III. 27)

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where

$$\langle \mathbf{i}, \alpha_{\mathrm{L}} | \mathbf{j}, \alpha_{\mathrm{R}} \rangle = \int_{-\infty}^{\infty} \mathrm{dx} \, \Psi_{\mathbf{i}}(\mathbf{x}, \alpha_{\mathrm{L}}) \, \Psi_{\mathbf{j}}(\mathbf{x}, \alpha_{\mathrm{R}}) \quad .$$
 (III. 28)

In the limit that $\gamma \rightarrow 0$, (III. 27) simplifies due to the orthogonality of the eigenfunctions:

$$Q(\alpha_{\rm L}, \alpha_{\rm R}) \simeq \frac{1}{1 - z_{\rm L} g \omega(\alpha_{\rm L})} \frac{\omega(\alpha_{\rm R}) - \omega(\alpha_{\rm L})}{\alpha_{\rm L} - \alpha_{\rm R}} \frac{1}{1 - z_{\rm R} g \omega(\alpha_{\rm R})}$$
(III. 29)

The result for the inverse is thus formally the same as (III.8) except now p(z) satisfies a different type of equation not derivable in a MPM.

To investigate what might happen in such a model in more detail, we shall make use in Section IV of a simplifying approximation, known in statistical mechanics as the mean field approximation (MFA). The approximation consists of replacing the long range part of the non-nearest neighbor interactions by its average value. In the KUH model, the limit $\gamma \rightarrow 0$ is meant to simulate a long range potential. For n produced particles, the mean field approximation replaces (III. 18) by

$$\frac{d^{n}Q_{n}}{dy_{1} \dots dy_{n}} = \prod_{i=1}^{n-1} \theta (y_{i+1} - y_{i} - b) e^{an^{2}/Y}$$
(III. 30)

since the average potential is the number of pairs times $\frac{1}{Y} \int_0^Y dy V(y)$, or an^2/Y . But (III. 30) is almost in the form of the Chew-Snider model, the "hard core" NNMM. The result is that the cross section $Q_{n_L}{}^{n_R}(Y_L, Y_R)$ is given by (III. 13) times the extra weight $e^{an^2/Y}$. In the next section, this model is compared with certain aspects of the data. For completeness we give here the parameter values used, and the prescription used for including charged and neutral particles. A fit to data^{3e} with this model favors

$$a = \frac{27}{64}$$
, $b = \frac{1}{8}$, $g = 4 e^{-7/4}$; $2\alpha - 2 = -1$ (III. 31)

which corresponds to a Van der Waals fluid at the critical point. The distinction between neutral and charged particles is accomplished by introducing a binomial weight factor [as in Ref. 3e]

$$\binom{{}^{n}_{R}}{{}^{n}_{R}}\binom{1}{3}{}^{n}_{R}^{-n}_{R}^{ch} \qquad \binom{n}{R}^{ch}_{R}$$

for right movers, and a corresponding weight factor for left movers, into the expression (III.13). Calling $Q_{n_R n_L}$ in (III.13), $Q_{n_R n_L}^{h.c.}$, the present model is

$$\sigma_{\substack{n_{R}^{ch}n_{L}^{ch}n_{R}^{n}n_{L}^{0}}} = e^{-Y} \begin{pmatrix} n_{R}^{ch} + n_{R}^{0} \\ n_{R}^{ch} \end{pmatrix} \begin{pmatrix} \underline{g} \\ 3 \end{pmatrix}^{n_{R}^{0}} \begin{pmatrix} \underline{2g} \\ 3 \end{pmatrix}^{n_{R}^{ch}} \\ \begin{pmatrix} \underline{2g} \\ 3 \end{pmatrix}^{n_{R}^{ch}} \\ \begin{pmatrix} \underline{2g} \\ 3 \end{pmatrix}^{n_{R}^{ch}} \end{pmatrix} \begin{pmatrix} \underline{g} \\ 3 \end{pmatrix}^{n_{L}^{ch}} e^{an^{2}/Y} Q_{n_{L}^{n_{R}^{ch}}}^{h.c.} (Y_{L}, Y_{R})$$

(III. 32)

with (III. 31) giving the parameter values.

C. Pomeron plus Regge Model

In applications of the types of model discussed in Sections III. A – III. B, it is usually assumed that diffractive effects are small. For this reason the elastic process is excluded in these analyses. To include diffraction explicitly in the multiperipheral picture, two exchanges have to be included in the input; a Pomeranchuk pole and a secondary meson pole representing an average of all lower lying trajectories. ¹⁶ If a matrix formalism is adopted, the formulas are similar in form to those in the NNMM of Section III. A. In place of (III. 1) we have

$$\frac{d^{n}\sigma_{n}}{dy_{1}\cdots dy_{n}} = G_{a}^{T} F(y_{1}-y_{a}) GF(y_{2}-y_{1}) G\cdots F(y_{n}-y_{n-1}) GF(y_{b}-y_{n})G_{b}$$
(III.33)

where the kernel is now an n × n diagonal matrix if there are n poles exchanged. (We do not take out the leading behavior in the definition of F as was done in (III. 1). Also we include in F the flux and phase space factor e^{-2y} .) G is an n × n matrix giving the couplings between the ith Regge pole and the jth Regge pole; G_a and G_b are column vectors of dimension n which give the coupling of the external lines to the exchanged Regge poles. Analogous to (III.5) we have the double transform of $Q_{n_L n_R}(Y_L, Y_R)$ [here $\sigma_0(Y) \equiv 1$]

$$Q_{n_{L}n_{R}}(\alpha_{L},\alpha_{R}) = G_{a}^{T}(\overline{F}(\alpha_{R})G)^{R} \frac{\overline{F}(\alpha_{L}) - \overline{F}(\alpha_{R})}{\alpha_{R} - \alpha_{L}} (G\overline{F}(\alpha_{L}))^{n_{L}} G_{b}.$$
(III. 34)

Summing over all $n_L^{n,n}$, the analog to (III.6) is

$$Q(\alpha_{\rm L}, \alpha_{\rm R}) = G_{\rm a}^{\rm T} (\mathbf{I} - \mathbf{z}_{\rm L} \mathbf{\overline{F}}(\alpha_{\rm L})G)^{-1} \frac{\mathbf{\overline{F}}(\alpha_{\rm L}) - \mathbf{\overline{F}}(\alpha_{\rm R})}{\alpha_{\rm R} - \alpha_{\rm L}} (\mathbf{I} - \mathbf{z}_{\rm R}G\mathbf{\overline{F}}(\alpha_{\rm R}))^{-1} G_{\rm b}.$$
(III. 35)

For comparison, we note that the Laplace transform of the total cross section is

$$Q(\alpha) = G_a^T \overline{F}(\alpha) \left[I - z G \overline{F}(\alpha) \right]^{-1} G_b \quad . \tag{III.36}$$

We see that the leading behavior of the cross sections are governed by the zeroes of

$$\det (\mathbf{I} - z \, \mathbf{G} \,\overline{\mathbf{F}}(\alpha)) = 0 \quad . \tag{III. 37}$$

In the general case, the leading behavior of the generating function $Q(Y_L, Y_R)$ will again be factorizable.

As mentioned in III. A, the above formalism is general and can be used to treat the exchange of several secondary trajectories, or the effects of isospin conservation. ²² Our interest here is in the two-pole exchange model, Pomeron plus Reggeon. The treatment has to be somewhat roundabout since an input pole at $\alpha = 1$, treated in the multiperipheral formalism, is well known to generate an output pole larger than unity. ²⁵ The difficulty remains in the inclusive picture and is connected with the vanishing of the triple Pomeron coupling or the decoupling of the Pomeron from physical processes. ²⁶ The standard treatment of these problems for phenomenology is to allow the Pomeron to be exchanged only once on the argument that the Pomeron coupling is very small. For the coupling and kernel matrix we have

$$\overline{F}(\alpha) = \begin{pmatrix} \frac{1}{\alpha} & 0 \\ & \\ 0 & \frac{1}{\alpha+1} \end{pmatrix}; \qquad G = \begin{pmatrix} 0 & \gamma \\ & \\ \gamma & g \end{pmatrix}$$
(III. 38)

assuming the Pomeron to be at unity and the secondary exchange to be at onehalf. These intercepts are taken to be only approximations and might be varied. For calculating multiplicities, the relevant quantity is

$$\left(\overline{\mathbf{F}}(\alpha)\mathbf{G}\right)^{\mathbf{n}} \cong \begin{pmatrix} \delta_{\mathbf{n}0} , & \frac{\gamma}{\alpha} \left(\frac{\mathbf{g}}{\alpha+1}\right)^{\mathbf{n}-1} \theta\left(\mathbf{n}-1\right) \\ & \\ \frac{\gamma}{\alpha+1} \left(\frac{\mathbf{g}}{\alpha+1}\right)^{\mathbf{n}-1} \theta\left(\mathbf{n}-1\right) , & \frac{\mathbf{g}}{\alpha+1} \left(\frac{\mathbf{g}}{\alpha+1}\right)^{\mathbf{n}-1} \end{pmatrix}$$
(III. 39)

After some algebra, and neglecting terms of order γ^2 and higher, we find (III. 34) has the form

The significance of these terms and a comparison of this type of model with data has been done by Snider.²⁷ We summarize briefly the physics of (III. 40). The first two terms are the elastic scattering given by Pomeron exchange $(\gamma_e^a \gamma_e^b)$ and the multiperipheral contribution to $\sigma(n_L, n_R)$ ($g_a g_b$). The latter is here given by the product of two Poisson distributions. The last two terms represent low and high mass diffraction. For example the terms proportional to $\gamma_e^b g_a$ consist of (1) Pomeron exchange extending across the left hemisphere to a low mass state in the right hemisphere — low mass diffraction; (2) Pomeron exchange which remains in the left hemisphere to a high mass state — high mass diffraction. We note that the inverse transform of these diffraction parts

$$\frac{1}{2\pi i} \int_{c-i\infty}^{c+i\infty} d\alpha \left[e^{\alpha Y} \right] \frac{\gamma g^{n-1} \theta (n-1)}{\alpha (\alpha+1)^n} = \gamma g^{n-1} \left[1 - e^{-Y} \sum_{k=0}^{n-1} \frac{Y^k}{k!} \right]$$
(III. 41)

which is an incomplete gamma function, has been discussed previously with regard to multiplicity data.²⁸

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The factorization properties of (III.40) are most easily seen by writing the equation in the form

$$Q_{n_{L}n_{R}}(\alpha_{L},\alpha_{R}) = A_{n_{R}}^{a}(\alpha_{R}) A_{n_{L}}^{b}(\alpha_{L}) + B_{n_{R}}^{a}(\alpha_{R}) B_{n_{L}}^{b}(\alpha_{L})$$
(III. 42a)

$$A_{n}^{i}(\alpha) = \gamma_{e}^{i} \frac{\delta_{n0}}{\alpha} + g_{i} \gamma \frac{g^{n-1} \theta(n-1)}{\alpha(\alpha+1)^{n}}$$
(III.42b)

$$B_{n}^{i}(\alpha) = \frac{g_{i}g^{n}}{(\alpha+1)^{n+1}} + \frac{\gamma_{e}^{i}\gamma g^{n-1}\theta(n-1)}{\alpha(\alpha+1)^{n}} \quad . \tag{III.42c}$$

Both terms in (III. 42a) give leading behavior in the approximation that $(\gamma Y)^2$ is negligible, and so the generating function $Q(z_L, z_R)$, which will have the form

$$Q(z_{L}, z_{R}) = A^{a}(z_{R}) A^{b}(z_{L}) + B^{a}(z_{R}) B^{b}(z_{L})$$
 (III. 43)

will not factorize. This lack of consistency with the starting assumptions is clearly related to our being forced for other reasons to keep only one Pomeron exchange in the approximation. To compare the basic features of this picture with data we simplify (III. 40) somewhat while retaining the basic form (III. 42). This will be explained in more detail in Section IV.

D. Inclusive Picture

It is interesting to counterbalance the above microscopic, or exclusive, models for left-right multiplicities with a simple Mueller-Regge model for the inclusive multiplicity moments $f_{n_L n_R}$. We choose for this purpose the model of Frazer, Peccei, Pinsky and Tan (FPPT).^{3f} It is assumed that the inclusive correlation functions are dominated by their values in the central region and further that these values can be well approximated by the Regge form even down to small rapidity differences. The diagrams considered are shown in Fig. 2. Formally the contributions of Fig. 2 have the MPM form (III. 33) for

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 \mathbf{Q}_n and (III.34) for $\mathbf{Q}_{n_{\rm L}n_{\rm R}}$ with G and F 2×2 matrices

$$G = \begin{pmatrix} g_{PP} & g_{PM} \\ g_{PM} & g_{MM} \end{pmatrix}; \qquad \overline{F}(\alpha) = \begin{pmatrix} \frac{1}{\alpha} & 0 \\ 0 & \frac{1}{\alpha + \frac{1}{2}} \end{pmatrix}$$
(III. 44)

and the external couplings two dimensional column vectors

$$G_{a} = G_{b} = \begin{pmatrix} 1 \\ 0 \end{pmatrix} .$$
 (III. 45)

We assume the meson trajectory is at $\alpha = \frac{1}{2}$, the Pomeron is $\alpha = 1$, and the particles are in the central region so that the external meson coupling can be ignored. To compute the integrated correlation functions f_n , we use the identities

$$\sum_{n=0}^{\infty} Q_n (z-1)^n = \exp \left[\sum_{n=1}^{+\infty} f_n \frac{(z-1)^n}{n!} \right]$$
(III. 46)

$$\sum_{n_{\rm L}, n_{\rm R}=0}^{\infty} Q_{n_{\rm L}n_{\rm R}} (z_{\rm L}^{-1})^{n_{\rm L}} (z_{\rm R}^{-1})^{n_{\rm R}} =$$

$$= \exp\left[\sum_{n_{L},n_{L}=1}^{\infty} f_{n_{L}n_{R}} \frac{(z_{L}-1)^{n_{L}}}{n_{L}!} \frac{(z_{R}-1)^{n_{R}}}{n_{R}!}\right]$$
(III. 47)

so that the results for $Q(\alpha_L, \alpha_R)$, $Q(\alpha_L)$ and det $(I - zG\overline{F}(\alpha)) = 0$ are the same as (III.35) - (III.37) with all z's replaced by z-1. The zeroes of

det $(I - (z-1)G\overline{F}(\alpha))$

$$0 = \frac{1}{\alpha \ \alpha + \frac{1}{2}} \left[\left(\alpha - (z - 1)g_{PP} \right) \left(\alpha + \frac{1}{2} - (z - 1)g_{MM} \right) - (z - 1)^2 g_{PM}^2 \right]$$
(III. 48)

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determine the leading behavior of the generating function $Q(\alpha)$ and $Q(\alpha_L, \alpha_R)$. For Q(Y) the result is

$$Q(Y) \simeq \frac{p_{+}(z) + \frac{1}{2} - (z-1) g_{MM}}{p_{+}(z) - p_{-}(z)} \exp(p_{+}(z)Y)$$
(III. 49)

 where^{29}

$$p_{\pm}(z) = \frac{-\frac{1}{2} + (z-1) \left[g_{MM} + g_{PP}\right]}{2} \\ \pm \frac{1}{2} \left[\left(\frac{1}{2} + (z-1) \left[g_{PP} - g_{MM}\right]\right)^2 + 4 (z-1)^2 g_{PM}^2 \right]^{1/2}$$
(III. 50)

are the solutions for (III.48). The result for $Q(z_L, z_R)$ is

$$Q(z_{L}, z_{R}) \cong \frac{(p_{+}(z_{L}) + \frac{1}{2} - (z_{L}^{-1}) g_{MM})(p_{+}(z_{R}) + \frac{1}{2} - (z_{R}^{-1}) g_{MM}) + (z_{L}^{-1}) (z_{R}^{-1}) g_{PM}^{2}}{\lfloor^{p_{+}(z_{L}) - p_{-}(z_{L})}] [p_{+}(z_{R}) - p_{-}(z_{R})]}$$
$$exp \left[p_{+}(z_{L}) Y_{L} + p_{+} (z_{R}) Y_{R} \right] . \qquad (III.51)$$

We see now explicitly that the left-right correlations $p_c(z_L, z_R)$ go to zero like 1/Y as $Y \rightarrow \infty$, which further illustrates the points made in Section II. Additionally we see the formalism of the Mueller-Regge picture closely parallels that of the MPM, justifying our including them in the same general picture. Finally, we note that it is a simple generalization to include isospin constraints explicitly in the above formalism. The coupling and kernel matrices then have a larger dimension.

IV. IMPLICATIONS FOR THE DATA

With the continuing analysis of bubble chamber³⁰ and two arm spectrometer³¹ data on high energy pp collisions, we are beginning to obtain the experimental values for the left-right charged cross sections $\sigma(n_L, n_R)$, the average multiplicities $\langle n_L \rangle vs n_R$, the mean-squared charge transfer ($\langle u^2 \rangle$) and the other coarse-grained quantities we have been discussing. The available current amount of data on these quantities is not large and there is still a lot of important information which can be obtained from analyses of experiments <u>below</u> NAL energies. For example, the question of the energy dependence of the experimental quantities in the AGS-ZGS energy range.

One thing which can be done with the data currently available is to examine the problem of separating that portion of the inelastic cross section which might be labelled diffractive fragmentation in a hybrid or 2-component theory. This problem is important because it strikes at the viability of the two component concept. If the diffractive and short-range components cannot be separated then it's not clear whether we can attach much significance to the two-component models currently being discussed.

A. <u>The Two-Component Concept and $\sigma(n_1, tot)$ </u>

The idea that a certain fraction of inelastic events in asymptotic protonproton collisions contain a quasi-elastically scattered proton in one hemisphere and an excited "fireball" which decays into two or more particles in the other hemisphere has received a good deal of attention. ¹⁵ Only recently has it become apparent that these types of events do not dominate production processes at current energies. ⁹ The experimental observation of a low-mass enhancement in the inclusive cross section $pp \rightarrow p + MM$ at NAL and ISR, ³² however, provides sufficient evidence that there are some of these events.

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Currently, the most frequently discussed possibility is a hybrid model for production processes containing both a fragmentation component and a shortrange order component. There remain several important unanswered questions in the two-component approach such as whether the dominant excitation of the "fireball" at current energies is due to Pomeron exchange. These questions can only be answered cleanly if there is some way of looking at the data which separates the two components so that each can be discussed separately.

In some versions of the two-component model, the charged prong multiplicity distribution would, at infinite energy develop a dip structure, the diffractive component contributing only to low multiplicity. If this characteristic were present in the data, it would be possible to make the separation into two components comparatively easily. The dip structure is certainly not present in the data at NAL energies and so a separation on the basis of total charged multiplicities is not possible. In addition there is certain theoretical uncertainty about whether any dip in the multiplicity distribution is really expected. A "self-consistent" model due to Ball and Zachariesen, ²⁸ for example, has several features in common with the other versions of the two-component approach except that the "diffractive" component contributes to cross sections out to the mean multiplicity. In this version of the two-component model no dip would develop but only a broad plateau. A related approach due to Frazer, Snider and Tan¹⁶ also predicts no dip.

We have to look at slightly more complicated objects than the total charged prong distribution in order to isolate the diffractive component. Since the Pomeron has isospin zero, it will contribute asymptotically only to those cross sections where the charge transfer between hemispheres is zero. In protonproton collisions, this means for example that the total charge in the forward

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hemisphere due to a diffractive event must be +1. If we look at the distribution $\sigma(n^{}_{\rm L},\, tot)$ vs $n^{}_{\rm L}$ then diffractive processes will contribute only to those cross sections with $n_{I_{L}}$ odd. It is not hard to convince one self that the most plausible situation is for a substantial portion of the diffractive cross section to be in $\sigma(1, tot)$. This is the cross section which contains those events in which there is a quasielastic proton in the left hemisphere. By symmetry, there are an equal number of events containing a quasielastic proton in the right hemisphere. In the plot of $\sigma(n_1, \text{tot})$ these show up in $\sigma(1, \text{tot})$, $\sigma(3, \text{tot})$, $\sigma(5, \text{tot})$, etc. The graph in Fig. 3 shows preliminary experimental data on $\sigma(n_L, tot)$ at $p_{lab} = 205 \text{ GeV/c.}^{30}$ The theoretical curves represent parametrizations based on models discussed in Section III and are explained more fully in the figure caption. The data completely support the qualitative expectations of the two component picture. The presence of the expected large surplus of events in $\sigma(1, \text{tot})$ followed by a dip in $\sigma(2, \text{tot})$ is reinforced by the sawtooth behavior of the rest of the distribution showing the tendency to prefer odd multiplicities (and, presumably, $\Delta Q = 0$) over even multiplicities. The presence of the dip allows us to extrapolate roughly the nondiffractive component and estimate

$$\sigma(1, \text{tot})_{\text{diff}} \cong 3.5 \text{ mb} \tag{IV. 1}$$

It might be noted parenthetically that this distribution is the first <u>inelastic</u> multiplicity distribution to show any sort of dip structure.

B. Estimates of the Diffractive Component from $\sigma(n_{L}, n_{R})$

To get more information about the amount of fragmentation or diffraction we can look at the distribution of $\sigma(n_L, n_R)$ directly. As pointed out by Nussinov, Quigg and Wang,⁵ it is possible to distinguish between the extremes of pure fragmentation models and pure short-range order models in a single experiment

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at a fixed energy by looking at the shape of $\sigma(n_L, N-n_L)$ as a function of n_L . It is easy to see how this can be if we assume exact factorization

$$\sigma(n_{L}, N-n_{L}) = a(n_{L}) a(N-n_{L}) . \qquad (IV.2)$$

Recall now that the short-range order concept discussed in Section II would require that both $a(n_L)$ and $a(N-n_L)$ be peaked around some central value $\langle n \rangle \propto \ln s$ and that they fall off rapidly for values of n larger or smaller than $\langle n \rangle$. We can see that in a short-range order picture $\sigma(n_L, N-n_L)$ will be maximum near the symmetric distribution $n_L = N-n_L$. The most popular "fragmentation" models would favor either n_L or $N-n_L$ small so that the antisymmetric distribution is preferred. The specific examples of $a(n_L)$ being Poisson or $1/n_L^2$ are discussed by Nussinov, Quigg and Wong.

The general feature of short-range order contributing mainly near $\langle n_L \rangle$ and "fragmentation" or "diffraction" contributing to small n_L should be true as well in two component models so that looking at $\sigma(n_L, N-n_L)$ should aid in the <u>separation</u> of fragmentation and short-range order. Preliminary data at 205 GeV/c are shown in Fig. 4 compared with a critical point fluid model and a two-component model based on the discussion in Section III. A model independent statement can also be made. Except for the four-prong events, the data favor the symmetric distribution of left-right multiplicities implied by the short-range order concept. The apparent surplus of events in $\sigma(3, 1)$ can be understood if the "fragmentation" component is entirely "diffractive" in that I=0 Pomeron exchange dominates. In the same experiment a more detailed analysis of the four-prong events has been performed. Figure 5 shows the missing mass spectrum of pp \rightarrow pX in the different four-prong charge configurations. ³² Not surprisingly, we see a low mass enhancement in $\sigma(1, 3)$ not present in either $\sigma(2, 2)$ or $\sigma(3, 1)$. If we identify the experimental cross

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section for producing this enhancement

$$\sigma_{\rm D}(1,3) = 0.82 \pm 0.07 \text{ mb}$$
 (IV.3)

with the diffractive component we then conclude that the diffraction and shortrange order components in $\sigma(1, 3)$ are about equal at this energy.

Snider²⁷ has done an analysis of the same data based on an assumption of a factorizable Pomeron. He obtains a larger diffractive component

$$\sigma_{\rm D}(1,3) = 1.35 \text{ mb}$$
 (IV.4)

than the experimental low mass enhancement. In his approach, however, it is quite reasonable to expect a large component of diffraction in the high missing mass as well as a diffractive component in $pp \rightarrow (n\pi^+)X$ which are difficult to separate from the short range order.

Equations (IV. 1) and (IV. 3) are consistent if we assume that $\sigma(1, 1)$ is predominantly diffractive and that the amount of diffraction in $\sigma(1, 5)$, etc. is small.

The obvious test for the presence of diffraction is, of course, the energy dependence of the cross sections. Figure 6 shows what data there is on the energy dependence of the four-prong data. In this plot $\sigma(1,3)$ is compared with $\sigma(0,4) + \sigma(2,2)$. Since Pomeron exchange should not contribute asymptotically to $\Delta Q=1$ the latter cross section should be free of diffraction. The experimental errors are certainly large but there is not much indication for a different dependence. The energy dependence alone is consistent with the diffractive component in $\sigma(1,3)$ being zero but it can be as large as that given by (IV.3).

A related test for the presence of diffraction involves plotting the data as a function of the maximum rapidity gap in the event and the charge transferred across that gap. Data on this is available at $p_{lab} = 12$ and 24 GeV. Since this quantity can in principle separate diffractive events with large missing mass data at higher energies would provide a check on the considerations presented here.

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C. Average Multiplicities and Factorization

As indicated in the discussion of Section III. B concerning the two-polenearest-neighbor-multiperipheral model, the standard formulation of the hybrid or two-component model is explicitly not factorizable in left-right multiplicities. We can define the hybrid model by writing the total cross section in the form

$$\sigma_{\text{tot}} = \sigma^{D} + \sigma^{S}$$
 (IV.5)

where σ^{D} is a diffractive piece and σ^{S} a short-range order piece. A large fraction of the elastic cross section is usually grouped with the diffractive portion. The arguments presented in Section II. C concerning the short-range order component indicate that it separately should asymptotically factorize with respect to left-right multiplicities. If the Pomeranchuk is a hard singularity as discussed in Section II. D then the diffractive piece associated with single Pomeranchuk exchange should also factorize. At high energy, it may therefore make sense to write

$$\sigma(n_{\rm L}, n_{\rm R}) \cong a^{\rm D}(n_{\rm L}) a^{\rm D}(n_{\rm R}) + a^{\rm S}(n_{\rm L}) a^{\rm S}(n_{\rm R})$$
 (IV.6)

where we have to keep track of the fact that the diffractive piece contributes only to those distributions where no charge is exchanged between left and right hemispheres. In terms of the two-pole-nearest-neighbor-multiperipheral model discussed in Section III, it is interesting to note that we recover the factorization properties (IV.6) if we group "high-mass diffraction" with the short-range component.

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The consequences of the factorization properties of (IV.6) are most clearly seen in the dependence with energy of the quantities

$$\langle n_{L} \rangle_{n_{R}} \cong \frac{\sum_{n_{L}} n_{L} \left[a^{D}(n_{L}) a^{D}(n_{R}) + a^{S}(n_{L}) a^{S}(n_{R}) \right]}{\sum_{n_{L}} \left[a^{D}(n_{L}) a^{D}(n_{R}) + a^{S}(n_{L}) a^{S}(n_{R}) \right]}$$
 (IV.7)

with the usual assumption that at high energy $a^{S}(n_{R})$ is peaked at $<n_{R}> \approx \ln s$ while $a^{D}(n_{R})$ is energy independent and peaked at small n_{R} ($n_{R}=1$ in pp collisions). We therefore get

$$< n_L >_{n_R} = 1 \cong \frac{\sum_{n_L} n_L a^D(n_L)}{\sum_{n_L} a^D(n_L)} \rightarrow \text{const}$$
 (IV.8)

For charge configurations in which diffraction cannot contribute such as $n_R^{=2}$ in pp collisions which must have charge transfer, or for right multiplicities large we obtain

$$\langle n_L \rangle_{n_R=2} \cong \langle n_L \rangle_{n_R} \simeq \langle n_R \rangle \cong \frac{\Sigma n_L a^S(n_L)}{\Sigma a^S(n_L)} \propto \ln s$$
 (IV.9)

There is some evidence for this distinction between $<n_L>n_R=1$ and other values of n_R in the ANL-NAL data³⁰ at 205 GeV/c. The fact that data from ISR does not show this distinction can perhaps be understood from the fact that the spectrometers in the ISR experiment do not detect events in a cone of about 2^o around the forward direction. They therefore miss most of the diffractive events and are insensitive to the possibility of two production mechanisms. Data from the NAL experiment are shown in Fig. 7. The theoretical curves are again the KUH fluid model and a version of the hybrid model.

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Note that if we write asymptotically an explicitly factorized form for $\sigma(n_L, n_R)$

$$\sigma(\mathbf{n}_{\mathrm{L}},\mathbf{n}_{\mathrm{R}}) \cong \left[\gamma b^{\mathrm{D}}(\mathbf{n}_{\mathrm{L}}) + b^{\mathrm{S}}(\mathbf{n}_{\mathrm{L}})\right] \left[\gamma b^{\mathrm{D}}(\mathbf{n}_{\mathrm{R}}) + b^{\mathrm{S}}(\mathbf{n}_{\mathrm{R}})\right]$$
(IV. 10)

we can still incorporate the idea of two types of production mechanisms. It is not easy to disentangle the various terms of this kind of <u>ad hoc</u> factorizing hybrid model to each $\sigma(n_L, n_R)$. If we neglect terms of $0(\gamma^2)$ the structure is quite similar to the usual hybrid model. In the factorized form we have explicitly the result

$${}^{}n_{R}=1 \simeq {}^{} \propto \ln s$$
 (IV.11)

Since many current approaches to the two-component model present it as consisting of the first two terms in a power series expansion in something like g_{PPP} ln s where g_{PPP} is a triple-Pomeron-coupling, it is perhaps premature to rule out the possibility that at superasymptotic energies we recover factorization in left-right multiplicities. It is not clear, therefore, that the prediction (IV.8) is an absolute test of these ideas. However, it is clear that measurements of $\langle n_L \rangle_{n_R}$ at several energies through the AGS-ISR range can supply a great deal of information with little fuss.

The factorization test (II.28) can also be used to investigate the possibility of a two-component structure such as (IV.6). Keeping in mind the fact that $a^{D}(n_{L})$ will only contribute to odd n (u=0) we would predict that ratios of cross sections involving only even multiplicities would be much closer to one than those involving odd multiplicities. In the 205 GeV/c data we have, for example,

$$\frac{\sigma(2,2) \sigma(4,4)}{\sigma(2,4) \sigma(2,4)} = 1.3 \pm .3$$
 (IV. 12)

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$$\frac{\sigma(3,3) \sigma(5,5)}{\sigma(3,5) \sigma(3,5)} = 1.8 \pm .6$$
 (IV. 13)

If this tendency is observed at other energies it would be an important confirmation of these ideas.

D. Data on Charge Transfer

The energy behavior of the mean-squared charge transfer between hemisphere has been proposed as a feature of the data which will distinguish shortrange order and fragmentation models.⁶ We define the charge transfer

$$\ll u^2 >> = \sigma_{tot}^{-1} \sum_{n_L, n_R, u} u^2 \sigma(n_L, n_R, u)$$
 (IV. 14)

In models with short-range order it is well known that

$$\ll u^2 >>_{S.R.O.} \sim const$$
 . (IV. 15)

The simple arguments given in Section II in terms of the generating functional formalism confirm this result. See, for example, Eqs. (II. 15') and (II. 17).

The arguments for the behavior of $\ll u^2 >>$ in the diffractive fragmentation picture are slightly more complicated. In these models we have for large s

$$\sigma(\mathbf{n}_{\mathrm{T}},\mathbf{n}_{\mathrm{T}},\mathbf{u}) \to 0 \quad , \quad \mathbf{u} \neq 0 \tag{IV. 16}$$

but the summation over ${\bf n}_{\rm L}, {\bf n}_{\rm R}$ in (IV. 14) diverges at large ${\bf n}$ to give

$$\ll u^2 \gg_{\text{Chou-Yang}} \propto s^{1/2}$$
 . (IV. 17)

This is another example of the general feature of diffractive fragmentation models of large correlations in the central region from high multiplicity events. The hybrid model in this case does not predict a behavior for the charge transfer intermediate between the short-range order and fragmentation pictures. To

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see this we write

$$\sigma(\mathbf{n}_{\mathrm{L}},\mathbf{n}_{\mathrm{R}},\mathbf{u}) = \sigma^{\mathrm{D}}(\mathbf{n}_{\mathrm{L}},\mathbf{n}_{\mathrm{R}},\mathbf{u}) + \sigma^{\mathrm{S}}(\mathbf{n}_{\mathrm{L}},\mathbf{n}_{\mathrm{R}},\mathbf{u})$$
(IV. 18)

and input the usual assumptions about the diffractive component

$$\sigma^{\mathrm{D}}(\mathrm{n_{L},n_{R},u}) \rightarrow 0 \qquad u \neq 0$$
 (IV. 19)

$$\sigma^{\rm D}(n_{\rm L}, n_{\rm R}, u) = 0 \qquad n_{\rm L} + n_{\rm R} > N$$
 (IV. 20)

where N is some (approximate) energy independent integer. In the hybrid picture no energy dependent divergence of the sum (IV.1) takes place and

$$\ll u^2 >>_H \sim \sigma_{tot}^{-1} \sum_{n_L, n_R, u} u^2 \sigma^S(n_L, n_R, u) \sim \text{const}$$
, (IV.21)

the mean squared charge transfer behaves asymptotically just as it does in a short-range order model.

The available data on charge transfer is shown in Fig. 7. The fact that we see little experimental evidence for the approach to a constant asymptotically indicates that the correction terms omit to factorization of the generating function $p_c(z_L, z_R s)$ discussed in Section II, may be large at available energies. This does not mean that the ideas presented in Section II should not be investigated experimentally. The suggestion is that the correction terms are large and this means that the energy dependence of quantities such as R(N, M), Eq. (II. 32) will provide nontrivial information.

V. SUMMARY AND CONCLUSIONS

The analysis of data on hadronic production processes is becoming more sophisticated as some of the simple early questions have been answered. We are still at a primitive stage in our understanding, however, so that it is important to be flexible in our approach. In this paper we have considered what can be learned from left-right multiplicities. Data on left-right multiplicities are easy to obtain and can provide important dynamical information both at intermediate and high energies. Along with other coarsegrained or integrated quantities left-right multiplicities illustrate the advantages of dealing with exclusive and inclusive data on an equivalent footing.

As an example of the dynamic questions which left-right multiplicities can provide we have seen how the short-range order concept leads to factorization. This left-right factorization can be interpreted as a consequence of the presence of a leading Regge pole in the short-range order framework. Left-right factorization is not a universal property of models for the production process but, due to the fact that unitarity¹⁸ and other constraints³³ favor approximately factorizable J-plane singularities, we discuss a series of tests for this property. Data at low energies is needed to see if there is an approach to factorization. The indication from the behavior of charge transfer is that correction terms may be large.

Another feature of left-right multiplicities is that they are sensitive to the presence of a diffractive component in the production mechanism. Preliminary data at 205 GeV/c give a good measure of the amount of diffraction in $\sigma(1, 3)$ and $\sigma(1, \text{tot})$.

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VI. ACKNOWLEDGEMENTS

We are grateful for early conversations with Lloyd Hyman and Rich Singer concerning the ANL-NAL data which stimulated our interest in this project. We thank Carl Bromberg and Tom Ferbel for sending us a preliminary analysis of the 102 GeV/c data and Brian Musgrave for providing us with $pp \rightarrow pX$ missing mass data at 205 GeV/c. For many important conversations, we are indebted to R. Arnold and Dale Snider.

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- d. "Two component models," J. D. Jackson and C. Quigg, NAL preprint (1972).
- e. Bjorken-Feynman-Wilson fluid analog, R. C. Arnold and G. H.
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- f. Mueller-Regge models, W. R. Frazer, R. D. Peccei, S. S. Pinsky, and C. -I. Tan, Phys. Rev. <u>D7</u>, 2647 (1973). These authors consider also a two component model, and give further references concerning such models.
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- 30. Analysis of bubble chamber data to obtain information on left-right multiplicities has been performed by: an ANL-NAL collaboration at 205 GeV/c. Data presented by J. Whitmore at the Nashville Conf. Bonn-Hamburg-Munchen collaboration at 12 and 24 GeV/c. V. Blobel <u>et al.</u>, Hamburg preprint (1973). 102 GeV/c. T. Ferbel, C. Bromberg, private communication.
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TABLE I

Table of Measurable Experimental Quantities

	"Coarse Grained" Integrated Quantities	ned'' I Increasing Momentum Resolution \rightarrow s		"Fine Grained" Differential Quantities
Inclusive Quantities	σ _{tot}	<n_l>,<n_l<sup>nR></n_l<sup></n_l>		ρ(y ₁) dQ/dy
	,	> «u ² >>		C(y ₁ ,y ₂)
Mixed Semi-inclusive Quantities	· ·	<u<sup>2>_{n_R}</u<sup>		
	• • •	^{<n< sup="">L^{>}n_R</n<>}		$\left[\begin{array}{c} \rho(\mathbf{y}_1) \right]_{n} \mathbf{ch}$
	<n0<sup>>n_,<nk<sup>>n_</nk<sup></n0<sup>	σ(n _L , tot)		
	σ(n _{ch})			
		$\sigma(n_L, n_R)$	$\sigma(n_L, n_C, n_R)$	
		σ(n _L , n _R , u)	for example, 3 momentum bins	
Exclusive Quantities	$\sigma(\mathbf{n}_{+},\mathbf{n}_{-},\mathbf{n}_{0})$	$\sigma_{n_{+L},n_{-L},n_{0L};}$ n_{+R},n_{-R},n_{0R}		$\frac{d\sigma(y_1 \cdots y_n)}{dy_1 \cdots dy_n}$
				// hima hh m

1 momentum bin 2 momentum bins \dots # bins >> n_{max}

"Coarse Grained" quantities are integrals ("sums") over fine grained quantities, averages over kinematic variables.

"Inclusive" quantities are sums over exclusive quantities, averages over (unseen particles) channels.

Everything tied together by generating functional.

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FIGURE CAPTIONS

- 1. The multiperipheral ladder diagram for the production process $ab \rightarrow a'b'$ $n_{L}(\pi's) + n_{R}(\pi's)$. The phase space is assumed to be one-dimensional and the x, shown are rapidity differences.
- 2. The multiperipheral Mueller diagram which contributes to the $n_L + n_R$ particle inclusive correlation function.
- 3. Preliminary experimental data on $\sigma(n_L, tot)$ from the ANL-NAL collaboration at $p_{lab} = 205 \text{ GeV/c}$.³⁰ The two theoretical curves are motivated from models discussed in Section III. The smooth curve is the prediction of a critical point fluid model given by Eq. (III. 32) with the parameters (III. 31). The dashed curve is a two component model with

$$\sigma(n_{\rm L}, n_{\rm R}) \propto \left(\frac{1}{n_{\rm L}^2} \frac{1}{n_{\rm R}^2} + 4.8 \frac{e^{-4} {}^{(4)} {}^{\rm R}}{n_{\rm L}!} \frac{e^{-4} {}^{(4)} {}^{\rm R}}{n_{\rm R}!}\right)$$

This parametrization retains the basic features of (III. 43).

- 4. Data on $\sigma(n_L, n_R)$ at 205 GeV/c from the ANL-NAL collaboration³⁰ presented at fixed prong number $N=n_L+n_R$ as a function of n_L . The theoretical curves are from the same models as in Fig. 3.
- 5. The missing mass spectrum off the slow proton in the four-prong events at 205 GeV/c for different left-right charge configurations. ³²
- 6. The energy dependence of the cross sections $\sigma(1,3) = \sigma_{Q=0}$ and $\sigma(0,4) + \sigma(2,2) = \sigma_{\Delta Q=1}$ as a function of energy. The data at 12 and 24 GeV/c is from the Bonn-Hamburg -Munchen collaboration. The 205 GeV/c data is from the ANL-NAL collaboration.³⁰
- Data on <n_L >n_R from a bubble chamber experiment at NAL.³⁰
 Data on ≪u²>> vs √s.³⁰

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FIG. 2











FIG. 6



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FIG. 8