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# TOWARDS A UNIFIED PARTON SHOWER DESCRIPTION OF MULTIPARTICLE PRODUCTION PROCESSES

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Since 1984, remarkable regularities related to negative binomial (NB) properties have been found in the experimental multiplicity distributions of $e^+e^-$ annihilation, leptoproduction and soft hadronic collisions at high energies. From a "Monte Carlo experiment" based on the Lund version of the QCD parton shower model for $e^+e^-$ annihilation, we have recently shown that the same NB regularities are predicted for multiparticle production by quark-antiquark and gluon-gluon systems up to 2 TeV, not only for the final hadrons but also for the final partons before hadronization. In addition, we have found an approximate, but physically intuitive, description of the parton shower (independent emission of selfsimilar gluon jets) and of the hadronization (a specific form of local parton-hadron duality). Motivated by the observed universality of the NB regularities, we now extend our approximate parton shower description to leptoproduction and soft hadronic collisions. We show that it provides a method of calculating partonic multiplicity distributions from the experimental hadronic ones, and we give applications to muon-proton and proton-(anti)proton collisions. In the soft hadronic case, we propose a dynamical justification for a perturbative parton shower description and we indicate consequences for ultra-relativistic nuclear collisions and the problem of quark-gluon plasma formation.

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## 1. *Introduction: Simplified Parton Shower Description* (*SPSD*)

Recent work on the QCD interpretation of multiparticle production in hard processes ($e^+e^-$ annihilation, leptoproduction) has shown that at very high energies the best descriptions are given by parton shower models derived from perturbative QCD and supplemented by a non-perturbative hadronization prescription.

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Among the various versions studied so far, the preferred one now appears to be a coherent parton shower model of Marchesini-Webber type (which simulates destructive coherence effects by an angular ordering prescription), with Lund-type string fragmentation for the hadronization phase [1][1]. Using this so-called Lund Shower Model with main parameters $Q_0 = 1$ GeV, $\Lambda = 0.4$ GeV [1b], we recently studied [2] the multiplicity distributions (MD) which it predicts for $e^+e^-$ annihilation at centre-of-mass energies $\sqrt{s} = 22$–$2000$ GeV, and we compared them with those which it predicts when the initial partonic system, instead of being a $q\bar{q}$ pair as in $e^+e^-$ annihilation, is composed of two gluons (gg system). We studied the MDs for the (meta)stable charged hadrons (mostly pions) and for the final partons of the shower (mostly gluons), in full phase space and in symmetric rapidity windows $|y| < y_0$ (the c.m. rapidity $y$ was defined with respect to the linear sphericity axis for the final partons).

As reported in Ref. [2], this "Monte Carlo experiment" revealed that with very few exceptions, the partonic and the hadronic MDs considered are very well described by negative binomial (NB) fits, with the following properties for the NB parameters $\bar{n}_p$, $k_p$ (partonic MD) and $\bar{n}_h$, $k_h$ (hadronic MD).

a) All four parameters increase for a growing window $|y| < y_0$ at fixed $\sqrt{s}$. For increasing $\sqrt{s}$ at fixed $y_0$, the mean multiplicities $\bar{n}_p$, $\bar{n}_h$ increase whereas $k_p$, $k_h$ decrease.

b) The partonic and hadronic NB distributions have approximately[2] the same $k$ and constant ratio of the $\bar{n}$:

$$k_h \simeq k_p, \qquad \varrho \equiv \bar{n}_h/\bar{n}_p \simeq 2. \tag{1}$$

This can be understood as a far-reaching manifestation of local parton-hadron duality (see Ref. [2], Section 3), far-reaching because it links the complete MDs, i.e., all their moments.

c) For fixed windows $|y| < y_0$, the quantities

$$\bar{N}_p = k_p \ln[1+(\bar{n}_p/k_p)], \qquad \bar{N}_h = k_h \ln[1+(\bar{n}_h/k_h)] \tag{2}$$

are approximately independent of $\sqrt{s}$, this independence being more precise for $\bar{N}_p$ than for $\bar{N}_h$. The growth of multiplicities with $\sqrt{s}$ is therefore mainly due to the increase of the quantities

$$\bar{n}_{cp} = \bar{n}_p/\bar{N}_p, \qquad \bar{n}_{ch} = \bar{n}_h/\bar{N}_h. \tag{3}$$

d) The following approximate relations hold between the MDs produced by the $q\bar{q}$ and the gg systems:

$$\bar{N}_p(gg) \simeq 2\bar{N}_p(q\bar{q}), \qquad \bar{n}_{cp}(gg) \simeq \bar{n}_{cp}(q\bar{q}). \tag{4}$$

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[1] See Ref. [1a] for the model with its various options and [1b] for the determination of the best option in $e^+e^-$ annihilation; the authors of the first Ref. [1b] coined the name Lund Shower Model. Reference [1c] concerns leptoproduction.

[2] The relation $k_h \simeq k_p$ was found to hold at $\sqrt{s} \geqslant 200$ GeV, not at $\sqrt{s} \leqslant 30$ GeV. In the latter case, however, the distributions are close to Poisson and the deviations are controlled by $k_h^{-1}$, $k_p^{-1}$ which are both very small, $k^{-1} \simeq k_p^{-1} \simeq 0$ (the NB distribution reduces to a Poissonian for $k = \infty$).

Similar relations hold for $\bar{N}_h$ and $\bar{n}_{ch}$, but with larger errors.

In Ref. [2], we gave a simple interpretation of these properties by using what we called [3] the *clan structure* of NB distributions, i.e., the mathematical fact that a NB of parameters $\bar{n}$, $k$ can always be generated by independent emission of $N$ "clans" of multiplicity $n_c$ with, for $N$, a Poisson distribution of average $\bar{N} = k \ln[1+(\bar{n}/k)]$, and for the *average* clan a logarithmic multiplicity distribution of average $\bar{n}_c = \bar{n}/\bar{N}$. These averages occur in Eqs. (2)–(4) for the partonic and hadronic levels.

At the partonic level we were led to interpret the clans within a given window $|y| < y_0$ to be "bremsstrahlung gluon jets" (BGJ), i.e., a simple type of jets mainly composed of gluons, emitted by an initial "skeleton" part of the partonic shower in bremsstrahlung fashion (Poisson distribution), and falling at least partly in the window. Within the latter, each BGJ is assumed to have a geometric distribution of final partons:

$$P(n_{cp}) = (\nu_p - 1)^{n_{cp}-1}/\nu_p^{n_{cp}}, \qquad n_{cp} \geqslant 1, \qquad \langle n_{cp} \rangle = \nu_p. \tag{5}$$

This distribution is reasonable because it corresponds to the simplest type of self-similar cascade processes; it is therefore expected to be a good approximation for QCD gluon jets, which contain only a few $q\bar{q}$ pairs. If the mean BGJ multiplicity $\nu_p$ is distributed as $d\nu_p/\nu_p$ over an interval $1 < \nu_p < 1+(\bar{n}_p/k_p)$, the average clan has a logarithmic MD. If, in addition, the number $N_p$ of BGJs is Poisson distributed with average $\bar{N}_p$ given by Eq. (2), the resulting partonic MD is a NB with parameters $\bar{n}_p$, $k_p$. The hadronic NB distribution of parameters $\bar{n}_h$, $k_h$ given by (1) then follows from the local parton-hadron duality principle as formulated in Section 3 of Ref. [2]. We shall use the name "simplified parton shower description" (SPSD) for this approximate description of the Lund Shower Model, which we proposed on the basis of the NB shapes of partonic and hadronic MDs supplemented by the properties a) to d).

Our SPSD can of course only be an approximation to the original shower model. Among other imperfections, conservation of energy, momentum and charge are not taken into account, and our formulation of parton-hadron duality neglects the resonances present in the hadronic final state. But the SPSD has the advantage of being much simpler than the full-fledged Lund Shower Model which, by the way, involves itself many approximations. The most attractive feature is that the SPSD describes two intermediate stages of the multiparticle production process, the bremsstrahlung gluon jets and the final parton distributions, relating their properties to the observed hadronic distributions by very simple equations. In its present state, the SPSD is still very incomplete (it is certainly not yet a model), but the above feature allows us to develop it step by step with guidance from experimental data.

The hadronic level results of Ref. [2], while being quantitatively close to the available MD data for $e^+e^-$ annihilation (Ref. [4], $\sqrt{s} = 29$ GeV), are qualitatively very similar to what is observed for the MDs of soft (low $p_T$) hadron-hadron collisions and deep inelastic muon-proton scattering (leptoproduction). We refer to Ref. [5] for a recent review of the experimental evidence. This similarity strongly suggests that one should try to extend the Lund Shower Model not only to leptoproduction, as discussed in Ref. [1c], but also to soft hadronic collisions (the present Lund Model for the latter class of collisions, called Fritiof [6], has no shower character). The question of extension also arises for our SPSD, and the principle of an answer was already briefly formulated at the end of Ref. [2]. We show in the present paper how it can be used to interpret variations observed experimentally in the MDs of leptoproduction and soft hadronic collisions.

Our method is explained in the next section, which also presents its application to the deep inelastic muon-proton data. Section 3 deals with the case of low $p_T$ proton-(anti)-proton collisions. We address in Section 4 the more general question of justifying on theoretical grounds the extension of a perturbative parton shower description to the low $p_T$ hadronic collisions which are usually regarded as soft. We argue that the answer is to be found in the occurrence of a large longitudinal excitation of the incident hadrons in an early phase of the collision. This can be considered as a hard subprocess, offering the possibility of a *unified parton shower description* for all multiparticle production processes. Section 5 applies our method to proton-nucleus collisions and presents implications for the problem of quark-gluon plasma formation. The paper ends with a summary and concluding remarks.

## 2. *Energy dependence of multiplicity distributions and application to leptoproduction*

Careful consideration of results a), b) and c) of our Monte Carlo experiment [2] shows that they cannot all be exact. This is seen by substituting Eq. (1) in Eq. (2) and considering a fixed rapidity window $|y| < y_0$ for increasing energy $\sqrt{s}$. Since $\bar{n}_p$ and $k_p^{-1}$ increase, $\bar{N}_p$ and $\bar{N}_h$ cannot both remain constant. Although both change very little over the wide $\sqrt{s}$-range covered in Ref. [2], $\bar{N}_p$ turns out to be distinctly closer to constant[3]. By assumption, we shall therefore take $\bar{N}_p$ to be independent of $\sqrt{s}$ at constant $y_0$ and make this a defining feature of our SPSD.

With property c) of Section 1 now taken to state that $\bar{N}_p$ is independent of $\sqrt{s}$, and with properties b) and d), we can study the applicability of the SPSD to other multiparticle production processes. The simplest case is leptoproduction. Here, as for $e^+e^-$ annihilation, it is reasonable to assume that the parton shower is initiated from (anti)quarks excited by the original electroweak interaction, without contribution from originally excited gluons [but the excited (anti)quark can emit initial state radiation, which is not the case in $e^+e^-$ annihilation]. We therefore expect that in leptoproduction the source of BGJ emission is of quark origin. For short we shall say that it is a *quark source*. For a fixed window $|y| < y_0$, the resulting $\bar{N}_p$ is predicted to be independent of $\sqrt{s}$.

In purely hadronic collisions, on the other hand, the original interaction (probably gluon exchange) can excite gluons as well, so that the source of BGJ emission can be partly of quark and partly of gluon origin. As we shall show in the next section, even an energy-

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[3] See Figs. 1 and 2 of Ref. [2]. As all the figures of Ref. [2] illustrate, there is here a general trend; the Monte Carlo results are systematically simpler and smoother for the partonic than for the hadronic MDs. This is probably due to the non-perturbative hadronization prescription, the most uncertain part of all QCD-based models.

-independent source can then give an $\bar{N}_p$ increasing with $\sqrt{s}$, because its gluon fraction can grow at the cost of its quark fraction and a gluon source produces more BGJs than a quark one [property d)].

Given the experimental fact that the charged hadron MDs in rapidity windows $|y| < y_0$ are of NB shape, the actual analysis is extremely simple. The NB fit to the data gives $\bar{n}_h$ and $k_h$ for each rapidity window. Property b) then tells us that the partonic MD in the same window is the NB distribution of parameters

$$\bar{n}_p = \bar{n}_h/2, \qquad k_p = k_h. \tag{6}$$

We identify its clans with the BGJs of the SPSD, their number being Poisson distributed with average

$$\bar{N}_p = k_p \ln\left[1+(\bar{n}_p/k_p)\right] \equiv \bar{n}_p/\bar{n}_{cp}. \tag{7}$$

We predicted above that in the case of leptoproduction this number should only depend on $y_0$ and not on energy, more precisely not on $W$, the total energy of the hadronic system in its own c.m. frame. We have tested this for deep inelastic muon-proton scattering,

TABLE I

The clan structure parameters at hadronic and partonic level, Eqs. (2) and (3), for deep inelastic muon-proton scattering, in various rapidity windows and hadronic energy intervals, from Ref. [7]

| | $W$ | $\bar{N}_h$ | $\bar{N}_p$ | $\bar{n}_{ch}$ | $\bar{n}_{cp}$ |
|---|---|---|---|---|---|
| | 6÷8 GeV | 1.40±0.02 | 0.72±0.01 | 1.049±0.001 | 1.025±0.003 |
| $\|y\| < 0.5$ | 12÷14 GeV | 1.35±0.02 | 0.71±0.01 | 1.127±0.003 | 1.065±0.001 |
| | 18÷20 GeV | 1.36±0.02 | 0.73±0.01 | 1.158±0.003 | 1.081±0.001 |
| | 6÷8 GeV | 2.65±0.04 | 1.37±0.02 | 1.076±0.002 | 1.039±0.001 |
| $\|y\| < 1.0$ | 12÷14 GeV | 2.57±0.03 | 1.38±0.02 | 1.170±0.003 | 1.087±0.001 |
| | 18÷20 GeV | 2.52±0.03 | 1.39±0.01 | 1.237±0.005 | 1.122±0.001 |
| | 6÷8 GeV | 3.89±0.07 | 2.00±0.03 | 1.059±0.002 | 1.030±0.001 |
| $\|y\| < 1.5$ | 12÷14 GeV | 3.72±0.04 | 2.01±0.02 | 1.174±0.004 | 1.085±0.001 |
| | 18÷20 GeV | 3.60±0.03 | 2.00±0.02 | 1.268±0.005 | 1.139±0.002 |

using the data of the European Muon Collaboration [7]. The results, shown in Table I, strikingly confirm the constancy of $\bar{N}_p$ as $W$ increases. The growth of multiplicity with $W$ is entirely due to the increase of $\bar{n}_{cp} = \bar{n}_p/\bar{N}_p$, which is the average multiplicity of the BGJs. In fact, $\bar{n}_{cp}$ is a double average; each BGJ has a mean multiplicity $\nu_p$, and $\nu_p$ itself is distributed over the interval $1 < \nu_p < 1+(\bar{n}_p/k_p)$ with weight $d\nu_p/\nu_p$ [4].

[4] As explained in Ref. [2], Section 5, this distribution of $\nu_p$ is a good approximation when the rapidity window is not too large. For large $y_0$, one finds a surplus of small clans, i.e., BGJs of low $\nu_p$, but this does not distort significantly the NB shape of the overall distribution.

### 3. *Soft hadronic processes. Appearance of the gluon source term*

We performed the same calculations for pp collisions at $\sqrt{s} = 22$ GeV (European Hybrid Spectrometer data of NA 22 Collaboration, Ref. [8]) and p$\bar{\text{p}}$ collisions at $\sqrt{s} = 200$–$900$ GeV (p$\bar{\text{p}}$ collider data of UA5 Collaboration [9]; we shall mainly use the 546 GeV data which have the highest statistics). Since the UA5 streamer chamber gave only pseudorapidity measurements, we did the analysis for pseudorapidity windows $|\eta| < \eta_0$. The results are listed in Table II for $\sqrt{s} = 22$ and 546 GeV. In contrast with

TABLE II

The clan structure parameters at hadronic and partonic level, Eqs. (2) and (3), for pp at $\sqrt{s} = 22$ GeV [8] and p$\bar{\text{p}}$ at $\sqrt{s} = 546$ GeV [9], in various pseudorapidity intervals. At 22 GeV, the values for $|\eta| < 0.2$ have been obtained from the published data for $|\eta| < 0.25$ by multiplying $\bar{n}_h$ by 0.8 and keeping $k_h$ unchanged. The 22 GeV data for $|\eta| < 2.5$ and 3.0 are derived from the even multiplicity events only

| | $\sqrt{s}$ | $\bar{N}_h$ | $\bar{N}_p$ | $\bar{n}_{ch}$ | $\bar{n}_{cp}$ |
|---|---|---|---|---|---|
| $\|\eta\| < 0.2$ | 22 GeV | $0.52 \pm 0.01$ | $0.27 \pm 0.01$ | $1.129 \pm 0.004$ | $1.066 \pm 0.001$ |
| | 546 GeV | $0.87 \pm 0.02$ | $0.49 \pm 0.02$ | $1.336 \pm 0.010$ | $1.175 \pm 0.003$ |
| $\|\eta\| < 0.5$ | 22 GeV | $1.18 \pm 0.03$ | $0.66 \pm 0.01$ | $1.278 \pm 0.010$ | $1.144 \pm 0.003$ |
| | 546 GeV | $1.72 \pm 0.03$ | $1.07 \pm 0.01$ | $1.743 \pm 0.025$ | $1.399 \pm 0.009$ |
| $\|\eta\| < 1.0$ | 22 GeV | $2.18 \pm 0.05$ | $1.26 \pm 0.02$ | $1.395 \pm 0.016$ | $1.207 \pm 0.005$ |
| | 546 GeV | $2.67 \pm 0.04$ | $1.79 \pm 0.02$ | $2.299 \pm 0.047$ | $1.714 \pm 0.016$ |
| $\|\eta\| < 1.5$ | 22 GeV | $3.05 \pm 0.08$ | $1.81 \pm 0.03$ | $1.499 \pm 0.028$ | $1.264 \pm 0.009$ |
| | 546 GeV | $3.45 \pm 0.05$ | $2.40 \pm 0.03$ | $2.748 \pm 0.069$ | $1.970 \pm 0.028$ |
| $\|\eta\| < 2.0$ | 22 GeV | $3.96 \pm 0.08$ | $2.35 \pm 0.03$ | $1.508 \pm 0.020$ | $1.269 \pm 0.007$ |
| | 546 GeV | $4.14 \pm 0.06$ | $2.95 \pm 0.03$ | $3.095 \pm 0.083$ | $2.171 \pm 0.034$ |
| $\|\eta\| < 2.5$ | 22 GeV | $4.95 \pm 0.15$ | $2.94 \pm 0.07$ | $1.508 \pm 0.032$ | $1.269 \pm 0.011$ |
| | 546 GeV | $4.72 \pm 0.06$ | $3.41 \pm 0.04$ | $3.370 \pm 0.096$ | $2.329 \pm 0.040$ |
| $\|\eta\| < 3.0$ | 22 GeV | $5.54 \pm 0.14$ | $3.26 \pm 0.06$ | $1.474 \pm 0.026$ | $1.250 \pm 0.009$ |
| | 546 GeV | $5.32 \pm 0.06$ | $3.88 \pm 0.04$ | $3.536 \pm 0.092$ | $2.425 \pm 0.039$ |

leptoproduction, we now find a systematic increase of $\bar{N}_p$ with $\sqrt{s}$ [5]. As we shall show, it can be interpreted along the lines indicated in Section 2 in terms of variable quark and gluon fractions of the source. In the spirit of our SPSD, the most natural assumption is that the overall source strength is energy-independent. We express this by the formulae

$$\bar{N}_p = \alpha \bar{N}_g + (1-\alpha)\bar{N}_q \tag{8}$$

$$0 \leqslant \alpha \leqslant 1 \tag{9}$$

with $\bar{N}_g$ and $\bar{N}_q$ the mean numbers of BGJs emitted by pure gluon and quark sources

[5] Interestingly, it is more uniform for $\bar{N}_p$ than for $\bar{N}_h$; see the footnote at the beginning of Section 2 for similar trends in the Monte Carlo experiment.

respectively (in the window $|\eta| < \eta_0$ under consideration), taken to be independent of $\sqrt{s}$. By invoking property d), first Eq. (4), we write for all $\eta_0$

$$\bar{N}_g = 2\bar{N}_q. \tag{10}$$

The second Eq. (4) means that for given $\eta_0$ and $\sqrt{s}$, the BGJs emitted by the two types of source have the same multiplicity distributions in the window.

Let us first assume that the gluon part of the source is negligible at $\sqrt{s} = 22$ GeV (we shall prove later that it is $\lesssim 10\%$). This is expressed by putting $\alpha = 0$ in (8) and we therefore write

$$\bar{N}_p(22) = \bar{N}_q. \tag{11}$$

The value $\bar{N}_p$ (546) of $\bar{N}_p$ at $\sqrt{s} = 546$ GeV corresponds to a positive $\alpha$ which is determined by combining (8), (10) and (11):

$$\bar{N}_p(546) = (1+\alpha)\bar{N}_q = (1+\alpha)\bar{N}_p(22) \tag{12}$$

which gives for $\alpha$

$$\alpha = \Delta/\bar{N}_p(22), \qquad \Delta \equiv \bar{N}_p(546) - \bar{N}_p(22). \tag{13}$$

The resulting values for various pseudorapidity windows are given in Table III. As a first consistency test, we see that (9) is always satisfied. It remains satisfied for the UA5 data at 900 GeV, the limit $\alpha = 1$ being reached within errors for $|\eta| < 0.2$. Our assumptions imply that the growth of $\bar{N}_p$ in this very small window should therefore stop at $\sqrt{s} \sim 1$ TeV.

A second test is provided by the expectation that the gluon part of the source should be more central than the quark part, which implies that $\alpha$ should decrease for a growing window $|\eta| < \eta_0$. Table III shows that this is verified. One can even say more in this respect. The gluon part in Eq. (8) is

$$\alpha\bar{N}_g = 2\alpha\bar{N}_q = 2\alpha\bar{N}_p(22) = 2\Delta. \tag{14}$$

TABLE III

The parameters $\alpha$ and $\Delta$ of Eq. (13) are calculated from Table II. For $|\eta| < 2.5$ at 22 GeV, the experimental value of $\bar{n}_h$ shows a rather large fluctuation compared to a smooth interpolation between $|\eta| < 2.0$ and 3.0. This has been removed in the present Table

| | $\alpha$ | $\Delta$ |
|---|---|---|
| $\|\eta\| < 0.2$ | $0.803 \pm 0.134$ | $0.220 \pm 0.025$ |
| $\|\eta\| < 0.5$ | $0.626 \pm 0.053$ | $0.413 \pm 0.017$ |
| $\|\eta\| < 1.0$ | $0.423 \pm 0.042$ | $0.532 \pm 0.025$ |
| $\|\eta\| < 1.5$ | $0.326 \pm 0.040$ | $0.594 \pm 0.044$ |
| $\|\eta\| < 2.0$ | $0.253 \pm 0.032$ | $0.596 \pm 0.048$ |
| $\|\eta\| < 2.5$ | $0.213 \pm 0.041$ | $0.600 \pm 0.073$ |
| $\|\eta\| < 3.0$ | $0.188 \pm 0.035$ | $0.613 \pm 0.073$ |

The last column of Table III shows that this quantity increases up to $\eta_0 = 1.5$ and then becomes practically constant. The physical meaning is that the gluon part of the source is almost entirely concentrated in the window $|\eta| < 1.5$[6].

We now consider the possibility that the source can already be partly gluonic at 22 GeV. Equation (11) is then replaced by (8) with a value $\alpha_0$ for the parameter

$$\bar{N}_p(22) = (1-\alpha_0)\bar{N}_q + \alpha_0 \bar{N}_g = (1+\alpha_0)\bar{N}_q. \tag{15}$$

Combining this with the first Eq. (12), we obtain

$$(1+\alpha_0)/(1+\alpha) = \bar{N}_p(22)/\bar{N}_p(546). \tag{16}$$

Since $\alpha \leqslant 1$, this gives an upper limit for $\alpha_0$:

$$\alpha_0 \leqslant [2\bar{N}_p(22)/\bar{N}_p(546)] - 1. \tag{17}$$

For the smallest window $|\eta| < 0.2$, Eq. (17) gives $\alpha_0 \leqslant 0.1$. Just as was the case at 546 GeV for $\alpha$, one expects $\alpha_0$ to decrease for a growing window. The bound $\alpha_0 \leqslant 0.1$ therefore applies throughout, and it is a good approximation to take the source to be of quark origin at $\sqrt{s} = 22$ GeV.

The interpretation of the growth of $\bar{N}_p$ with $\sqrt{s}$ which we have presented seems to be the most natural one in the spirit of our SPSD, but there are other possibilities (experiment will have the last word, of course). One that comes immediately to mind is to assume a source composed of an energy-independent quark part $\bar{N}_q$ and an increasing gluon part $\beta\bar{N}_g = 2\beta\bar{N}_q$. This replaces Eq. (8) by

$$\bar{N}_p = \bar{N}_q + \beta\bar{N}_g = (1+2\beta)\bar{N}_q. \tag{18}$$

In this case there is no upper bound for $\beta$. Hence one cannot derive from the data that the gluon part of the source is small at 22 GeV. If one assumes $\beta = 0$ at 22 GeV, the resulting value of $\beta$ at 546 GeV is simply $\alpha/2$ with $\alpha$ given as in Table III, and one can again conclude that the gluonic part of the source is concentrated in the pseudorapidity window $|\eta| < 1.5$.

Although the hadronic collisions just studied show an increase of $\bar{N}_p$ with $\sqrt{s}$, in contrast with the energy-independence found in leptoproduction and for the $q\bar{q}$ and gg systems of Ref. [2], one should stress that the effect is small, especially for the larger pseudorapidity windows. As is seen in Table II, the bulk of the multiplicity increase is due to the growth of $\bar{n}_{cp}$, i.e., the multiplicity growth of the BGJs. In other words, as $\sqrt{s}$ increases, the BGJs grow more massive. But also their $p_T$ can be expected to grow with $\sqrt{s}$, suggesting the possibility that the minijets observed at the CERN $p\bar{p}$ collider are BGJs in the multi-GeV tail of their $p_T$ distribution[7].

[6] Interestingly, this simple behaviour does not appear in the hadronic quantities. As seen from Table II, the hadronic analogue of $\Delta$ which is $\bar{N}_h(546)-\bar{N}_h(22)$ first increases to $\sim 0.5$ and then decreases to $\sim -0.2$ as $\eta_0$ increases from 0.2 to 3.0.

[7] M. Jacob, private communication.

*4. Large longitudinal mass excitation in soft collisions*

It is now generally admitted that in low-$p_T$ hadronic collisions, a large longitudinal rapidity separation occurs between incident partons. This is true in the Dual Parton Model (DPM, Ref. [10]), because the two leading chains, attached at one end to the leading (di)quarks with large momentum fraction $x$, have their other end attached to low-$x$ quarks (weight function $\propto x^{-\frac{1}{2}}$). It is also true in the Lund Fritiof Model [6] because each incident hadron is assumed to get excited into a string ending on one of its valence partons with very low $x$ (weight function $\propto x^{-1}$). In both cases a large excitation occurs between the unseparated and separated states of the incident hadrons, giving them high effective masses.

In Fritiof, which kinematically is a double diffraction dissociation model, the excited states of the incident hadrons have large masses by assumption. In the DPM, the masses are large because the low-$x$ partons on which the leading chains end get a sizeable transverse mass $m_T$ in the gluon exchange processes which initiate the reaction (this can be seen from the derivation of the DPM by means of the $1/N_c$ expansion of the Reggeon and Pomeron calculus, $N_c$ being the number of colours, $N_c = 3$ in actual QCD). In fact, just a single gluon exchange involving an incident low-$x$ parton and giving it a sizeable $m_T$ leads to an excited system of large mass

$$M \simeq m_T(x'/x)^{1/2}, \tag{19}$$

where $x'$ is the momentum fraction of the leading (di)quark in the same hadron ($x \ll x'$).

In both cases the excitation is a longitudinal one because the separated partons have $m_T$'s of normal size ($\sim$ 0.5–1 GeV), but largely different longitudinal rapidities [8]. We therefore conclude that low-$p_T$ hadronic collisions begin with a large longitudinal mass excitation, and we propose that this "hard subprocess" plays a rôle similar to the large momentum transfers occurring in the traditional hard collisions. It is then natural to assume, as we did in the previous section, that a perturbative parton shower can form in low-$p_T$ hadronic collisions with properties analogous to those of hard processes, the result being a unified QCD-based approach to all multiparticle production phenomena.

It is interesting to note that trends toward shower-related unification are appearing in the thinking of other workers. Thus, the Lund group put growing emphasis on gluon radiation, also for low $p_T$ hadronic collisions [11], exploiting the fact that such gluon emission effects can be regarded as an alternative description of QCD cascades [12]. One can therefore understand that with the inclusion of semi-hard gluon effects, the Fritiof model does quite well in reproducing the observed multiplicity distributions, and an essentially similar situation prevails for the Lund description of leptoproduction data.

A quite different example of the above-mentioned trends is given by Hwa and collaborators, who advocate stochastic branching as the unifying feature in multiparticle production, although they maintain that QCD partons, while relevant for the hard processes, are irrelevant for the soft ones [13].

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[8] Note that if one wants to apply (19) to the incident hadrons before they interact, their low-$x$ partons should be given very small $m_T \propto x^{1/2}$ in order to keep the overall hadron mass small.

## 5. *Remarks on nuclear collisions and quark-gluon plasma formation*

As found by the NA5 experiment (Ref. [14], $p_{\mathrm{lab}} = 200$ GeV/$c$), the multiplicity distributions of charged hadrons produced in proton-nucleus collisions have negative binomial shape in symmetric windows $|y| < y_0$ for $y_0 \lesssim 2$, and in forward and backward windows $0 < \pm y < y_0$ for all $y_0$, where $y$ is the rapidity in the proton-nucleon centre-of--mass system and is positive in the forward (proton) hemisphere. In Section 4 of Ref. [5], we performed the hadronic clan analysis of the NA5 data; we here reproduce the results in Figs. 1a and 2a. Proceeding as before, we now calculate the corresponding partonic

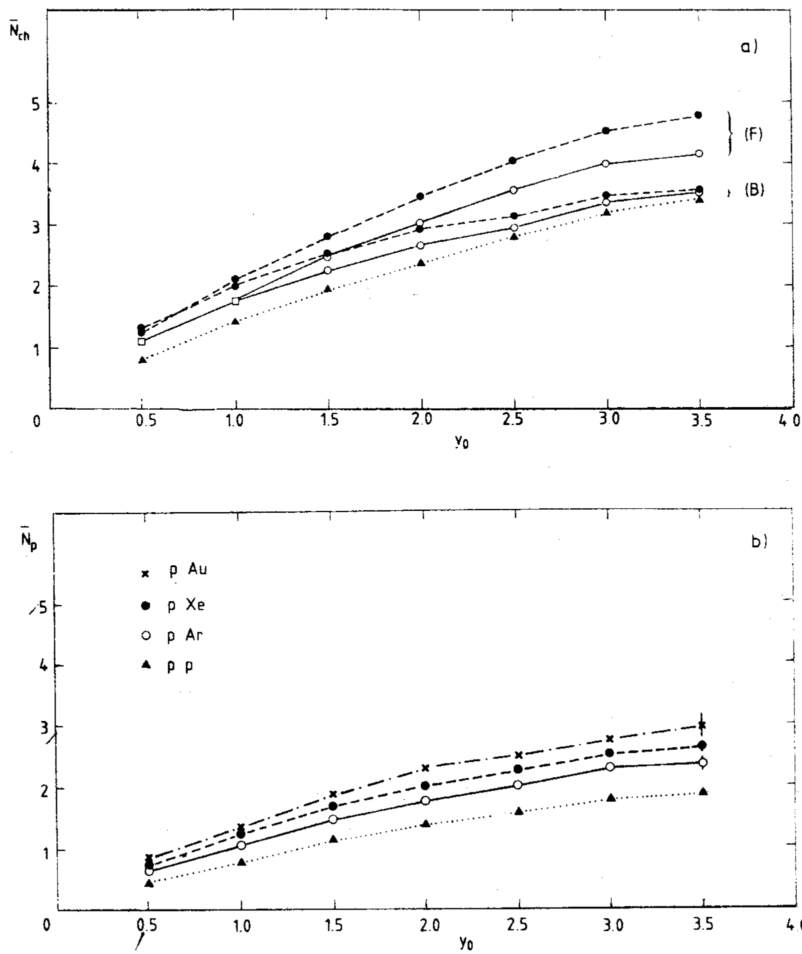


Fig. 1. The mean number of clans for collisions of 200 GeV/$c$ protons on proton, argon and xenon targets (data of Ref. [14]) and of 350 GeV/$c$ protons on a gold target (data of Ref. [15]), in various forward (F, $0 < y < y_0$) and backward (B, $-y_0 < y < 0$) rapidity intervals. a) Number of clans for charged hadrons. b) Number of clans at parton level. The forward and backward values are equal within errors. The symbols are defined in Fig. 1b

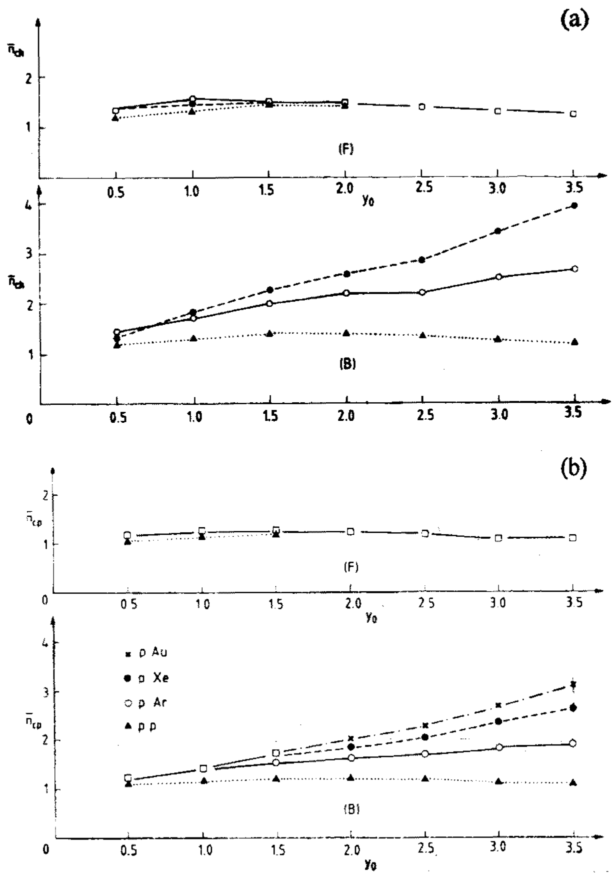


Fig. 2. The mean number of particles per clan for the same reactions as in Fig. 1, for charged hadrons (a) and for partons (b). The symbols are defined in Fig. 2b

clan parameters $\bar{N}_p$, $\bar{n}_{cp}$ by using Eqs. (6) and (7). They are shown in Figs. 1b and 2b, where we included very recent data for proton-gold at $p_{lab} = 350$ GeV/$c$ [15][9].

In contrast with the lepton-hadron and hadron-hadron collisions studied before, the simplifications obtained by going from hadronic to partonic level here turn out to be surprisingly large. The mean numbers of clans become equal within errors for the forward and backward hemispheres (their common values and typical errors are shown in Fig. 1b), and the $A$-dependence, which was quite complicated at hadronic level (Fig. 1a), now becomes very simple. These partonic properties are therefore much easier to analyze than the hadronic ones.

The emission of partonic clans, i.e., of BGJs, is not only about the same forward and backward, which suggests that it is controlled by the central region, but it has also a very weak dependence on $A$. The increase of the backward-hemisphere multiplicities with $A$

[9] The proton-aluminium data of Ref. [15] are compatible within their large errors with the more precise proton-argon data of Ref. [14].

is mainly caused by the multiplicity growth of the backward BGJs shown in Fig. 2b, B. This strong nuclear effect on the size of the backward BGJs could be due to larger virtualities of their initial gluons and/or a rapid shower development including cascading.

It is interesting to discuss the $A$-dependence of $\bar{N}_p$ and $\bar{n}_{cp}$ more quantitatively. The absence of $A$-dependence for the forward $\bar{n}_{cp}$ is strikingly shown by Fig. 2b, F. For the backward $\bar{n}_{cp}$ (Fig. 2b, B), the pp-pXe comparison (for which the statistics is highest), when parametrized as an $A^{\alpha}$-dependence, gives $\alpha$ growing smoothly from $\sim 0.0$ at $y_0 = 0.5$ to $\sim 0.2$ at $y_0 = 3.5$ (the pAr-pXe comparison gives a growth from $\sim 0.0$ to $\sim 0.3$). In contrast, an $A^{\alpha}$-parametrization of $\bar{N}_p$ (Fig. 1b) gives a value of $\alpha$ decreasing very slowly from $\sim 0.1$ at $y_0 = 0.5$ to $\sim 0.07$ at $y_0 = 3.5$ for pp versus pXe (for pAr versus pXe, $\alpha$ increases from $\sim 0.11$ at $y_0 = 0.5$ to $\sim 0.13$ at $y = 1.0$ and decreases to $\sim 0.09$ at $y_0 = 3.5$, but the errors are too large to regard this as a significant effect).

The weak $A$-dependence of $\bar{N}_p$ suggests a discussion of the more sensitive quantity

$$R = [\bar{N}_p(\mathrm{pA})/\bar{N}_p(\mathrm{pp})] - 1 \tag{20}$$

which is plotted in Fig. 3. $R$ measures the relative increase of $\bar{N}_p$ due to multiple collisions in the target nucleus. It is natural to compare its $A$-dependence with that of the mean number $\nu'$ of "downstream" collisions suffered by the projectile after its first collision with a target nucleon. Using Poisson statistics for the successive collisions of the projectile, we find

$$\nu' = [\bar{\nu}/(1-e^{-\bar{\nu}})] - 1, \qquad \bar{\nu} = A\sigma_{\mathrm{inel}}(\mathrm{pp})/\sigma_{\mathrm{inel}}(\mathrm{pA}), \tag{21}$$

where $\bar{\nu}$ is the mean number of inelastic projectile-nucleon collisions. With $\bar{\nu} = 2.4$, 3.3 and 3.9 for argon, xenon and gold, we find $\nu' = 1.6$, 2.4 and 3.0 respectively. Within their rather large errors, the numbers of Fig. 3 are proportional to these $\nu'$-values, with some trend for $R$ to grow slightly faster. If confirmed with higher statistics data, this trend would indicate that after the first projectile-nucleon collision, the "wounded" projectile has a slightly larger inelastic cross-section. The most important finding, however, is the

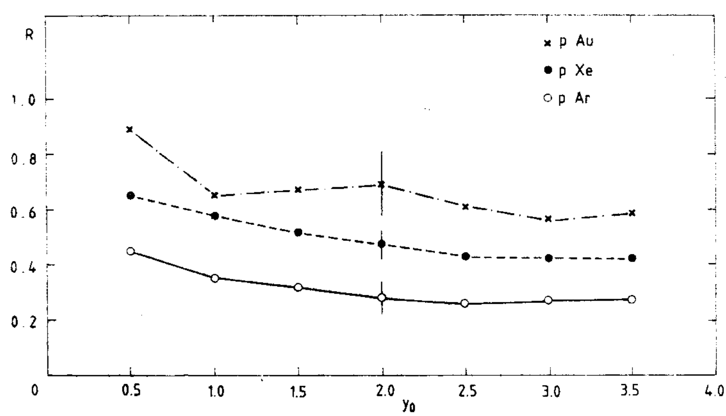


Fig. 3. The quantity $R$ defined in Eq. (20), extracted from the data in Fig. 1b

small value of $R/v'$, which decreases from $\sim$ 0.27 at $y_0 = 0.5$ to $\sim$0.17 at $y_0 = 3.5$. It demonstrates *a strongly reduced capability of the wounded projectile for further emission of BGJs.*

A last remark on the parton-level analysis of proton-nucleus collisions is that, for $y_0 \leqslant 2.5$, the $\bar{N}_p$ values in Fig. 1b agree, within errors, with the hadronic-level values $\bar{N}_h$ calculated for negative hadrons from Table II of Ref. [14]. Differences occur at larger $y_0$, and for $\bar{n}_c$ at all $y_0$. We have for the moment no explanation to offer for this empirical observation.

We expect that the extension of the preceding considerations to nucleus-nucleus collisions will be of interest for the experimental programme recently started with ultra-relativistic ion beams [16]. The most important new feature is that the distribution of the number $\nu$ of nucleon-nucleon collisions is very much wider than for proton-nucleus collisions, where it is always narrow. While the multiplicity distributions for small intervals of $\nu$ (as occur for a given impact parameter) are again expected to be of NB type, the averaging over all $\nu$ radically changes their shape. The analysis of these problems is beyond the scope of the present paper.

We end the present section by drawing attention to the implications of our Simplified Parton Shower Description for another aspect of ultra-relativistic ion collisions, to wit, the possible formation of a quark-gluon plasma. Evidently, in any QCD parton shower description of high energy collisions, there is first the *shower process* creating a system of quarks and gluons (let us call it the quark-gluon system, QGS), and then the *hadronization process* transforming the QGS into the system composed of the final hadrons. Does that mean that every collision produces a quark-gluon plasma which then makes a phase transition to hadrons? The answer is "No", firstly because the quarks and gluons resulting from the shower process are virtual (with masses $\sim Q_0 = 1$ GeV in the model calculation of Refs. [1b] and [2]), and also because they have a considerable degree of ordering in phase space, correlated with their colour states. They mostly fly away from each other and their growing separation rapidly gives rise to colour flux tubes (strings), whose breaking produces the hadrons. This is not what is meant by a plasma making the QCD phase transition. For the QGS to be considered a plasma, there should be considerable randomness in the phase space distribution of the partons (although full local thermalization need not be attained), and the lifetime of the QGS should be sufficiently long for them to reach virtual masses of order $\Lambda \sim 0.4$ GeV (in a dense medium, particles are never exactly on their mass shell).

More randomness in phase space is certainly expected for the QGS formed in central ultra-relativistic collisions of sufficiently heavy nuclei, and the higher parton density should go in the direction of giving more time before string formation. But it is difficult to define a sharp distinction between collisions with or without plasma formation, a difficulty also existing in other models, in particular the "colour rope" models invoking a very strong colour field as the source of parton creation [17].

The matter is of relevance for those experimental tests of plasma formation which refer to early times in the collision process, in particular the signals based on dilepton and J/ψ production. It should be realized that a sufficiently dense QGS at early times would create dileptons and would suppress J/ψ production in much the same way as usually discussed for a plasma [18], even if the partons have virtual masses $\sim$ 1 GeV and non--random motion. The situation is different for late-time tests based on the expectation that the plasma-to-hadron transition can last $\gtrsim$ 20 fm/$c$ because of the long time needed to liberate the large latent heat of the plasma. We want to stress, however, that early-time signals like J/ψ suppression remain of great interest, because they give information on the density and space-time development of the partonic system, even if the latter does not have all the characteristics of a plasma.

## 6. *Summary and concluding remarks*

Multiplicity distributions (MD) of charged hadrons in full phase space and in symmetric rapidity windows $|y| < y_0$ have been extensively studied in recent years. The same *negative binomial* (*NB*) *regularities* have been found experimentally in the charged hadron MDs of leptonic, semi-leptonic and hadronic collisions (see Ref. [5] for a review), and theoretically in the MDs predicted by the QCD parton shower model which best fits the $e^+e^-$ annihilation data [1a, 2]. In the latter case, similar NB regularities were found in the MDs of the final partons of the shower, and the partonic and hadronic NB parameters $\bar{n}$ and $k$ turn out to be approximately related by remarkably simple equations, which suggest a specific form of local parton-hadron duality [2].

Using the NB regularities and the fact that their parameters have particularly simple properties when expressed in terms of the "clan structure" of NB distributions, we were led in Ref. [2] to give an approximate but physically intuitive description of the parton shower as an independent emission of gluon jets with geometric MDs characteristic of self--similar cascades. We call them *bremsstrahlung gluon jets* (BGJ).

In the present paper, we show how this *simplified parton shower description* (SPSD) can be extended to leptoproduction and purely hadronic collisions. By assuming for the partonic NB regularities and for the parton-hadron duality relations the same degree of universality as observed experimentally for the hadronic MDs, we can calculate the partonic MDs from the measured hadronic ones. We then analyze their behaviour in the light of the properties found in Ref. [2] for showers initiated by a $q\bar{q}$ system (as in $e^+e^-$ annihilation) or a gluon-gluon (gg) system. Among the latter properties, the most unexpected and remarkable ones are the energy independence of the Poissonian MDs of BGJs in fixed rapidity windows $|y| < y_0$ over the wide range $\sqrt{s}$ = 22–2000 GeV, with about the same BGJs for both systems, but with the average numbers of BGJs about twice as large for the gg as for the $q\bar{q}$ system.

For leptoproduction, keeping in mind that the parton shower must originate from (anti)quarks since only these can be excited by the initial electroweak interaction, we predict that the multiplicity of BGJs must have a Poisson distribution independent of the hadronic energy $W$. As shown in Section 2 and Table I, this prediction is strikingly confirmed.

In Section 3 we performed the same analysis for low-$p_T$ proton-(anti) proton collisions, considering pseudorapidity windows $|\eta| < \eta_0$ for which data are available over the wide range $\sqrt{s}$ = 22–900 GeV, again with all NB properties required for our SPSD to apply.

Here the average number $\bar{N}_p$ of BGJs in fixed $\eta$-windows is found to increase slowly with $\sqrt{s}$, although the overall increase of multiplicity is controlled by the much larger growth of the BGJs themselves (Table II). We attribute the increase of $\bar{N}_p$ to the fact that in purely hadronic collisions both types of partons can be excited by the initial interaction. Using the above-mentioned property that BGJs are about twice as numerous in gluon-initiated showers than in quark-initiated ones, we successfully account for the data by assuming that the only energy variation is an increase with $\sqrt{s}$ of the fraction $\alpha$ of collisions which excite gluons. For the smallest window, $|\eta| < 0.2$, $\alpha$ is found to grow from $\lesssim 0.1$ at $\sqrt{s} = 22$ GeV to close to its maximum value of 1 at the upper energies. For larger windows $|\eta| < \eta_0$, making the reasonable assumption that $\alpha$ remains $\lesssim 0.1$ at 22 GeV, we calculate that the $\alpha$ values for the upper energies decrease rapidly with $\eta_0$ and that the gluonic contribution is concentrated in $|\eta| \lesssim 1.5$ (Table III).

This paper discusses the energy variations of the quantities in Tables I and II in the framework of our SPSD. The actual values and their comparison with the corresponding numbers for the $q\bar{q}$ and gg systems of Ref. [2] should also contain significant information for the further development of the still very incomplete SPSD. While we left this topic for later study, we addressed in Section 4 a question of principle raised by our approach to low-$p_T$ hadronic collisions, i.e., the identification of a large energy-momentum transfer justifying the formation of a perturbative parton shower. We propose that this hard element is to be found in high mass excitations caused by the longitudinal rapidity separation of partons in the early collision phase. On this basis, we believe that the NB regularities and the parton-hadron duality offer good prospects for developing step by step a unified parton shower description of all multiparticle production processes. As shown in Section 5, our method also has interesting implications for proton-nucleus collisions, where the partonic structure turns out to be much simpler than the hadronic one, and for the problem of quark--gluon plasma formation.

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