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CERN-EP 89-159

NUCLEAR RESEARCH

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CM-P00060910

CERN/EP 89-159
8 December 1989

HADRONIC PRODUCTION OF PARTICLES AT LARGE TRANSVERSE MOMENTUM:
ITS RELEVANCE TO HADRON STRUCTURE, PARTON FRAGMENTATION AND SCATTERING

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ABSTRACT

Soon after the discovery of scaling structure functions in deep inelastic lepton-nucleon scattering and its theoretical explanation based upon the parton model, a surprisingly large yield of hadrons with high transverse momenta was observed at the ISR. This had been suggested to be a manifestation of pointlike parton processes in purely hadronic interactions. An impressive body of detailed data has been collected so far by experiments investigating events including hadrons of high transverse momentum. The interpretation of the measurements in the framework of Quantum Chromo-Dynamics is rather far advanced, even on a quantitative level. Colliding hadrons are effectively "wide band" beams of various constituents, whose pointlike, strong interactions have been studied in unprecedented detail: Fragmentation mechanisms of various (systems of) final state partons were scrutinized in a way often complementary to lepton induced interactions. With this knowledge at hand, conclusions can be drawn concerning the effective composition of the colliding hadrons and the dynamics of constituent scattering. In this spirit the present review emphasizes both experimental and theoretical work done during the second decade of physics with single particles at large transverse momentum.

Submitted to Physics Reports

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1. INTRODUCTION AND HISTORY

The start of the experimental program at the CERN Intersecting Storage Rings (ISR) [ISR 71] coincided with a change of conception of hadron structure [PANO 68, BREI 69, BJOR 69a] and of strong interaction dynamics [FEYN 69a,b]. The study of particle production at large transverse momentum had a major impact on the further evolution of this understanding in the framework of Quantum Chromo-Dynamics (QCD).

This introductory section gives a historical survey of the major steps leading to this evolution from the point of view of participating experimentalists. Theoretical efforts will be discussed which were particularly stimulating to the experimental program. Especially those experimental results will then be considered which were essential for the survival of QCD based models as the most promising theoretical framework for the description of these reactions.

1.1 Deep inelastic scattering and the rise of the parton model

The quark model found its first successful application in the description of the spectroscopy of hadrons, their static electromagnetic properties and their weak decays [KOKK 69]. Its application to hadron dynamics was hampered by the complexity of the observed phenomena and the missing consistent theoretical basis for such an approach. The first spectacular breakthrough on the way to this goal was the observation of scaling behaviour of the structure functions in deep inelastic electron scattering [PANO 68, BREI 69] and its interpretation by Feynman and Bjorken [BJOR 69a, FEYN 69a, FEYN 72]. It was demonstrated that a hadron can be viewed – in a frame where its momentum is infinite – as a system composed of independent pointlike scattering centers. This lead to speculations about deep inelastic hadronic processes [BERM 70]. An explicit connection of the parton model with the quark model was formulated by Bjorken and Paschos [BJOR 69b].

Encouraged by this success Berman, Bjorken and Kogut [BERM 71] extended the parton model to describe the production of hadrons C (with transverse momentum p_T far in excesss of 1 GeV/c) in inclusive hadronic processes

$$A + B \rightarrow C + X$$

(*) "Phenomena concede views of the invisible" (Anaxagoras).

at very high energies. From photon exchange between quarks of the incoming hadrons the existence of "multiple cores" or "jets" originating from hadronization of the affected quarks was predicted. Moreover they stressed that, in addition to this electromagnetic mechanism, "elementary parton processes might be conceivably present associated with spin 1 gluon exchange. The expected cross section should exceed the electromagnetic one by a factor $\leq 10^4$ ". Assuming scaling behaviour of structure functions, as observed in deep inelastic lepton scattering, and conjecturing scaling of fragmentation functions in the relative momentum component along the jet axis (as well as a strong damping of momenta perpendicular to it), the invariant inclusive cross section was expected to exhibit scaling behaviour at fixed polar angle θ :

$$E \frac{d^3\sigma}{dp^3} (A + B \rightarrow C + X) = \frac{1}{p_T^4} f(x_T, \theta).$$

Here $(E, \vec{p})$ is the four momentum of hadron C, $x_T = 2p_T/\sqrt{s}$ is the transverse scaling variable, and $\sqrt{s}$ denotes the center of mass energy.

The phase space configuration of the final state was predicted to differ characteristically from the bulk of hadronic interactions: the large transverse momentum particles were expected to belong to two jets essentially coplanar with and at large angles to the incident beam. These sideways jets are due to the "hard" scattered quarks. The spectator quarks which do not participate in the hard interaction are expected to produce jets along the direction of the incoming hadrons. Hence these reactions should exhibit a four jet structure.

1.2 Observation of hadrons with large transverse momentum

An experimental program along this line of reasoning was not uncontroversial at that time since many theorists thought that "it is more important to understand processes with large cross sections because there can be very many competing small effects which would become important in calculating a cross section of small size".

The observation by three ISR experiments [ALPE 73, BANN 73, BÜSS 73] of a yield of large p_T particles much larger than expected from a naive extrapolation of the exponential behaviour observed at low p_T raised the

hope that the parton model might indeed be applicable to hadronic interactions. This stimulated an extensive study of large p_T particle production. Unfortunately, the detectors available for the first generation experiments at the CERN ISR and at Fermilab were not specifically designed for this kind of physics, hence decisive measurements were not performed for some time.

The cross section [BÜSS 73] for the process $pp \rightarrow \pi^0 + X$ (fig. 1.1) measured at a polar angle $\theta = 90^\circ$ in the center of mass system (c.m.s.) for different energies shows the expected scaling behaviour as a function of p_T . The measured exponent n of a power law in p_T ($n = -8.24 \pm 0.05$) differs, however, from the naive parton model expectation $n = -4$. This experimental result triggered a boom of theoretical speculation in the field.

Bjorken [BJOR 73] proposed to measure the parton-parton cross sections directly in a calorimeter experiment in order to avoid complications due to parton fragmentation.

A very detailed early discussion of how to observe directly parton-parton scattering in purely hadronic interactions is due to Ellis and Kisslinger [ELLI 74]. They considered first order scattering of quarks and gluons assuming scaling behaviour of the relevant distribution functions. Based on vector gluon exchange they predicted a four-jet structure. Moreover they pointed out that a particle at large p_T - called "trigger particle" for experimental reasons - is the leading fragment in the corresponding jet and should contain the scattered quark as a valence quark. Ellis and Kisslinger indicated the need for more sophisticated structure and fragmentation functions in order to reproduce the measured inclusive p_T spectra.

Emphasizing the importance of the difference between observed and predicted shape of the inclusive cross section at large p_T , Blankenbecler and Brodsky [BLAN 74] proposed quark-hadron scattering as an alternative dominant source of large p_T particles. Their predictions rely heavily on counting rules, which relate the shape of measured invariant inclusive cross section at given polar angle θ to the number n_a of "active" quarks q in the hard scattering process:

$$E \frac{d^3\sigma}{d^3p} (A + B \rightarrow C + X) \sim p_T^{-2(n_a - 2)} f(x_T, \theta) .$$

For large x_T the scaling function $f(x_T, \theta)$ should be given by a simple power law in x_T . One expects $n_a = 6$, i.e. $n = 8$, if $q + \pi^0 \rightarrow q + \pi^0$ is the dominant subprocess of high p_T π^0 production in pp collisions. This is in agreement with the measurements in fig. 1.1. This model, called constituent interchange model (CIM), was extremely successful in predicting the p_T dependence of the measured single particle spectra. Since cross sections represented the most detailed experimental information for a few years, the CIM dominated for some time the interpretation of large p_T data [SIVE 76, ELLI 77].

The matrix elements for parton-parton scattering in lowest order QCD were derived by Combridge, Kripfganz and Ranft [COMB 77], and by Cutler and Sivers [CUTL 77,78].

1.3 Event structure

The experimental study of the phase space structure of events selected by a large p_T trigger particle allowed to differentiate between the different theoretical approaches. The observation of secondaries in a "trigger jet", which were not due to two-body decays of a resonance yielding the trigger as well [BÜSS 76, DARR 76, DELL 76a], and the detection of a jet structure [DELL 76b] in the azimuthal angular region opposite to the trigger particle ("away region") disfavoured the CIM and revived the interest in investigations [FOX 76] based upon the quark parton model.

In a series of papers Feynman, Field and Fox [FIEL 77, FEYN 77, FEYN 78] showed how a QCD inspired parton model predicted the full phase space structure of events with a triggering particle at large p_T . Three experimental results strongly influenced the development: The observed symmetry around c.m.s. rapidity $y = 0$ of the "away" particle distribution [DELL 76a], even for large p_T triggers at an extreme forward polar angle θ ($\theta = 10^\circ$), constrained strongly the basic two-body parton scattering cross sections [FOX 76, KRIP 77]. Moreover, the observed azimuthal asymmetries of leading spectator fragments [DELL 75], and the fact that away jets are less coplanar [DELL 77] than predicted by parton models, lead various authors [FOX 76, DELL 77, FEYN 78] to introduce a large effective transverse momentum of partons inside hadrons.

With this important basic knowledge at hand a new generation of experiments with single hadron triggers and further theoretical analyses have focussed on many more details of hard scattering processes. The results of these more recent investigations will be discussed in the following. Contributions published after early summer 1989 are not considered.

For more complete presentations and discussions of earlier results one should consult some of the previous excellent reviews [ELLI 77, JACO 78, DARR 80, MCCU 81, MOLL 84].

1.4 Jet triggers

An in many aspects complementary method of investigating hard parton-parton scattering dynamics was pursued in parallel making use of calorimetric detection of large angle jets comprising all parton fragments as suggested in ref. [BJOR 73]. Hence inclusive production of jets is probed, which corresponds to larger cross sections than the production of those particular jets which fragment into a fast leading fragment, i.e. the trigger hadron. This reduction of cross section, called "trigger bias", was investigated in detail by Jacob and Landshoff [JACO 76].

In fixed target experiments with beams of a few hundred GeV/c this approach turned out to be unsuccessful. Obviously, cross sections for isotropic high multiplicity events are dominating jet-jet cross sections. Later on, ISR experiments employing calorimeter triggers yielded clear evidence for large angle jets, reflecting a stronger increase with $\sqrt{s}$ of cross sections for hard parton-parton scattering than for high multiplicity events ($\sim \ln s$). The history of this type of experimentation is very lucidly summarized by Tannenbaum [TANN 89].

With the advent of the SppS collider calorimeter triggers were producing such clear jet structures that the interpretation in terms of hard parton-parton scattering in the framework of QCD is on extremely solid ground. A thorough review of the more recent relevant collider results was undertaken by DiLella [DILE 85]. This subject is therefore not addressed in detail in this report which focusses on the physics with individual hadrons at large transverse momentum.

* * * * *

The main theoretical concepts needed to understand deep inelastic hadronic processes will be given in the next section.

In sect. 3 simple average properties of events triggered by large p_T particles will be presented with emphasis on the four-jet pattern.

Sect. 4 is then devoted to a discussion of parton fragmentation, i.e. of momentum and quantum number structures of all 4 jets. The concept of flavour tagging of parton jets is inferred from the data and, finally, the question of Lorentz non-invariance of current fragmentation schemes is addressed.

Valence and sea quarks as well as gluons and diquarks are shown in sect. 5 to be the hadron constituents actively involved in hard processes. Included in this section is also a different aspect of hadron structure, i.e. the transverse motion of partons.

Hard parton-parton interactions, are considered in sect. 6. Experimental information on the angular dependence of the relevant matrix elements is summarized. Furthermore, methods to separate experimentally parton subprocesses are described. The section is concluded by a presentation of inclusive cross sections.

By structuring the present review in this way it was attempted to split the large body of data into nearly self-contained parts. It is obvious that experimentally one observes a convolution of all effects discussed; hence a separation into independent sections is not strictly possible. As a consequence, various arguments may be repeated for better understanding; in other cases arguments presented in one section will be refined and completed later on. This seems to be unavoidable.

Conclusions are drawn in sect. 7, where also a list of topics awaiting further study is given.

Definitions of the symbols used are given in the Appendix which includes also formulae for the matrix elements for various hard parton processes. Throughout, the transverse momentum of the trigger particle is denoted by " p_T ", while " p_t " is used for associated hadrons in the same event; consequently the symbol " p_T " is also applied for inclusive spectra.

FIGURE CAPTION

Fig. 1.1 Invariant cross section for the reaction $pp \rightarrow \pi^0 + X$ as function of the transverse momentum for five centre-of-mass energies [BÜSS 73].

Fig. 1.1

2. THEORETICAL BACKGROUND

From scale invariance observed in lepton-nucleon scattering it was inferred that hadrons may be represented by a set of point-like constituents. Similarly, in hadron physics, the observation that the inclusive cross section exceeds at large p_T by orders of magnitude the standard e^{-6p_T} shape at low p_T was interpreted as an indication of granularity inside the hadrons. This was shown to be well described in the following picture:

- Each hadron may be considered as dissociated into constituents, which behave as free states during the short interaction time.
- At least 2 constituents then suffer a "hard" large angle scattering.
- Eventually all four parton states - scattered constituents and spectator's "remnants" - materialize into hadrons.

The underlying idea is to describe the global process as a succession of steps, well enough separated in time, such that it may be represented by the product of the relevant probability densities (factorization hypothesis). The resulting inclusive cross section turns out to behave as a power law in p_T ; although small, it should eventually dominate over the exponential component, and does so for $p_T > 2 \text{ GeV}/c$.

It is the purpose of this section to review briefly the phenomenological consequences that one may extract from this hard scattering picture. The first sub-section defines the global formalism and the following ones describe the individual sub-processes. Quantitative predictions are compared to the data throughout the other sections.

There is currently about one dozen of algorithms in the framework of QCD, some of which are updated at least once per year. A complete survey is not attempted in this report. Additional details may be found in ref. [WEBB 86]. However, the main ingredients of earlier versions are described and potential improvements are pointed out.

2.1 General framework

In various models, the constituents involved in the hard scattering were identified with hadrons or partons. In the following the second possibility will be considered which is most probably dominant. More

precisely, the Quark Parton Model (QPM) will be discussed, where charged point-like constituents (partons) are identified with quarks [FEYN 72]. One notices that neutral partons are needed as well (see sect. 2.2). These are the gluons which are naturally introduced through Quantum Chromodynamics (QCD) [FRIT 73]. Unless otherwise mentioned, the implicit basis for what follows is the QPM in the QCD framework.

A direct application of the probabilistic picture sketched above gives the single particle differential cross-section for the inclusive reaction $A + B \rightarrow C + X$:

$$d\sigma = \sum G_{A \rightarrow a}(x_a) dx_a \cdot G_{B \rightarrow b}(x_b) dx_b \cdot \left[\frac{d\sigma}{d\hat{t}} \right] d\hat{t} \frac{d\hat{\phi}}{2\pi} \cdot D_{c \rightarrow C}(z) dz$$

(see [SIVE 76] for a derivation).

Here (A, B) are the initial hadrons, (a, b) are the corresponding partons which participate in the hard scattering, (c, d) are the scattered partons and C is the hadron observed at large p_T , usually called the "trigger" particle. The structure functions G and the fragmentation functions D describe the transitions, respectively, from hadrons to partons and, conversely, from partons to hadrons.

$G_{A \rightarrow a}(x_a) dx_a$ represents the number of partons of type "a" found inside the initial hadron A, at a fractional longitudinal momentum x_a within dx_a . The scattering cross section for the process $a + b \rightarrow c + d$ is represented by $[d\sigma/d\hat{t}]$. The azimuthal angle $\hat{\phi}$ is taken about the direction of the partons a, b in their c.m.s. and $\hat{t}$ is the squared 4-momentum transfer between partons a and c. $D_{c \rightarrow C}(z) dz$ represents the number of hadrons of type C found at a fractional longitudinal momentum z within dz from the hadronization of the parton c. The sum $\sum$ runs over the possible parton species a, b, c, d, compatible with the nature of A, B, C. The variables x and z should be estimated in infinite momentum frames. It is strictly equivalent to consider ratios of light cone variables ($p_{\pm} = E \pm p_1$). This exhibits Lorentz invariance under longitudinal boosts.

This expression must be recast to exhibit the invariant phase-space element d^3p_C/E_C . Using $\hat{t} d\hat{\phi} dz = (2/z) d^3p_C/E_C$, one gets:

$$\frac{d\sigma}{d^3p_C/E_C} = \sum [G(x_a) dx_a \cdot G(x_b) dx_b \cdot \frac{1}{\pi} \left[\frac{d\sigma}{df} \right] D(z)/z]$$

Hadronization of the parton d and of the "remnants" of the initial hadrons integrate out due to implicit summation over all possible final states. At large p_T the cases where the hadron C is formed from one of the "remnants" should be negligible.

The formula may be extended to take into account transverse momentum components $\vec{k}_{ta}$ and $\vec{k}_{tb}$ of partons inside the initial hadrons, and as well the transverse momentum $\vec{q}_t$ of the hadron C relative to its parent parton c. For this, a term $w(\vec{k}_{ta}) w(\vec{k}_{tb}) w'(\vec{q}_t) d\vec{k}_{ta} d\vec{k}_{tb} d\vec{q}_t$ must be introduced under the integral, where w and w' represent the respective probability density functions.

The 2-particle cross section may be treated in a similar way if the two hadrons C_1 and C_2 originate from the partons c and d. The essential modification consists in replacing $D_{c \rightarrow C}(z)$ by $D_{c \rightarrow C_1}(z_1) \cdot D_{d \rightarrow C_2}(z_2)$ and by using a different Jacobian for the change of variables [ELLI 74]. In the cases where C_1 and C_2 come from the same parton one needs another fragmentation function, namely $D_{c \rightarrow C_1 C_2}(z_1, z_2)$. Further inclusive n-particle cross-sections require higher order fragmentation so that one will eventually resort to totally exclusive simulation methods.

In the following, more details of the different stages of the global process will be considered. In addition sub-section 2.4 will treat parton evolution, a complication of our simple picture due to possible transitions of the partons a, b, c, d (such as $q \rightarrow q + g$) as a necessary consequence of QCD.

2.2 Hadron structure

By isospin symmetry - local in x - the functions $G_{p \rightarrow a}$ and $G_{n \rightarrow a}$, for protons and neutrons are related. The density functions $G_{p \rightarrow a}$ are conventionally renamed using the symbol of the parton a.

$$u(x) = G_{p \rightarrow u}(x) = G_{n \rightarrow d}(x) \quad \bar{u}(x) = G_{p \rightarrow \bar{u}}(x) = G_{n \rightarrow \bar{d}}(x)$$

$$d(x) = G_{p \rightarrow d}(x) = G_{n \rightarrow u}(x) \quad \bar{d}(x) = G_{p \rightarrow \bar{d}}(x) = G_{n \rightarrow \bar{u}}(x)$$

$$s(x) = G_{p \rightarrow s}(x) = G_{n \rightarrow s}(x) \quad \bar{s}(x) = G_{p \rightarrow \bar{s}}(x) = G_{n \rightarrow \bar{s}}(x)$$

$$g(x) = G_{p \rightarrow g}(x) = G_{n \rightarrow g}(x)$$

These functions are constrained by conservation laws (sum rules) and may in principle be determined in electroweak interactions. The conservation of charge, Q , 3rd component of isospin, I_z , and strangeness, S , impose the following equations for protons:

$$Q = 1 = \frac{2}{3} \int_0^1 (u - \bar{u}) dx - \frac{1}{3} \int_0^1 (d - \bar{d}) dx - \frac{1}{3} \int_0^1 (s - \bar{s}) dx$$

$$I_z = \frac{1}{2} = \frac{1}{2} \int_0^1 (u - \bar{u}) dx - \frac{1}{2} \int_0^1 (d - \bar{d}) dx$$

$$S = 0 = \int_0^1 (s - \bar{s}) dx$$

One can formally distinguish the distributions of valence and sea components, u_v and u_s , by redefining:

$$u_v = u - \bar{u}, \quad \bar{u}_v = 0, \quad u_s = \bar{u}_s = \bar{u} \quad (\text{and similarly for } d \text{ and } s \text{ quarks}).$$

This is guided by the physical meaning one wants to attach to the concepts "valence" and "sea", but does not introduce any restrictive hypothesis. It then follows for protons that:

$$\int_0^1 u_v dx = 2; \quad \int_0^1 d_v dx = 1; \quad \int_0^1 s_v dx = 0$$

Notice that the functions u_v , d_v , s_v need not be positive semi-definite for all x -values, contrary to the case of the original functions u , d , s , since the latter were introduced by a probabilistic concept. In particular s_v need not be identical to 0 although this is usually assumed.

The total momentum carried by the partons must equal that of the proton, thus giving the following relation where the gluons turn out to be essential, and where heavier quarks are neglected:

$$\int_0^1 x(u + \bar{u} + d + \bar{d} + s + \bar{s} + g) dx = 1$$

2.2.1 Dependence on x

The inclusive charged lepton-nucleon scattering cross-sections may be expressed as [BJOR 69a]:

$$\frac{d^2\sigma}{d\Omega dE'} = \frac{4\alpha^2 E'^2}{Q^4} \{2W_1 \sin^2(\theta_L/2) + W_2 \cos^2(\theta_L/2)\},$$

where a lepton of initial energy E is scattered by a nucleon at rest (of mass M) through an angle θ_L to a final energy E' . The structure functions W_1 and W_2 depend on the type of leptons and nucleons involved and are functions of 2 variables, x and Q^2 :

$$Q^2 = 4EE' \sin^2(\theta_L/2), \quad x = Q^2/2M\nu, \quad \nu = E - E'$$

It was predicted [BJOR 69a] that νW_2 and W_1 should become functions of x alone for large Q^2 (and, of course, large ν). This was experimentally tested to a good first approximation (sect. 2.2.2). A third function W_3 needs to be included to describe neutrino-nucleon interactions.

The structure functions may be expressed in terms of the 6 quark densities (neglecting heavier quarks). For ep and en scattering one obtains:

$$F_2^{ep} = \nu W_2^{ep}/x = \frac{4}{9} (u + \bar{u}) + \frac{1}{9} (d + \bar{d} + s + \bar{s})$$

$$F_2^{en} = \nu W_2^{en}/x = \frac{4}{9} (d + \bar{d}) + \frac{1}{9} (u + \bar{u} + s + \bar{s})$$

Varying the types of projectile ($e, \mu, \nu, \bar{\nu}$) and target ($p, n, \dots$), one extracts the structure functions from which one can reconstruct quark densities. For strange quarks only the sum ($s + \bar{s}$) can be determined.

A compilation [ARNE 89] of measured structure functions $x \cdot u_v(x)$ is given in fig. 2.1. The quark densities diverge at $x \approx 0$ (but less strongly than $1/x$), and are proportional to $(1-x)^m$ at $x \approx 1$ with $m \approx 3-4$ for valence quarks and $m \approx 7-10$ for sea-quarks.

The total momentum carried by quarks and antiquarks is experimentally found to be about 50% of that of the initial nucleon (see e.g. [DEDE 75]). Consequently, this leaves an equal amount for neutral partons, probably gluons, which cannot be directly probed by electroweak interactions. A potential rôle of diquarks [SZCZ 89] is discussed in sect. 2.2.5.

2.2.2 Dependence on Q^2

When studied in more detail, the structure functions turn out to depend as well on Q^2 as exemplified for F_2 in fig. 2.2 [AUBE 87]. The main feature is, with increasing Q^2 , an increase at low x and a corresponding decrease at larger x . This may be explained in the framework of QCD. An increase in Q^2 implies that partons are hit harder, augmenting the amount of gluon bremsstrahlung, or, equivalently, are probed with a finer resolution. At any scale, quarks and gluons may be resolved into constituents, i.e. again into quarks and gluons, similarly to fractal curves. In other words, there is no Q^2 asymptotia of naked partons, where Q^2 scaling would occur.

The problem of extracting parton densities from experimental data in the context of QCD has been extensively studied in ref. [BURA 77, FOX 77, GLÜCK 82, DUKE 84], typical results can be found in ref. [ABRA 82,83]. Usually, one starts from parametrized expressions of the parton densities at some reference scale Q_0 , e.g. $Q_0 \approx 2$ GeV, then the QCD evolution equations (sect. 2.4) allow to calculate densities at Q , which are then compared to the data. The transitions $q \rightarrow q + g$ and $g \rightarrow q + \bar{q}$ give a handle on the gluon density, so that an estimation of both quarks and gluons distributions may be obtained.

Practical results are best presented as parametrizations whose coefficients are adjusted to the experimental data. For example [DUKE 84]:

$$\begin{aligned} x(u_v + d_v) &= N_{ud} x^{\eta_1} (1-x)^{\eta_2} (1 + x\gamma_{ud}) \\ xd_v &= N_d x^{\eta_3} (1-x)^{\eta_4} (1 + x\gamma_d) \end{aligned}$$

where $\eta_1 \dots \eta_4$ and γ_{ud}, γ_d are quadratic functions in $\ln\{\ln(Q^2/\Lambda^2)/\ln(Q_0^2/\Lambda^2)\}$ and N_{ud}, N_d are normalization constants determined by the equations $\int_0^1 u_v dx = 2, \int_0^1 d_v dx = 1$.

One should keep in mind that the parameters obtained do depend implicitly on the scale Λ (0.2 and 0.4 GeV were used in ref. [DUKE 84]). Thus they should not be replaced by one's favorite value. Figs 2.3 show quark and gluon structure functions using the parameters of ref. [DUKE 84] (set 1 with $\Lambda = 0.2$ GeV); the dependence on Q^2 is clearly visible. The ratio $d(x)/u(x)$ in fig. 2.4 is seen to decrease with x .

2.2.3 Dependence on transverse momentum of partons

Finally, parton densities depend also on the parton transverse momentum k_t . Two different sources can be identified. First, there is a contribution from internal motion of the partons inside the proton; it is assumed to factorize, i.e. $G(x, Q^2)$ is replaced by $G(x, Q^2) * w(\vec{k}_t)$.

The second source is characteristic of QCD. It originates from gluons emitted by the partons before a hard scattering (sect. 2.4). This component should depend both on x and Q^2 , since the larger Q^2 the more gluons are emitted.

In some studies, these 2 components are effectively described by one function $w(\vec{k}_t)$, so that factorization is assumed to be a fair approximation. Typically an exponential in k_t or k_t^2 is used for the density in k_t^2 , with a scale parameter adjusted such that $\langle k_t \rangle \approx 0.6-0.8$ GeV/c, possibly x -dependent. From model calculations of inclusive single particle cross sections, one finds that the transverse momenta of the two partons participating in the hard scattering align themselves preferentially with the trigger (sect. 5.6). Indeed, the high transverse momentum of the trigger is not only due to large angle scattering angle. Large values of the component of $\vec{k}_t$ of each parton along the trigger direction contribute as well. Quantitative predictions for the cross sections reflect strongly the choice of $w(\vec{k}_t)$ for $p_T \leq 4$ GeV/c, but this effect becomes negligible above 5 GeV/c (sect. 6.4).

2.2.4 Structure of pions

Up to now only the parton densities in protons and neutrons were considered which can be obtained from lepton scattering experiments. The case of antiprotons and antineutrons is easily dealt with using C-parity. For pions and kaons there is no such easy way. Estimates may be obtained from Drell-Yan experiments [KENY 82].

2.2.5 Diquarks

There is now ample evidence for the contribution of diquarks to hard scattering (sect. 5.5). These constituents do not have the pointlike structure of quarks and gluons. Their size may be parametrized by a form factor $F(Q^2) = 1/(1 + Q^2/M_0^2)$. The density of matter in the diquark, $f(r)$,

may be related to its form factor, $F(q)$, $q = \sqrt{Q^2}$, by a Fourier transform. In order to get a gross order of magnitude, one assumes spherical symmetry of $f(r)$ (the relation is now a Hankel transform):

$$f(r) = \sqrt{\frac{2}{\pi K}} \int_0^\infty F(q) \frac{q \sin(qr/K)}{r} dq ,$$

$$F(q) = \sqrt{\frac{2}{\pi K}} \int_0^\infty f(r) \frac{r \sin(qr/K)}{q} dr .$$

One may then estimate the mean squared radius of the diquark by the relation

$$\langle r^2 \rangle = E(r^2) = -6K^2 \left. \frac{d \ln F}{dq^2} \right|_{q=0}$$

which, for the form factor considered, becomes $\langle r^2 \rangle = 6K^2/M_0^2$.

Such a Q^2 dependence might explain, at least in part, the non-scaling of structure functions. Many studies have been devoted to this subject but no consensus has been reached. However, one should keep in mind that the fraction of the nucleon momentum carried by gluons may be significantly lower than the claimed value of 50%.

The main source of this estimation comes from the integral of the structure function $F_2 = \nu W_2/x$ through observations of inclusive cross sections induced by neutrinos on isoscalar targets. In the absence of diquarks, one expects $\int_0^1 (F_2^{\nu p} + F_2^{\nu n}) dx$ to be proportional to $\int_0^1 (u + \bar{u} + d + \bar{d} + s + \bar{s}) dx$. Observations of the integral of F_2 [DEDE 75] allow then to deduce the fractional momentum of gluons, which is found to be close to 50%. In case diquarks exist as constituents of the nucleon, they contribute to the integral of F_2 . Unfortunately, they do not only contribute through their density, but a multiplicative form factor is involved. It follows that there does not exist any more a simple relation between the integral of F_2 and the gluon momentum contribution. Strictly speaking, the gluon might carry either less or more momentum when diquarks are allowed. However, it seems probable that it carries less.

An extreme case was studied [FRED 82, FRED 83] where there are no gluon constituents inside the nucleon though evidently gluons are still QCD mediators (in this case they are used to bind the diquark). Irrespective of the merits of this model, one should notice that the absence of gluon constituents in the nucleon does not violate sum rules.

A much less ambitious goal was tackled [BREA 85b, LOHS 85], purely aiming at an understanding of high p_T proton production. One considers the proton to exist in two non-interfering states: a conventional non-diquark state $|u, u, d, g, \text{sea}\rangle$ and another state $|\tilde{u}\tilde{d}, \tilde{u}, \tilde{g}, \tilde{\text{sea}}\rangle$ with respective probabilities β and $1 - \beta$. The conventional parton distributions were kept unmodified, with their standard fitted shapes. When in the diquark state, one assumes for the proton constituents the following simple shapes (a tilda denotes such distributions), expressed in terms of "conventional" distributions (without tilda):

$$\tilde{u}\tilde{d} = \frac{1}{2} u + d; \quad \tilde{g} = g; \quad \tilde{d} = 0; \quad \tilde{u} = \frac{1}{2} u$$

In addition, two variants of this scheme were used, one where $\tilde{g} = 0$ and one where β is a linear function of x . A description of the comparison with data will be found in sect. 5.5.

This introduction of diquarks as nucleon constituents gives rise to a contribution of the gluons to the total momentum varying from 0 to 14% as compared to the canonical 50%. It should be noticed, however, that no comprehensive study including all "relevant" data was performed so far.

2.3 Parton-Parton scattering

Now the heart of the high p_T process is addressed, i.e. the hard interactions between two partons. In principle, one may compute the 2-parton $\rightarrow$ N-parton cross sections from Feynman diagrams. The smallness of the coupling constant allows to restrict the calculations to lowest orders. Practically, one considers only the simplest case of diagrams with 2 vertices. There are calculations on $2 \rightarrow 3$ transitions [BERE 81, ELLI 85], however, there is no clear-cut distinction from parton evolution (sect. 2.4) and consequently some risk of double counting [GOTT 85].

2.3.1 Scattering amplitude

The parton-parton differential cross section is written as:

$$\frac{d\sigma}{d\hat{\phi} d\hat{t}} = \frac{\alpha_s^2}{2 \hat{s}^2} |\Gamma|^2$$

where $\hat{s}$, $\hat{t}$, $\hat{u}$ are the Mandelstam variables and Γ is the matrix element.

2.3.2 Matrix elements

Appendix 2 gives the formulae for $|\Gamma|^2$ for most of the relevant subprocesses together with graphs of some distributions as function of angle. Although their vector nature is now well established, formulae are given for both scalar and vector gluons. The color factors are exhibited in such a way that they can be set to 1 in a colorless model. The partons are assumed to be massless in these calculations inducing well known divergences at $\hat{t} = 0$ and $\hat{u} = 0$.

The case of scalar diquarks is included as well (sect. 5.5).

2.3.3 Color lines

Although color is not observable, one has to take into account color correlations between the parton states participating in the scattering. Hadronization of color neutral parton systems is currently formulated in a non Lorentz-invariant way, thus definite reference frames must be chosen. In case of the Lund model (sect. 2.5.1.2), it is postulated that the reference frame is that of the color string which connects 2 partons. Whatever its justification, this hypothesis poses the problem of the association of "color lines" to "parton lines". Unfortunately this cannot be done in a unique way. One usually attaches the labels of long strings, or short strings, to the configurations where the 2 partons effective mass is large, or small, respectively. Using for colors c the familiar notation $c \in \{r, g, b\}$, there are 6 combinations $c\bar{c}'$ ($c' \neq c$), which represent gluons. In addition one can make 3 orthogonal linear combinations of the $c\bar{c}$ associations. The symmetric combination $r\bar{r} + g\bar{g} + b\bar{b}$, a color singlet, does not represent a gluon but the 2 other combinations do. In the case of quark-quark scattering (fig. 2.5), one may draw the color lines in two different ways. One corresponds to long strings, (fig. 2.5(a)). The other, where the gluon is in one of the two states mentioned earlier, corresponds to short strings (fig. 2.5(b)). For $N_c = 3$ colors, the ratio of cross sections is $N_c^2 = 9$, in favor of the long string configuration. In case all colors are identical, there is an interference term which cannot be treated directly. A prescription used in the Lund model is to share the contribution of this term between long strings and short strings configurations, keeping the cross section, i.e. the sum of the three terms, fixed.

2.3.4 Q^2 scale

The strong coupling constant, α_s , depends on a scale Q^2 which characterizes the "hardness" of the parton-parton interaction: the larger Q^2 , the harder the interaction. QCD is an asymptotically free theory in the sense that partons may be considered as free states at infinite Q^2 , where the coupling constant α_s is equal to 0. For large values of Q^2 , though experimentally accessible, $\alpha_s(Q^2)$ may be small enough so that perturbation methods may be applied. In hadronic interactions, only the Leading Log Approximation (LLA, sect. 2.4) has been considered up to now in most cases. Up to 2-loop order, $\alpha_s(Q^2)$ may be written as:

$$\alpha_s(Q^2) = \alpha_o(Q^2) \cdot [1 - \rho \ln \ln(Q^2/\Lambda^2)] + O[(\alpha_o)^3]$$

with the following notations:

$$\alpha_o = \frac{4\pi}{\beta_o \ln(Q^2/\Lambda^2)}, \quad \beta_o = (11/3) N_c^2 - (2/3) N_f$$

$$\rho = \frac{\beta_1}{4\pi \beta_o}, \quad \beta_1 = (34/3) N_c^2 - (10/3) N_c N_f - N_f(N_c^2 - 1)/N_c$$

with $N_f = 3$ flavours when only light quarks are considered.

In the case of lepton-nucleon scattering, Q^2 is defined by the four-momentum transfer squared of the lepton. The parameter Λ is to be determined experimentally, as a measure of the scale-breaking. Notice that this parameter will depend on the order of the calculations, in particular $\alpha_s(Q^2) = \alpha_o(Q^2)$ should be used if LLA is assumed.

The case of hadronic interactions is more complex, since there is no unique definition of Q^2 . Different formulations based on the Mandelstam invariants of parton-parton scattering have been tried, such as $Q^2 = \hat{s}$. The choice of $Q^2 \sim p_T^2$ (where p_T is characteristic of the scattered partons) has the advantage that it is more directly connected to observables. The most popular definition was proposed by [FEYN 78]:

$$Q^2 = \frac{2 \hat{s} \hat{t} \hat{u}}{\hat{s}^2 + \hat{t}^2 + \hat{u}^2}$$

Quoted values for Λ range from $\Lambda = 0.2$ to 0.6 GeV [VOSS 88]. Notice that the value of Λ depends on the choice of the definition of Q^2 and, in addition, on the order of approximation.

2.3.5 Dependence on p_T

Early calculations [ELLI 74] predicted that the inclusive single particle cross section should behave as:

$$\frac{d\sigma}{dp_T^2} \xrightarrow{s \rightarrow \infty} \frac{f(x_T, \theta)}{s^2}$$

where f is an a priori unknown function of $x_T = 2 p_T / \sqrt{s}$ and of θ . This results solely from the hypothesis of scattering by vector gluon exchange, given that structure and fragmentation functions essentially scale. This can be checked experimentally by measuring ratios of cross sections at different energies $\sqrt{s}$ for fixed x_T and θ . At fixed x_T , s^2 is proportional to p_T^4 hence one arrives at $d\sigma/dp_T^2 \sim p_T^{-n}$, $n = 4$. All measurements, however, gave values of $n \approx 6-8$, first at the ISR and later at the SPS collider (sect. 6.4).

In fact, many corrections must be applied to this naive expectation. The decrease of $\alpha_s(Q^2)$ with increasing Q^2 will naturally induce a larger value for n . The Q^2 variation of both structure and fragmentation functions has a similar tendency. The transverse momenta k_t and q_t have also some influence up to $p_T \approx 5$ GeV/c. The largest effect comes from gluon bremsstrahlung occurring in parton evolution (sects 2.4 and 6.3.3.1).

Therefore, in hindsight, the early ISR results on the shape of inclusive spectra were giving very first experimental evidence for scale breaking effects in the framework of QCD.

2.3.6 Jet cross section

The formula (sect. 2.1) representing single particle inclusive cross section can be trivially reduced to the case where jets are considered:

$$\frac{d\sigma}{d^3p_{\text{jet}}/E_{\text{jet}}} = \sum \int (G(x_a) dx_a \cdot G(x_b) dx_b \cdot \frac{1}{\pi} \left[\frac{d\sigma}{d\hat{t}} \right])$$

This is obtained through the identification of the jet of hadrons with their parent parton so that the fragmentation function may be replaced by $\delta(z - 1)$.

2.4 Parton evolution

A detailed study of parton evolution is now addressed, in view of a representation in terms of exclusive parton cross sections.

The existence of gluons, as postulated by QCD, implies transitions between quarks and gluons, namely:

$$q \rightarrow q + g, \quad g \rightarrow g + g, \quad g \rightarrow q + \bar{q}$$

These branching processes may be described by the Altarelli-Parisi relations [ALTA 77] which, given a probabilistic interpretation, represent parton evolution. One considers first the case of the quark branching process $q \rightarrow q + g$, then extensions needed for incident gluon are indicated. A discussion of the possible choices of variables describing the branching process and of kinematical constraints will follow. Finally, evolutions of time-like and space-like partons will be compared.

2.4.1 Master equations

Consider a parton of virtuality t (defined, for example, by its invariant mass-squared) which splits into 2 partons. The distribution of energies between the 2 outgoing partons is represented by a variable z , e.g. the ratio of the longitudinal momenta of the outgoing and incoming partons in the process $q \rightarrow q + g$. Other choices of the variables t and z will be discussed later. At this level it is assumed that these variables are sufficient to define the kinematics of the branching process, once the energy of the incoming parton and the masses of the outgoing partons are given. The basic equation describes the probability that a transition occurs and the distribution of z :

$$dT(t, z) = \alpha_s(t) (dt/t) f(z) dz$$

where $f(z)$ depends on the process considered:

$$f_{qq}(z) = C_2(R) (1 + z^2)/(1 - z) \quad \text{for } q \rightarrow q(+g),$$

$$f_{qg}(z) = (1/2)[z^2 + (1-z)^2] \quad \text{for } g \rightarrow q (+\bar{q}),$$

$$f_{gg}(z) = 2C_2(G) [z/(1-z) + (1-z)/z + z(1-z)] \quad \text{for } g \rightarrow g (+g).$$

The factors C_2 are related to the number N_c of colors by $C_2(R) = (N_c^2 - 1)/2N_c$ and $C_2(G) = N_c$.

Notice the standard notation $f_{qg}(z)$ describing the process $g \rightarrow q (+\bar{q})$, where $z = z_q$ corresponds to the parton q . Using the identity $z_q + z_{\bar{q}} = 1$, one may deduce e.g. that $f_{qg}^-(z) = f_{qg}(1-z)$ and, in particular, that $f_{gq}(z) = f_{qg}(1-z) \neq f_{qg}(z)$.

Values of z are confined to the interval $[0, 1]$, but one finds logarithmic divergences at $z = 1$ for f_{qg} , and at $z = 0, 1$ for f_{gg} . Such infrared divergences, also existing in QED, correspond to the radiation of gluons of vanishing momentum, which evidently may occur arbitrarily frequently. They will require an ad hoc treatment such as a reduction of the range in z , or better a regularisation using parton masses.

Starting from a parton of virtuality t_0 , e.g. $t_0 \sim Q^2$, one needs to express the probability $dP(t, z|t_0)$ that, conditional to the fact that the parton kept its integrity in the range between t_0 and t (excluding t), a branching process occurs at t with energy sharing z . Labelling the first component by $E(t|t_0)$ (evolution from t_0 to t) and using for the second the basic definition of $dT(t, z)$ one gets:

$$dP(t, z|t_0) = E(t|t_0) dT(t, z) .$$

One writes $dT(t) = \int [dT(t, z)/dz] dz$, and similarly

$$dP(t|t_0) = \int \{dP(t, z|t_0)/dz\} \cdot dz .$$

Eventually one uses

$$P(t|t_0) = \int_{t_0}^t dP(t|t_0) = 1 - E(t|t_0),$$

from which one gets

$$E(t|t_0) = \exp[T(t_0) - T(t)]$$

Transposing to the easier case of the decay of an unstable particle, $dT(t) = dt/\tau$, where t represents time and τ mean life, one gets the familiar formula $E(t|t_0) = \exp[(t_0 - t)/\tau]$. Here z , characteristic of the kinematics of the decay, has been integrated out and $P(t|t_0) = 1 - \exp[(t_0 - t)/\tau]$.

Notice that $P(t|t_0)$ is a probability, but $T(t)$ is not (the former is bounded by 1, the latter is not). As shown, the existence of $E(t|t_0)$ follows trivially from probability definitions (or, equivalently, from unitarity). However, this term, usually called "Sudakov factor" [ALTA 82, SUDA 56], was studied on more general grounds, where its exponential shape follows from non-trivial summations.

In addition to singularities of $f(z)$, the relation $dT(t) \sim dt/t$ gives rise to a divergence at $t = 0$. Neither of these divergences poses a problem for the functional definition of the probability $dP(t, z|t_0)$. However, there are physics reasons which restrict its validity. The formula defining the running coupling constant $\alpha_s(t)$ is meaningless for $|t| \leq \Lambda^2$, reflecting the fact that perturbative calculations become unjustified. Practically, one should stay at $|t| > t_{\min}$ and describe hadronization by phenomenological fragmentation functions. In principle, one could consider refinements to this Leading Log Approximation, such as $2 \rightarrow 3$ processes, but no general formalism exists yet for the study of exclusive reactions through simulation programs.

In the case of gluon branching ($g \rightarrow g + g$ and $g \rightarrow q + \bar{q}$), one must generalise the formalism. It is sufficient to consider different splitting functions $f_i(z)$ and, correlatively, $dT_i(t, z)$, $i = 1, 2$. As usual the probability of a branching process is based upon the sum over the different channels i contributing proportionally to $dT_i(t)$. There are no fundamental difficulties, however, explicit implementation in a simulation program requires some care as will be apparent in the following section.

2.4.2 Towards next order

It should be clear that the LLA is a first order approximation, justified only asymptotically. The merit of different schemes of parton evolution, whose differences are asymptotically vanishing, may therefore not easily be classified theoretically.

It should be emphasized that in the previous probabilistic approach t was implicitly a monotonic function of time. The usefulness of such a treatment was justified on fairly general grounds, the so-called jet-calculus [KONI 78]. It shows that the graphs involved are essentially dominated by phase-space regions where t is strongly ordered. This means that successive processes correspond, on the average, to decreasing values of $|t|$. In addition this decrease of $|t|$ is rapid. Note that the probabilistic interpretation requires a time ordering, a notion of "before"/"after", but this does not necessitate an ordering in $|t|$. For time-like partons successive processes occur by construction at decreasing values of t , when it is identified with a mass-squared.

In other schemes the virtuality is given by the transverse momentum of the outgoing parton relative to the incoming one or by the opening angle of the outgoing partons. Here ordering is not enforced by the choice of the variable t , but the strong ordering of t should eventually occur on the mean.

The variable z may as well be defined differently, inasmuch as a ratio of energies or longitudinal momenta or their sums are asymptotically equivalent. There exists no a priori good definition in the absence of a "next to leading order" calculation.

The formulation based on angle [MARC 84, WEBB 84] is a typical example of a non-trivial choice of variables. One uses for z the ratio of energies namely $z = E_j/E_i$ (in a process labelled by " $i \rightarrow jk$ ") and for the evolution variable t the expression $\xi = (E_j E_k - \vec{p}_j \cdot \vec{p}_k)/E_j E_k$. For small virtualities of the outgoing partons, $\xi \approx 1 - \cos\theta$, where θ is the opening angle. It has been shown that there are strong destructive interferences in the soft region (z close to 0 or 1). It turns out that such phase-space regions are eliminated more effectively when one imposes strong ordering on the variable $t = \xi$ than by the choice $t = p_i^2$. All exclusive models rely on strong ordering to justify the neglect of interference. As argument of α_s one chooses $2\xi(1 - z)p_i^2$ which is approximately equal to the square of the relative transverse momentum of the outgoing partons. Such a scheme is now used, with different variations, in many simulation programs.

Another complication arises from the requirements of exact four-momentum conservation. The prediction for the occurrence of a branching process is based upon the integration over the range of t and z . Limits are needed in order to control divergences at $z = 0$ or/and $z = 1$. Such divergences vanish when outgoing partons are massive, which indeed is the case when t is the square of the parton mass. The kinematics of a branching process can be only defined after later processes are known. More difficult is the fact that the probability for a process depends critically on the range of z . This involves rather intricate coupled integral equations [KAJA 80] when t is identified with parton mass-squared.

It is not clear whether such complications are relevant. They may, however, give rise to significant effects, in particular in terms of parton multiplicity. Alternatively, one may utilize these effects as an estimate of systematical errors.

2.4.3 Space-like evolution

Partons emerging from the hard scattering are time-like. They evolve from a virtuality t , comparable to the scale Q^2 characteristic of the scattering, down to a cut-off value below which LLA is unjustified. A similar situation occurs before the scattering, but the space-like nature of the partons introduces further complications. Partons evolve from a (negative) virtuality $t = -Q_0^2$ (e.g. $Q_0 \approx 1$ GeV) down to $t \approx -Q^2$, where the hard scattering occurs. At each branching process, a space-like parton gives rise to one space-like and one time-like parton. Eventually two lines of space-like partons connect the original beam particles to the scattering process.

Apart from variations on the choice of the relevant variables, current simulation programs can be classified as follows: (a) forward evolution, where branching processes are generated first long before the scattering and continue up to the scattering, and (b) backward evolution, where the reverse order of generation is adopted. Both schemes are equally justified, they differ only in technical aspects and, eventually, in generation efficiency. The forward evolution algorithm has, however, a clear drawback: the two partons involved in the hard scattering are

evolved independently of each other so that unrealistic kinematical conditions, e.g. $\hat{s} < Q^2$, may occur. Such cases are eventually discarded, but may impair the generation efficiency. A decision must be taken at each branching process as to which parton is space-like. For forward evolution the space-like parton should, in the mean, be the hardest so that enough energy is available for the hard scattering. For backward evolution the last parton reached should be of a flavor which reflects the hadron structure function. The problem evidently lies in the requirement that strong boundary conditions exist for both ends of the chain.

The structure functions of the beam particles may be defined at some Q_0^2 , the scale of the last space-like parton for backward evolution (or the first parton in the forward evolution scheme) and the Q^2 dependence will follow from the generation method. Indeed the Altarelli-Parisi evolution equations used are precisely what describes the non-scaling effects, i.e. the Q^2 dependence. But which structure function should one use? It was shown [GOTT 85] that the phase-space region sampled in the space-like parton evolution is different from that one in DIS experiments so that structure functions estimated from the latter should not be used in the former.

2.5 Fragmentation

In the preceding section the evolution of the outgoing partons from a scale characteristic of the hard scattering to some scale Q_0 was considered. Now one needs a description of the transition from partons to hadrons. Confinement, according to which partons cannot asymptotically exist as free states, is up to now only conjectured to be a property of QCD. Thus, a fortiori, hadron formation is not understood and one must resort to phenomenology. There are in principle two approaches: Cascade models and Colorless Cluster models; more recent refinements which did not influence the data analyses at ISR energies are not mentioned.

2.5.1 Cascade models

The formalism of cascade models was abstracted from multiperipheral models [KRZY 72, FINK 72, JENG 73], and later quantitatively applied to experiments in refs [FIEL 77, FIEL 78a, FEYN 77,78], collectively referred to later as [FF]. The ideas developed in the Lund model [ANDE 83] started from a simplified, exactly calculable case of one space and one time coordinate. This was then heuristically generalized to 4 dimensions.

The following is a short résumé from ref. [KRZY 78] introducing the cascade concept. The term "jet" will designate the collection of hadrons resulting from the fragmentation of a parton, and $J(P)$ stands for the complete set of jets of total four-momentum P . This set is characterized by the multiparticle momentum spectra which in turn define the multiplicity distribution (all hadrons are assumed identical for simplicity).

The basic postulate is that in each jet a particle may be formally distinguished in a way, not an arbitrary one, such that the collection of the other particles constitutes a jet as well. To be explicit, let $J(P_0)$ be the set of jets of momentum P_0 . From each jet, the distinguished particle, of momentum p_d , is removed resulting in a collection of particles of total momentum $P_0 - p_d$. One may group all such collections with a common momentum $p_d = p_1$ into a set $S(P_1 = P_0 - p_1 | P_0)$ which, a priori, depends both on P_1 and P_0 . The postulate states that this set S does not depend on P_0 and is identical to the complete set $J(P_1)$. The collections of particles mentioned above are not arbitrary collections, they are jets, and all sets $J(P_0)$ generate the same set $J(P_1)$, provided $P_0 \geq P_1$. It is then sufficient to make a hypothesis on the invariant spectrum of p_1 , $h(p_1 | P_0)$, to calculate all statistical distributions of the jet. It should be stressed that this postulate imposes strong constraints on the structure of jets (see below). Collections of particles with small relative transverse momentum do not necessarily verify this postulate.

For example, neglecting masses and transverse momenta, and assuming scale invariance, i.e. $h(p|P) = h(z)$ with $z = |\vec{p}|/|\vec{P}|$ in the infinite momentum frame, one obtains the one particle invariant inclusive spectrum $g(z)$:

$$g(z) = h(z) + \int_0^{1-z} h(v) \frac{dv}{v} \cdot g\left(\frac{z}{1-v}\right)$$

One recognizes the contribution $h(z)$ of the distinguished particle. Various results based on this postulate may be found in refs [FIEL 78a] or [KRZY 78], in particular the development of a plateau in rapidity and, consequently, a mean multiplicity proportional to the height of this plateau and to the logarithm of the invariant mass of the jet. The normalization of the invariant spectra $h(z)$ and $g(z)$ correspond to $\int_0^1 h(z) \frac{dz}{z} = 1$ (one distinguished particle), $\int_0^1 g(z) \frac{dz}{z} = \langle N \rangle$ (mean

multiplicity) and $\int_0^1 g(z)dz = 1$ (total momentum). The notations of [FIEL 78a] are slightly different: $h(z)$ is replaced by $f(1 - z) = h(z)/z$ and $g(z)$ by $F(z) = g(z)/z$, because distributions in z rather than invariant spectra are considered. This function f is used in the next two sections.

While the practical implementation will be described in the next section, a few comments on the postulate should be made here. A jet is decomposed into a distinguished particle and a remnant jet, which is decomposed again and the process repeated until all particles are considered. By this method, one can order the particles by the "rank" in the decomposition where they are distinguished (rank 1 in the original jet, rank 2 in the next jet ...). The next section will present a simulation algorithm where particles are generated precisely in this order. This does not correspond to the order of times at which particles are created. The rank order is Lorentz invariant, the time order is not. In a given Lorentz frame, the first particles produced are on the mean the slowest. In the c.m.s. of a parton-parton interaction, the fragmenting partons are fast. The function $h(z)$ which approximates best the experimental data favours high z . It follows that, on the mean, the rank 1 particle is fast and created last. This is the "in-out" cascade [BJOR 76] where particles strongly correlated with the fragmenting parton, i.e. low rank, are produced last. Notice that the rank of a particle is in general not observable. One may however construct functions $h(z)$ where the rank is observable, e.g., if $h(z)$ is equal to zero in the range $z \in [0, 0.5]$. It remains true that the fastest particle of a jet in the parton-parton c.m.s. frame (often the "trigger particle"), is of rank 1 with a high probability.

The next subsections will exhibit strong correlations of quantum numbers between a particle of rank 1 and its parent parton. This is the basis of qualitative inferences made in the other sections. When relevant, quantitative corrections will be estimated from model calculations.

2.5.1.1 Field-Feynman model

The cascade concept is now extended to take into account masses and transverse momenta, parton flavor and the quark composition of hadrons.

Let q, q' represent quarks, $\bar{q}, \bar{q}'$ the corresponding anti-quarks and h denote a hadron. The algorithm of ref. [FF] describes the fragmentation of the quark q : $q \rightarrow q' + h$ where $h = q\bar{q}'$ is the hadron formed. This may be viewed as a vacuum excitation into a $q'\bar{q}'$ pair followed by a $q\bar{q}'$ recombination into a hadron h . This scheme is successively applied to the remaining quark q' , and over again until no further hadron is produced (through a decision to be defined below). Thus, each parent quark gives rise to a cascade of hadrons and a left over quark. The statistical distribution of the hadron produced in such a process is assumed to depend only on its parent quark, not on the prior processes. This is an explicit realisation of the scheme described in 2.5.1; the hadrons produced successively are the ordered distinguished particles.

Here, only the excitation of light quark pairs is considered, with the relative frequencies $s\bar{s}/u\bar{u}/d\bar{d} = \lambda/1/1$, with a free parameter λ , the strangeness suppression factor. The quarks q' and $\bar{q}'$ have opposite transverse components $\vec{k}_t$; the distribution of $\vec{k}_t$ is of Gaussian shape with $\langle k_t^2 \rangle \approx (0.44 \text{ GeV}/c)^2$. The longitudinal distribution is given by $f(1-z) = (1-a) + 3a(1-z)^2$ (sect. 2.5.1) with $a \approx 0.77$. The variable z is expressed here as the ratio of light cone variables $E + p_1$ rather than of longitudinal momentum p_1 which is difficult to handle in the infinite momentum frame. The longitudinal axis is taken along the momentum of the quark which originates the chain. The hadrons defined by their quark content $q\bar{q}'$ are assumed to be produced equally frequently independent of their spin, whereas spin arguments would lead to a ratio 1 to 3 for pseudoscalar and vector mesons.

In case of gluons originating the chain, a similar scheme is adopted, in which gluons correspond to either $u\bar{u}$ or $d\bar{d}$ or $s\bar{s}$ pairs with relative frequencies $1/1/\lambda$, such that one studies the process $q\bar{q} \rightarrow q'\bar{q}' + h$ with $h = q\bar{q}'$. The function f is now taken [FEYN 77] to be $f(1-z) = (1-z)^3$. There exists a variant of this scheme in which the gluon is explicitly split into two quarks which fragment independently. The fraction ρ of the gluon energy taken by a quark is then distributed in an ad-hoc manner either uniformly in $[0, 1]$, or most asymmetrically, i.e. proportional to $0.5 * [\delta(\rho) + \delta(1-\rho)]$. The latter scheme replaces effectively the gluon by either a quark or an antiquark on a random basis.

The rapidity spectrum of hadrons resulting from this scheme shows a plateau extending towards $y \rightarrow -\infty$. Large negative rapidities, i.e. $z \approx 0$ correspond to unphysically large energies. Thus one must reject hadrons with longitudinal momenta lower than some threshold. In ref. [FF] a cut-off at $p_1 = 0$ was applied but this may be considered to be adjustable. Notice that this threshold does define the end of the fragmentation chain by the condition that further transitions would produce hadrons necessarily below this cut. These "rejected hadrons" give rise to non-conservation of energy-momentum as well as of quantum numbers. In addition, each fragmentation chain leaves a quark (or a quark pair in the case of gluons) unassociated with hadrons. One may balance the global interaction in various ways but such refinements are difficult to verify from the data.

The partons initiating the fragmentation chains have been assumed to have evolved down to a scale Q_0 (sect. 2.4) where the splitting function $f(1 - z)$ was assumed to be reasonably adjusted. The Q^2 dependence, which cannot be neglected at energies e.g. in the ISR range, was taken into account in [FF] using the QCD moments method. Alternatively, one may consider an effective splitting function taking a power law $f(1 - z) = (b + 1) * (1 - z)^b$ where b is Q^2 dependent. The following choice [LOHS 82] gives reasonable results at ISR energies, adjusted to match the data on K^- production:

$$b = 1 + 0.64 \tau \text{ for quarks,}$$

$$b = 2.4 \tau \text{ for gluons,}$$

with $\tau = \ln[\ln(Q^2/\Lambda^2)/\ln(Q_0^2/\Lambda^2)]$ and $Q_0 = 2 \text{ GeV}$.

2.5.1.2 Lund model

In the model of ref. [FF], one assumed that partons fragment independently of each other. In such an Independent Fragmentation (IF) model, Lorentz invariance is violated. However, it is postulated that, in the overall c.m.s., the slow hadrons, which essentially cause this problem, may not legitimately be associated with the fragmentation of a specific parton. In contrast, the Lund model [ANDE 83] is more ambitious. It insists on Lorentz invariance which imposes restrictions and in turn provides predictions.

After parton evolution, one may draw color lines between partons in a nearly unique way. These lines connect color triplets (quarks) and color antitriplets (antiquarks or quark pairs), and may join gluons in between. An equivalence is made between such color lines and strings in a two-dimensional model (E and $|\vec{p}|$) where one may show the existence of stable states and study their creation time and dynamical properties. A heuristic step is then made towards strings in four-momentum space where stable states are identified with hadrons. Only the salient points of this model are recalled here. Details are given in [ANDE 83, BENG 85, SJOS 86].

The splitting function is constrained in this model by requiring symmetry when exchanging the ends of the fragmentation chains. One is eventually led to the shape $f(z) \sim 2/z (1 - z)^a e^{-bm_T^2/z}$, where m_T is the transverse mass of the hadron formed, a and b are adjustable parameters (though not exploited yet, the parameter a might be flavor dependent). In the model of ref. [FF], the splitting function $f(z)$ was adjusted to the data, and kept independent of the hadron mass for simplicity. In the Lund model, a mass dependence emerges naturally.

By construction, this model is Lorentz invariant. Apart from being conceptually appealing, this gives rise to definite predictions for 3-jet events in e^+e^- interactions. In such a case, one expects a depletion of the density of hadrons in the $q-\bar{q}$ wedge opposite to the gluon jet, and this is compatible with observations. In the case of deep inelastic hadronic interactions, the original string picture gave rise to difficulties. The simplest case is that of two scattered quarks and two spectator systems. Connecting these four objects by color strings, one finds two configurations of defined relative importances. One has "short" strings and "long" strings, i.e. strings of low and high invariant effective masses. It turns out that the resulting predictions are in disagreement with the data (sect. 4.4). This simple picture of four fragmenting partons has been extensively modified by the introduction of parton evolution prior to the fragmentation process. A consequence is now a variable number of gluons so that the extreme configurations of short and long strings have virtually disappeared. Such extensions of the Lund model are not discussed here.

2.5.2 Colourless clusters

Here it is assumed that the scale Q_0 is low enough so that hadrons may be formed directly. As was seen in sect. 2.5.1.2, a color line may be associated to each parton and the parton lines displayed on a plane (planar diagrams) so that they may be grouped into colorless clusters. In the general case, these clusters are assumed to decay directly into hadron pairs using phase space and spin factors (tensor hadrons are included). Some critical mass is considered below which close clusters are merged, and above which clusters are treated in a manner similar to that of the Lund model. A more detailed description of this type of model may be found in [GOTT 83].

2.5.3 Lorentz frame

The cascade models in sect. 2.5.1 are not Lorentz invariant, thus one must define a frame in which the algorithm is to be applied. In the model of ref. [FF], there is no evident best choice: generally one uses the overall c.m.s. frame. In the case of two collinear jets, a Lorentz boost along their common direction would give rise to more restrictive cuts for one jet and conversely less restrictions on the other jet. However, due to non linearity, these effects would not compensate. Differences would appear for slow hadrons but those cannot be safely predicted. Indeed, the model is based on asymptotic properties and slow particles may not legitimately be associated to either of the two jets. Even if the model does make such an association, results sensitive to this aspect are therefore not to be trusted.

By contrast, the Lund model was claimed to be predictive in its definite choice of Lorentz frames, especially for angular densities from e^+e^- 3-jet events. Straight application of this model, in lowest order quark-quark scattering, to 4-jet events observed in pp interactions is in clear disagreement with the data (sect. 4.4). In the current implementation of this model, parton evolution (sect. 2.4) generates now a larger number of strings and, ultimately, Lorentz transformations from string to c.m.s. frames are definitely less important than in lowest order. Whatever the reasons of the compatibility of the model with e^+e^- data, it is hoped that its current, more complex, string structure will give specific predictions for pp interactions which may be experimentally tested.

2.5.4 Spectators

Hard interaction occur frequently between valence quarks, leaving diquarks as spectators in case of pp collisions. In the case of hard gluon or sea quark scattering, the spectators are 3- or 4-quark systems. The fragmentation of such 2-, 3- 4-quark systems into hadrons is treated in various ways. In the Lund model, 3- and 4-quark systems are degraded immediately into a leading baryon, leaving no or 1 quark. Thus, one is left with 0-, 1-, 2-quark systems at the end of a string. A different scheme consists in using a phenomenological description from hadronic low- p_T data. This inherently includes leading baryon effects. Finally, spectators may be treated in analogy to cascade models [YEUN 87]. Leading baryons are introduced through a special choice of the splitting function to be applied at the first step in the chain. Notice that the fragmentation of scattered diquarks should be treated in a way similar to that of equivalent spectators.

2.5.5 Summary

Empirical fragmentation schemes provide the necessary link between the present level of perturbative QCD predictions and experimental data. In this phenomenological approach, actual data are used with variable emphasis: the more they are used, the less can be predicted genuinely. The faith into QCD evolution varies inversely. In other words, the intermediate Q^2 range is either attributed to parton evolution or to hadronization. This is partly subjective and only data can indicate which direction is to be followed.

FIGURE CAPTIONS

Fig. 2.1 A compilation of measured structure functions $xu_v(x)$ [ARNE 89].

Fig. 2.2 The structure function F_2 as a function of Q^2 for different bins of x [AUBE 87].

Fig. 2.3 Dependence on x and Q^2 of the structure functions, (a) for valence up quarks, (b) for valence down quarks, (c) for gluons and (d) for sea quarks $\bar{q}_s$ ($\bar{q}_s = \bar{u} = \bar{d} = \bar{s}$).

Fig. 2.4 Display of the ratio of the structure functions of down and up quarks.

Fig. 2.5 Gluon exchange graphs which correspond, in the case of forward scattering, to (a) "long strings" and (b) "short strings".

Fig. 2.2

Fig. 2.1

Fig. 2.3

Fig. 2.4

(a)

(b)

Fig. 2.5

3. JET STRUCTURE AND GLOBAL EVENT PROPERTIES

Within the parton model, hard scattering of constituents from two colliding hadrons is expected to produce a characteristic four-jet structure [BERM 71] (fig. 3.1). Two jets transverse to the beam direction result from the fragmentation of the hard scattered partons. The spectator jets originate from the fragments of the proton remnants - and are therefore directed approximately along the beams. In this section the available experimental evidence for this expected four-jet structure is presented and global event properties are discussed.

3.1 Four-jet structure

The most direct evidence [BREA 84a] of a four-jet structure is shown in fig. 3.2, where the normalized momentum flow F per event projected onto the scattering plane is plotted as a function of the polar angle θ . The scattering plane is defined by the direction of the incoming protons and the large p_T trigger particle. The normalization was done by scaling the projected momentum q_i of particle i in a bin $\Delta\theta$ by the corresponding average momentum per particle $\langle q \rangle_{in}$ determined from normal inelastic hadronic interactions in the same θ interval:

$$F(\theta) = \frac{1}{\langle q \rangle_{in}} \frac{1}{N_{ev}} \sum_{j=1}^{N_{ev}} \sum_{i=1}^{n_j} q_i$$

$$\langle q \rangle_{in} = \frac{1}{N'_{ev}} \sum_{j=1}^{N'_{ev}} \frac{1}{n_j} \sum_{i=1}^{n_j} q_i$$

N_{ev} and N'_{ev} are the total number of large p_T and normal inelastic events, respectively, and n_j is the number of charged tracks in the angular interval $\Delta\theta$ in event j . $F(\theta)$ is displayed for normal inelastic events in fig. 3.2(a). It represents, up to a normalization factor, the inclusive θ distribution. A clear two-jet structure of the beam fragments is observed.

In fig. 3.2(b) the momentum flow for an event sample with a trigger particle at $\theta = 50^\circ$ with $p_T > 6$ GeV/c is shown. A pronounced increase of the momentum flow is observed around the position of the trigger particle. Note that the triggering particle is excluded from $F(\theta)$. This increased momentum flow is due to a clustering of particles in a

narrow cone around the direction of the trigger particle. The spectator jets produce a distinct large momentum flow in a narrow angular interval around the direction of the beams at $\theta \approx 0^\circ$ and $\theta \approx 180^\circ$, respectively.

No obvious jet-like behaviour is observed opposite to the trigger particle ("away-side"), merely a general increase of the momentum flow can be recognized. It extends over a wide range of polar angles. Exactly this behaviour is expected, since the broad momentum distribution of partons within the incoming protons causes the c.m.s. polar angle of one scattered parton, and therefore of the corresponding jet, to vary widely from event to event, once the direction of the other scattered parton is approximately fixed by the trigger particle. The small value of $F(\theta)$ around $\theta = 270^\circ$ reflects incomplete cancellation of acceptance effects.

In order to exhibit the presence of the fourth jet more clearly, a similar condition to that on the trigger side is applied. Only events with a further large p_t particle in a limited θ interval on the "away side" are selected. In fig. 3.2(c) such an "away trigger" with $p_t > 2$ GeV/c is demanded in the "back to back" configuration ($230^\circ \leq \theta \leq 260^\circ$). As for the primary trigger particle also this away trigger particle is not included when the flow $F(\theta)$ is evaluated. A jet-like structure of width similar to the one of the trigger jet is observed. The corresponding plot with an away trigger in the "back to antiback" configuration ($280^\circ \leq \theta \leq 310^\circ$) is shown in fig. 3.2(d). Even in this extreme case the "away side" secondaries cluster around the direction of the away trigger.

Such a four-jet structure is observed for triggering $\pi^\pm$ and $K^\pm$ mesons as well as for protons [BREA 84a, BREA 87c].

Given this clear evidence of a four-jet structure it is possible to read off details of parton scattering dynamics. Only global properties are discussed here, details are given in the following sections.

3.2 Particle density distributions

The density distribution of secondary particles in the azimuthal angular region of the trigger- and the away-jet reveals first features of parton dynamics in large p_T processes. In fig. 3.3 the rapidity distribution of fast secondary particles in the away-side region is

displayed for large p_T trigger particles at three different polar angles θ [BREA 87d]. The most characteristic property of these data is the near symmetry around $y = 0$, as already observed by [DELL 76a], although the polar angle of the trigger particle varies between the extreme forward ($\theta \approx 10^\circ$) and the central region ($\theta \approx 50^\circ$).

This observation poses severe constraints on the shapes of scattering amplitudes [FOX 76, KRIP 77]. An "isotropic" scattering amplitude results - specifically for the $\theta = 10^\circ$ data - in a "back to antiback" configuration for the range of transverse momenta of the trigger particle considered. The argument goes as follows: Scattering angles around 90° in the parton-parton rest frame would be preferred since they correspond to the lowest value of Bjorken x at given p_T . Due to the generally unequal relative momenta x of two incoming partons, there is a Lorentz boost along the beam axis resulting in a configuration with nearly the same rapidity of the trigger- and away-side jet fragments. On the other hand, a scattering amplitude peaked at small angles allows to reproduce the observed rather symmetric rapidity distributions. In sect. 6.2 it will be shown that these types of measurements allow to derive the angular dependence of the scattering amplitude and to exclude spin 0 gluon exchange.

In fig. 3.4 the rapidity distributions of secondary particles with $p_t > 1$ GeV/c in an azimuthal angular region $\Delta\phi = \pm 30^\circ$ around the trigger particle (toward side) and at $\Delta\phi = 180^\circ \pm 40^\circ$ (away side) are presented [BREA 84a]. The data are given for three intervals of p_T of the trigger particle. One observes a shift of the average direction of the away particles from the symmetric configuration to a "back to back" configuration in rapidity as p_T increases. Besides a kinematical contribution due to momentum conservation [BREA 84a], this shift occurs as a consequence of the change of the relative importance of contributing parton subprocesses: whereas at moderate trigger momenta the away side particles are expected to be predominantly gluon fragments (sect. 5.2), the relative contribution from quark-quark scattering should increase with the trigger momentum. Fig. 3.4 therefore is suggestive of the presence of more than one parton subprocess. This observation will be discussed further in sect. 5.

Selecting events with an away trigger, one arrives at the distributions in fig. 3.5(a,b) [BREA 84a]. Away triggers with $-0.8 < y < -0.4$ and $x_E = p_t/p_T > 0.3$ are required in fig. 3.5(a), while fig. 3.5(b) comes from away triggers with $0.4 < y < 0.8$. Note that neither the trigger nor the away trigger particle were included in figs 3.5(a,b). Again one observes that the associated particles follow the trigger and away trigger, respectively. The width of the clustering in both cases is $\Delta y \approx 1.0$, as expected for parton fragmentation. Jets observed at the CERN SpPS collider [DILE 85] reveal (fig. 3.6), on a single event basis, similar features as those observed in the data discussed here.

The width of the transverse jets is not only limited in rapidity but also in the other dimension. This is shown for the trigger-jet in figs 3.7(a,b) where the particle density

$$\rho = \frac{1}{N_{ev}} \frac{\Delta^2 N}{\Delta\phi\Delta y}$$

is presented as function of the differences in rapidity (Δy) and in azimuthal angle ($\Delta\phi$) between secondaries and the trigger particle [BREA 84b]. A sharp correlation structure emerges, which appears nearly background free for secondaries with $p_t = 1.6 \div 2$ GeV/c (fig. 3.7(b)). In order to demonstrate particle clustering in the azimuthal angle for both sideways jets, the density of secondaries is plotted in fig. 3.8 as a function of both rapidity y and azimuthal angle ϕ for events with $p_T > 4$ GeV/c [BREA 84a]. The distributions were obtained by normalizing the particle densities (excluding the trigger) of large p_T events to the density measured for normal inelastic events. The density distribution is shown for three intervals of p_t . At low p_t only a minor deviation from uniformity is observed. With increasing p_t of the secondaries a clear particle clustering at $\phi = 0^\circ$ and $\phi = 180^\circ$ is observed. The normalized particle density is larger than 1 as indicated by the dotted lines. The width in ϕ at $\phi = 0^\circ$ and $\phi = 180^\circ$ is compatible with an average transverse momentum relative to the jet axis of a few 100 MeV/c. The width of the away-jet is measured to be larger than that on the trigger side. This observation is compatible with parton fragmentation convoluted with an additional transverse momentum of the incoming partons [DELL 77]. This will be discussed again in sect. 4.2.3.

Figs 3.8(a-c) demonstrate also the clear separation of spectator fragments from those of the sideways jets. Specifically at low p_t (fig. 3.8(a)), where most spectator fragments are produced, background from the tails of the sideways jets is low. Studying spectator fragments in the azimuthal angular region around $\phi \approx 90^\circ, 270^\circ$ yields therefore a data sample which is particularly free of background from sideways jets.

The clean separation of all 4 jets demonstrated in this section allows to study the fragmentation of partons yielding the different jets (sect. 4), to disentangle the contribution of different parton sub-processes (sect. 6.3), and to determine properties of the scattering amplitude (sect. 6.2).

3.3 Average multiplicities

It is intuitively expected that hard collisions of two hadrons lead on average to the production of more particles than normal inelastic reactions. This expectation was confirmed by measurements performed with a streamer chamber recording most of the charged particles emerging from proton-proton collisions [DERA 78]. The chamber was triggered by the presence of a large $p_T \pi^0$ emitted at $\theta \approx 90^\circ$ or 53° . It was found that the average number $\langle n \rangle$ of charged secondaries is increasing with the transverse momentum of the π^0 up to $p_T \approx 2 \text{ GeV}/c$ where it levels off (fig. 3.9).

A more detailed study of multiplicity distributions was performed for pp collisions in which a large p_T charged hadron was produced at $\theta \approx 45^\circ$ [SOSN 83]. The results are shown in figs. 3.10 and 3.11. Here again one notices that $\langle n \rangle$ increases with increasing p_T , reaches a maximum and then slowly decreases. On the other hand the dispersion D of the multiplicity distribution does not depend significantly on p_T . It is, however, smaller than the dispersion from soft collisions with the same value of $\langle n \rangle$.

One can explain qualitatively the dependence of $\langle n \rangle$ on p_T and the low value of D if one assumes that, in hard collisions, there are more independent sources of particle production than in soft inelastic reactions. In fact, it is sufficient to assume that in hard collisions particles originate from two independently fragmenting systems, each very similar to

the one emerging from a soft hadronic collision. This conjecture may be justified by the observation in the previous section that hard collisions show a four-jet structure in contrast to the two-jet structure of soft reactions. According to this finding the total c.m.s. energy $\sqrt{s}$ of a hard collision splits up into the rest energies $\sqrt{s_1}$ and $\sqrt{s_2}$ of two 2-jet systems. If these systems do not move much relative to each other, one has $\sqrt{s} \approx \sqrt{s_1} + \sqrt{s_2}$.

Within this simple model, the multiplicity of hard collisions is then given by

$$\langle n \rangle = \langle n(s_1) \rangle + \langle n(s_2) \rangle$$

with two contributions $\langle n(s_1) \rangle$ and $\langle n(s_2) \rangle$ of logarithmic energy dependence (see e.g. [MATT 83]). By changing the value of p_T one modifies the splitting of $\sqrt{s}$. Hence, the average multiplicity of particles produced in hard collisions depends on p_T . However, a calculation of the two contributions using the parametrization of [MATT 83] leads to $\langle n \rangle = 23$ at $\sqrt{s} = 63$ GeV and $p_T > 4$ GeV/c, whereas the data shown in fig. 3.10 hardly exceed 18. One reason is that the large p_T particle carries nearly 40% of the available energy, e.g. of $\sqrt{s_1}$ (sect. 4.1), and therefore the system produces about $\langle n(0.36 s_1) \rangle + 1$ particles, a number smaller than $\langle n(s_1) \rangle$. In addition, some fraction of the total energy $\sqrt{s}$ is taken by the relative motion of the two systems.

The same simple model which was used to explain the $\langle n \rangle$ - p_T correlation is also able to reproduce the small dispersion of the multiplicity distribution from particles produced in hard collisions. For inclusive particle production in proton-proton collisions the dispersion D as function of the average charged particle multiplicity $\langle n \rangle$ is shown in fig. 3.12 [BREA 84d]. They are related by a linear dependence [WROB 73]. The measurement obtained for proton-proton collisions triggered by a large p_T hadron is included as well. The dispersion is in this case much smaller than the one measured inclusively for the same value of $\langle n \rangle$. This is again a consequence of the multiple source structure of large p_T events. For instance, if two fragmenting systems, each with $\langle n \rangle = 9$ and therefore $D = 4.8$, constituted large p_T events with $\langle n \rangle = 18$, the dispersion for such events would be $D = \sqrt{2} \cdot 4.8$ due to the summation of squared dispersions. This gives $D \approx 6.8$, a value close to that observed for large p_T events.

One concludes that the shape of the multiplicity distribution from high p_T events is sensitive to the number of fragmenting systems, their subenergies, and consequently their relative motion. Unfortunately, no detailed predictions from the various fragmentation schemes (sect. 2.5) exist so far.

In the next section detailed properties of all 4 individual jets are worked out.

FIGURE CAPTIONS

Fig. 3.1 Sketch of the characteristic four jet topology of hard proton-proton interactions.

Fig. 3.2 Projected, normalized momentum flow (3.1) in the scattering plane as a function of the polar angle θ for (a) normal inelastic events, (b) high p_T events triggered by a particle with $p_T > 6$ GeV/c at $\theta \sim 50^\circ$. (c,d) same as (b) with the additional requirement of a charged particle with $p_t > 2$ GeV/c at $\langle\theta\rangle = 295^\circ$ (c) and $\langle\theta\rangle = 245^\circ$ (d). The momentum of the pseudo-trigger is not included. The position of the trigger is indicated by the arrow [BREA 84a].

Fig. 3.3 Rapidity distribution of fast secondary particles ($1 \text{ GeV/c} < p_t < 1.5 \text{ GeV/c}$) in the azimuthal angular region opposite to the trigger particle [BREA 87d]. The polar angle of the trigger particle is (a) $\theta = 45^\circ$ ($\langle p_T \rangle \approx 4.2 \text{ GeV/c}$), (b) $\theta = 20^\circ$ ($\langle p_T \rangle \approx 2.8 \text{ GeV/c}$), (c) $\theta = 10^\circ$ ($\langle p_T \rangle \approx 2.2 \text{ GeV/c}$).

Fig. 3.4 Particle density as a function of the rapidity for 3 different transverse momenta of the trigger $p_T > 2, 4$ and 6 GeV/c [BREA 84a]. The rapidity of the trigger is indicated by an arrow. The density of particles in the trigger jet (excluding the trigger particle) is plotted upwards, that in the away-jet downwards; only particles with $p_t > 1 \text{ GeV/c}$ are included in the plot.

Fig. 3.5 Particle densities as a function of rapidity in events with $p_T > 4 \text{ GeV/c}$ [BREA 84a]. A pseudotrigger T' with large p_t is required in the away region. Only associated particles with $p_t > 1 \text{ GeV/c}$ and $0 < \phi < 25^\circ$, or $155^\circ < \phi < 205^\circ$ are considered.

Fig. 3.6 Distribution of transverse energy in phase space for events with total transverse energy larger than 100 GeV [DILE 85].

FIGURE CAPTIONS (Cont'd)

Fig. 3.7 Particle density as function of Δy and $\Delta\phi$ for $p_T = 4.6$ GeV/c [BREA 84a]. (a) Secondaries with $0.4 < p_t < 0.8$ GeV/c and (b) with $1.6 < p_t < 2$ GeV/c are considered. The trigger particle itself is excluded.

Fig. 3.8 Particle density observed in hard hadronic interactions normalized to normal inelastic events as a function of rapidity y and azimuthal angle ϕ [BREA 84a]. The trigger particle ($p_T > 4$ GeV/c) is not included. The transverse momenta of the secondaries considered are indicated; the dashed line indicates a unit ratio.

Fig. 3.9 Mean charged multiplicity at $\sqrt{s} = 53$ GeV as a function of the transverse momentum of the trigger particle [DERA 78].

Fig. 3.10 Mean charged multiplicities as a function of p_T at $\sqrt{s} = 44$ and 63 GeV [SOSN 83]. The values determined for non-diffractive soft hadronic interactions are indicated by arrows.

Fig. 3.11 Dispersion D of the multiplicity distribution as function of p_T ; the arrow indicates the value of D for non-diffractive hadronic interactions [SOSN 83].

Fig. 3.12 Dispersion D of the multiplicity distribution of charged hadrons produced in soft hadronic interactions as function of the average multiplicity $\langle n \rangle$. The measurement for high p_T events is included [SOSN 83].

Fig. 3.1

Fig. 3.2

(c)

(d)

Fig. 3.2

Fig. 3.3

Fig. 3.4

Fig. 3.5

Fig. 3.6

Fig. 3.7

Fig. 3.8

Fig. 3.9

Fig. 3.10

Fig. 3.11

Fig. 3.12

4. FRAGMENTATION

Knowledge on the interaction of coloured partons is always extracted from the observed jets of colourless hadrons. A good understanding of fragmentation is therefore mandatory to extract dynamical features of parton processes. In addition fragmentation in itself is of course, interesting since it reflects large distance properties of strong forces. Some selected recent data are presented in the following separately for sideways and longitudinal jets. However, it is a priori not clear (see sect. 4.4) whether it is correct to consider individual jets.

4.1 Toward jet

The jet including the large p_T trigger particle is the most untypical one in hard hadron-hadron interactions: The jet axis is roughly collinear with the momentum of the trigger particle which carries most of the jet energy. It will turn out that this unusual feature opens the door to flavour tagging of parton jets and hence to rather detailed tests of QCD predictions. In lepton induced reactions jets of this type are usually not studied due to their small cross sections.

4.1.1 Longitudinal structure

From e^+e^- interactions the Mark I Collaboration [HANS 82] has collected a large number of jets containing a leading fragment ("trigger") at $x_{//}^{\max} = 2 p_{//}^{\max} / \sqrt{s} > 0.5$; $p_{//}$ denotes the momentum component along the jet axis. The density of associated charged particles in those jets is strongly suppressed compared to the inclusive yield as shown in fig. 4.1(a). Experimentally, one finds a strong correlation between normalization and slope of $dN/dx_{//}$ and $\langle x_{\max} \rangle$, where $\langle x_{\max} \rangle$ is the fraction of jet momentum carried by its leading fragment. Hence, $\langle x_{\max} \rangle = \langle z_{TR} \rangle$, can be derived from such distributions.

The quantitative determination of $\langle z_{TR} \rangle$ for toward jets was based upon the following method [FISC 81a, BUCH 82]: The probability $P(z_{TR})$ for a trigger particle to carry a fraction z_{TR} of the parent parton's momentum is approximated by a bell-shaped curve with a maximum near $\langle z_{TR} \rangle$. This shape can be understood as follows: Due to steep fragmentation functions, $P(z_{TR})$ decreases for $z_{TR} > \langle z_{TR} \rangle$. At fixed p_T , $z_{TR} < \langle z_{TR} \rangle$ would correspond to larger momenta of the scattered partons. Since the cross section

for the production of a parton in hadronic collisions decreases strongly with its momentum, $P(z_{TR})$ decreases also for $z_{TR} < \langle z_{TR} \rangle$. Effectively, this empirical shape of $P(z_{TR})$ reflects the dynamics of hard parton scattering. With a fixed parametrization of $P(z_{TR})$ and the fragmentation scheme of [FF] for the remaining fraction $1 - z_{TR}$ of the parton energy, the density of associated particles can be obtained. A comparison with measurements of dN/dz_f ($z_f = p_{//}/|\vec{p}_{trig}|$) then yields estimates of $\langle z_{TR} \rangle$ with a systematic error of about 3%, which are shown in fig. 4.2(a) [BUCH 82]. One finds that $\langle z_{TR} \rangle$ is an increasing function of only x_T with $\langle z_{TR} \rangle > 0.7$ for $x_T \geq 0.05$.

The Mark I data are again displayed in fig. 4.1(b) [GEIS 84] as function of z ($= x_{//}$). Included in fig. 4.1(b) are densities of charged particles from jets which contain a large p_T hadron as function of $z = p_t/p_T \cdot \langle z_{TR} \rangle$. The toward jet densities from pp collisions at the ISR [BREA 84b] are free of background from spectator jets for $p_t > 1.0$ GeV/c. They usually exceed the inclusive yield as also demonstrated by the collider data [ARNI 82b] in fig. 4.3. The hadronic high p_T data in fig. 4.1(b) match the trend of the Mark I results. Assuming the same type of jets, this confirms that triggering hadrons with $p_T = 4 \div 8$ GeV/c carry most of the toward jet momentum, the fraction increasing with p_T .

These measurements of $\langle z_{TR} \rangle$ agree also with predictions in ref. [ELLI 74] and approximations [ALON 80] based on data of ref. [ALBR 78b]. On the other hand, they are in disagreement with analytic approximations [JACO 76] which give $\langle z_{TR} \rangle > 0.9$. Within the statistical accuracy, $\langle z_{TR} \rangle$ is independent of the trigger flavour. In [ANGE 82] one finds also that $\langle z_{TR} \rangle$ depends only on x_T ; however, larger values of $\langle z_{TR} \rangle$ were derived (fig. 4.2(b)). This is probably a systematic effect, since the jet energy was calculated from particles with $p_t > 0.3$ GeV/c only. Thus one tends to underestimate the jet energy.

From the experimental finding that $\langle z_{TR} \rangle < 0.9$ one may conclude that simple CIM processes [BLAN 77], where the trigger particle carries all the energy transferred in the hard collision, do not dominate.

For a more detailed study of the toward jet structure, a momentum component u perpendicular to the trigger momentum (defining $u = 0$) is

defined [BREA 84b] in the scattering plane (the plane which contains the momenta of the beam and the trigger, fig. 4.4). At low z_f rather large densities are found for $u < -1$ GeV/c due to background from spectator jets at $\theta \approx 0^\circ$. Comparing the measured distributions and spectator densities estimated from average inelastic events one infers a background of $\approx 10\%$ for $z_f < 0.4$.

The background subtracted distributions of toward jet fragments in fig. 4.5 are asymmetric around $u = 0$. This indicates that toward jet particles compensate a transverse momentum of the trigger particle relative to the common jet (i.e. parton) axis. The average angle $\Delta\theta$ between trigger momentum and parton axis was extracted [BREA 84b] from these measurements; it is always smaller than a few degrees and decreases with p_T as shown in fig. 4.6.

Given a Rutherford-type of parton scattering with a maximum amplitude at $\theta \approx 0^\circ$, it is plausible that trigger particles at large θ originate from scattered partons at a smaller polar angle $\theta' = \theta - \Delta\theta$. The trigger particle acquires an additional angle $\Delta\theta$ from its transverse momentum relative to the jet axis. Monte-Carlo calculations support these simple ideas ([BREA 84b], fig. 4.6). The values of $\Delta\theta$ obtained are about as large as the angles α between comparable fragments and the jet (thrust) axis in e^+e^- collisions [ALTH 84]. For, e.g., $0.6 \leq x_{//} \leq 0.8$ one finds a value of $\langle\alpha\rangle$ which is about equal to the one of $\langle\Delta\theta\rangle$ determined at $p_T \approx 2.3$ GeV/c ($\langle z_{TR} \rangle \approx 0.7$).

It should be noted that the dependence of $\Delta\theta$ on p_T in these inclusive processes implies a further "scaling-violation" of the respective cross section; i.e. a stronger decrease than p_T^{-4} (sect. 2.3.5).

Knowing $\langle u \rangle$, toward jet densities are now determined for the half-jet away from the spectator jet ($u > \langle u \rangle$); they are virtually free of background from spectator fragments. For charged trigger hadrons the resulting z_f distributions are shown in fig. 4.7 [BREA 84b] to be of about exponential shape and to depend only on x_T . This is equivalent to a scaling of $\langle z_{TR} \rangle$ with x_T (see above). As will be shown in sect. 4.1.3, QCD predictions are consistent with these data.

Investigations of the kind presented here cannot easily be done at small θ , since toward and spectator jets tend to overlap.

4.1.2 Transverse structure

Transverse momentum distributions are displayed in fig. 4.8 [BREA 84b] as function of $q_t^2 (= p_{\text{perp}}^2)$, which is measured relative to the toward jet axis given by $\langle u \rangle$. One finds roughly exponential distributions up to $q_t^2 \approx 1 \text{ (GeV/c)}^2$. The slope parameter decreases with increasing $\langle z_f \rangle$; this is also found for inclusive jet production in e^+e^- collisions [ALTH 84]. In fig. 4.9 it is furthermore shown that $\langle q_t \rangle$ is a strong function of $p_{\parallel}$ only, like in e^+e^- annihilation [ALTH 84].

4.1.3 Quantum number correlations within the toward jet: Flavour tagging

Since ν ($\bar{\nu}$) transform d (u) quarks into u (d) quarks, the x_F distributions in fig. 4.10 [ALLE 83] correspond to u and d quark fragmentation functions into pions "containing" and "not containing" the fragmenting quark. One finds that for $x_F > 0.7$ about 80% of the pions do contain a quark with the flavour of the parent quark. Since high p_T hadrons in hadronic collisions carry more than 70% of the jet energy (sect. 4.1.1), they are therefore expected to tag the flavour of the scattered parent parton nearly on an event-by-event basis.

The difference ΔQ between the charge of the parent parton and of its leading fragment depends on the flavour of the parent parton. For those π^+ (π^-) mesons which are leading fragments of u (d) quarks one has $\Delta Q = -1/3$ ($2/3$). If these leading fragments are trigger particles in deep inelastic hadron-hadron (hh) collisions, and if the charge ΔQ is compensated by associated toward jet particles, one expects characteristic patterns of charge densities in toward jets. Figs 4.11(a,b) show background free distributions dN/dz_f for positive and negative particles ($u > \langle u \rangle$) associated with π^+ and K^+ triggers [BREA 84c] from pp interactions. The distributions are compatible, as expected if both trigger species came from the same type of scattered partons, i.e. from u -quarks in pp collisions. They differ from the ones obtained with π^- triggers (fig. 4.11(c)) for which a larger charge ($\Delta Q = 2/3$) has to be compensated if they are leading fragments of d -quarks. Finally, the largest charge asymmetry is found for K^- triggers (fig. 4.11(d)), [BREA 85a]) which suggests that K^- come from a

neutral parent parton. Superimposed are QCD predictions [YEUN 87, BREA 85a] incorporating the fragmentation scheme of ref. [FF]. While the general trends are well reproduced, the predicted normalization differs from the data by up to 50%. This may not be too surprising in view of the involved experimental procedure and since these distributions depend sensitively on $\langle z_{TR} \rangle$ (see fig. 4.3). On the other hand, the z_f dependence of the charge ratios (not shown) is well reproduced [YEUN 87]. These QCD calculations suggest that about 70% (60%) of all π^+ (K^-) triggers come from scattered u-quarks (gluons) in this kinematical configuration. Charge compensation mechanisms were also studied in ref. [AKES 84] for toward jets; usually fair agreement with predictions from the Lund scheme was found.

Reconstruction of K_s^0 opens the door to a measurement of strangeness production in the toward jets. The average number of K_s^0 per event with $z_f > 0.2$ is shown in fig. 4.12 for pion and K^+ triggers [BREA 84c]. For K^+ triggers the observed rate of K_s^0 is about a factor 2 larger than for pion triggers. This is qualitatively expected from local strangeness compensation in case of K^+ triggers being leading fragments of u-quarks. For K^- triggers with $p_T > 3$ GeV/c an increase of the K_s^0 flux in the toward jet of a factor 2.3 ± 0.3 was determined relative to pion triggers [BREA 85a]. Again, this is evidence that K^- triggers come from a flavour neutral parent parton, such that strangeness is conserved locally in the trigger jet.

First data on baryon and antibaryon production in toward jets accompanying meson and (anti-) proton triggers has been reported in [BREA 87c]. The most prominent feature (fig. 4.13) is a strong proton-antiproton correlation for triggers including (anti-)protons.

More precise measurements of baryon production correlated with different types of trigger particles are obviously needed.

4.1.4 SU(3) mixing

In the framework of the non-relativistic quark model, the observed mesons η and η' are linear combinations of octet (η_8) and singlet (η_1) states: $|\eta\rangle = |\eta_8\rangle \cos\psi - |\eta_1\rangle \sin\psi$ and $|\eta'\rangle = |\eta_8\rangle \sin\psi + |\eta_1\rangle \cos\psi$ with a mixing angle $\psi = (-20.2 \pm 2.0)^\circ$ [ALBR 87]. The eigenstates are given by

$|\eta_1\rangle = 1/\sqrt{3}(|u\bar{u}\rangle + |d\bar{d}\rangle + |s\bar{s}\rangle)$, $|\eta_8\rangle = 1/\sqrt{6}(|u\bar{u}\rangle + |d\bar{d}\rangle - 2|s\bar{s}\rangle)$ and $|\pi^0\rangle = 1/\sqrt{2}(|u\bar{u}\rangle - |d\bar{d}\rangle)$. Fragmentation of (anti-)quarks into π^0 , η , η' should therefore directly reflect SU(3) mixing for large values of z , where masses and contributions from decays of heavier resonances (sect. 4.1.6) can be neglected. Experimentally, e.g. the ratio

$$r_D = D_{q \rightarrow \eta}(z)/D_{q \rightarrow \pi^0}(z) = [\langle q\bar{q} | \eta \rangle / \langle q\bar{q} | \pi^0 \rangle]^2$$

with $q = (\bar{u}, \bar{d})$ is accessible; it is given by the η/π^0 ratio at high p_T ($z > 0.7$): $\eta/\pi^0 \approx r_D$. This near equality can be justified since both η and π^0 should be produced by hard collisions of the same type of non-strange quarks (sect. 5.3) such that the same structure functions and matrix elements are involved and therefore cancel in the η/π^0 ratio. These considerations lead to $\eta/\pi^0 \approx 0.65$. Experiments typically yield $\eta/\pi^0 \approx 0.4 \div 0.5$ [BÜSS 75, DONA 78a, KOUR 79a, AKES 85, ANTI 87], as shown e.g. in fig. 4.14 [BONE 89] for pp collisions. No SU(3) mixing would imply a ratio of about 0.3. In contrast, experiment suggests that $\eta'/\eta \approx 1.7$ in pp collisions [DIAK 80] which is closer to the non-mixing case ($\eta'/\eta = 2$) than to SU(3) mixing ($\eta'/\eta = 0.8$). Potentially this apparent discrepancy may be resolved by a contribution from higher twist processes (sect. 6.3.4).

4.1.5 Strangeness suppression

In sect. 4.1.3 it was found that the charge of π^+ and K^+ triggers are compensated in the same way. It is therefore very likely that both are leading fragments of the same type of scattered quarks, i.e. u-quarks in case of pp scattering. The observed strong strangeness compensation for K^+ triggers suggests that K^+ contains a $\bar{s}$ quark from a $s\bar{s}$ pair created in the u quark fragmentation process. Hence the inclusive ratio K^+/π^+ reflects rather directly the strangeness suppression factor $\lambda = s\bar{s}/d\bar{d}$ at $z \approx \langle z_{TR} \rangle$. For $\sqrt{s} > 20$ GeV and $p_T > 2$ GeV/c, $K^+/\pi^+ > 0.4$ was found (fig. 4.15(a), [BREA 84f, JAFF 89]). With decreasing p_T the ratio drops toward the value 0.1 which is typical for average hh collisions [BAIL 87a, MALH 83]. Small values of K/π are also reported from e^+e^- [BUCH 87] and ℓN interactions [ADER 85]. In fig. 4.15(b) an attempt was made to compare data derived from ℓN and deep inelastic hadron collisions at $\sqrt{s} = 20$ GeV as function of the jet energy E_{jet} [GEIS 87]. For hadron collisions,

$E_{\text{jet}} \approx p_T / \langle z_{\text{TR}} \rangle$, with $\langle z_{\text{TR}}(x_T) \rangle > 0.7$ (fig. 4.2(a)). The K^+/π^+ ratio from ℓN interactions was determined for $z \geq 0.2$ [JONE 85]. A difference of about a factor 2 exists between both data sets for $E_{\text{jet}} \approx 3 \div 4$ GeV. The lepton data correspond to tagged d-quark jets (sect. 4.1). The hadron data come primarily from u-quark jets. The only conclusion to be drawn independently of Monte-Carlo simulations is that λ may increase between $z \geq 0.2$ (ℓN data) and $z = \langle z_{\text{TR}} \rangle > 0.7$ (hadron data).

Model calculations yield typically $\lambda \approx 0.3$ for lepton induced and average hadronic collisions [MALH 83]. The ratio $K^+/\pi^+ = 0.46$ at high transverse momentum corresponds to $\lambda = 0.55 \pm 0.05$ [BREA 84f]. The difference of 0.09 between the latter two values is due to K^* production (sect. 4.1.6).

4.1.6 Leading resonances

It has been known for quite some time that more than 70% of the inclusive stable meson yield comes from resonance decays. At high p_T a different situation is expected due to the strong p_T dependence of cross sections: A resonance contribution to the stable hadron yield at high p_T is only expected for the rare cases in which most of the resonance momentum is transferred to one decay particle. This seems to be supported qualitatively by the data in sect. 4.1.3 showing a different level of charge compensation in the toward jet for π^+ and π^- triggers. If both came mainly from ρ decay one should have seen the same amount of compensating charge on the average.

Just because only a small fraction of meson triggers may come from resonances, a measurement of relative inclusive resonance yields should be rather directly related to the ratio $V/P = (\text{number of direct vector mesons})/(\text{number of direct pseudoscalar mesons})$, which is a crucial ingredient of most fragmentation models. The measured resonance yields - and therefore the ratio V/P - correspond to large values of z , since one has $z(\text{resonance}) \geq z_{\text{TR}}$ and $\langle z_{\text{TR}} \rangle \geq 0.7$ (sect. 4.1.1). Evidence for resonance production at high p_T in pp collisions at the ISR was first reported in [ALBR 78b]. The most recent results are presented in the following:

4.1.6.1 Vector and tensor mesons

Significant signals of ρ^0 and $(\bar{K})^{*0}$ production were obtained [BREA 89a, MANK 88] by combining $\pi^\pm$ and $K^\pm$ triggers (at $\theta = 50^\circ$, $p_T \geq 4$ GeV/c) with charged secondary particles taken to be pions.

A new event mixing/subtraction scheme [DRIJ 84] together with a simulation of decay kinematics and of acceptance yield the relative production rates for ρ^0 as function of p_T in fig. 4.16(a). The ratio of the values found for ρ^0/π^- and ρ^0/π^+ is a reflection of the inclusive ratio π^+/π^- (sect. 5.1). The ratio $\rho^0/\frac{1}{2}(\pi^+ + \pi^-) = \rho^0/\pi^0$ is close to unity. Due to the resonance contribution to the π^0 yield one expects $V/P > \rho^0/\pi^0$. Based on the fragmentation scheme of [FF] the parameters V/P were derived from the measured ratios for strange and non-strange particles. In the case of non-strange vector mesons a determination of V/P depends on a further assumption concerning relative ρ^0 and ω production. Previous high p_T experiments measured $\omega/\rho^0 \approx 1$ [DIAK 80] and $\omega/\rho^0 \approx 0.44$ [DONA 80, POVL 83]. Recent data on current jet fragmentation from deep inelastic μ scattering is compatible with $\omega/\rho^0 \approx 1$ [ARNE 86a]. With $\omega/\rho^0 = 1$, $(V/P)_{u,d} = 1.66 \pm 0.18 \pm 0.63$ is obtained [BREA 89a, MANK 88] from the data in fig. 4.16. As can be concluded from a compilation in fig. 4.17 [MANK 88], most of the lepton induced reactions tend to yield $(V/P)_{u,d} \approx 1$. From $(\bar{K})^{*0}$ production at large p_T , $(V/P) = 0.93 \pm 0.17 \pm 0.25$ was derived [BREA 89a, MANK 88] whereas lepton experiments tend to favour larger ratios [DERR 85].

Inclusive production of ϕ mesons in pBe collisions at 400 GeV/c is reported in ref. [AKER 77]. The measurements are shown in fig. 4.16(b); no theoretical interpretation has been attempted so far.

In ref. [BREA 89a] also tensor meson production was observed; for $\theta \approx 10^\circ$ and $\langle p_T \rangle = 2.8$ GeV/c one finds $f_2/\rho^0 = 0.30 \pm 0.08$.

4.1.6.2 Production of Δ^{++}

For proton triggers at $\theta \approx 10^\circ$ the experimental procedure mentioned above yields significant Δ^{++} signals and a ratio $\Delta^{++}/p(\text{direct}) = 0.29 \pm 0.04$ at $p_T \geq 2$ GeV/c [BREA 87c]. This ratio is not likely to

reflect decuplet suppression in quark or gluon jets, since most of the trigger protons are probably leading fragments ($z \geq 0.7$) of hard scattered diquarks (sect. 5.5). The rather substantial values of Δ^{++}/p hint, in this framework, for a non-negligible rôle of hard scattered uu -systems in addition to the probably more frequent ud scattering processes (sect. 6.2). Evidence for Δ^{++} production at high p_T was first found in [ALBR 78b].

4.2 Away-jet

Generally, neither the away jet axis nor its parton flavour are constrained by the requirement of a leading fragment. Hence inclusive jet fragmentation can be studied in principle.

Measurements are usually done relative to the experimentally observed scattering plane given by the incoming beam and the momentum of the trigger. This scattering plane is not necessarily identical to the parton-parton scattering plane since (a) the trigger is produced at a finite angle relative to the momentum of the scattered parton (sect. 4.1.1) and (b) partons carry an effective transverse momentum k_t (sect. 4.2.3). For away jet fragments, the momentum component p_{out} perpendicular to the experimental scattering plane is larger than expected from the transverse momentum q_t (relative to the parton direction) generated in the fragmentation process. The measured values of p_{out} may also be affected by the presence of gluon bremsstrahlung off away partons. Wide distributions in p_{out} are equivalent to large apparent widths in azimuth ϕ (sect. 3.2).

The polar angle distribution of away-jet particles is essentially given by the direction and flavour of the trigger particle, due to the QCD matrix elements for the various subprocesses, as well as by parton structure functions and, potentially, by effects of gluon bremsstrahlung.

Hence, the relative motion of the hadron-hadron and parton-parton rest frames causes rather wide distributions in azimuth around $\phi \approx 180^\circ$ and in y ($|y| < 2.5$) as actually observed (sect. 3.2).

Consequently, density distributions of away-jet particles are biased if kinematic cuts restricting the range in ϕ and/or $|y|$ are applied to suppress contributions from spectator fragmentation. Such cuts may also

influence the relative weight of the various partons contributing to away jet fragmentation. Absolute away-jet properties are therefore not easily extractable in the ISR energy range without major Monte-Carlo efforts.

On the other hand, the dependence of average away jet properties, e.g. of flavour composition, on transverse momentum, polar angle, and flavour of the trigger can be well measured virtually free of model dependent corrections. One can therefore establish rather directly differences between quark and gluon jets, if their contribution to away jets can be controlled in a well defined way. This will be shown below.

4.2.1 Longitudinal structure; gluon fragmentation

Due to the lack of precise knowledge of the away jet momentum, the fragmentation variable z is usually approximated by $x_E = p_t/p_T \approx z/\langle z_{TR} \rangle$. These relations become exact for $z_{TR} = 1$, $k_t = 0$, and if experimental scattering plane and parton-parton scattering plane coincide. Like toward jets at low z_f , also away jets at low x_E are affected by background from spectator fragmentation. Usually average events are used to approximate spectator jet fragmentation taking into account a potential finite angle between the beam direction and the spectator jet axis (sect. 5.6). Background from spectator jets at a level of $\geq 15\%$ for $x_E < 0.3$ is typically estimated [HERR 85].

From the beginning, x_E distributions were compared to equivalent distributions obtained from e^+e^- experiments and compatibility was found [DARR 80]. Whereas these findings were comforting - one was obviously talking about the same type of jets, i.e. partons - this is a fortunate coincidence: The flavour composition of e^+e^- jets is determined by the parton charges whereas it is given by an interplay of hadron structure and colour factors in the case of deep inelastic hh scattering. In addition, corrections for kinematic cuts applied to hadroproduced jets are not straightforward. Therefore, no attempt is made here to compare the two sets of data.

As can be seen in fig. 4.18 [BREA 87d], dN/dx_E is of about exponential shape above $x_E \approx 0.3$ and seems to scale, i.e. the distributions are independent of p_T above $p_T > 3$ GeV/c. Earlier detailed work [ALBR 78b] is

also providing evidence for scaling within the statistical precision and systematic uncertainties. Fig. 4.19 shows a measured x_E distribution together with a set of model predictions [LOHS 82]. One concludes that, to come close to the data, one needs (a) a gluon contribution to the away jets, (b) scalebreaking effects such as Q^2 dependence of all quantities entering the calculation, and (c) $k_t > 0$; $\langle k_T \rangle = 848 \text{ MeV/c}$ [FEYN 78] was used for fig. 4.19. These theoretical considerations imply non-scaling x_E distributions at some level. To verify this more precise experimental and theoretical analyses are needed.

In fig. 4.19 there remains, however, a normalization problem, which cannot be cured easily even by varying all model parameters to their reasonable limits. From a comparison of the calculations with the data, it looks like the effective away-jet energy is smaller than predicted. This would happen in case of final state gluon bremsstrahlung off the away parton. Actually, very recent quantitative calculations taking into account gluon bremsstrahlung [YEUN 87] tend to support this idea as will be discussed in sect. 6.3.3.

The Monte-Carlo predictions in fig. 4.19 incorporate a contribution of both gluons and quarks to the away-jets. The mixture should depend on the parton-parton c.m.s. energy $\sqrt{\hat{s}} = \sqrt{x_1 x_2 s}$ and on the flavour of the trigger particle (i.e. of its parent parton). The fractional parton momenta x_i are basically given by the kinematic configuration of the trigger, i.e. by θ and p_T , and can be calculated with negligible model dependence. Experimentally one can therefore try to enhance the contribution of away quark and gluon jets by varying θ at fixed $\sqrt{\hat{s}}$ and by varying also $\sqrt{\hat{s}}$. This is sketched in fig. 4.20 for π^+ triggers [BREA 87a]. In fig. 4.21(a,b) [BREA 87a] one finds that dN/dx_E , which is always measured in the same region of phase space, i.e. in the same section of the detector, does not strongly depend on θ for $\sqrt{\hat{s}} < 15 \text{ GeV}$. For $\langle \sqrt{\hat{s}} \rangle = 18 \text{ GeV}$ there is, however, a significant difference between dN/dx_E measured for $\theta \approx 45^\circ$ and $\theta < 20^\circ$ (fig. 4.21(c)). Since decreasing angles θ correspond to increasing values of $\langle x \rangle$ of the parton yielding the trigger, the away parton becomes softer ($\sqrt{\hat{s}}$ is fixed) such that the contribution of gluons to away jets is expected to rise (see also sect. 4.2.2). The measurements in fig. 4.21(c) support therefore the expectation that gluon fragmentation is softer than quark fragmentation [GROS 73]; this conclusion is of minimal model

dependence, mainly based on kinematical arguments. Weaker and more indirect indications of differences between quark and gluon jets were obtained in [BAGN 84b, ARNI 86, OLSE 88, BRAU 89]. The experimental data presented above imply, of course, some scale breaking of dN/dx_E due to the gluon contribution to the away-jets and its dependence upon $\sqrt{s}$, i.e. upon p_T .

4.2.2 Particle composition

In addition to charged fragments one expects also π^0 production from away jet fragmentation. This component was measured with a liquid argon calorimeter simultaneously to the charged particles [DORO 84]. Fig. 4.22 shows dN/dx_E for π^0 and charged particles. No deviation from isospin symmetry is observed.

Typical charge structures of away jets are reflected in the ratio R of densities of positive and negative secondaries given as function of x_E in fig. 4.23 [AKES 84]. Generally, one finds a rise from $R \approx 1$ at $x_E \approx 0$ to $R \geq 1.5$ at $x_E \geq 0.5$. This signals the presence of scattered valence quarks which are dominated by u quarks. Again R should depend in principle on the kinematical configuration of both sideways jets due to varying relative contributions from different subprocesses. This was investigated using π^+ triggers [BREA 87a]. The data in fig. 4.24 confirm that the away jets carry less net charge, i.e. contain more gluons (sect. 4.2.1), for $\langle\theta\rangle \approx 18^\circ$ than for $\theta \approx 45^\circ$. Hence, this supports the idea that one can change the flavour composition of away partons in a controlled way by varying direction and momentum of a flavour identified trigger particle (sect. 4.2.1). The dependence of R on the rapidity of the away jet fragments will be discussed in sect. 5.2.2.

A component of strange particles (K^-) was measured with high statistics for jets recoiling against π^+ and K^+ triggers with $p_T = 3 \div 4$ GeV/c as shown in fig. 4.25 [FISK 78]. A detailed interpretation of these data would require a comparison with measurements of dN/dx_E for other particle types and/or a Monte-Carlo comparison. Limited measurements are available for π^+ triggers [FISK 78]; for $x_E > 0.75$ one finds $K^+/\pi^+ = 0.45 \pm 0.10$, $K^-/\pi^- = 0.16 \pm 0.04$ and $\bar{p}/p = 0.1 \pm 0.03$. These ratios are also found for inclusive production at the corresponding values of transverse momenta [BREA 84f], i.e. $p_T \geq 4 \div 5$ GeV/c.

Significant production at $p_t \geq 2$ GeV/c of ρ , ϕ and f recoiling against a single particle with $p_T > 0.9$ GeV/c was recently found in $\pi^- N$ interactions at 300 GeV/c [BENA 87a].

It would be desirable to have better data on away strange meson ($K^{\pm 0}$) and baryon (p , $\bar{p}$, Λ , $\bar{\Lambda}$) production for detailed tests of fragmentation models.

4.2.3 Transverse structure

The transverse profile of away jets is usually analyzed in terms of the transverse momentum component p_{out} . As mentioned above, $\langle p_{out} \rangle$ receives contributions from fragmentation (q_t), intrinsic transverse momentum (k_t), from the finite angle $\Delta\theta$ between trigger and trigger-jet axis and from gluon bremsstrahlung effects. The latter two features were not considered in the early analyses.

First measurements [DELL 77, KOUR 80] showed that $\langle p_{out} \rangle$ increases linearly with x_E and that, at given x_E , $\langle p_{out} \rangle$ is independent of p_T above 2 GeV/c. The x_E dependence suggested a rather large value of $\langle k_t \rangle$ (> 1 GeV/c).

Figs 4.26 summarize more recent detailed measurements [ANGE 80] of $\langle |p_{out}|^2 \rangle$ as function of p_T and x_E . With $\langle |p_{out}|^2 \rangle = 2\langle |k_{t0}|^2 \rangle x_E^2 + \langle |q_{t0}|^2 \rangle (1 + x_E^2)$ [FEYN 78], where k_{t0} and q_{t0} are measured out of the scattering plane, $\langle q_t \rangle \approx .56$ GeV/c and $\langle k_t \rangle > 1$ GeV/c was derived. Such a large value of $\langle k_t \rangle$, if confirmed by more detailed and complete analyses including the effect of $\Delta\theta$, seems to be too large to be attributed to primordial k_t . Qualitatively, however, one expects an increase of the "effective" $\langle k_t \rangle$ due to gluon bremsstrahlung. A different method to determine $\langle k_t \rangle$ will be presented in sect. 5.6. The value of $\langle q_t \rangle \approx 0.56$ GeV/c determined in this analysis is larger than what is usually found (sect. 4.3.2).

From the e^+e^- data in fig. 4.27 [ALTH 84] a potential bias can be guessed which may be introduced by determining $\langle q_t \rangle$ from a measurement of transverse momenta relative to the trigger momentum instead of the jet-axis (thrust axis, shaded area). In fig. 4.27 one has $x_{//}(\text{trig}) = 2p_{//}(\text{trig})/\sqrt{s} > 0.25$ for the "trigger" in the "trigger-jet"; this corresponds to an average

angle $\langle \alpha \rangle \approx 18^\circ$ between trigger and jet direction [ALTH 84]. For both jets one finds differences between $\langle q_t \rangle$ measured relative to the trigger axis and the values (shaded) which were determined relative to the jet axis. Larger discrepancies appear at larger $x_{//}$.

One anticipates therefore from fig. 4.27, that $\langle q_t \rangle$ tends to be overestimated in the analysis of the data in fig. 4.26. The bias should be weaker than in fig. 4.27, since one has $\Delta\theta < 2^\circ$ for $p_T > 4 \text{ GeV}/c$ (sect. 4.1.1) which is smaller than $\langle \alpha \rangle \approx 18^\circ$. This effect should die out with increasing p_T , i.e. decreasing $\Delta\theta$ (sect. 4.1.1).

4.3 Spectator jets

The two high transverse momentum jets, the most characteristic feature of deep inelastic hadronic collisions, are nearly always produced in the central rapidity region. Most of the experimental effort went therefore into the study of this region of phase space. Simple reasoning suggests, however, that quite substantial information can be obtained from measurements at large longitudinal momenta: the incoming hadrons virtually dissociate into the partons to be scattered at large angle and the non-active "spectator" parton systems fragmenting into jets close to the beam direction. Hence, in the case of e.g. q-g scattering, one spectator jet should consist of 3 valence quarks, the other one of two valence quarks. For $\theta \ll 90^\circ$ the probability is high that trigger jet and spectator jet in the same rapidity hemisphere originate from the same incident hadron. Hence their flavours must add up to the flavour of the incoming hadron, or with other words: a study of the quantum numbers of the spectators reveals the nature of the scattered parton.

Furthermore, fragmentation of parton systems consisting of 3, 2 or 1 valence quark can be studied in principle. In pp interaction one expects e.g. ud, uud and u spectator partons in case of hard scattering of u-quarks, gluons and ud-diquarks, respectively.

Finally, the spectator partons are supposed to compensate the transverse momentum k_t of the scattered parton. A measurement of a finite angle of spectator jets relative to the beam direction should therefore contribute to a deeper understanding of the transverse Fermi motion of partons (sect. 5.6).

4.3.1 Longitudinal structure

The first data on the x_F dependence of positive particle yields in the spectator jets were presented in terms of correlation functions to minimize acceptance problems. For trigger particles at $\theta = 90^\circ$ the densities of positive particles, normalized to fully inclusive yields, were determined, and the deviation from unity is given in fig. 4.28 for various values of p_T [ALBR 78a]. One finds a stronger suppression for larger p_T , especially at large x_F , potentially due to energy-momentum conservation. This effect is most pronounced for positive triggers which may be a consequence of charge conservation. These data indicate that diffractive processes which yield positive particles at large x_F become suppressed in high p_T processes (sect. 6.2).

Absolute densities of identified $\pi^\pm$ and p in events triggered by positive and negative hadrons at $\theta = 30^\circ$ with $p_T = 2.5 \div 4$ GeV/c are shown in fig. 4.29 as function of $1 - z$ [HANN 81]. Here, z is the fragmentation variable for the spectator quark system; it was estimated by normalizing the particles' longitudinal momenta p_L to an average spectator momentum derived from simplified kinematics. Whereas a strong dependence on z , nearly independent of the trigger charge, is found for mesons from spectator jets, flatter spectra with significant trigger flavour dependence are observed for protons at large p_L . These distributions cannot easily be turned into (multi)quark fragmentation functions, since the rest system of both spectator jets moves relative to the pp cms with a momentum given by the momentum imbalance of the two sideways jets. From the data of fig. 4.29 it was concluded [BEAV 81] for the case of valence quark scattering that, (a) the diquark spectators fragment as an entity (hence the similarity of meson densities for positive and negative triggers), (b) the diquark fragmentation functions into pions are compatible with deep inelastic lepton-nucleon data, (c) uu and ud diquarks fragment into protons in a similar way. For gluon scattering a triquark fragmentation into protons (pions) of the shape

$$D_{uud \rightarrow p} = \frac{0.19}{z} \quad (D_{uud \rightarrow \pi^+} = 0.15 \frac{(1 - z)}{z})$$

is needed to fit the data [BEAV 81].

A complete analysis of fragmentation of the various (multi)quark spectator systems produced in high p_T processes is still missing.

4.3.2 Transverse structure

A measurement of $d\sigma/dp_t^2$ for particles with $x_F > 0.5$ in high p_T events and normal events is given in ref. [BREA 84h]. Nearly identical exponential distributions were obtained; the measured values of $\langle p_t \rangle$ differ by about 7%. In fig. 4.30 $d\sigma/dp_t^2$ is shown for high p_T events as function of p_t^2 for particles with $2 < p_L < 10$ GeV/c produced within $\pm 40^\circ$ about the normal to the scattering plane [BREA 84h]. The latter cut was chosen for high p_T events in order to minimize the influence of sideways jets. The curves in fig. 4.30 [YEUN 87] incorporate a convolution of an exponential in q_t ($\langle q_t \rangle \approx 0.4$ GeV/c) and a Gaussian in k_t with $\langle k_t \rangle = 0.6$ GeV/c. Good agreement for $p_t^2 < 2(\text{GeV/c})^2$ is seen. The need for a non-vanishing value of $\langle k_T \rangle$ reflects a recoil of spectator jets against the scattered partons. This creates a finite angle between the spectator axis and the beam direction and widens the transverse momentum distribution measured relative to the beam axis.

For a more detailed analysis, $\langle p_t^2 \rangle$ is given as function of p_L in fig. 4.31(a) for normal events and events triggered by a high p_T particle at 20° , and in fig. 4.31(b) for events from 45° triggers [BREA 84h]. One clearly sees (a) the seagull effect, (b) larger values of $\langle p_t^2 \rangle$ for high p_T events, and (c) asymmetries of $\langle p_t^2 \rangle$ around $p_L = 0$ which decrease with increasing θ . A Monte Carlo simulation shows that the data for $|p_L| < 2$ GeV/c reflect directly properties of the fragmentation process (i.e. of dN/dq_t^2) whereas the remaining data can be used to establish bounds on dN/dk_t^2 . With $\langle q_t \rangle = 0.4$ GeV/c and $\langle k_t \rangle = 0.6$ GeV/c, the data are reproduced reasonably well over most of the kinematic range (fig. 4.31(b)). As in sects 4.2.1 and 4.2.3 one finds therefore evidence for Fermi motion of partons, here even in the direction normal to the scattering plane. Transverse momentum components in the scattering plane are more directly affected by $\langle k_t \rangle$; this will be discussed in sect. 5.6.

4.3.3 Two-particle correlations

Previously, it was argued that spectator jets and the longitudinal jets in normal events may differ in their valence quark and, maybe, gluon content. It is therefore interesting to determine two-particle correlations within spectator jets. The measured charge density balance and two-particle correlation functions [DRIJ 80] are displayed in fig. 4.32;

the fixed rapidity of one of the two particles is indicated. The results are identical to those obtained from non-diffractive events (solid lines, [DRIJ 79a]). This suggests a kind of universality of jets from different light quark systems. Similar conclusions hold for light quark-jets produced in e^+e^- annihilations [AIHA 84].

4.4 Fragmentation scheme

So far, jets were analyzed independently of each other. Already fig. 4.1. shows, however, that the 2 jets produced in e^+e^- annihilation are not completely independent.

According to QCD, hard scattered partons carry colour such that they must be correlated with other coloured partons before breaking up into hadrons. From detailed investigations of 3-jet events in e^+e^- annihilation one has actually disfavoured independent fragmentation. The data are generally found to be compatible with QCD shower calculations and with string models [SAXO 85]. In the latter picture strings are stretched between coloured quarks with gluons causing kinks. In 3-jet e^+e^- events the finite angles between all jets introduce a relative motion between the string rest system, in which fragmentation is assumed to occur, and the overall rest system. This results (fig. 4.33, [AIHA 86]) in a depletion of energy flow between the two quark jets and in an increased energy flow between the gluon and each quark jet (always understood relative to 2-jet events where an additional single photon simulates 3-jet kinematics). Due to the small gluon energy, only moderate Lorentz transformations are involved, hence the observed effects are subtle.

In deep inelastic hadron collisions at the ISR with 2 large angle partons all four jet momenta and also the angles between them are usually rather large. Potentially, there may be therefore large Lorentz transformations between colour neutral systems and the overall rest system. This promising situation is probably obscured in reality, at least by the many possible string configurations in first order QCD (fig. 4.34(a), [BENG 84]) whose relative weights are functions of $\sqrt{s}$, i.e. p_T and θ , and also of the trigger flavour.

In fig. 4.34(b), the Lorentz factors $\beta_{//}$ and β_T (parallel and perpendicular to the incident beam) are given [BREA 84i] for $q'q$ "forward"

(curve c) and backward (curve d) scattering corresponding to the first and second string configuration in fig. 4.34(a), respectively. Large values of $\beta_{//}$, i.e. $|\beta_{//}| > 0.5$, are found suggesting large effects on particle densities. For the azimuthal hemisphere of the trigger, measured densities of particles with $p_t > 1$ GeV/c in a wedge around the scattering plane are displayed in fig. 4.35 [BREA 84i]. Model predictions including no or only small Lorentz transformations (scheme a: independent fragmentation, scheme b: toward-away jet and spectator-spectator fragmentation) reproduce the data rather well. In contrast, predictions from string configurations corresponding to lowest order $q'q$ scattering (schemes c and d) are in obvious disagreement with the data. They differ from the data and from independent fragmentation by factors up to 5. In fig. 4.36 [BREA 84i] the particle density close to the trigger rapidity does not depend strongly on the rapidity of the away jet in contrast to what is expected from the string configurations for $q'q$ forward and backward scattering. Finally, u distributions (sect. 4.1.1) are given in fig. 4.37 [BREA 84i] for different trigger flavours. The particle density near the trigger ($u = 0$), is nearly the same for all trigger flavours. From the string model one expects differences in the case of K^- triggers which tag gluon jets (sect. 5.2) and therefore enhance different string configuration (fig. 4.34(a)).

To summarize, the data presented in this section do not seem to correspond to predictions from a string picture applied to first order QCD $q'q$ scattering. One may conclude that either (a) qg string configurations, which were not simulated, are more important than expected in current formulations of first order QCD, or (b) a strong contribution from higher order QCD diagrams gives rise to "shorter" strings via additional bremsstrahlung gluons (sect. 6.3.3), or (c) the type of strings considered does not exist.

It is emphasized again that the potential existence of large Lorentz transformations makes more detailed analyses of the type outlined above very attractive. As concluded in sect. 3.3 also the shape of the multiplicity distribution may be sensitive to fragmentation schemes. It should be pointed out that neither higher order nor QCD shower analyses of all data available were performed so far.

FIGURE CAPTIONS

Fig. 4.1 (a) Corrected particle-density distributions $(1/\sigma)d\sigma/dx_{//}$ vs $x_{//}$ for various $x_{\max}$ cuts for $7.0 < \sqrt{s} < 7.8$ GeV. $x_{\max}$ refers to the highest- x particle and is not included. The jet direction is oriented so that $x_{\max}$ is at positive $x_{//}$. The distributions are normalized to the cross sections for jets with $x_{\max}$ within the specified range [HANS 82].

(b) Inclusive z -distribution of charged jet fragments obtained at $\sqrt{s} = 6.8$ GeV [HANS 82]. Also given are corresponding distributions for charged particles from trigger jets using single particle triggers with $p_T > 4$ and 6 GeV/c at $\sqrt{s} = 62$ GeV [GEIS 84].

Fig. 4.2 (a) Values of $\langle z_{TR} \rangle$ obtained from the data as function of x_T , compared to a prediction [BUCH 82].

(b) Measurements of $\langle z_{TR} \rangle$ ($= \langle z_{trig} \rangle$) versus x_T ($= x_t$ (trig)) for inclusive π^0 production at $\sqrt{s}$ of 31, 44.8, and 62.4 GeV [ANGE 82].

Fig. 4.3 p_t spectrum of towards secondaries within $\Delta y = \pm 0.5$ and $\Delta\phi = \pm 30^\circ$ of the trigger particle, compared to the minimum-bias inclusive p_T spectrum [ARNI 82b].

Fig. 4.4 Definition of momentum components relative to trigger direction and event plane.

Fig. 4.5 Secondary particle density as a function of u (fig. 4.4), for different z_f intervals [BREA 84b]: (a) for high p_T events with $\langle p_T \rangle = 4.6$ GeV/c and $\sqrt{s} = 62$ GeV; (b) for minimum bias events (background); (c) for background subtracted high p_T events.

Fig. 4.6 Angle difference between trigger direction and jet axis (extracted from u -asymmetry) as function of p_T ($= p_T^{TR}$), compared to model calculation [BREA 84b].

FIGURE CAPTIONS (Cont'd)

- Fig. 4.7 z_f distributions, after background subtraction, for different $\langle p_T \rangle$ ($= \langle p_T^{TR} \rangle$) values at $\sqrt{s} = 62$ GeV, and for the "half cylinder" $u > \langle u \rangle$ [BREA 84b].
- Fig. 4.8 Distributions of q_t^2 ($= p_{\text{perp}}^2$) for different x_f intervals, q_t is the transverse momentum of secondaries in the trigger jet relative to the jet axis; $\langle p_T \rangle = 4.6$ GeV/c [BREA 84b].
- Fig. 4.9 Average transverse momentum $\langle q_t \rangle$ ($= \langle p_{\text{perp}} \rangle$) as function of $p_{//}^{\text{sec}}$, the momentum of secondaries parallel to the trigger particle, for various values of $\langle p_T \rangle$ ($= \langle p_T^{TR} \rangle$) [BREA 84b].
- Fig. 4.10 Corrected normalized x_F distributions $D^{\pi^\pm}(x_F) = (1/N_{\text{ev}}) dN^{\pi^\pm}/dx$ of π^+ (full circles) and π^- (open circles) in the forward ($x_F > 0$) c.m.s. hemispheres from νp and $\bar{\nu} p$ events. The distributions are arranged such that simple isospin relations of the QPM can be tested. The error bars include the statistical errors and the uncertainty in the correction functions [ALLE 83].
- Fig. 4.11 Particle density in the trigger jet versus z_f for particles with charges equal (circles) or opposite (dots) to the charge of the trigger ((a) π^+ ; (b) K^+ ; (c) π^- ; (d) K^-); $p_T = 4 \div 6$ GeV/c [BREA 84c, 85a]. Monte Carlo predictions are given by solid lines [BREA 89b, YEUN 87].
- Fig. 4.12 Average number of reconstructed K_S^0 with $z_f > 0.2$ per event for different triggers [BREA 84c].
- Fig. 4.13 Fraction of (a) protons and (b) antiprotons among all positive (a) and negative secondary (b) particles in the trigger jet for various trigger particles at $\theta \approx 45^\circ$ and $p_T \geq 4$ GeV/c. The kinematic cuts for the secondaries are indicated [BREA 87c].
- Fig. 4.14 Ratio of inclusive cross sections for η and π^0 production as function of p_T [BONE 89].

FIGURE CAPTIONS (Cont'd)

Fig. 4.15 (a) K^+/π^+ production ratios versus x_T at $\sqrt{s} = 62$ and 45 GeV for $\theta \approx 50^\circ$ and at $\sqrt{s} = 27$ and 23 GeV for $\theta \approx 90^\circ$. The dashed line comes from a QCD calculation with $\lambda = 0.5$ [BREA 84f].

(b) The ratios K^0/π^- (full dots) and K^+/π^+ as functions of jet energy E_{jet} [GEIS 87].

Fig. 4.16 (a) Cross section ratios ρ^0/π^+ and ρ^0/π^- at $\theta = 45^\circ$ as a function of p_T . The full lines represent model predictions for $V/P = 1.66$ [BREA 89a, MANK 88].

(b) Ratio ϕ/π^- from pBe collisions at 400 GeV/c at $x_F \approx -0.08$ as function of p_T ($= p_\perp$) [AKER 77]. Upper limits from ref. [APPE 75] and the relative yield of μ^- from ϕ decay are also included.

Fig. 4.17 $\alpha_V = V/(V + P)$ as function of the mass M_V of vector mesons relative to pseudoscalars (M_P). High p_T data are compared with results from HRS, TPC, JADE and TASSO [MANK 88].

Fig. 4.18 Number of tracks per event in the away jet as function of x_E for pion triggers [BREA 87d].

Fig. 4.19 See fig. 4.18, but for π^+ triggers with $p_T > 4$ GeV/c. Monte Carlo predictions based upon various assumptions are displayed as well [LOHS 82].

Fig. 4.20 Fractional contributions of different subprocesses to the away jet (with $x_E > 0.5$) as function of $\langle\sqrt{s}\rangle$ for two different values of θ , as predicted by model calculations [BREA 87a].

Fig. 4.21 Normalized x_E distributions for events with a π^+ trigger. All particles with $-2 < y < +2$ in the away wedge are included. The distributions are shown for two different trigger angles and three values of $\langle\sqrt{s}\rangle$ [BREA 87a].

FIGURE CAPTIONS (Cont'd)

Fig. 4.22 x_E distribution of reconstructed π^0 mesons in comparison with the distribution of charged pions [DORO 84].

Fig. 4.23 Number of positive particles relative to negatives in the away jet as function of x_E and results from the Lund model (hatched) [AKES 84].

Fig. 4.24 Away side charge ratio, R , for π^+ triggered events; (a) as a function of y for all particles with $0.3 < x_E < 1.0$ averaged over three different regions of $\sqrt{s}$; (b) as a function of $\langle\sqrt{s}\rangle$ for all particles $-2 < y < +2$ and $0.3 < x_E < 1.0$; (c) as a function of x_E for all particles with $-2 < y < +2$ averaged over three different regions of $\langle\sqrt{s}\rangle$. Data from trigger particles at $\langle\theta\rangle = 47^\circ$ and $\langle\theta\rangle = 18^\circ$ are shown separately [BREA 87a].

Fig. 4.25 The differential conditional probability dn/dx_E for finding a K^- in the away jet for events triggered by π^+ and K^+ . These are proton-beryllium data with $3 \leq p_T (= p_{\text{trig}}) \leq 4$ GeV/c [FISK 78].

Fig. 4.26 $\langle |p_{\text{out}}| \rangle^2$ versus x_E^2 for different data samples. The crosses are the points used in the fit. The straight lines are obtained keeping the intercept the same for all nine data samples [ANGE 80].

Fig. 4.27 The average transverse momentum relative to the trigger particle as a function of $x_{\parallel}$ for the trigger side and the away side from e^+e^- collisions at $W = 34$ GeV. A particle with momentum greater than 4 GeV/c was required as the trigger. The shaded band corresponds to measurements where the sphericity axis was used as a reference axis [ALTH 84].

FIGURE CAPTIONS (Cont'd)

Fig. 4.28 The dependence of the correlation function R on p_T ($= p_{\text{trig}}$) for positive particles in three intervals of $|x_F|$ ($= |x|$). Data are shown separately for positive and negative triggers. Also shown (crosses) is the mean value of R for minimum-bias events when a low- p_T ($p_T < 0.5$ GeV/c) track is required in the central region ($|y| < 0.5$) [ALBR 78a].

Fig. 4.29 Dependence of forward particle production on z , the fragmentation variable for spectator jets, for triggering pions with $2.5 < p_T < 4.0$ GeV/c [HANN 81] and $\theta \approx 30^\circ$ at $\sqrt{s} = 63$ GeV.

Fig. 4.30 Distribution of transverse squared momentum, $d\sigma/dp_t^2$, for secondaries with longitudinal momentum between 2 and 10 GeV/c and with azimuthal angles between 70° and 110° or between 250° and 290° , is given together with Monte Carlo calculation (sect. 6.33, [YEUN 87]) [BREA 84h].

Fig. 4.31 (a) Dependence of average transverse momentum squared (relative to beam axis) on the longitudinal momentum for minimum bias events (triangles) and events with a large p_T hadron emitted at 20° (circles). Only secondaries with azimuthal angles between 50° and 130° or between 230° and 310° are included [BREA 84h].

(b) The average value of the transverse momentum squared versus the longitudinal momentum for events with a 45° trigger hadron. The smooth curves show results of a Monte Carlo simulation for different values of $\langle k_T \rangle$, and for $\langle q_T \rangle = 0.4$ GeV/c [BREA 84h].

FIGURE CAPTIONS (Cont'd)

Fig.4.32 (a) Charge density balance $\Delta q(y|y^{(s)})$ of particles associated with a selected spectator fragment at $y^{(s)} = 2$ and $y^{(s)} = 4$ respectively, from high p_T events. The full line reproduces the corresponding results of non-diffractive inelastic events. The fragments of the trigger jet and the away jet are localized at negative rapidities [DRIJ 80].

(b) Conditional two-particle density in high p_T events with a selected particle in the spectator fragmentation region. The full line reproduces the corresponding result of non-diffractive inelastic events. The fragments of the trigger jet and the away jet are localized at negative rapidities [DRIJ 80].

Fig. 4.33 Directivity diagram of charged-hadron flow in the event plane, $(1/N_{\text{event}})(dn/d\phi)$, as a function of the azimuthal angle ϕ , for $q\bar{q}g$ events (solid circles) and $q\bar{q}\gamma$ events (open circles). In $q\bar{q}g$ events, jet 3 is typically the gluon, whereas jet 1 and jet 2 are q or $\bar{q}$ jets. Solid lines, asymptotic QCD predictions for the flow of soft gluons, for $\phi_q = 0$, $\phi_{\bar{q}} = 153^\circ$, $\phi_{\lambda,q} = 231^\circ$; the normalization is arbitrary, but identical for $q\bar{q}g$ (solid) and $q\bar{q}\gamma$ (dashed). Radial scale is logarithmic [AIHA 86].

Fig. 4.34 (a) Graphs for hard subprocesses and the corresponding string configurations [BENG 84].

(b) Distribution of longitudinal and transverse Lorentz boosts in schemes (a) to (d) for $p_T = 4-6$ GeV/c [BREA 84c].

Fig. 4.35 Rapidity density distributions for secondaries with $p_t > 1$ GeV/c in an azimuthal wedge of $\pm 25^\circ$ about the trigger particle for $p_T = 4-6$ GeV/c, and a comparison to simulations in the framework of schemes (a) to (d) [BREA 84i].

FIGURE CAPTIONS (Cont'd)

Fig. 4.36 Rapidity distributions for the azimuthal trigger hemisphere for different selections of recoil ("away") jets [BREA 84i].

Fig. 4.37 u-distributions at z_f between 0.4 and 0.6 for different types of the trigger particles with $p_T (= p_T^{TR}) = 4 \div 6$ GeV/c [BREA 84i].

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Fig. 4.32

Fig. 4.33

Fig. 4.34(a)

Fig. 4.34(b)

Fig. 4.35

Fig. 4.36

Fig. 4.37

5. THE STRUCTURE OF HADRONS

As already discussed in previous sections, at the origin of large p_T particles emitted in collisions of two high-energy hadrons is supposedly a pointlike interaction of hadron constituents. A study of the production of large p_T particles should therefore reveal properties of the substructure of hadrons.

From the analysis of deep inelastic collisions of electrons, muons and neutrinos, it is known (sect. 2.2.1) that about one half of the momentum of proton is carried by charged partons, i.e. quarks in the QPM. They interact with charged leptons because of their electric charge and with neutrinos because of their "weak charge". However, there is still the remaining fraction of the proton momentum which can be tracked down only via strong interactions. It may be carried by neutral partons which can be detected when they scatter off other strongly interacting partons. Hadroproduction of high p_T particles should therefore be investigated to determine their properties.

Obvious candidates for neutral partons are gluons - quanta of the strong force - which, according to QCD, keep together hadron constituents. Often, only the presence of quarks and gluons is assumed when interpreting the production of large p_T particles in hadronic collisions. Are, however, only these partons involved in hard collisions? A maybe more complete list of potential candidates includes:

- valence quarks,
- gluons,
- sea quarks,
- diquarks,
- virtual mesons,
- supersymmetric partners of conventional partons,
- constituents of quarks.

It may still be extended due to new theoretical ideas or experimental discoveries. Experiments concerning the last two types of candidates would probably require much higher energies than those accessible with present accelerators. On the other hand the contribution to hard scattering of the remaining types of partons has been investigated experimentally and will be discussed in the subsequent sections.

5.1 Valence Quarks

Unfortunately, one cannot observe partons directly as the confinement forces transform a scattered parton into a stream of hadrons - a jet. However, the fastest jet fragment contains often the original parton (sect. 4.1.3).

One therefore expects that scattered u-quarks are the dominant source of large transverse momentum positive pions whereas scattered d-quarks lead to negative pions. In a proton there are twice as many u-quarks as d-quarks. Consequently, in proton-proton collisions π^+ mesons should occur more frequently at large p_T than π^- mesons do. The ratio of cross sections for π^+ and π^- production at the ISR [BREA 84e] is shown in fig. 5.1. Indeed, one finds that π^+ mesons dominate π^- at large p_T . However, the ratio deviates from 2, probably because the fragmentation of scattered valence quarks is not the only source of large p_T pions. Additional charge symmetric processes, e.g. gluon or sea quark scattering, may lower this ratio especially at smaller p_T .

On the other hand, at transverse momenta close to the kinematic limit, the ratio of the two cross sections can be significantly larger than two. This is due to the difference between structure functions of quarks. A u-quark carries often a larger fraction of the proton momentum than a d-quark does (sect. 2.2.1), giving rise to a larger effective parton-parton energy. After scattering, this corresponds to a larger transverse momentum which is transferred to a leading π^+ . Fig. 5.2 shows indeed that the π^+/π^- ratio at large $x_T = 2 p_T/\sqrt{s}$ obtained from protons at 70, 200, 300 and 400 GeV exceeds a value of 2 [ABRA 85, ANTR 79]. For $x_T \approx 0.5$, i.e. for mesons carrying as much as a half of the maximum transverse momentum, the π^+/π^- ratio reaches a value of 2.0. Contrary to p-p collisions, the production rates of positive and negative pions emitted in proton-neutron interactions are equal (fig. 5.2). This is to be expected because the initial p-n system is symmetric in u- and d-quarks.

Consider now the production of high p_T positive pions and kaons in proton-proton collisions. Both mesons contain a u-quark so either one may be produced when a u-quark is scattered. In this case the K^+/π^+ ratio

should be independent of kinematic variables. This is illustrated in fig. 4.15 (and was discussed in sect. 4.1.5) again supporting the idea of valence quark scattering.

A similar situation occurs when comparing production of π^0 and η at large p_T . In collisions of two protons both mesons may originate from scattered u or d quarks. In this case one expects again that the η/π^0 ratio should not depend on kinematical quantities, as confirmed by the data shown in fig. 4.14 (sect. 4.1.4). The η/π^0 ratio does not depend on p_T and is about the same for different c.m.s. energies.

More evidence for the rôle of valence quarks comes from a comparison of proton-proton and pion-proton collisions. A proton is composed of three valence quarks whereas a pion is made up by a quark-antiquark pair. Hence, the momentum of an incoming proton is shared between more partons than that of an incoming pion. Therefore, quarks and antiquarks inside a pion carry larger momenta than quarks inside a proton in the c.m.s. of a pion-proton collision. This should lead to an asymmetric distribution, e.g. of large p_T π^0 in this reference system. Such an asymmetry is in fact observed experimentally [DONA 78b, BONE 87, DEMA 87b]. Fig. 5.3(a) shows the invariant inclusive cross section as function of x_F for neutral pions from pion-proton collisions [BONE 87]. Its maximum is shifted along the direction of the incoming pion. The corresponding distribution from proton-proton collisions in fig. 5.3(b) [BONE 88b] does not show any asymmetry.

Since (anti-) quarks from pions carry on average more momentum than quarks from protons, the effective quark-quark collisions energy is larger for pion beams than for proton beams of the same momentum. This should result in more abundant π^0 production at large p_T in pion-proton collisions. Ratios of cross sections for π^0 production in pp and π^-p interactions were measured by various experiments [DONA 76, BADI 86, DEMA 87b, BONE 88b]. An example [BONE 88b] is shown in fig. 5.4 as function of p_T ; it confirms these expectations.

The quantum numbers of a scattered parton and of the remaining spectator system must add up to the quantum number of the incident

hadron. Hence, a study of spectator jets yields valuable information on the scattered partons. Analyses along these lines have been performed and support the idea of valence quark scattering in the case of $\pi^\pm$ and $K^\pm$ production at high p_T in pp collisions (sect. 6.2).

5.2 Gluons

Gluons may carry about one half of the proton momentum (sect. 2.2.5). The scattering of gluons might contribute, therefore, substantially to the production of high transverse momentum secondaries in hadron-hadron collisions. A verification of gluon scattering is based upon the fact that gluons are flavour neutral and that their fragments may carry any flavour allowed by energy conservation.

5.2.1 Production of K^-

It was argued previously that scattering of valence u- and d-quarks contributes to $\pi^\pm$ and $K^\pm$ production at large p_T . In contrast, in pp collisions K^- at large p_T can hardly come from scattered valence quarks because of their $\bar{u}s$ composition. In principle it may happen that a scattered quark is not contained in the fastest fragment of the jet. The leading fragment could then carry any flavour. However, the probability of such a "rank reordering" is low, otherwise the charge compensation of K^- triggers (fig. 4.10) would be equal to the one of $\pi^\pm$ and $K^\pm$ triggers. It is therefore justified to assume that the production of K^- mesons is not dominated by the scattering of valence quarks.

Next, a consequence of the gluon structure function of protons is discussed. It decreases with x much faster than the structure functions of quarks (sect. 2.2). So, on average, a gluon carries less momentum than any of the valence quarks. This corresponds to smaller parton-parton sub-energies $\sqrt{\hat{s}}$ and, in turn, to smaller transverse momenta after scattering. Consequently, mesons produced from gluon fragmentation should have smaller transverse momenta than those originating from valence quarks. Fig. 5.5 shows the K^-/π^- ratio [BREA 84f]; it decreases with increasing p_T in contrast to the K^+/π^+ ratio (see also [JAFF 89]). This supports the expectation that K^- 's are mainly fragments of relatively slow partons. In principle, scattering of rather slow s-quarks from the sea could also result in a decreasing K^-/π^- ratio. However, it will be argued in sect. 5.3 that sea quarks do not contribute substantially to K^- production at large p_T .

Similar trends should be observable for heavier flavours; first data on $D^{(-)}$ -meson production at $p_T > 2$ are becoming available [ADAM 89].

5.2.2 Parton-parton kinematics

In order to obtain further evidence for gluon scattering, simple kinematics of parton-parton collisions is considered. When two partons scatter, their incoming momenta generally differ in the proton-proton c.m.s. Therefore, the momenta of the two scattered partons are frequently distinctly non-collinear in the overall c.m.s. This situation occurs in particular when a gluon scatters off a valence quark. The incoming gluon momentum is usually smaller than that of an incoming valence quark. After scattering both partons move therefore along the direction of the incoming quark (fig. 5.6(a)) in the pp reference frame. However, the momenta of incident valence quarks may often be of similar magnitude which gives rise to a collinear configuration after the scattering as shown in fig. 5.6(b).

A large $p_T \pi^+$ meson at $\theta \approx 45^\circ$ (i.e. $p_L > 0$) comes usually from a scattered u-quark. If the scattering occurred off another quark, both final state quarks would be most often in the back-to-back configuration sketched in fig. 5.6(b). The quark recoiling against the triggering π^+ is on average positively charged because there are more u- than d- quarks in a proton. The fragments of the recoiling quark should be produced predominantly at $p_L < 0$ and generate an excess of positive over negative charges. On the other hand, when a u- quark scatters off a gluon the back-to-antiback configuration shown in fig. 5.6(a) is more likely. The gluon fragments would have mainly positive longitudinal momenta and no average electric charge. Relevant measurements have been performed at the ISR [BREA 84c]. The ratio of densities of positive and negative away jet fragments recoiling against a large p_T meson at $\theta \approx 45^\circ$ was measured as function of their scaled transverse momenta x_E . It is displayed in fig. 5.7 for fragments with $p_L < 0$ (a) and for those with $p_L > 0$ (b). The ratio rises to a value 2 for the back-to-back configuration ($p_L < 0$) and is significantly larger than that for the back-to-antiback case ($p_L > 0$). This reflects a higher contribution of neutral partons in the back-to-antiback configuration as expected.

A complementary kinematical method to identify gluon scattering was developed in sect. 4.2.1.

5.2.3 Direct photons

It was found experimentally that, among secondaries emitted with large transverse momenta in hadronic collisions, the production of photons exceeds the yield due to decays mainly of π^0 , η and η' . The rate of such "direct" photons provides an argument for a gluon component inside hadrons. It is much higher than expected from a purely electromagnetic mechanism: Large transverse momentum photons can be produced, e.g., via Compton scattering of a photon from the electromagnetic field of one hadron off a quark from the other hadron (fig. 5.8(a)). The production cross section is in this case proportional to the square of the electromagnetic coupling constant α . However, it may also happen that, in a "quark-gluon Compton process", a gluon from one hadron is absorbed by a quark from the other hadron with a subsequent emission of a photon (fig. 5.8(b)); finally photons may be produced by quark-antiquark annihilation (fig. 5.8(c)). Now the cross section is more than one order of magnitude larger as it is proportional to $\alpha \cdot \alpha_s$. Measured cross sections for inclusive production of direct photons in pp collisions [CAVA 89] are shown in fig. 5.9. In this figure theoretical predictions [AURE 88] for photon production, dominated by quark-gluon fusion (fig. 5.8(b)), are included as well. The observed agreement between calculated and measured cross sections proves that a gluon component in hadrons is essential for the production of large p_T direct photons. Having verified the existence of direct photons one can then proceed and try to extract the gluon structure function from the differential production cross sections [AKES 87b] in a way complementary to the "standard" procedure (sect. 2.2).

Due to parton-parton kinematics (as in the case of high p_T π^0 (sect. 5.1)) direct photons are (a) more abundantly produced in π^+p collisions than in pp collisions [BONE 88b, DEMA 87a] and (b) predominantly emitted with $\langle p_L \rangle > 0$ ($\langle p_L \rangle = 0$) in π^+p (pp) collisions [BONE 88a,b]. For a comprehensive review on direct photons it is referred to refs [FERB 84, CAMI 88, FERB 88]; some additional recent data are supplemented in sect. 6.3.2.

5.3 Sea quarks

In the previous section it was argued that the production of large p_T K^- is predominantly due to fragmentation of scattered gluons. However,

quarks and antiquarks from the "sea" are in many respects equivalent to gluons. They are neutral on average, have much lower momenta than valence quarks, and their fragments may carry any energetically allowed flavour. Further experimental evidence is presented now that sea quarks do not dominate the production of K^- at large p_T .

Sea quarks come in pairs. When a s -quark from the sea is scattered it turns up in one of the fragments of a transverse jet whereas its anti-quark is contained in a fragment of the spectator jet. Contrary to that, when an s - $\bar{s}$ pair is created in the fragmentation of a scattered parton both will appear in the same transverse jet.

Experimental data on strangeness compensation in the trigger jet show that the latter occurs for K^- triggers (sect. 4.1.3). One concludes that K^- mesons at large p_T are due to scattering of nonstrange partons.

A further observation rules out a dominance of strange or non-strange sea-quark scattering in K^- production at large p_T . The compensation in the trigger jet of the charge of the trigger particle is much stronger for K^- triggers than for π^+ and K^+ triggers (sect. 4.1.3). This shows that K^- triggers are not produced by scattered strange or non-strange (anti-) quarks.

The reason why scattering of sea (anti-) quarks does not contribute strongly to the yields of large p_T particles is their low average longitudinal momentum, i.e. their soft structure functions. Therefore they cannot attain a significant transverse momentum when scattered off other partons.

Nevertheless, there is one experimental observation indicating that at lower values of p_T strange sea quarks do play a rôle: If a $\bar{s}$ -quark is scattered the associated s quark remains in one spectator jet. There it may form a Λ hyperon. Fig. 5.10 [SMIT 87] shows the yields of Λ hyperons produced along the direction of an incoming proton for proton-proton collisions with a trigger particle with $p_T > 1.15$ GeV/c at $\theta = 90^\circ$. The yield of Λ is much larger for events associated with K^+ triggers than for those which were selected by triggering on protons or pions.

5.4 Virtual mesons

Although scattering of individual sea (anti-) quarks does not contribute strongly to the production of large p_T particles they may still be involved in such processes. A sea antiquark (quark) can form a $q\bar{q}'$ -pair with a valence quark (antiquark) which, after a scattering, carries enough transverse momentum to produce large p_T fragments or to appear directly as a large p_T meson [GUNI 72]. The formation of such a pair which can be considered as a virtual meson, is due to the attractive quark-antiquark interaction which is responsible for the existence of real mesons.

The importance of $q\bar{q}'$ pairs as a source of large p_T mesons can be tested in collisions of $\pi^- (= \bar{u}d)$ with protons. Due to the presence of $\bar{u}$ valence quarks it is possible to form $\bar{u}d$ pairs which would scatter and emerge as large p_T π^- 's. There is, however, no $\bar{d}$ valence antiquark to create $\bar{d}u$ pairs; therefore π^+ at large p_T should be produced at a much lower rate. In other words, the π^-/π^+ ratio should be large. Experiment shows [FRIS 83] that the ratio is close to unity (fig. 5.11) which rules out the scattering of $q\bar{q}'$ pairs as a dominant mechanism. The trend of the data can be explained by recalling that there are three ways to produce a large p_T π^- from valence quark scattering ($\bar{u}$, d from a pion and d from a proton) while there are only two for π^+ . The observed ratio is a little lower than this simple prediction corresponding to a ratio 3:2. However, d -quarks from protons are rather soft (sect. 2.2.2) and contribute thus less to the production of π^- at large p_T . This and a potential contribution of charge symmetric mechanisms, e.g. gluon scattering, will lower the π^-/π^+ ratio. Hence, the observed value is in qualitative agreement with a dominant valence (anti-) quark contribution to the production of large p_T pions, leaving little room for virtual mesons.

5.5 Diquarks

There is hardly any experimental evidence suggesting that, e.g., a valence quark and a sea antiquark of an incoming hadron attract each other strongly enough to form a system which is scattered as a single entity. The forces between two quarks are predicted to be even weaker [DOSC 89]. Consequently such a pair, called "diquark", is not expected to act as a single parton. Contrary to this consideration, experimental evidence will be presented here that diquarks do exist in protons such that their

scattering leads to production of large p_T baryons. The focus will be on proton-proton collisions. It should be pointed out, however, that diquarks have been suspected to play a non-negligible role in deep inelastic lepton-proton scattering [FRED 82, FRED 83] as well as in large angle elastic pp scattering [ANSE 87] and other exclusive processes [KROL 88].

A rather "conventional" mechanism for the production of large p_T protons may be valence quark scattering. The scattered quark has to pick up two sea quarks to form a proton instead of fragmenting into a meson by attaching one antiquark. It is generally assumed that fragmentation of a quark factorizes to a good approximation from the scattering process (sect. 2). In this framework the yield of protons relative to positive pions and kaons should therefore be independent of transverse momentum, collision energy and production angle. Obviously the data in figs 5.12 and 5.13 disagree with this concept [BREA 84g, BREA 85b]. The p/π^+ and p/K^+ ratios do depend on transverse momentum, i.e. they decrease with increasing p_T for $\sqrt{s} \geq 20$ GeV (see also [JAFF 89]); at $\sqrt{s} = 11.5$ GeV one finds an increase (fig. 5, 13(b), [ABRA 85]). The yield of protons relative to pions also decreases with increasing polar angle θ and with collision energy (figs 5.14 and 5.15, [BREA 84g, BREA 85b]). From these measurements one concludes that quark fragmentation cannot be a dominant mechanism of proton production in the kinematic domain considered.

Protons might also be produced at large p_T by scattered gluons fragmenting into proton-antiproton pairs. However, protons and anti-protons should be produced at equal rate by this mechanism, which is in violent disagreement with the data shown in figs 5.12 and 5.13. Hence, scattering and fragmentation of gluons cannot dominate the production of large p_T protons.

Finally, the hypothesis of diquark scattering must be considered. In analogy to single quarks fragmenting into leading mesons, diquarks are expected to hadronize into leading baryons. Obviously, a diquark is a composite object which may break up when it receives a large momentum transfer $\sqrt{Q^2}$. A form factor, e.g. of the dipole form $F(Q^2) = 1/(1 + Q^2/M_0^2)$, suppresses scattering at large Q^2 . The mass parameter M_0 describes the strength of diquark binding (sect. 2.2.5); Q^2 characterizes an elementary

diquark-parton interaction which is not directly observable. Even if Q^2 is therefore not uniquely defined, it must become large with increasing p_T , θ and/or $\sqrt{s}$. Consequently, the data on relative proton yields in figs 5.12-5.15 are compatible with a diquark form factor. In fact, calculations assuming diquark scattering [BREA 85b, EKEL 84, LARS 84, KIM 88, LAPE 82] reproduce reasonably well all these experimental results, as seen in figs 5.12-5.15 [BREA 85b]. The interpretation of the increasing ratio p/π^+ in fig. 5.13(b) requires a diquark structure function which does not decrease as strongly as those of valence quarks such that the suppression due to the form-factor is compensated [KIM 88]. These observations suggest that a diquark mechanism is responsible for the production of most of the large p_T protons in the kinematic range discussed.

An independent argument supporting this idea has been obtained from a measurement of baryon production in spectator jets [SMIT 87]. The data will be discussed in detail in sect. 6.3.

The question now arises whether uu or ud pairs are the prevailing diquarks. In the framework of $SU(6)$, for example, scalar ud diquarks are more abundant than uu vector diquarks by a factor 2. In fact, hadro-production of large p_T baryons suggests that there is non-negligible uu diquark scattering. It is known (sect. 4.1.6) that a significant fraction of protons comes from Δ^{++} decay ($\Delta^{++} = uuu$). While it is hard to imagine that fragmentation of ud diquarks will lead to Δ^{++} formation, Δ^{++} resonances are the natural leading fragments of scattered uu diquarks. The relative importance of ud and uu diquarks will be investigated in more detail in sect. 6.2.2.

The data presented here yield very compelling and direct experimental evidence for a diquark substructure of protons. It remains to be shown how this relates to the extraction of structure functions from lepton-nucleon collisions.

5.6 Transverse momentum of partons

The analysis of deep inelastic scattering of leptons off nucleons has demonstrated that a nucleon behaves like a collection of independent pointlike partons. These partons move approximately along the direction of the nucleon and each of them carries a fraction x of the proton momentum.

In previous sections it was shown that quarks which are at the basis of deep inelastic lepton-nucleon scattering, are also agents in hard collisions of hadrons. It seems therefore reasonable to assume identical structure functions for deep inelastic scattering of leptons and nucleons off nucleons. Nevertheless, one should realize that this is not necessarily so (sect. 2.4.3). So far, the rather good agreement between model calculations and experimental data indicates that this is at least a very useful approximation of the physical reality.

In addition to its longitudinal momentum component a parton may have also some initial transverse momentum k_t before a hard scattering. In fact, the uncertainty relation implies that the average transverse momentum $\langle k_t \rangle$ of a parton, and the dimension R of a volume in which partons are contained should fulfil the inequality

$$\langle k_t \rangle \cdot R \geq \hbar.$$

Taking for R the radius of a proton one obtains $\langle k_t \rangle \geq 250 \text{ MeV}/c$. This so-called primordial contribution to the transverse momentum is a consequence of the finite size of the parent hadron. In reality the transverse momentum of a parton may include also some dynamical part and may therefore be significantly larger than about $250 \text{ MeV}/c$. One assumes usually that gluon bremsstrahlung is responsible for the dynamical component of k_t (see sect. 6.3.3.1).

The transverse momenta k_t of partons influence the four-jet structure of hard hadronic collisions. In case of scattered partons with $k_t = 0$ the spectator partons follow exactly the direction of the incoming hadrons. The situation is different for $k_t \neq 0$. The spectator partons have to have a transverse momentum which compensates at least the primordial part of $\langle k_t \rangle$. Therefore, the spectator jets will not be any longer collinear with the incoming hadrons (fig. 5.16). Circumstantial experimental evidence for a non-collinearity was already discussed in sect. 4.3.2.

A more transparent method to determine the non-collinearity of spectator jets in events triggered by a large p_T particle is presented in the following. It is particularly "economic" for a trigger hadron to acquire a large transverse momentum, if the transverse momenta k_t of both interacting partons are parallel to the transverse momentum generated

by the scattering process. At fixed p_T , this corresponds to smaller parton-parton energies $\sqrt{\hat{s}}$ and momentum transfers $\sqrt{\hat{t}}$ than in the case of $k_t = 0$ (sect. 2.3.1). Hence, the spectator jets recoil with transverse momenta compensating those of the scattered partons. Fig. 5.17 supports this idea [DELL 77]. Azimuthal distributions of spectator fragments are displayed for proton-proton collisions at $\sqrt{s} = 52.5$ GeV with a large p_T hadron at $\theta \approx 20^\circ$ and $\phi = 0^\circ$. Clearly, the spectator fragments show an excess at $\phi = 180^\circ$. This proves that the spectator jets have their largest transverse momentum components antiparallel to the transverse momentum of the trigger particle. These data are compatible with $\langle k_t \rangle \approx 600$ MeV/c (solid line, [DELL 77]) as are the data presented in sect. 4.3.2.

A further interesting trend emerges from the data in fig. 5.17. The ϕ asymmetry is more pronounced for the spectator jet in the longitudinal hemisphere containing the trigger particle than for the other spectator jet. Kinematic considerations in the QPM framework show that the parton yielding the trigger particle carries, on average, a larger fraction x of the beam momentum than the other scattered parton. Hence, based upon the association of trigger jet and away jet with the respective spectator system, the difference of the ϕ asymmetries of both spectator systems was translated into a dependence of $\langle k_t \rangle$ on x (dotted line, [DELL 77]). This was also expected on a theoretical basis [LAND 77].

The same recoil effect was investigated in refs [YEUN 87, BREA 89b, GESS 85] for p-p collisions at $\sqrt{s} = 63$ GeV using the average transverse momentum component $\langle p_x \rangle$ in the scattering plane. The measured values of $\langle p_x \rangle$ deviate significantly from zero for all longitudinal momenta. Since particles with large longitudinal momenta are fragments of spectator jets, this reflects a finite recoil angle of the spectator partons relative to the beam direction which corresponds to a transverse momentum component of about 450 MeV/c.

Another consequence of finite values of $\langle k_t \rangle$ is an acoplanarity of the four jets produced in hard hadronic collisions. As discussed in sect. 4.2.3, the away jet has usually a non vanishing momentum component perpendicular to the scattering plane which cannot be determined uniquely so far. Nevertheless, evidence was obtained for a rather large average effective transverse momentum, $\langle k_t \rangle \approx 1$ GeV/c, which is also required by the absolute normalization of away jet densities in sect. 4.2.1.

Last, but not least, a further consequence of non-vanishing transverse momenta k_t should be mentioned. It was already pointed out above that, for given p_T , the invariant parton-parton energy $\sqrt{\hat{s}}$ and momentum transfer $\sqrt{\hat{t}}$ become smaller with increasing $\langle k_t \rangle$. Since the parton-parton cross section is proportional to $\hat{s}^{-2}$ and contains terms proportional to $\hat{t}^{-l}$, $l = 1, 2$ (sect. 2.3.1), this corresponds to an increase of the inclusive cross section. It was actually from an analysis of the shape and normalization of inclusive cross sections for pions with $p_T \geq 2$ GeV/c from pp collisions at $\sqrt{s} = 52.5$ GeV that first phenomenological evidence for $\langle k_t \rangle \gg 250$ MeV/c was obtained [DELL 77]; a later analysis [FEYN 78] used $\langle k_T \rangle = 848$ MeV/c. This will be again discussed in sects 6.3.3.1 and 6.4.

Summing up this subsection, one can state that very different measurements - of inclusive cross sections and of momenta of away jets - point towards values of $\langle k_t \rangle$ which are much larger than expected from the size of the parent hadrons: $\langle k_T \rangle \approx 0.6 \div 1$ GeV/c. The recoil of spectator jets is consistent with smaller values of $\langle k_t \rangle$: $\langle k_T \rangle \approx .45 \div 0.6$ GeV/c. The numerical value of $\langle k_T \rangle$ depends, of course, to some extent on the particular method by which it was extracted.

In the framework of QCD, a natural explanation of large values of $\langle k_T \rangle$ is based on gluon bremsstrahlung. A thorough theoretical treatment of the large body of data accumulated is really needed, but barely beginning (sect. 6.2.4, [YEUN 87, BREA 89b]). Further detailed analyses should yield a deeper understanding of the primordial and dynamical components of parton transverse momenta.

FIGURE CAPTIONS

Fig. 5.1 The π^+/π^- production ratios versus x_T for $\sqrt{s} = 45$ and 62 GeV and $\theta \approx 50^\circ$. Measurements at $\sqrt{s} = 27$ GeV and $\theta \approx 90^\circ$ are also shown. The cross hatched band displays predictions from a QCD calculation [BREA 84e].

Fig. 5.2 Ratio of π^+/π^- production versus x_T measured at $\theta = 90^\circ$ for (a) pp and pn collisions [ANTR 79], and (b) for pp collisions at 70 and 400 GeV/c [ABRA 85].

Fig. 5.3 (a) Invariant cross section for π^0 production as a function of x_F in the three p_T bins $4.0 < p_T < 4.5$ GeV/c, $4.5 < p_T < 5.0$ GeV/c, $5.0 < p_T < 6.0$ GeV/c. The error bars are statistical only. The error obtained by combining statistical and systematic errors in quadrature is also indicated. Superimposed on the data is the result of the phenomenological fit [BONE 87].

(b) See (a), but for pp collisions [BONE 88b].

Fig. 5.4 Ratio $\sigma(\pi^+p \rightarrow \pi^0X)/\sigma(pp \rightarrow \pi^0X)$ as a function of p_T for $|x_F| < 0.15$. The errors are statistical only; the systematic errors are estimated to be smaller than the statistical errors [BONE 88b].

Fig. 5.5 K^+/π^+ production ratios in proton-proton collisions at $\sqrt{s} = 62$ and 45 GeV for $\theta \approx 50^\circ$, plotted as function of x_T . The solid line is a fit to the K^+/π^+ ratio [BREA 84f].

Fig. 5.6 Typical parton kinematics in the pp c.m.s. for partons of (a) equal and (b) different momenta.

Fig. 5.7 Ratio of particle densities for positive and negative away jet secondaries versus x_E for (a) back-to-back and (b) back-to-antiback configurations [BREA 84c].

Fig. 5.8 Some Feynman diagrams for direct photons in hard hadron interactions.

FIGURE CAPTIONS (Cont'd)

Fig. 5.9 Inclusive cross sections for direct photons produced in pp collisions at $\sqrt{s} = 63$ GeV as a function of p_T [CAVA 89]. The lines display results of calculations from ref. [AURE 87a].

Fig. 5.10 Number of forward π^+ , π^- , p, Λ^0 per large p_T trigger particle at $\theta_T = 90^\circ$ with $1.15 \text{ GeV/c} < p_T < 2.25 \text{ GeV/c}$ [SMIT 87].

Fig. 5.11 The π^-/π^+ ratio versus p_T in 200 and 300 GeV π^-p collisions. The same ratio measured in 200 GeV pp collisions is shown for comparison. The lines represent theoretical predictions [FRIS 83].

Fig. 5.12 Ratios p/K^+ and $\bar{p}/K^-$ of production rates in pp collisions at $\theta = 45^\circ$. The lines represent model predictions [BREA 85b].

Fig. 5.13 (a) Production ratios p/π^+ and $\bar{p}/\pi^+$ as function of p_T . Theoretical predictions are also included [BREA 85b].
(b) Ratio p/π^+ as function of x_T for $p_{\text{beam}} = 70 \text{ GeV/c}$ [ABRA 85].

Fig. 5.14 Ratio p/π^+ and $\bar{p}/\pi$ of production rates at $p_T = 3.5 \text{ GeV/c}$ as a function of the polar angle. The curves represent model predictions [BREA 85b].

Fig. 5.15 Ratio p/π^+ of production cross sections in pp collisions at $x_T = 0.15$ as a function of the centre-of-mass energy $\sqrt{s}$; also included are model predictions [BREA 85b].

Fig. 5.16 Kinematics of parton-parton scattering in the presence of transverse momenta of the incoming partons.

Fig. 5.17 Azimuthal distribution of positive particles with $|p_L/p_{\text{max}}| > 0.5$ (leading particles) and with $p_t > 0.2 \text{ GeV/c}$: (a) for particles in the rapidity hemisphere opposite to that of the trigger and (b) for particles in the same rapidity hemisphere as the trigger [DELL 75]. The triggering particle is at $\theta = 20^\circ$ and at ϕ as indicated by the arrow. The lines show predictions of models, normalized to the observed number of leading particles [DELL 77].

Fig. 5.1

Fig. 5.2(a)

Fig. 5.2(b)

Fig. 5.3

Fig. 5.4

Fig. 5.5

Fig. 5.6

Fig. 5.7

Fig. 5.8

Fig. 5.9

Fig. 5.10

Fig. 5.11

Fig. 5.12

Fig. 5.13(a)

Fig. 5.13(b)

Fig. 5.14

Fig. 5.15

Fig. 5.16

Fig. 5.17

6. PARTON INTERACTIONS

The previous sections 5 and 4 focussed on partons prior to a hard interaction and on their fragmentation thereafter. Valence quarks, gluons and diquarks turned out to be essential. In the present section a closer look is attempted at the dynamics of hard scattering of hadron constituents. Of special relevance is the fact that in hadronic collisions gluons appear in the Born terms on equal footing with quarks.

In order to arrive at an unambiguous picture of the interactions of hadron constituents, a potential background process, i.e. high mass diffraction, is discussed first in section 6.1.

The matrix elements for most relevant subprocesses are predicted (sect. 2.3.2) to exhibit about the same dependence on the scattering angle $\hat{\theta}$ (for $\hat{\theta} \lesssim 90^\circ$) as a consequence of the spin of the exchanged colored gluons. Hence one expects that the shape of the matrix elements can be determined experimentally to a good approximation, without identifying the participating partons, just by a measurement of their momenta. This is described in sect. 6.2 together with a new approach to investigate backward ($\hat{\theta} > 90^\circ$) scattering.

To separate individual lowest order subprocesses the quantum numbers of active partons must be inferred from the detected particles. The feasibility of such a program will be demonstrated in sect. 6.3, where also first results on non-leading order processes will be summarized.

All subprocesses contribute to inclusive distributions, which represent therefore the most global aspects of constituent dynamics as discussed in sect. 6.4.

6.1 Diffraction production

Conceivably, high p_T hadrons at $\theta \ll 90^\circ$ may come from decays of high mass systems X produced diffractively (fig. 6.1) by Pomeron exchange [GOUL 83]. If this were the case one would expect (a) an enhanced phase space density near the rapidity of the high p_T hadron (trigger) due to the other decay products of X , and (b) the appearance of an isolated

"diffractive" proton with $|x_F| > 0.9$ in the longitudinal hemisphere opposite to the trigger. This type of process can be modelled according to published data: the production of the system X is described by the differential cross section $(s/\pi) d^2\sigma/dtdM_X^2 \sim 1/M_X^2 \exp(bt + ct^2)$ [ALBR 76], where t is the square of the four-momentum transfer, $b = 6.5 \text{ GeV}^{-2}$ and $c = 1.5 \text{ GeV}^{-4}$. The multiplicity of charged particles from decays of M_X follows a gaussian distribution with $\langle n_{ch} \rangle$ depending logarithmically on M_X , and with a dispersion $D = \langle n_{ch} \rangle / 2$ [GOUL 83]. The experimental fact that the decays are not isotropic [BERN 85, SMIT 86] is to be taken into account; double diffraction can be treated based upon factorization.

Densities of charged particles with $|\phi - \phi(\text{trig})| > 90^\circ$ produced in events with forward trigger protons are shown in fig. 6.2, where they are compared to expectations from (double) diffractive production [BREA 88]. Obviously, there is strong disagreement yielding an upper limit of at most a few per cent for any diffractive component.

A more direct estimate of a diffractive contribution is based upon a detection of diffractive protons with $|x_F| > 0.9$ in events with high p_T hadrons. Such a measurement was done recently at the ISR [BREA 88] exploiting a new technique to correct for geometrical acceptance and finite momentum resolution for tracks of large momentum [SCHM 87].

The resulting fractions of events with diffractive protons are always smaller than a few per cent in fig. 6.3, where they are given as function of the transverse momentum of trigger particles at $\theta \approx 10^\circ, 20^\circ$, and 45° [BREA 88]. These data confirm more qualitative measurements done long ago [ALBR 78a, DELL 75]. Hence, diffractive production of high p_T particles can be neglected for a discussion of parton subprocesses.

However, it is tempting to speculate about the production mechanism of these rare events. Starting from experimental evidence for pointlike Pomeron-proton interactions [SMIT 85], one may expect large angle scattering of the Pomeron or its constituents off proton constituents [INGE 85]. With first measurements at $\sqrt{s} = 560 \text{ GeV}$ supporting this idea [BONI 88], the ISR data may serve as a useful reference at lower masses M and lower energy $\sqrt{s}$.

6.2 Angular dependence

Gluons are the gauge bosons of QCD. First experimental evidence for their vector nature was obtained from e^+e^- annihilation:

- (a) The angular distribution of the thrust axis in 3-jet events [BRAN 80, BERG 80, DORF 81], when interpreted as manifestation of gluon radiation, favors a spin-1 assignment.
- (b) The decay $Y \rightarrow ggg$ is compatible with vector-gluons [BERG 81a, NICZ 81a, NICZ 81b], while abelian gluons [WALS 80] and scalar gluons [KOLL 79] are excluded.

Already in 1971 [BERM 71] it was pointed out that a study of deep inelastic hadronic interactions would reveal the spin of the exchanged particle. This is obvious from the matrix elements given in Appendix 2. For scalar gluons there is a very weak dependence on the parton-parton scattering angle $\hat{\theta}$, whereas spin-1 exchange corresponds approximately to the familiar Rutherford shape, $\sin(\hat{\theta}/2)^{-4}$, i.e. to a dominance of small angle scattering (fig. A2.1(a), Appendix 2). These trends were exploited in [DRIJ 83] to determine experimentally the spin of the exchanged quantum, i.e. the gluon. The basic idea can be grasped by considering simplified kinematics on the parton level. The rapidities $\hat{y}_T$ and $\hat{y}_A$ of the "trigger" and "away" parton are correlated through conservation laws:

$\hat{y}_A = \hat{y}_T + 2 \ln|\text{tg}(\hat{\theta}/2)|$, which holds for a given transverse momentum and for $k_T = 0$. For $\hat{y}_T > 0$ one expects therefore more away partons at $\hat{y}_A < 0$ for vector gluons than for scalar gluons since the matrix elements favor small values of $\hat{\theta}$. Experimentally, $\hat{y}_T$ is close to the rapidity of the trigger particle (sect. 4.1); in a similar fashion fast away jet fragments can be used to estimate $\hat{y}_A$. Consequently, the rapidity distribution for away jet particles with $x_E > 0.5$ is given in fig. 6.4, together with a full simulation of all lowest order parton diagrams including parton fragmentation [DRIJ 83]. One finds that the hypothesis of vector-gluon exchange is clearly favored.

To actually determine the shape of the matrix elements independently of models one has to know the momenta of both scattered partons in their common rest system on an event-by-event basis. In this case one obtains directly the distribution of scattering angles $\hat{\theta}$ averaged over all contributing subprocesses, which are of nearly identical shape, provided

the scattered partons are not identified. If, e.g. gluons and quarks are not separated, one observes a symmetric distribution also in the case $gq \rightarrow gq$. The approximate universality of angular distributions is called "generalized factorization property" [COHE 83, COMB83, HALZ 83].

Based upon the fact that single hadrons with $p_T \geq 2$ GeV/c approximate the jet (i.e. parton) axis to better than a few degrees (sect. 2.1), distributions of scattering angles $\hat{\theta}$ were derived from events with two neutral pions with $p_T \geq 2.5$ GeV/c in opposite azimuthal hemispheres [ANGE 82]. By transforming into a reference frame with no net longitudinal momentum of the $di-\pi^0$ system, which is close to the parton-parton rest frame, one obtains an approximate parton-parton scattering angle $\hat{\theta}$ for each event. The results are shown in fig. 6.5(a) for various $di-\pi^0$ masses $m(\pi^0\pi^0) \approx \sqrt{s}$. As expected, a dominance of small angle scattering is observed. At fixed $di-\pi^0$ masses, these distributions were fitted to the following formula:

$$d\sigma/d \cos\hat{\theta} \sim \left[\frac{1}{(1 - \cos\hat{\theta})^a} + \frac{1}{(1 + \cos\hat{\theta})^a} \right];$$

This can be reformulated using Mandelstam variables:

$$d\sigma/d \cos\hat{\theta} \sim [1/\hat{t}^a + 1/\hat{u}^a].$$

Independent of $m(\pi^0\pi^0)$, $a = 2.7-2.9$ was found in good agreement with a calculation from lowest order QCD diagrams, which yields also the relative contributions of dominant subprocesses as function of $m(\pi^0\pi^0)$ (fig. 6.5(b), [ANGE 82]). While a somewhat weaker dependence on $\hat{\theta}$ is expected for the process $gg \rightarrow gg$, this is not borne out by the measurement at small $di-\pi^0$ masses. Most probably, this is a reflection of the change of Q^2 with $\hat{\theta}$ at a fixed $di-\pi^0$ mass, which results in an effectively $\hat{\theta}$ -dependent gluon structure function. It is also interesting to note that the dependence of α_s on Q^2 contributes about 0.3 units to a .

Similar analyses were later performed at the SppS collider [BAGN 84a, ARNI 85]; again agreement with QCD predictions was found.

Of special interest is the process $gq \rightarrow gq$ which is not symmetric about $\hat{\theta} = 90^\circ$, if quarks and gluons are distinguished (fig. A2.1(a), Appendix 2). If one managed furthermore to suppress experimentally

symmetric processes, (e.g. $gg \rightarrow gg$ and $qq \rightarrow qq$) and to become sensitive to the region $\hat{\theta} > 90^\circ$, a more subtle check of QCD amplitudes would become feasible. In previous sections it was shown that $\pi^\pm$ and $K^\pm$ triggers come predominantly from scattered valence quarks; hence these triggers suppress gluon-gluon scattering. They do enhance $qg \rightarrow qg$ subprocesses due to their large cross sections unless one considers away jets in the back-to-back configuration (sect. 6.3.1) where recoiling valence quarks become dominant. With other words, events with $\pi^\pm$ and $K^\pm$ triggers (at $y > 0$) and away jets in the back-to-antiback ($y < 0$) configuration should correspond to a rather clean sample of $qg \rightarrow qg$ scatterings (see also sect. 6.3.1).

Some kinematical considerations on the parton level will now illustrate how the shape of amplitudes can be estimated also for backward scattering [BREA 89c]: At fixed invariant parton-parton energy $\sqrt{\hat{s}}$ the angular distribution for two partons with rapidities $\hat{y}_1$ and $\hat{y}_2$ is given by:

$$\frac{d^3\sigma}{d\sqrt{\hat{s}} d\hat{y}_1 d\cos\hat{\theta}} = \frac{\pi}{\hat{s}^{3/2}} \cdot \alpha_s^2 G(x_1)G(x_2) |\Gamma(\cos\hat{\theta})|^2$$

with $\hat{x}_{1,2} = \sqrt{(\hat{s}/s)} \exp(\pm (\hat{y}_1 + \hat{y}_2)/2)$ and $\cos\hat{\theta} = \tanh((\hat{y}_1 - \hat{y}_2)/2)$. Choosing now 2 configurations with $\hat{y}_1 + \hat{y}_2 = \hat{y}'_1 + \hat{y}'_2$, which ensures that $\hat{x}_{1,2} = \hat{x}'_{1,2}$, the ratio A of their differential cross sections is given by the following expression:

$$A = \frac{d^3\sigma}{d\sqrt{\hat{s}} d\hat{y}'_1 d\cos\hat{\theta}'} \bigg/ \frac{d^3\sigma}{d\sqrt{\hat{s}} d\hat{y}_1 d\cos\hat{\theta}} = \frac{|\Gamma(\cos\hat{\theta}')|^2}{|\Gamma(\cos\hat{\theta})|^2}$$

By a Taylor expansion in $1/\Delta$, $\Delta = \coth(\hat{y}_1 - \hat{y}'_1)$, and using the relation $\cos\hat{\theta} = (1 + \Delta \cos\hat{\theta}')/(\cos\hat{\theta}' + \Delta)$ one arrives at:

$$\begin{aligned} A &= \left| \Gamma\left(\frac{1 - \Delta \cos\hat{\theta}}{\cos\hat{\theta} - \Delta}\right) \right|^2 / |\Gamma(\cos\hat{\theta})|^2 \\ &= 1 - \frac{\sin^2\hat{\theta}}{\Delta} \frac{d}{d\cos\hat{\theta}} \log |\Gamma(\cos\hat{\theta})|^2 + O\left(\frac{1}{\Delta^2}\right) \end{aligned}$$

A is therefore a measure of the slope of the angular distribution.

A first attempt to apply this procedure was made recently [BREA 89c]. Events with triggering $\pi^\pm$ and $K^\pm$ mesons with $p_T > 4$ GeV/c at $\theta \approx 45^\circ$ were

selected such that two samples with $\Delta \approx \coth(\langle y_T \rangle - \langle y_T' \rangle) = 6.4$ were obtained. Furthermore, only events with fast away-jet fragments ($x_E > 0.4$) were retained for the analysis. One can therefore approximate the parton rapidities by the trigger rapidity (y_T or y_T') and by the rapidity of the away fragment, respectively. Since $\sqrt{s}$ is related to the invariant mass $\sqrt{s}$ of the trigger particle and the away-jet fragment, i.e. $\hat{s} \approx \tilde{s}/(\sqrt{x_E} \langle z_T \rangle)$, parton-parton kinematics are well known. The ratios A derived according to the formula above are shown as function of $\cos\hat{\theta}$ in fig. 6.6 together with QCD predictions. While for $\cos\hat{\theta} \gtrsim 0.2$ all QCD subprocesses are of similar shape (as discussed above), one finds experimentally evidence for a different shape of the process $ug \rightarrow ug$ - mediated by gluon exchange - for backward scattering ($\cos\hat{\theta} < 0$), as expected (fig. A.2.1(a), Appendix 2). The events contributing here come from the back-to-antiback configuration, where away jets are dominated by gluons (sect. 6.3.1.2). It is also important to note that the normalization of A for $\cos\hat{\theta} \geq -0.4$ excludes abelian QCD and scalar gluons. These measurements of the shape of the angular distribution provide therefore probably the most convincing evidence for the non-abelian nature of QCD, up to date.

Investigations of this type require "flavour tagging" of parton jets and can therefore not be easily realized with jet triggers. A particularly clean reaction, the production of single photons ($ug \rightarrow \gamma u$) was not analyzed yet in this framework.

6.3 Separation of subprocesses

It goes without saying that QCD would be seriously tested by a detailed investigation of individual subprocesses as already indicated in the previous section. While this is a real experimental challenge, considerable progress was made during the past few years. Lowest order quark and gluon interactions were identified, and there is accumulating evidence for gluon bremsstrahlung, direct hadron production via higher twist processes, simultaneous multiple hard scattering, and even successive multiple parton scattering in nuclear matter.

6.3.1 Subprocesses with two partons in the final state

6.3.1.1 Identification of one final state parton

As shown in sect. 4.1.3, the quantum numbers of a scattered parton can be tagged by an appropriate trigger hadron. This finding is mainly based upon the fact that the trigger particle is the leading fragment of a scattered parton, and on the measured quantum number correlations between the trigger particle and the remaining jet fragments. The analysis of inclusive cross section ratios from pp collisions finally lead to the conclusion that high p_T $\pi^\pm$ and $K^\pm$ mesons come from scattered valence quarks (u,d), while K^- are due to gluons, and that protons originate from hard diquark scattering. The relative contributions of different parent partons in table 6.1 were calculated in a QCD framework for trigger hadrons at two values of $\sqrt{s}$ and x_T [ABRA 85]; there is a surprisingly good qualitative agreement with similar calculations done for $\sqrt{s} = 62$ GeV [LOHS 86].

Further qualitative support of this picture comes from the following analysis of data on spectator fragments, which was suggested and applied first in ref. [DRIJ 79b].

The flavor of a scattered parton yielding a trigger hadron at $\theta \ll 90^\circ$ and the flavor of the spectator jet at $\theta \approx 0^\circ$ are correlated, if both originated from the same incoming hadron. Hence, the quantum number composition of the spectator system should differ in a characteristic way for different types of trigger particles (fig. 6.7).

For pp collisions this is investigated here using the ratios $R_S^\pm(h_1/h_2) = \rho^\pm(h_1)/\rho^\pm(h_2)$ of the densities of positive and negative secondaries with $p_t > 0.3$ GeV/c in events triggered by large p_T hadrons h_i , $i = 1, 2$. Acceptance effects tend to cancel for ratios defined in this way. Measurements of $R_S^\pm$ are displayed in fig. 6.8 as function of x_F for $h_1 = \pi^-, \pi^+ + K^+, p$ and $h_2 = \pi^+, K^+, K^-, \pi^-$ [BREA 84c, 85a]. For $x_F < 0$, i.e. in the longitudinal hemisphere not containing the trigger, $R_S^\pm$ is generally consistent with unity. The ratios $R_S^\pm$ obtained from π^+, K^+ and π^- triggers produced at $\theta \approx 45^\circ$ and $p_T = 4 \div 6$ GeV/c in fig. 6.8(a,b) indicate that $\rho^\pm(\pi^+) = \rho^\pm(K^+)$. This confirms the conclusion reached earlier (sect. 4.1.3) that both π^+ and K^+ triggers come from the same

scattered partons, i.e. u-quarks. The measured relation $R_s^+((\pi^+ + K^+)/\pi^-) > 1 > R_s^-((\pi^+ + K^+)/\pi^-)$ in fig. 6.8(a,b) reflects the smaller positive charge of the spectator system for π^+ or K^+ triggers relative to π^- triggers, as expected from d-quark (or gluon) scattering in the case of π^- triggers. Absence of Δ^{++} production in spectator jets associated with π^+ triggers and its observation for π^- triggers (not shown, [FISC 81b]) also favors this interpretation.

Relative spectator densities for K^- triggers at $\theta \approx 45^\circ$ in fig. 6.8(c) differ significantly from those measured for π^+ or K^+ triggers. As the latter originate predominantly from u-quark scattering, high p_T K^- cannot be predominantly higher rank fragments of u-quarks (and, therefore, of the less abundant d-quarks). The measured values of $R_s^\pm$ for $x_F > 0$ are consistent with neutral partons as the dominant source of K^- triggers. In such a scheme the remaining uud spectator system should produce more positive hadrons than the ud-spectator left behind by π^+ triggers.

All measured ratios in fig. 6.8 are in rather good agreement with predictions from lowest order QCD [YEUN 87].

Production mechanisms for protons at high transverse momentum are considered next. Proton production by hard scattered valence quarks implies $R_s^\pm(p/\pi^+) = R_s^\pm(p/K^+) = 1$ (u-quark scattering), or $R_s^\pm(p/\pi^-) = 1$ (d-quark scattering) independent of x_F . On the other hand, $R_s^\pm(p/\pi^\pm)$ should decrease with increasing $x_F > 0$ if trigger protons were diquark fragments. In this case a single spectator quark may produce fast positive fragments less often (fig. 6.7) than a di-quark spectator (π^- triggers) which is an efficient source of fast protons. The experimental ratios of corrected particle densities are displayed in fig. 6.9 for triggering particles at $\theta \approx 10^\circ$ and $p_{lab} > 13$ GeV/c [BREA 87c]; the ratios are significantly smaller than 1 for $x_F > 0.2$ ruling out a dominance of valence quark scattering. Similar correlations were already found much earlier [ALBR 78a, DRIJ 79b]. The deviations from unity become smaller as θ increases (not shown, [BREA 87c]). This reflects the decreasing contribution of diquark scattering, due to the diquark form factor.

A most direct proof of diquark scattering would be a measurement of the absence of baryon number in spectator jets. For trigger protons produced at $\theta \approx 90^\circ$ and $p_T \approx 1.2 \div 2.2$ GeV/c a significant suppression of protons relative to π^+ triggers is shown in fig. 6.10 [SMIT 87]; the data are given as function of $x'_F = x_F/(1 - 2p_T/\sqrt{s})$, a variable which approximately takes into account the reduced momentum range of the spectator system.

While it is now rather clear that hard diquark scattering is at the origin of proton production of high p_T , the relative contribution of ud and uu diquarks is unknown as yet. Actually, observation of Δ^{++} production at large transverse momentum (sect. 4.1.6) suggests a non-negligible contribution of hard scattering of uu diquarks. In order to estimate the relative importance of ud and uu scattering, the acceptance corrected spectator ratio $R'_s(h) = \rho^-/\rho^+$ is studied in fig. 6.11 for trigger hadrons h emitted at $\theta \approx 10^\circ$ and $p_T > 2$ GeV/c [BREA 87c]. Indeed, for π^+ (π^-) triggers which correspond to ud (uu) spectator jets, one finds good agreement in fig. 6.11(a) with charge ratios from ud (uu) target jets from $\bar{\nu}p$ (νp) scattering [ALLE 83, ARNE 85]. The charge ratio $R'_s(p)$ measured for proton triggers is substantially larger. A quantitative comparison with predictions from single quark spectators is shown in fig. 6.11(b). For a numerical calculation data from current jets [ALLE 83, ARNE 85] were used (fig. 4.10). Kinematics dominate the relation between the fragmentation variable z (ν data) and x_F (pp collisions): $\langle z \rangle \approx 0.2 + 1.3 \cdot x_F$ [BREA 87c]. If ud and uu diquarks were contributing to the proton wave function according to SU(6) one would expect 67% spectator u-quarks and 33% spectator d-quarks; the data are well below this prediction. On the other hand, if only ud systems would scatter at large angles, R'_s should be reproduced entirely by spectator u-quark fragmentation. For proton triggers the spectator ratio $R'_s(p)$ is in between the two extremes considered above, which is suggestive of a non-negligible contribution of hard uu scattering. This is also in line with the measured relative yield of Δ^{++} at high p_T (sect. 4.1.6) which tends to be larger than the corresponding ratio measured for target jets (mainly ud) in deep inelastic μp scattering [ARNE 86b].

So far it was shown that in the case of a single identified hadron at high p_T only a very limited number of hard subprocesses can contribute; this was achieved by "flavor tagging" of one of the scattered partons. The way this conclusion is reached is summarized in table 6.2.

6.3.1.2 Dominance of a single subprocess

In a final step it will be attempted now to demonstrate that sets of data can be obtained which are dominated by a single subprocess. In order to do so the identity of the parton which is at the origin of the away-jet recoiling against the trigger particle, must be determined at least in a statistical sense. Also its kinematical configuration relative to the trigger must be measured. The argumentation will be based upon the acceptance corrected charge ratios $R_a = \rho^+ / \rho^-$ of densities of away-jet particles with $|\phi - \phi(\text{trig})| > 145^\circ$. For trigger particles at $\theta \approx 45^\circ$, R_a was measured as function of $x_E = p_t / p_T$ separately for $y < 0$ (back-to-back) and $y > 0$ (back-to-antiback).

The strong increase with x_E of R_a for π^+ and K^+ triggers in fig. 5.7(a) [BREA 84c] for the back-to-back configuration is consistent with a dominance of valence quark fragmentation with a large contribution from u-quarks. Keeping in mind that π^+ and K^+ triggers are mainly due to scattered u-quarks this means that the back-to-back configuration is dominated by "classical" quark-quark interactions with a large contribution from the subprocess $uu \rightarrow uu$.

For π^+ and K^+ triggers R_a is rather small in the back-to-antiback configuration (fig. 5.7(b)) indicating a larger contribution of neutral partons. Because of kinematics (sect. 5.2.2) these partons have also a softer structure function than valence quarks. Consequently, the subprocess $ug \rightarrow ug$ is very likely to dominate events in which both π^+ (or K^+) triggers and away-jets are produced in the same longitudinal hemisphere. Due to the rather soft d-quark structure function (sect. 2.2) a larger fraction of high p_T π^- mesons may be due to scattered gluons (table 6.1). This would explain why all features of R_a discussed above are less pronounced for π^- triggers than for π^+ and K^+ triggers.

Charge ratios R_a for K^- triggers are shown in fig. 6.12 [BREA 85a] together with the ratios obtained from π^+ and K^+ triggers. For the back-to-back configuration (fig. 6.12(a)), the measurements agree with those for π^+ and K^+ triggers, while they exceed the ratios R_a from π^+ and K^+ triggers for $y > 0$ (fig. 6.12(b)). These data imply a dominating subprocess $gu \rightarrow gu$ with gluons yielding the K^- triggers and away jets due to u-quarks in both configurations.

Comparing away jets recoiling against proton and positive meson triggers at $\theta \approx 10^\circ \div 45^\circ$ one finds similar charge ratios when integrated over x_E and y . Hence, diquarks scatter off quarks and gluons in much the same way as quarks [BREA 87c].

In conclusion, there are experimental means to select samples of events which are dominated by the hard subprocesses $uu \rightarrow uu$, $du \rightarrow du$, $ug \rightarrow ug$ and diquark-parton $\rightarrow$ diquark-parton. With a two-arm spectrometer capable of particle identification one would be able to study nearly on an event-by-event basis many individual subprocesses in an appropriate kinematic range such as (for pp collisions): $uu \rightarrow uu$ by triggering on two π^+ mesons, $gu \rightarrow gu$ (K^- , π^+), $gg \rightarrow gg$ (K^- , K^-), diquark-gluon $\rightarrow$ diquark-gluon (p , K^-) or diquark-diquark $\rightarrow$ diquark-diquark (p , p). In a similar fashion subprocesses from meson-nucleon collisions could be analyzed. Measurements of this type would be useful to e.g. get a handle on the diquark structure function in protons, on structure functions of kaons, or to study fragmentation of flavor tagged spectator systems in color singlet or octet states.

6.3.2 Direct photons

Whereas the final state particles discussed in the previous section are fragments of scattered partons, the photons considered now are directly produced in hard parton-parton interactions, i.e. they are not part of a jet. A more direct access to the basic scattering mechanisms is therefore anticipated. Because of recent excellent reviews of this field [FERB 84, CAMI 88, FERB 88], only those experimental results are covered here which provide evidence for individual subprocesses.

The pointlike coupling of photons to electric charges is well understood and has been exploited for decades as a probe of hadron structure. Whereas in deep inelastic scattering of charged leptons one is sensitive to quarks inside nucleons, the observation of direct photons in the final state of deep inelastic hadronic collisions provides information on gluons inside hadrons (Compton diagram, fig. 5.8(b)) and on gluon fragmentation (annihilation, fig. 5.8(c)). In collisions of charged mesons and nucleons, both subprocesses are dominated by u ($\bar{u}$)-quark scattering due to its larger absolute charge. Consequently, rather clean

u-quark jets recoil against single photons, e.g. in pp and π^+p collisions, where annihilation is negligible. A measurement of the charge ratio R_a of these jets (fig. 6.13) is qualitatively compatible with expectation [CASS 87]. The charge ratio from recoiling jets in π^-p collisions is smaller, due to a large gluon contribution from annihilation process.

Of special interest relative to the annihilation scheme are projectiles with valence $\bar{u}$ -quarks, e.g. π^- or $\bar{p}$. While there is no significant difference between inclusive direct photon production at $p_T \leq 5$ GeV/c in pp and $p\bar{p}$ collisions at $\sqrt{s} = 53$ GeV [AKES 85], π^- are more efficient in producing single photons than π^+ [BONE 88a, DEMA 87a]. The difference between inclusive cross sections for direct γ production in π^- and π^+ induced interactions is essentially due to annihilation (fig. 5.8(c)), since Compton contributions tend to cancel. A measurement of this difference at a beam momentum of 280 GeV/c is displayed in fig. 6.14 as function of p_T for $-0.45 < x_F < 0.55$ [BONE 88a]. It compares well with a next-to-leading order QCD calculation [AURE 87], using optimized perturbation theory [STEV 86] and a QCD parameter $\Lambda = 200$ MeV together with a soft gluon structure function [OWEN 84] (solid line); the dashed line corresponds to $\Lambda = 400$ MeV and a hard gluon structure function [OWEN 84]. With a good control of the annihilation diagram a reliable determination of the gluon structure functions from the dominating Compton process should then become feasible. This is very promising especially, if further kinematic constraints are provided by measurements of away jets.

It is worth pointing out that single photon events provide a source of flavor-tagged spectator jets for fragmentation studies.

Production of pairs of direct photons (fig. 6.15) may proceed via (a) an annihilation subprocess, (b) a higher order mechanism involving a virtual quark-loop (box diagram), or (c) bremsstrahlung. The contribution from bremsstrahlung can be minimized experimentally by removing events with hadrons close to any of the two photons ("isolation cut"). Although the experiments are by no means trivial, evidence for $\gamma\gamma$ production was indeed obtained [KOUR 82, AKES 86b]. A more recent measurement provides a signal of 139 ± 22 events (after background subtraction) in the reaction $\pi^-p \rightarrow \gamma\gamma X$ at 280 GeV/c [BONV 89]. Photons were detected at $p_T > 2.75$ GeV/c and $-1 < y < 1.25$; a contribution from bremsstrahlung is included. The

cross section for photons from this data sample is given in fig. 6.16 as function of p_T . It agrees with predictions [AURE 85] based upon "beyond leading logarithm" QCD using set I structure functions of ref. [OWEN 84] and $\Lambda = 200$ MeV. In the low p_T range the box diagram contributes about 10% to the cross section.

As in the case of single direct photons it would be instructive to measure differences between beams of particles and antiparticles.

6.3.3 Higher orders in α_s

6.3.3.1 Gluon bremsstrahlung

In sect. 5.6 it was concluded that, in the framework of lowest order QCD, an initial parton transverse momentum $\langle k_T \rangle \lesssim 600$ MeV/c is required to reproduce data on spectator recoil whereas the azimuthal distributions of away jets tend to imply $\langle k_T \rangle \approx 1$ GeV/c. Hence, $\langle k_T \rangle$ is larger than expected from the proton size. Starting from this observation, possible dynamical contributions to $\langle k_T \rangle$ due to parton branching processes - gluon bremsstrahlung - will be investigated here. In such an approach a hard parton-parton process would look like in fig. 6.17. On a qualitative basis one anticipates therefore, that gluon bremsstrahlung influences spectator jets less than inclusive cross sections and away jet properties, both of which are directly modified due to gluon radiation off the relevant partons.

An analysis along those lines was performed by comparing predictions of the Monte-Carlo program JETS [YEUN 87], which incorporates initial and final state branching processes, to measurements performed at the ISR. As an example, fig. 6.18 illustrates a distribution of primordial transverse momentum with an average of 450 MeV/c evolving due to initial state parton branching to a broad distribution of initial transverse momentum k_T at the time of hard scattering. One has therefore reconciled large values of $\langle k_T \rangle$ with a relatively small value of the primordial component. In this approach inclusive cross sections for meson production in the ISR energy range are well reproduced down to $p_T \approx 4$ GeV/c with $\langle k_T \rangle = 1$ GeV/c (not shown, [YEUN 87]).

A problem left unsolved in sect. 4.2.1 was the discrepancy by about a factor 1.5 between measured and predicted particle densities from away jets. This is illustrated again in figs 6.19 and 6.20; where a rapidity distribution for $0.5 \leq x_E \leq 1$ and an x_E distribution for $|y| < 2$ are displayed for π^+ triggers with $p_T \geq 4$ GeV/c [BREA 89b, YEUN 87]. Predictions from the JETS code [YEUN 87] with and without gluon bremsstrahlung are included. Remarkable agreement is achieved in the latter case. Even the p_T distribution from spectator fragments is rather well reproduced after inclusion of gluon bremsstrahlung as shown in fig. 4.30. Gluon bremsstrahlung would also explain the lack of success of the simple string model presented in sect. 4.4.

Encouraged by these results one may ask if there is more direct experimental evidence for gluon bremsstrahlung, i.e. for the appearance of additional jets, as e.g. in the case of three-jet events in e^+e^- annihilation. A search for events with 3 large angle jets was done at the ISR [ANGE 81] in events triggered by π^0 with $\langle p_T \rangle \approx 6$ GeV/c. Using a cluster algorithm, events of this topology were found at a level predicted on the basis of QCD. These findings were confirmed more recently by experiments using a calorimetric jet trigger [AKES 86a] and a neutral transverse energy trigger [ANGE 88]. They are generally well described in the framework of QCD incorporating higher order processes.

6.3.3.2 Double parton scattering

Finally, experimental evidence for a more exotic higher order process will be considered here. Production of deuterons as well as of antideuterons [CRON 75, ABRA 87] and tritium [ABRA 87] at large p_T was measured in pp collisions. The ratios of inclusive cross section for d and p production in fig. 6.21 [ABRA 87] are in the order of $4 \cdot 10^{-3}$ at 70 GeV/c and of 10^{-2} at 300 GeV/c [CRON 75]; for $p_T \geq 2$ GeV/c no strong dependence on transverse momentum is observed. Guided by ideas developed for deuteron production at lower beam energies and at smaller momenta, an interpretation based on a fusion model was proposed [SCHW 63, ABRA 87]. In this picture a proton and a neutron of momenta $\vec{p}/2$ fuse to give a deuteron of momentum $\vec{p}$: $\sigma_d(\vec{p}) = C(\vec{p}) \cdot \sigma_p(\vec{p}/2) \sigma_n(\vec{p}/2) = C(\vec{p}) \sigma_p^2(\vec{p}/2)$. The coefficient $C(\vec{p})$ incorporates the p-n correlation and the probability for fusion; its measured value is consistent with being independent of p_T and of $\sqrt{s}$. All these features are compatible with a fusion model, as well as the data on $\bar{d}$ production (not shown, [ABRA 87]). Also the data on $\bar{d}$

production in e^+e^- annihilation [ALBR 85] tend to be consistent with a $\bar{p}\bar{n}$ fusion mechanism. A complete understanding of the data depends then on a model for simultaneous production of 2 (anti-)nucleons yielding (anti-) deuterons in pp collisions. The least adhoc mechanism seems to be one in which two simultaneous diquark-parton (parton-parton) scatterings occur in one pp interaction (fig. 6.22) producing two baryons (antibaryons) [EFRE 87]. In principle, this rather simple scheme could be verified experimentally by looking for the absence of spectator jets in case of deuteron production. Also two rather soft away jets, which probably overlap partially and compensate the transverse momentum of the trigger nucleus should be detected in contrast to harder single away jets balancing the momentum of meson and nucleon triggers.

While this scheme of double parton scattering may seem a bit far fetched at first, it should be pointed out that events with 2 balanced jet-jet systems were observed recently at the ISR [AKES 87a]. Also these topologies were attributed to double-parton scattering. There is no experimental evidence for a contribution from double-parton scattering to high p_T production of pseudo-scalar mesons.

6.3.4 Higher twists

All measurements of inclusive production of pseudoscalar mesons and (anti-)protons at high transverse momentum considered so far were successfully reproduced quantitatively in the framework of QCD where subprocesses involving a minimal number of single "free" partons are assumed (minimal twist processes). Consequently, these hadrons appear in the final state as (leading) fragments of jets. In this spirit also proton production from diquarks may be attributed to a minimal twist process. There is a priori no reason for the absence of hard processes in which pairs of (anti-) quarks or gluons from one hadron are involved directly in a coherent fashion as illustrated in fig. 6.23(a). In this case systems with more than the minimal number of partons participate, these mechanisms are therefore labelled "higher twist" processes. From this point of view diquark scattering is a particular higher twist process, as is double parton scattering.

As mentioned in sect. 5.4 there is no experimental evidence from charged pion production at high p_T in π^-p collisions for higher twist

effects due to quark pairs in the initial state. On the other hand a contribution to Drell-Yan processes of initial state quark pairs was established experimentally (fig. 6.23(b), [CONW 88]). Some evidence of higher twist effects due to the structure of incident hadrons comes from π^-p collisions with two very energetic transverse jets (fig. 6.23(c), [NAUD 86]), in which no forward spectator seems to be present.

Here only the direct involvement of parton pairs in the final state, i.e. direct production of single mesons, will be further discussed. Cross sections for higher twist processes are expected generally to decrease more strongly with the momentum transfer in the pointlike interaction than those for minimum twists. However, since the direct hadrons carry all of the momentum transfer $\sqrt{Q^2}$, they may dominate at large p_T over hadrons originating from single parton scattering. The latter retain only a fraction of $\sqrt{Q^2}$ due to parton fragmentation. A comprehensive theoretical treatment can be found in [BENA 87d] for mesons; direct baryons are considered in [GROZ 87].

First experimental evidence for direct meson production comes from an experiment investigating ρ^0 production at large transverse momentum in π^-Be collisions at 300 GeV/c [BENA 87a]. Minimal twist and higher twist mechanisms may contribute as illustrated in fig. 6.24. To enrich a data sample with hard processes, ρ^0 with $p_t > 2$ GeV/c recoiling against hadrons with $p_T > 1$ GeV/c were searched for. The ρ^0 production rate was then determined as function of $x_E = p_t(\rho^0)/p_T$ from invariant mass distributions (fig. 6.25). Between $x_E = 1 \div 2$ and $x_E > 2$ an increase of about 38% of the fraction of ρ^0 -resonances among all $\pi^+\pi^-$ combinations was found. Predictions based upon the LUND Monte-Carlo (lowest order QCD, [BENG 85]) reproduce the background under the ρ^0 signal well, but fail to give the correct dependence on x_E of ρ^0 production. Of course, the predictions depend on the ρ^0 fragmentation function. However, a modification of the shape of this function is claimed to lead to inconsistencies with other data. A modified LUND Monte-Carlo which incorporates higher twist processes [INGE 86] agrees well with the data. It suggests that about 55% of all ρ^0 mesons at $x_E \geq 2$ are produced directly.

Motivated by this positive result, data on η' production at high p_T were investigated in ref. [BENA 87b]. In sect. 4.1.4 it was mentioned that for inclusive production $\eta'/\eta \approx 1.7$ is measured, while $\eta'/\eta \approx 0.8$

would result from single quark scattering and SU(3) mixing. A combined analysis of π^0 , ω , η and η' production at high transverse momentum leads to the conclusion that a higher twist contribution, especially of the gluon-gluon fusion type (fig. 6.26) accounts fully for η' production with negligible change of the inclusive ratios η/π^0 and ω/π^0 .

It may also be useful to analyse the data on inclusive ϕ production in pBe collisions at 400 GeV/c [AKER 77] in this framework.

More precise data on ratios of inclusive cross sections as well as a detailed investigation of event topologies would certainly be welcome.

According to recent ideas, also gluonia (glueballs) are produced via higher twist mechanisms (fig. 6.27, [BENA 87c]) in addition to lowest order QCD processes. A search for the glueball candidates ι (iota) and θ recoiling against single particles with $p_T > 0.9$ GeV/c yielded upper limits only [BENA 87c]; the upper limit for θ production is a factor 14 below the QCD prediction incorporating higher twists, while upper limit and the predicted signal are compatible for ι production.

6.3.5 Nuclear effects

Inelastic hadron-nucleus (hA) collisions involve basically only surface nucleons due to strong absorption, i.e. large cross sections. Hence, the inelastic cross section $\sigma_{\text{in}}(\text{hA})$ is proportional to the number N_s of surface nucleons: $N_s \sim R_A^2 \sim A^{2/3}$; $R_A \sim A^{1/3}$ is the radius of a nucleus of baryon number A. Inclusive differential cross sections are found to be proportional to A^α with $\alpha = \alpha(x_F) \simeq 0.45 \div 0.75$ [BART 83]. For hard scattering, i.e. particle production at high p_T , $\alpha = 1$ was naively anticipated, since the incident partons would "see" all target partons due to the relatively small QCD cross sections. Surprisingly, $\alpha > 1$ was measured for inclusive production of $\pi^\pm$, $K^\pm$ and $(\bar{p})$ at $p_T > 1$ GeV/c as shown in fig. 6.28 [CRON 75]. Such a "cooperative" phenomenon was not understood well theoretically for quite some time.

With accumulating confidence into the reliability of QCD predictions for hadronic hard processes, these data were reanalyzed in the framework of lowest order QCD subprocesses [LEV 83]. Detailed calculations assuming multiple parton-parton scattering were done for pion and kaon production, good agreement with the data (fig. 6.29) for $p_T \geq 1$ GeV/c was achieved.

This can easily be understood qualitatively: The QCD subprocesses resemble Rutherford scattering with large cross sections for small angle scattering. Hence, it is more "economic" for an incident parton to undergo multiple scattering in nuclear matter than to generate a high p_T large angle hadron by a single hard impact in the case of a proton target.

In this scheme also π^0 production at high p_T in $\alpha\alpha$ collisions was predicted [LEV 83]. Recent data in fig. 6.30 [ANGE 87] on the relative inclusive yields of π^0 in $\alpha\alpha$, and pp collisions are compatible with these calculations. One may therefore conclude that the picture of multiple parton scattering is compatible with observed features of pion (and kaon) production off nuclear targets.

It should, however, be pointed out that (a) $\alpha(p_T \leq 8 \text{ GeV}/c) \approx 1$ for pion production in αp and dd collisions [FISC 88] and (b) proton yields may be more strongly enhanced by nuclear targets than those of mesons and antiprotons [BREA 86]. The relevant data are shown in figs 6.31 and 6.32, respectively. While there is no detailed calculation yet for the latter feature, one may speculate that the diquark form factor is at its origin.

Corroboration of these data and of its detailed understanding as well as an "application" of high p_T processes as a diagnostic tool of high-energy nuclear collisions [GEIS 88] is left for the future.

6.4 Single particle inclusive cross section

Inclusive single particle cross sections summarize rather global aspects of constituent scattering, hence they are much harder to understand than the detailed correlation data discussed in the foregoing sections. Being easily measurable, the inclusive cross sections, however, and the proposed interpretations have influenced strongly the development of the field.

In fig. 6.33 the inclusive cross section for the production of π^0 and charged hadrons is displayed as a function of the transverse momentum p_T for centre-of-mass energies $5 \text{ GeV} < \sqrt{s} < 1800 \text{ GeV}$ [BEIE 78, LOHS 82, ARNI 82a, ABE 88]. The compilation of lower energy data [BEIE 78] in fig. 6.33(a) includes an empirical parametrization of the invariant cross section containing two terms:

$$A \exp(-B\sqrt{(p_T^2 + M_1^2)})g(x_T - x_0) \text{ and } C(1 - x_T)^m (p_T^4 + M_2^4)^{-n/4}$$

with free parameters A , B , C , M_1 , M_2 , m , n and x_0 . The first term reflects basically the familiar exponential shape for $p_T < 1$ GeV/c. It is modified by the function $g(x_T - x_0)$ describing a suppression of the cross section for $x_T > x_0 = 0.5$, presumably due to phase space effects. This is clearly borne out by the data at 12 and 24 GeV/c. The second "inverse power" term is often used to empirically describe high p_T data. Historically, the observed strong excess over the expected exponential behaviour (fig. 1.1) triggered all investigations summarized in this review. Even more striking today is the strong increase of the inclusive cross section with $\sqrt{s}$ at fixed p_T . It amounts to more than three orders of magnitude at $p_T = 3$ GeV/c between the lowest ISR and the Tevatron collider energy (fig. 6.33). Qualitatively this behaviour had been predicted as a consequence of hard parton-parton scattering [BERM 71]. On the other hand the observed p_T dependence was stronger than expected in the framework of the naive quark-parton model which presumed scaling of the structure and fragmentation functions [ELLI 74].

If one considers all subprocesses in first order QCD and takes into account scale breaking effects [OWEN 78], data and predictions in fig. 6.34 are in reasonable agreement for $p_T \geq 5$ GeV/c; especially the shape is well reproduced. At smaller transverse momenta the measured cross section is larger than the predicted one by factors up to ten. This observation has been interpreted to be indicative of an additional scale breaking effect - not considered in ref. [OWEN 78] - of particular importance at intermediate p_T . The effective transverse momentum k_T of the incoming partons (sect. 5.6) is a natural source of this additional scale breaking effect [DELL 77, FEYN 78]. As discussed in sect. 6.3.3.1 a convolution of the primordial transverse momentum of partons and of gluon bremsstrahlung gives rise to the large values of $\langle k_T \rangle$ ($\langle k_T \rangle \approx 1$ GeV/c) which are actually needed to reproduce the cross section.

The onset of the additional scale breaking effects at intermediate p_T is most apparent if one parametrizes the measured cross section according to

$$E \, d^3\sigma/dp^3 = p_T^{-n} f(x_T, \theta)$$

and determines n locally. The result of such an analysis is shown in fig. 6.35 [GEIS 84]. At $x_T \geq 0.25$ the exponent derived from the inclusive single particle spectra coincides with the corresponding result derived from large transverse energy jet cross sections. The asymptotic values tend to agree with lowest order QCD predictions neglecting k_T effects [FIEL 78b]. The estimated increase of $n(x_T)$ at small x_T due to the influence of the effective transverse momentum k_T of the primary partons is indicated by the full (broken) line in fig. 6.35. One further observation is of interest: the exponent for high p_T proton production is larger than the corresponding one for mesons. This additional scale breaking effect for high p_T baryon production may be due to the diquark form factor discussed in sect. 5.5.

This interpretation of the measured exponent n assumes implicitly that a potential K factor does not depend on $\sqrt{s}$, x_T (and θ); an experimental K factor is given by the ratio between data and predictions in lowest order QCD. Such a K factor cannot easily be determined for parton-parton scattering yielding single particles at high p_T since rigorous simulations of lowest order processes require excessive computational efforts. Basically, this is due to the hadronization process. Introduction of approximations may lead to uncertainties of the overall scale of the order of 2 [YEUN 87]. Nevertheless a K -factor of about 2 seems to be required.

As demonstrated by fig. 6.36 [AKES 82] the inclusive parton (i.e. jet) cross section at the highest ISR energy is a few orders of magnitude larger than that for single particles; this results from the fact that no specific (and rare, $\langle z_{TR} \rangle > 0.7$, sect. 4.1.1) fragmentation mode is selected.

It follows also from fig. 6.34 that the observed inclusive cross section cannot be traced back to a single constituent scattering process; rather a sum of different processes, whose relative strength varies with $\sqrt{s}$, x_T and θ , has to be considered [GLÜC 81]. Differences of inclusive cross sections may be dominated by one subprocess as pointed out in ref. [ANSE 83].

It has been shown [FEYN 78] that the shape of longitudinal distributions depend sensitively on the strength of the different subprocesses

involved as well as on the momentum dependence of the respective structure- and fragmentation-functions. The latter are, of course, to a large extent fixed by deep inelastic lepton-nucleon and e^+e^- data as discussed in sect. 2. A recent experimental result covering a large fraction of longitudinal phase space is shown in fig. 5.3(b). The cross section for production of π^0 at fixed p_T is approximately constant in the central region ($x_F < 0.1$) while a stronger decrease is observed for large x_F .

The dependence on x_F for $|x_F| < 0.1$ is weaker than for fully inclusive production of π^- in pp collisions [BAIL 87b]. This finding corresponds to the very gentle dependence of differential cross sections at high p_T on θ for $90^\circ \leq \theta \leq 20^\circ$ ($|y| \leq 1.7$) observed earlier [LLOY 80].

Another clear demonstration of the sensitivity of longitudinal distributions is provided by the measurement of (relative) rates for the production of positive and negative particles as function (of p_T and) of x_F , as shown in fig. 6.37 [BREA 87b]. While the fraction of π^- increases with x_F , probably due to the disappearance of K^- (similar to the K^-/π^- ratio as function of p_T , sect. 5.2.1), the relative yield of π^+ decreases, probably to a growing contribution from high p_T protons (see also sect. 5.5). From a comparison to theoretical predictions [BREA 87b] it is obvious that the data for positive particles cannot be reproduced by the scattering of quarks and gluons alone. The inclusion of hard diquark scattering (sect. 6.3.1), however, reconciles measurements for positive particles and theory.

The success of QCD inspired models is convincingly demonstrated by the quantitative agreement between data and calculations [BREA 84c, YEUN 87], keeping in mind that the inclusive cross section varies as function of p_T and $\sqrt{s}$ by many orders of magnitude. A calculation of π^0 and single photon production at $\sqrt{s} = 540$ GeV in a QCD framework [HORG 80] is shown in fig. 6.38(a). The model reproduces the data obtained at ISR energies and was then applied successfully (fig. 6.38(b), [JACO 87]) to predict inclusive cross sections from proton-antiproton collisions at collider energies.

TABLE 6.1

Relative contribution (in %) to the inclusive cross section for production of the meson (marked h) at 90° in the c.m.s. in pp collisions from fragmentation of the gluon, quarks, and antiquarks.
Calculations were done in a QCD framework [ABRA 85]

$P_{\text{beam}},$ GeV/c	$P_{\text{T}},$ GeV/c	x_{T}	h	g	u	d	s	$\bar{u}$	$\bar{d}$	$\bar{s}$
70	2.0	0.35	π^+	25	65	5.8	0.4	0.4	3.0	0.4
			π^-	35	18	42	0.6	3.3	0.8	0.6
			K^+	23	66	5.0	0.4	0.3	0.3	4.8
			K^-	54	19	8.3	11	5.6	1.1	0.8
70	3.5	0.61	π^+	6.7	88	3.2	0.1	0.1	1.7	0.1
			π^-	14	19	64	0.2	2.7	0.3	0.2
			K^+	6.4	88	2.7	0.1	0.1	0.1	2.7
			K^-	36	30	10	15	7.6	0.7	0.5
400	3.0	0.22	π^+	20	63	7.2	1.0	1.0	6.4	1.0
			π^-	26	17	46	1.2	7	1.5	1.2
			K^+	20	61	6.1	0.8	0.6	0.7	11
			K^-	39	16	8.7	22	11	2.0	1.7
400	5.0	0.36	π^+	3.6	88	4.3	0.3	0.3	3.5	0.3
			π^-	6.7	19	6.7	0.5	5.3	0.7	0.5
			K^+	3.9	86	3.6	0.3	0.2	0.2	5.6
			K^-	18	29	12	26	13	1.5	1.2

TABLE 6.2

Summary on flavor tagging

Hypotheses:		Experimental tools:				
Trigger: parton		Inclusive ratio	Charge ratio, trigger jet	K_S^0 , trigger jet	Correl. with spectator jet	Charge ratio, away jet
$\pi^+/K^+ \left\{ \begin{array}{l} : \bar{d} \\ \pi^- : \bar{u} \end{array} \right.$		No	No	Yes	No	?
$\pi^+/K^+ \left\{ \begin{array}{l} : g \\ \pi^- : g \end{array} \right.$		No	No	Yes	No	No
$\pi^+/K^+ \left\{ \begin{array}{l} : g \\ \pi^- : d \end{array} \right.$		No	No	Yes	Yes	No
$\pi^+/K^+ \left\{ \begin{array}{l} : u \\ \pi^- : g \end{array} \right.$		?	Yes	?	Yes	?
$\pi^+/K^+ \left\{ \begin{array}{l} : u \\ \pi^- : d \end{array} \right.$		Yes	Yes	Yes	Yes	Yes
$K^- \left\{ \begin{array}{l} : \bar{u} \\ \pi^+/K^+ : u \end{array} \right.$		Yes	No	Yes	Yes	No
$K^- \left\{ \begin{array}{l} : s \\ \pi^+/K^+ : u \end{array} \right.$		Yes	Yes	No	Yes	No
$K^- \left\{ \begin{array}{l} : u, d \\ \pi^+/K^+ : u \end{array} \right.$		Yes	Yes	Yes	Yes	No
$K^- \left\{ \begin{array}{l} : g \\ \pi^+/K^+ : u \end{array} \right.$		Yes	Yes	Yes	Yes	Yes
$p \left\{ \begin{array}{l} : u, d \\ \pi^+ : u \end{array} \right.$		No	?	-	No	?
$p \left\{ \begin{array}{l} : g \\ \pi^+ : u \end{array} \right.$		No	?	-	No	?
$p \left\{ \begin{array}{l} : uu, ud \\ \pi^+ : u \end{array} \right.$		Yes	?	-	Yes	?

Definition: No: Data qualitatively inconsistent with hypothesis.
Yes: Data qualitatively consistent with hypothesis.
?: Numerical calculations are needed for reasonable conclusion.

FIGURE CAPTIONS

Fig. 6.1 Schematic representation of high p_T particle production by single diffraction dissociation.

Fig. 6.2 Corrected densities of secondary particles with $0.3 \leq p_t \leq 1$ GeV/c (a) and $1 \leq p_t \leq 2$ GeV/c (b) at $\Delta\phi \geq 155^\circ$ relative to proton triggers with $p_T \geq 2$ GeV/c at production angles $\theta \sim 10^\circ$ compared to model calculations for single and double diffractive processes [BREA 88].

Fig. 6.3 Contribution from single diffraction dissociation to high p_T production as a function of the average trigger transverse momentum. The brackets display the linear sums: measurement $\pm$ stat. $\pm$ syst. error. The solid line corresponds to a simple exponential fit [BREA 88].

Fig. 6.4 Rapidity distribution of particles with $x_E > 0.5$ in the away jet ($\Delta\phi = 180^\circ \pm 40^\circ$). Distributions from events with a π^+ trigger with $4 \text{ GeV/c} < p_T < 6 \text{ GeV/c}$ are displayed along with "standard" assumptions for vector and scalar gluons [DRIJ 83].

Fig. 6.5 (a) Angular distributions as a function of $\cos\theta$ for fixed $\pi\pi$ -mass bins at $\sqrt{s} = 62.4 \text{ GeV}$ [ANGE 82].
(b) Relative contributions of valence quark-valence quark (qq-qq), gluon valence-quark (gq-gq) and gluon-gluon (gg-gg, gg-qq) scattering. Curves labelled (1) and (2) use gluon fragmentation functions proportional to $(1-z)^2$ and $(1-z)^3$, respectively [ANGE 82].

Fig. 6.6 The quantity A , a measure of the local slope of the square of the scattering matrix element, as function of $\cos\hat{\theta}$ ($= \cos\theta^*$) for various trigger particles with $p_T > 4 \text{ GeV/c}$ at $\theta \sim 45^\circ$. Predictions based upon (non)-abelian QCD and a scalar gluon theory are shown as well [BREA 89c].

Fig. 6.7 Dominating scattering mechanisms for the production of some high p_T hadrons.

FIGURE CAPTIONS (Cont'd)

Fig. 6.8 Density of positive and negative particles versus x_F in events with π^- triggers (T) relative to that for π^+ triggers (a) and K^+ triggers (b), and for the combination of π^+ and K^+ triggers (T) relative to K^- triggers (c) [BREA 84c, 85a]. Theoretical predictions based upon the JETS code are shown as well [YEUN 87].

Fig. 6.9 Ratios of positive spectator secondaries associated with proton and π^+ triggers at $\theta \approx 10^\circ$ with $p_T \geq 13$ GeV/c. The shaded regions at large x_F are predictions based upon deep inelastic neutrino scattering data [BREA 87c].

Fig. 6.10 Relative yields as function of x_F of protons in events with p and π^+ triggers; $x_F' = 1$ corresponds to the estimated kinematical limit in events with two sideways jets [SMIT 87].

Fig. 6.11 (a) Charge density ratio ρ^-/ρ^+ as function of x_F for pion and proton triggers. The shaded regions are predictions based upon deep inelastic neutrino and antineutrino scattering data ((ud) and (uu) target jets) [BREA 87c].

(b) Charge density ratio ρ^-/ρ^+ as function of x_F for proton triggers, compared to predictions based upon deep inelastic neutrino scattering for pure u-quark fragmentation and for a mixture of 33% d-quark and 67% u-quark fragmentation [BREA 87c].

Fig. 6.12 Ratio of densities of positive and negative away jet secondaries versus x_E for (a) back-to-back and (b) back-to-antiback configurations [BREA 84c, 85a].

Fig. 6.13 Ratio of densities of positive and negative secondaries as function of x_E ($= z$) for π^+p and pp collisions with a high p_T single photon [CASS 87].

FIGURE CAPTIONS (Cont'd)

Fig. 6.14 The difference between invariant cross sections for the reaction $\pi^- p \rightarrow \gamma X$ and $\pi^+ p \rightarrow \gamma X$ as a function of p_T for the full x_F range. The solid and dashed lines correspond to the predictions of second order QCD calculations using optimized scales with $\Lambda = 200$ MeV and $\Lambda = 400$ MeV, respectively [BONE 88a].

Fig. 6.15 Production of pairs of direct photons via annihilation, and higher order processes.

Fig. 6.16 Differential cross section for photons from 2-photon events [BONV 89].

Fig. 6.17 Dynamical (bremsstrahlung) contribution to k_T in hard parton-parton processes.

Fig. 6.18 The evolution of the intrinsic k_T distribution of protons with mean value of 0.45 GeV/c (dotted line) into an effective transverse momentum distribution with mean value of > 1 GeV/c (solid line) [YEUN 87].

Fig. 6.19 Corrected rapidity distribution for away jet fragments with $0.5 \leq x_E \leq 1.0$ and $\phi = (180 \pm 25)^\circ$ using π^+ triggers at $\phi \approx 0^\circ$ and parton model predictions. Spectator background is subtracted [BREA 89b, YEUN 87].

Fig. 6.20 Number of away jet tracks per event as function of x_E ($\phi = (180 \pm 25)^\circ$ and $-2 \leq y \leq 2$) and parton model predictions [BREA 89b, YEUN 87].

Fig. 6.21 Ratio d/p of inclusive cross sections for deuteron and protons as a function of p_T in pp interactions at 70 GeV/c [ABRA 87].

Fig. 6.22 Schematic representation of production of n-p and n- $\bar{p}$ pairs at high p_T .

FIGURE CAPTIONS (Cont'd)

Fig. 6.23 (a) Generic higher twist diagrams.

(b) Higher twist contribution to the Drell-Yan process.

(c) Suppression of forward spectator jets in a higher twist framework.

Fig. 6.24 (a) Example of a higher twist diagram for the direct production of a ρ^0 meson. (b) Example of a diagram for the indirect production of a ρ^0 meson, in the reaction πN , via quark scattering and subsequent fragmentation.

Fig. 6.25 Effective ($\pi^+\pi^-$) mass distribution for $1. < x_E < 2.$; solid line: fitted mass spectrum, dashed line: fitted values of the background function. The inserts at the bottom represents the reflection of the $\phi \rightarrow K^+K^-$ (dashed-dotted line) and $K^{*0} \rightarrow K^+\pi^-$ (dotted line) decay. (b) As (a) but for $x_E \geq 2$. [BENA 87a].

Fig. 6.26 Gluon-gluon fusion into an isosinglet $q\bar{q}$ pair giving the dominant contribution to the higher twist contribution of η . These diagrams are examples only.

Fig. 6.27 Example of a higher twist mechanism for the direct production of a glueball.

Fig. 6.28 The power α of the A dependence of the invariant cross section versus p_T for the production of hadrons by 400 GeV protons: (a) π^+ , (b) π^- , (c) K^+ , (d) K^- , (e) p, and (f) $\bar{p}$. Unless indicated, the errors are smaller than or equal to the size of the points [CRON 75].

Fig. 6.29 A subset of the data from fig. 6.28 and a comparison with predictions based upon multiple parton scattering [LEV 83].

Fig. 6.30 Ratio of the measured inclusive cross section for π^0 production (in $\alpha\alpha$, dd and pp collisions) to the pp fit. In addition to the statistical errors shown, there is a systematic error resulting from the $\pm 1 \sigma$ variation of the pp fit [ANGE 87].

FIGURE CAPTIONS (Cont'd)

Fig. 6.31 Ratio of normalized (by $A_1 \cdot A_2$) inclusive cross sections for pion production in $A_1 A_2$ ($A_1 = d, \alpha, A_2 = p, d, \alpha$) and pp collisions as function of p_T [FISC 88].

Fig. 6.32 Particle ratios R^+ (a) and R^- (b) for different c.m. polar angles θ as measured in pp, dd, and $\alpha\alpha$ interactions. The systematic errors are indicated in the figures. For comparison, pp data are also shown [BREA 86].

Fig. 6.33 (a) An early compilation of inclusive cross sections as function of p_T for π^- production in pp collisions at various energies. A fit of an empirical formula is superimposed [BEIE 78].

(b) A compilation of inclusive cross sections as function of p_T for pion production in the ISR to Tevatron energy range.

Fig. 6.34 Invariant cross section for π^0 production at $\sqrt{s} = 53$ GeV and a comparison with calculations in the framework of QCD [OWEN 78].

Fig. 6.35 A compilation of effective powers n obtained from the parametrization $E \frac{d\sigma}{dp} \sim p_T^{-n} f(x_T, \theta)$ [GEIS 84].

Fig. 6.36 Invariant cross section at $y = 0$ for inclusive production of jets and π^0 at $\sqrt{s} = 63$ GeV [AKES 82]; ISAJET predictions for jet production are included.

Fig. 6.37 Relative yields for positive and negative pions (π^+ /all positive and π^- /all negative) as function of Feynman x for fixed p_T at two ISR energies [BREA 87b].

FIGURE CAPTIONS (Cont'd)

- Fig. 6.38 (a) Inclusive cross sections predicted for π^0 production in pp collisions at 540 GeV assuming full scaling violation (lower solid curve), scaling violation of structure functions only (upper solid curve). Results from the standard parton model are given by the dashed curve. The prompt photon production is indicated as a dotted line [HORG 80].
- (b) Comparison of the predictions from fig. 6.38(a) with inclusive yields of π^0 and γ obtained by UA2; also inclusive jet cross sections are given [JACO 87].

Fig. 6.1

Fig. 6.2

Fig. 6.3

Fig. 6.4

Fig. 6.5(a)

Fig. 6.5(b)

Fig. 6.6

Fig. 6.7

Fig. 6.8

Fig. 6.9

Fig. 6.10

Fig. 6.11

Data Analysis

Fig. 6.12

Fig. 6.13

Fig. 6.14

Fig. 6.15

Fig. 6.16

Fig. 6.17

Fig. 6.18

Fig. 6.19

Fig. 6.20

Fig. 6.21

Fig. 6.22

Fig. 6.23

Fig. 6.24

Fig. 6.25

Fig. 6.26

Fig. 6.27

Fig. 6.28

Fig. 6.29

Fig. 6.30

Fig. 6.31

Fig. 6.32

Fig. 6.33(a)

Fig. 6.33(b)

Fig. 6.34

Fig. 6.35

Fig. 6.36

Fig. 6.37

Fig. 6.38(a)

Fig. 6.38(b)

7. SUMMARY AND OUTLOOK

The discovery at the ISR of an unexpectedly high yield of particles with large transverse momenta and the subsequent analysis of events triggered by the presence of such particles led, after a long venture, to a comprehensive understanding of parton-parton processes, in the framework of QCD, at a level of detail probably not anticipated. The more recent collider data have provided an unchallenged basis for such an interpretation.

Basically one owes this insight to the experimental observation that the trigger particle carries typically more than about 70% of the momentum of the scattered parton; hence, as the leading fragment, it plays a unique role in the hadronization process: its quantum numbers are strongly correlated to those of its parent parton. Therefore, triggering on large p_T flavour identified particles allows to tag parton species. This method, originally developed in the field of large p_T physics, is now a generally accepted tool in fragmentation studies [SAXO 85]. Furthermore, the unravelling of parton processes was based substantially on the observation of local quantum number correlations in the trigger jet as well as of long range quantum number correlations between jets.

The combination of these two powerful tools resulted in very restrictive constraints for the interpretation of the measurements in terms of the dynamical features of the basic scattering process and of the type of all participating partons. In many cases the conclusions drawn are implied directly by the data and are independent of any sophisticated Monte-Carlo studies which are, of course, welcome as a further helpful back-up.

The more important achievements are summarized in the following:

- (1) First, the parton composition of hadrons was probed. In hard proton-proton scattering processes u-, d-valence quarks, gluons and diquarks were identified as the dominating sources of the observed large p_T ($x_T \approx 0.1 \div 0.2$) particles. Triggering π^+ and K^+ mesons originate from the fragmentation of scattered u-quarks, while large p_T π^- mesons come very often from d-quark fragmentation. Large p_T K^- mesons are dominantly gluon fragments. This observation is not unexpected since incoming protons and K^- mesons have no valence quark in common. The proof of the latter claim follows from (a) the study of charge and strangeness correlations in the

trigger-jet, which characterizes the parent parton as flavour neutral, and (b) from the different dependences of the inclusive cross sections for K^- and π^- mesons on p_T , which demonstrates that the structure function of the parton yielding large p_T K^- mesons is softer than that of the π^- meson parent.

Evidence for gluons inside hadrons was corroborated by the detection of direct photons and of neutral jets recoiling against the triggering π^+ and K^+ mesons.

Studies of large p_T proton production show that the dominant source is the scattering of a pair of valence quarks, i.e. of a "diquark" scattering off valence quarks or gluons. Again this conclusion follows from extensive correlation studies and from the measurement of the proton production cross section in a broad kinematical range.

Not only the quantum numbers of the proton constituents participating in large p_T hadron production have been determined but also information on the parton wave function was derived. Already during the early stage of the experimental program it was shown that the partons must have a transverse momentum k_t before they scatter. The observed shape of the inclusive transverse momentum spectrum, the particle density of the away jet and the observed transverse momentum spectrum of the spectator fragments cannot be explained otherwise. Up to now it has not been possible to separate uniquely the primordial transverse motion of partons from a gluon radiation contribution to k_t , though there is compelling evidence that both components exist.

- (2) As a second major achievement the measurement of the angular dependence of parton-parton scattering cross sections has to be mentioned. It was shown that exchange of scalar gluons can be excluded as the basic scattering mechanism, while the data are compatible with spin-1 exchange.

As a result of gluon self-coupling qq and qg scattering contribute with comparable strength to large p_T processes. It turned out to be feasible to separate experimentally these scattering processes. One finds that the relative cross sections are compatible with what is expected in the framework of QCD. Hence one has obtained a consistent result both for normalization and angular dependence of parton processes.

The recent measurement of the angular dependence for backward scattering provides probably the most convincing experimental evidence for the non-Abelian nature of QCD so far.

- (3) The third important result concerns fragmentation of quarks and gluons. Hadroproduction of large p_T particles offers the unique chance to investigate with high statistics parton fragmentation at large values of the fractional momentum z . The fragmenting parton is far off shell, hence the influence of a different scale on SU(3) breaking can be studied in contrast to the fragmentation at low z values accessible so far in e^+e^- annihilation and lepton scattering experiments. Four relevant results should be mentioned: The observed rate of relative π^0 and η production is in accordance with expectations based on SU(3) mixing of pseudoscalar mesons. Secondly, at large z , the SU(3) breaking of strangeness production appears to be smaller than at small z . The mean relative vector and scalar meson production rates, on the other hand, are compatible with the ones at low z and deviate from spin weights. First evidence, independent of Monte-Carlo simulations, shows finally that gluon fragmentation is softer than quark fragmentation.

The results discussed prove the level of maturity reached in the understanding on the parton basis of large p_T physics. To a surprisingly large extent all observations can be described quantitatively in a unique and coherent way by scattering of valence quarks and gluons in the framework of QCD, a success which is due to the dedicated and careful work of a few experimental teams over more than a decade. The ISR energy range turned out to be especially suited for these investigations since the cross sections were large enough to allow for comfortable production rates. Moreover, the energy available to the system of interacting partons is high enough to allow large angle jets to evolve such that their separation from the bulk of secondary particles was possible. On the other hand, the available energy still allowed to identify single particles in jets. Ample evidence has been presented in this review for the importance of this experimental fact. At moderate energies multigluon effects are small and therefore do not obscure details of the 4-jet structure. Higher twist effects are usually negligible, unless enhanced by special triggers. Generally, these detailed results are not superseded by the advent of the

SppS collider data which have given unambiguous support to the concept of hard parton-parton scattering, but are typically very sensitive to distinctly complementary aspects of deep inelastic phenomena.

As usual in physics, once a satisfying level of understanding is reached, new questions evolve and stimulate new investigations. It is tempting to collect some of those problems which can now be attacked with a fair chance to find interesting answers.

- Measurements of the angular dependence of parton-parton scattering should be improved for the backward region. This necessitates a powerful separation of subprocesses.
- Determination of hadron (π , K, p) structure functions via tagged quarks and gluons in hard hadronic interactions. Emerging data can be compared in some cases to the structure functions measured in deep inelastic lepton scattering and Drell-Yan processes. It provides a test of recently developed ideas on parton evolution.
- A better separation of primordial and dynamical contributions to the transverse momenta of partons should be attempted.
- Inclusive fragmentation should be studied of quarks and gluons which are recoiling against flavour identified trigger particles and are not subject to kinematic constraints. The results presented here provide already evidence that gluon fragmentation is softer than quark fragmentation. Furthermore, the flavour structure of quark and gluon jets could be studied separately.
- Studies of the production of tensor mesons, hyperons, baryon resonances and glueballs at large z are possible. Important information on all stages of hadronization can be derived this way; it is hardly accessible in other types of experiments.
- The study of associated spectator fragmentation in events with a tagged gluon jet at large p_T allows to determine in a model independent way the fragmentation of quark systems in a color octet state. A comparison with gluon jet fragmentation gives important clues on the relevance of color and flavor quantum numbers in fragmentation processes.

- The rôle of polarization in (multi-) quark and gluon fragmentation has not been addressed yet.
- Distinct string configurations can easily be verified in different kinematical configurations. Their existence and fragmentation properties can be studied.
- Quantum number structures of direct photon events should be measured precisely.
- Higher twist processes, production of deuterons at high transverse momentum, and nuclear effects also deserve some more attention.

From this probably incomplete list one may conclude that there is still a potential to answer important physical questions in the ISR energy range $\sqrt{s} \sim 20 \div 60$ GeV now accessible at the Tevatron. The availability of different types of hadron beams for fixed target physics at the Tevatron is stimulating. One may also hope that some of the ideas developed at these energies can be adapted to analyses of collider data.

In summary, from the study of large p_T particle production, especially at the ISR, a consistent picture of hadronic hard scattering processes was developed revealing the applicability of QCD to this specific type of hadronic reactions. It is only here, so far, that the strong interactions between quasi-free quarks and gluons occur on equal footing in the Born term and can thus be probed rather easily. Having established a sound basis, these analyses open the door for a future physics program, which has the potential of interesting new insights into the constituent structure of hadrons and into constituent interactions at short and long distances.

Acknowledgements

This report was encouraged by M. Jacob, and is based upon numerous contributions to this field from experiments at Serpukhov, Fermilab and CERN. All of them helped to shape this presentation, even if they are not explicitly mentioned. This type of physics owes very much to the excellent performance of the CERN ISR and its staff, and to the enthusiastic physicists working there. We have benefitted from their experience through competition and discussions. Most of our thanks are, however, due to our colleagues of the Ames-Bologna-CERN-Dortmund-Heidelberg-Warsaw Collaboration (and its predecessor), who devoted their knowledge over long years to this subject, and who were pleasant colleagues to work with. In particular, we would like to acknowledge H.G. Fischer whose major role in all aspects of the collaboration's physics program cannot be sufficiently appreciated here. Last not least, we enjoyed the help of E. Delucinge, who prepared the manuscript with outstanding dedication and professional skill.

A. Putzer and D. Wegener were supported by a grant from the Bundesministerium für Forschung und Technologie of the Federal Republic of Germany.

APPENDIX 1

Most of the variables used in this report are defined here for easier reference. Only very specific kinematic quantities are explained in the text. The hadron-hadron centre of mass frame is assumed unless stated otherwise; the centre of mass energy is $\sqrt{s}$. Four-momenta are given by energy E and momentum $\vec{p}$; if needed, indices refer to beam, trigger particle, associated secondary particle, jets, or partons.

A1. KINEMATIC VARIABLES FOR HADRONS RELATIVE TO THE BEAM DIRECTION

$P_T (p_t)$	: Transverse momentum of trigger particle (of associated secondaries) or for inclusive spectra.
x_T	: $x_T = 2 P_T / \sqrt{s}$.
x_E	: $x_E = P_t / P_T$.
P_L	: Longitudinal momentum, i.e. parallel to the beam.
x_F	: $x_F = 2 P_L / \sqrt{s}$.
y	: rapidity.
P_{out}	: $P_{out} = \vec{p} \cdot \vec{v}$; $\vec{v} = (\vec{p}_{beam} \times \vec{p}_{trig}) / \vec{p}_{beam} \times \vec{p}_{trig} $ defines the normal to the "scattering plane".
ϕ, θ	: Azimuthal, polar angle.

A2. KINEMATIC VARIABLES FOR HADRONS RELATIVE TO A JET AXIS

$P_{//}$	: Component of $\vec{p}$ parallel to jet axis.
$x_{//}$	: $x_{//} = P_{//} / \vec{p}_{jet} $.
$\vec{q}_t$	: Components of $\vec{p}$ perpendicular to jet axis.
z	: Fragmentation variable.
z_f	: $z_f = P_{//} / \vec{p}_{trig} \approx P_t / P_T$.
$\langle z_{TR} \rangle$	: Average value of z for trigger particle.
u	: Component of $\vec{p}$ in scattering plane and perpendicular to $\vec{p}_{trig}$.

$\Delta\theta$: Angle between particle and jet axis (corresponds to α in e^+e^- collisions).

A3. KINEMATIC VARIABLES OF PARTONS

p_1 : Longitudinal momentum (component parallel to $\vec{p}_{\text{beam}}$).

x : Bjorken- x .

$\vec{k}_t$: transverse momentum of incoming partons relative to $\vec{p}_{\text{beam}}$.

$\hat{s}, \hat{t}, \hat{u}$: Mandelstam variables.

$\hat{\phi}, \hat{\theta}$: Azimuthal, polar angle in parton-parton c.m.s.

Q^2 : Square of 4-momentum transfer.

APPENDIX 2

The squared matrix elements $|\Gamma|^2$ for parton-parton scattering in lowest order are collected here. The color factors are enclosed in brackets "<...>": these factors should be set to 1 for an uncolored model. The variables $\hat{s}$, $\hat{t}$, $\hat{u}$ correspond to the Mandelstam invariants for the subsystem. The relation $\hat{s} + \hat{t} + \hat{u} = 0$ is implicitly used, so formulae published elsewhere might differ. The parton labels q , q' , $\bar{q}$, $\bar{q}'$, g and S correspond to (anti)quarks, gluons and scalar diquarks, respectively. When color is considered, gluons are octets, quarks are triplets and diquarks are anti-triplets. The formulae are given for both scalar and vector gluons.

Vector gluons:

$$\begin{aligned}
 q_i q_j \rightarrow q_i q_j & \quad \left\langle \frac{2}{9} \right\rangle 2 \frac{\hat{s}^2 + \hat{u}^2}{\hat{t}^2} + \delta_{ij} \left[\left\langle \frac{2}{9} \right\rangle 2 \frac{\hat{t}^2 + \hat{s}^2}{\hat{u}^2} + \left\langle -\frac{2}{27} \right\rangle \frac{4\hat{s}^2}{\hat{u}\hat{t}} \right] \\
 q_i \bar{q}_j \rightarrow q_i \bar{q}_j & \quad \left\langle \frac{2}{9} \right\rangle 2 \frac{\hat{t}^2 + \hat{u}^2}{\hat{s}^2} + \delta_{ij} \left[\left\langle \frac{2}{9} \right\rangle 2 \frac{\hat{u}^2 + \hat{s}^2}{\hat{t}^2} + \left\langle -\frac{2}{27} \right\rangle \frac{4\hat{u}^2}{\hat{s}\hat{t}} \right] \\
 gg \rightarrow q\bar{q} & \quad \left\langle \frac{1}{12} \right\rangle 2 \frac{\hat{u}^2 + \hat{t}^2}{\hat{u}\hat{t}} - \left\langle \frac{3}{16} \right\rangle 4 \left(1 - \frac{\hat{u}\hat{t}}{\hat{s}^2} \right) + \left\langle \frac{3}{32} \right\rangle > 4 \\
 q\bar{q} \rightarrow gg & \quad \left\langle \frac{16}{27} \right\rangle 2 \frac{\hat{u}^2 + \hat{t}^2}{\hat{u}\hat{t}} - \left\langle \frac{4}{3} \right\rangle > 4 \left(1 - \frac{\hat{u}\hat{t}}{\hat{s}^2} \right) + \left\langle \frac{2}{3} \right\rangle > 4 \\
 gq \rightarrow gq & \quad -\left\langle \frac{2}{9} \right\rangle > 2 \frac{\hat{u}^2 + \hat{s}^2}{\hat{u}\hat{s}} + (\left\langle 2 \right\rangle - 1) \frac{\hat{s}^2 + \hat{u}^2}{\hat{t}^2} \\
 gg \rightarrow gg & \quad (\left\langle 2 \right\rangle - 1) \frac{9}{2} \left(3 - \frac{\hat{u}\hat{t}}{\hat{s}^2} - \frac{\hat{t}\hat{s}}{\hat{u}^2} - \frac{\hat{s}\hat{u}}{\hat{t}^2} \right) \\
 qS \rightarrow qS & \quad -\left\langle \frac{2}{9} \right\rangle \frac{4\hat{u}\hat{s}}{\hat{t}^2} \\
 SS \rightarrow SS & \quad \left\langle \frac{2}{9} \right\rangle \left[\frac{(\hat{s} - \hat{u})^2}{\hat{t}^2} + \frac{(\hat{s} - \hat{t})^2}{\hat{u}^2} \right] - \left\langle \frac{2}{27} \right\rangle 2 \left(1 + \frac{2\hat{s}^2}{\hat{t}\hat{u}} \right) \\
 gS \rightarrow gS & \quad \left\langle \frac{1}{2} \right\rangle \frac{(\hat{s} - \hat{u})^2}{\hat{t}^2} + \left\langle \frac{7}{18} \right\rangle \\
 & \quad + \frac{4\hat{t}^2}{\hat{s}^2} \left(2\left\langle \frac{2}{9} \right\rangle - \left\langle \frac{7}{36} \right\rangle - \left\langle \frac{1}{4} \right\rangle \right) + \frac{2\hat{u}(2\hat{s} - \hat{t})}{\hat{s}^2} \left(\left\langle \frac{1}{2} \right\rangle - 2\left\langle \frac{1}{4} \right\rangle \right)
 \end{aligned}$$

Scalar gluons:

$$q_i q_j \rightarrow q_i q_j \quad \langle \frac{2}{9} \rangle + \delta_{ij} \{ \langle \frac{2}{9} \rangle - \langle -\frac{2}{27} \rangle \}$$

$$q_i \bar{q}_j \rightarrow q_j \bar{q}_i \quad \langle \frac{2}{9} \rangle + \delta_{ij} \{ \langle \frac{2}{9} \rangle - \langle -\frac{2}{27} \rangle \}$$

$$g g \rightarrow q \bar{q} \quad \langle \frac{1}{12} \rangle^2 \frac{\hat{t}^2 + \hat{u}^2}{\hat{t}\hat{u}} - \langle \frac{1}{96} \rangle$$

$$q \bar{q} \rightarrow g g \quad \langle \frac{16}{27} \rangle \frac{\hat{t}^2 + \hat{u}^2}{2\hat{t}\hat{u}} - \langle -\frac{2}{27} \rangle$$

$$g q \rightarrow g q \quad -\langle \frac{2}{9} \rangle \frac{\hat{u}^2 + \hat{s}^2}{\hat{u}\hat{s}} + \langle -\frac{1}{36} \rangle^2$$

$$g g \rightarrow g g \quad 0$$

In figs A2.1 and A2.2 squared matrix elements $|\Gamma|^2$ are displayed as function of the scattering angle θ for some dominating processes (in pp collisions) in the framework of standard QCD. Fig. A2.3 shows what would happen in the case of no color.

FIGURE CAPTIONS

- Fig. A2.1 (a) The dependence of the squared matrix elements, $|\Gamma|^2$, on the parton scattering angle $\hat{\theta}$ for interactions of quarks (q, q'), vector gluons (g), antiquarks ($\bar{q}, \bar{q}'$) in the framework of lowest order QCD.
- (b) See (a) but involving scalar diquarks (S).

Fig. A2.2 See fig. A.2.1(a) but for colorless partons (abelian theory).

Fig. A2.1(a)

Fig. A2.1(b)

Fig. A2.2

REFERENCES

- [ABE 88] F. Abe et al., Phys. Rev. Lett. 61 (1988) 1819.
- [ABRA 82] H. Abramowicz et al., Z. Phys. C12 (1982) 289.
H. Abramowicz et al., Z. Phys. C13 (1982) 199.
- [ABRA 83] H. Abramowicz et al., Z. Phys. C17 (1983) 283.
- [ABRA 85] V.V. Abramov et al., Sov. J. Nucl. Phys. 41 (1985) 445.
- [ABRA 87] V.V. Abramov et al., Sov. J. Nucl. Phys. 45 (1987) 845.
- [ADAM 89] M.I. Adamovich et al., CERN/EP 89-123.
- [ADER 85] M. Aderholz et al., Z. Phys. C27 (1985) 43.
- [AIHA 84] H. Aihara et al., Phys. Rev. Lett. 53 (1984) 2199.
- [AIHA 86] H. Aihara et al., Phys. Rev. Lett. 57 (1986) 945.
- [AKER 77] C.W. Akerloff et al., Phys. Rev. Lett. 39 (1977) 861.
- [AKES 82] T. Akesson et al., Phys. Lett. 118B (1982) 185.
- [AKES 84] T. Akesson et al., Nucl. Phys. B246 (1984) 408.
- [AKES 85] T. Akesson et al., Phys. Lett. 158B (1985) 282.
- [AKES 86a] T. Akesson et al., Z. Phys. C32 (1986) 317.
- [AKES 86b] T. Akesson et al., Z. Phys. C32 (1986) 491.
- [AKES 87a] T. Akesson et al., Z. Phys. C34 (1987) 163.
- [AKES 87b] T. Akesson et al., Z. Phys. C34 (1987) 293.
- [ALBR 76] M.G. Albrow et al., Nucl. Phys. B108 (1976) 1.
- [ALBR 78a] M.G. Albrow et al., Nucl. Phys. B135 (1978) 461.
- [ALBR 78b] M.G. Albrow et al., Nucl. Phys. B145 (1978) 305.
- [ALBR 85] H. Albrecht et al., Phys. Lett. 157B (1985) 326.
- [ALBR 87] H. Albrecht et al., Phys. Lett. 199B (1987) 457.
- [ALLE 83] P. Allen et al., Nucl. Phys. B214 (1983) 369.
- [ALON 80] J.L. Alonso et al., Z. Phys. C6 (1980) 21.
- [ALPE 73] B. Alper et al., Phys. Lett. 44B (1973) 521.
- [ALTA 77] G. Altarelli and G. Parisi, Nucl. Phys. B126 (1977) 298.
- [ALTA 82] G. Altarelli, Phys. Rep. 81 (1982) 1.
- [ALTH 84] M. Althoff et al., Z. Phys. C22 (1984) 307.

REFERENCES (Cont'd)

- [ANDE 83] B. Andersson, G. Gustafson, G. Ingelman and T. Sjöstrand, Physics Report 97 (1983) 31.
- [ANGE 80] A.L.S. Angelis et al., Phys. Lett. 97B (1980) 163.
- [ANGE 81] A.L.S. Angelis et al., Phys. Lett. 105B (1981) 233.
- [ANGE 82] A.L.S. Angelis et al., Nucl. Phys. B209 (1982) 284.
- [ANGE 87] A.L.S. Angelis et al., Phys. Lett. 185B (1987) 213.
- [ANGE 88] A.L.S. Angelis et al., Nucl. Phys. B303 (1988) 569.
- [ANSE 83] M. Anselmino, P. Kroll and E. Leander, Z. Phys. C18 (1983) 307.
- [ANSE 87] M. Anselmino et al., Z. Phys. C36 (1987) 89.
- [ANTI 87] J. Antille et al., Phys. Lett. 194B (1987) 568.
- [ANTR 79] D. Antreasyan et al., Phys. Rev. D19 (1979) 764.
- [APPE 75] J.A. Appel et al., Phys. Rev. Lett. 35 (1975) 9.
- [ARNE 85] M. Arneodo et al., Phys. Lett. 150B (1985) 458.
- [ARNE 86a] M. Arneodo et al., Z. Phys. C33 (1986) 167.
- [ARNE 86b] M. Arneodo et al., Nucl. Phys. B264 (1986) 739.
- [ARNE 89] M. Arneodo et al., Nucl. Phys. B321 (1989) 541.
- [ARNI 82a] G. Arnison et al., Phys Lett. 118B (1982) 167.
- [ARNI 82b] G. Arnison et al., Phys. Lett. 118B (1982) 173.
- [ARNI 85] G. Arnison et al., Phys. Lett. 158B (1985) 494.
- [ARNI 86] G. Arnison et al., Nucl. Phys. B276 (1986) 253.
- [AUBE 87] J.J. Aubert et al., Nucl. Phys. B293 (1987) 740.
- [AURE 85] P. Aurenche et al., Z. Phys. C29 (1985) 459.
- [AURE 87] P. Aurenche et al., Nucl. Phys. B286 (1987) 509.
- [AURE 88] P. Aurenche et al., Nucl. Phys. B297 (1988) 661.
- [BADI 86] J. Badier et al., Z. Phys. C31 (1986) 341.
- [BAGN 84a] P. Bagnaia et al., Phys. Lett. 144B (1984) 283.
- [BAGN 84b] P. Bagnaia et al., Phys. Lett. 144B (1984) 291.
- [BAIL 87a] J.L. Bailly et al., Phys. Lett. 195B (1987) 609.
- [BAIL 87b] J.L. Bailly et al., Z. Phys. C35 (1987) 309.
- [BANN 73] M. Banner et al., Phys. Lett. 44B (1973) 537.

REFERENCES (Cont'd)

- [BART 83] D.S. Barton et al., Phys. Rev. D27 (1983) 2580.
- [BEAV 81] D. Beavis and B.R. Desai, Phys. Rev. D23 (1981) 1967.
- [BEIE 78] E.W. Beier et al., Phys. Rev. D18 (1978) 2235.
- [BENA 87a] M. Benayoun et al., Phys. Lett. 183B (1987) 412.
- [BENA 87b] M. Benayoun et al., Phys. Lett. 192B (1987) 447.
- [BENA 87c] M. Benayoun et al., Phys. Lett. 198B (1987) 281.
- [BENA 87d] M. Benayoun et al., Nucl. Phys. B282 (1987) 653.
- [BENG 84] H.U. Bengtsson and G. Ingelmann, CERN-TH 3820 (1984).
- [BENG 85] H.U. Bengtsson and G. Ingelmann, Comp. Phys. Comm. 34 (1985) 251.
- [BERE 81] F.A. Berends et al., Phys. Lett. 103B (1981) 124.
- [BERG 80] Ch. Berger et al, Phys. Lett. 97B (1980) 459.
- [BERG 81] Ch. Berger et al, Z. Phys. C8 (1981) 101.
- [BERM 70] S.M. Berman and M. Jacob, Phys. Rev. Lett. 25 (1970) 1683.
- [BERM 71] S.M. Berman, J.D. Bjorken and J.B. Kogut, Phys. Rev. D4 (1971) 3388.
- [BERN 85] D. Bernard et al., Phys. Lett. 166B (1985) 459.
- [BJOR 69a] J.D. Bjorken, Phys. Rev. 179 (1969) 1547.
- [BJOR 69b] J.D. Bjorken and E.A. Paschos, Phys. Rev. 185 (1969) 1975.
- [BJOR 73] J.D. Bjorken; Phys. Rev. D8 (1973) 4098.
- [BJOR 76] J.D. Bjorken, SLAC-PUB-1756 (1976).
- [BLAN 74] R. Blankenbecler and S.J. Brodsky, Phys. Rev. D10 (1974) 2973.
- [BLAN 77] R. Blankenbecler et al., Phys. Rev. D18 (1977) 900.
- [BONE 87] A. Bonesini et al., Z. Phys. C37 (1987) 39.
- [BONE 88a] A. Bonesini et al., Z. Phys. C37 (1988) 535.
- [BONE 88b] A. Bonesini et al., Z. Phys. C38 (1988) 371.
- [BONE 89] A. Bonesini et al., Z. Phys. C42 (1989) 527.
- [BONI 88] R. Bonino et al., Phys. Lett. 211B (1988) 239.
- [BONV 89] E. Bonvin et al., Z. Phys. C41 (1989) 591.
- [BRAN 80] R. Brandelik et al, Phys. Lett. 97B (1980) 453.
- [BRAU 89] W. Braunschweig et al., DESY 89-032, and refs therein.

REFERENCES (Cont'd)

- [BREA 84a] A. Breakstone et al., Z. Phys. C23 (1984) 1;
- [BREA 84b] A. Breakstone et al., Z. Phys. C23 (1984) 9.
- [BREA 84c] A. Breakstone et al., Z. Phys. C25 (1984) 21.
- [BREA 84d] A. Breakstone et al., Phys. Rev. D30 (1984) 528.
- [BREA 84e] A. Breakstone et al., Phys. Lett. 135B (1984) 505.
- [BREA 84f] A. Breakstone et al., Phys. Lett. 135B (1984) 510.
- [BREA 84g] A. Breakstone et al., Phys. Lett. 147B (1984) 237.
- [BREA 84h] A. Breakstone et al., Contribution to the XXIIth Int. Conf. on High Energy Phys., Leipzig (1984).
- [BREA 84i] A. Breakstone et al., Contribution to the XXIIth Int. Conf. on High Energy Phys., Leipzig (1984).

- [BREA 85a] A. Breakstone et al., Z. Phys. C27 (1985) 205.
- [BREA 85b] A. Breakstone et al., Z. Phys. C28 (1985) 335.

- [BREA 86] A. Breakstone et al., Z. Phys. C30 (1986) 507.

- [BREA 87a] A. Breakstone et al., Z. Phys. C33 (1987) 475.
- [BREA 87b] A. Breakstone et al., Z. Phys. C35 (1987) 157.
- [BREA 87c] A. Breakstone et al., Z. Phys. C36 (1987) 567.
- [BREA 87d] A. Breakstone et al., unpublished.

- [BREA 88] A. Breakstone et al., Z. Phys. C40 (1988) 207.

- [BREA 89a] A. Breakstone et al., Z. Phys. C43 (1989) 185.
- [BREA 89b] A. Breakstone et al., Contribution to the XXIV Conf. on High Energy Physics, Munich (1988) and to be published.
- [BREA 89c] A. Breakstone et al., to be published.

- [BREI 69] M. Breidenbach et al., Phys. Rev. Lett. 23 (1969) 935.

- [BUCH 82] C.D. Buchanan, Proc. XVIIth Rencontre de Moriond, Les Arcs, France, vol. II (1982) 139.

- [BUCH 87] C.D. Buchanan and S.B. Chun, Phys. Rev. Lett. 59 (1987) 1997.

- [BURA 77] A.J. Buras and K.J.F. Gaemers, Nucl. Phys. B132 (1978) 249.

- [BÜSS 73] F.W. Büsser et al., Phys. Lett. 46B (1973) 471.

- [BÜSS 75] F.W. Büsser et al., Phys. Lett. 55B (1975) 232.

- [BÜSS 76] F.W. Büsser et al., Nucl. Phys. B106 (1976) 1.

- [CAMI 88] L. Camilleri, Proceedings of the XIXth International Symposium on Multiparticle Dynamics, Arles, France (1988).

- [CASS 87] A.J. Cass, Ph.D. Thesis, Univ. of Liverpool, September 1987.

- [CAVA 89] V. Cavasinni, CERN/EP 89-47, to appear in Riv. Nuovo Cim.

- [COHE 83] G. Cohen-Tannoudji et al., Phys. Rev. D28 (1983) 1628.

- [COMB 77] B.L. Combridge et al., Phys. Lett. 70B (1977) 234.

- [COMB 83] B.L. Combridge and C.J. Maxwell RL-83-095 (1983).

REFERENCES (Cont'd)

- [CONW 88] I. Conway and W. Louis, Proc. XXIIIrd Rencontre de Moriond, Les Arcs, France (1988) 369.
- [CRON 75] J.W. Cronin et al., Phys. Rev. D11 (1975) 3105.
- [CUTL 77] R. Cutler and Sivers, Phys. Rev. D16 (1977) 679.
- [CUTL 78] R. Cutler and Sivers, Phys. Rev. D17 (1978) 196.
- [DARR 76] P. Darriulat et al., Nucl. Phys. B107 (1976) 429.
- [DARR 80] P. Darriulat, Ann. Rev. Nuc. Sci. 30 (1980) 159.
- [DEDE 75] H. Deden et al., Nucl. Phys. B85 (1975) 269.
- [DELL 75] M. Della Negra et al., Phys. Lett. 59B (1975) 401.
- [DELL 76a] M. Della Negra et al., Nucl. Phys. B104 (1976) 365.
- [DELL 76b] M. Della Negra et al., CERN/EP/PHYS 76-35, Contribution to 18th Int. Conf. on High Energy Physics, Tbilissi, 1976.
- [DELL 77] M. Della Negra et al., Nucl. Phys. B127 (1977) 1.
- [DEMA 87a] C. De Marzo et al., Phys. Rev. D36 (1987) 8.
- [DEMA 87b] C. De Marzo et al., Phys. Rev. D36 (1987) 16.
- [DERA 78] I. Derado et al., Nucl. Phys. B143 (1978) 40.
- [DERR 85] M. Derrick et al., Phys. Lett. 158B (1985) 519.
- [DIAK 80] M. Diakonou et al., Phys. Lett. 89B (1980) 432.
- [DILE 85] L. DiLella, Ann. Rev. Nucl. Part. Sci. 35 (1985) 107.
- [DONA 76] G.J. Donaldson et al., Phys. Rev. Lett. 36 (1976) 1110.
- [DONA 78a] G.J. Donaldson et al., Phys. Rev. Lett. 40 (1978) 684.
- [DONA 78b] G.J. Donaldson et al., Phys. Lett. 73B (1978) 375.
- [DONA 80] G.J. Donaldson et al., Phys. Rev. D21 (1980) 828.
- [DORF 81] J. Dorfan et al, SLAC-PUB-2813 (1981).
- [DORO 84] K. Doroba et al., Z. Phys. C24 (1984) 31.
- [DOSCH 89] H.G. Dosch et al., Z. Phys. C42 (1989) 167.
- [DRIJ 79a] D. Drijard et al., Nucl. Phys. B155 (1979) 269.
- [DRIJ 79b] D. Drijard et al., Nucl. Phys. B156 (1979) 309.
- [DRIJ 80] D. Drijard et al., Nucl. Phys. B166 (1980) 233.
- [DRIJ 83] D. Drijard et al., Phys. Lett. 121B (1983) 433.
- [DRIJ 84] D. Drijard et al., Nucl. Instr. and Meth. 225 (1984) 367.

REFERENCES (Cont'd)

- [DUKE 84] D.W. Duke and J.F. Owens, Phys. Rev. D30 (1984) 49.
- [EFRE 87] A.V. Efremov and V.T. Kim, JINR, E2-87-74, Dubna (1987).
- [EKEL 84] S. Ekelin and S. Fredriksson, Phys. Lett. 149B (1984) 509.
- [ELLI 74] S.D. Ellis and M.B. Kislinger; Phys. Rev. D9 (1974) 2027.
- [ELLI 77] S.D. Ellis and R. Stroynowski, Rev. Mod. Phys. 49 (1977) 753.
- [ELLI 85] R.K. Ellis and J.C. Sexton, Fermilab-PUB-85/152-T (1985).
- [FERB 84] T. Ferbel and W.R. Molzon, Rev.Mod.Phys. 56 (1984) 181.
- [FERB 88] T. Ferbel, Proc. of Physics in Collision, October 1988, Capri, Italy.
- [FEYN 69a] R.P. Feynman, 3rd Int. Conf. on High Energy Collisions, Stony Brook, 1969.
- [FEYN 69b] R.P. Feynman, Phys. Rev. Lett. 23 (1969) 1415.
- [FEYN 72] R.P. Feynman, Photon-Hadron Interactions (Benjamin, Reading, MA) (1972);
- [FEYN 77] R.P. Feynman, R.D. Field and G.C. Fox, Nucl. Phys. B128 (1977) 1.
- [FEYN 78] R.P. Feynman, R.D. Field and G.C. Fox, Phys. Rev. D18 (1978) 3320.
- [FF] Stands for refs [FIEL 77, FIEL 78a, FEYN 77,78].
- [FIEL 77] R.D. Field and R.P. Feynman, Phys. Rev. D15 (1977) 2590.
- [FIEL 78a] R.D. Field and R.P. Feynman, Nucl. Phys. B136 (1978) 1.
- [FIEL 78b] R.D. Field, Phys. Rev. Lett. 40 (1978) 997.
- [FINK 72] J. Finkelstein and R.D. Peccei, Phys. Rev. D6 (1972) 2606.
- [FISC 81a] H.G. Fischer, Proc. 3rd Warsaw Symposium on Elem. Part. Phys., Nukleonika 26 (1981) 201.
- [FISC 81b] H.G. Fischer, Proc. Int. Conf. on High Energy Phys., Lisbon (1981) 297.
- [FISC 88] H.G. Fischer, Z. Phys. C38 (1988) 105, and refs therein.
- [FISK 78] R.J. Fisk et al., Phys. Rev. Lett. 40 (1978) 984.
- [FOX 76] G.C. Fox, Talk Brookhaven APS meeting, October 1986.
- [FOX 77] G.C. Fox, Nucl. Phys. B131 (1977) 107.
- [FRED 82] S. Fredriksson et al., Z. Phys. C14 (1982) 35.
- [FRED 83] S. Fredriksson et al., Z. Phys. C19 (1983) 53.
- [FRIS 83] H.J. Frisch et al., Phys Rev. D27 (1983) 1001.

REFERENCES (Cont'd)

- [FRIT 73] H. Fritzsch, M. Gell-Mann and H. Leutwyler, Phys. Lett. 47B (1973) 365.
- [GEIS 84] W.M. Geist, Proc. 4th Topical Workshop on Proton-Antiproton Collider Physics, Bern (1984).
- [GEIS 87] W.M. Geist, Phys. Lett. 196B (1987) 251.
- [GEIS 88] W.M. Geist, Phys. Lett. 211B (1988) 233.
- [GESS 85] L. Gesswein, Diploma thesis, Univ. of Heidelberg (1985).
- [GLÜC 81] M. Glück and E. Reya, Phys. Lett. 98B (1981) 444.
- [GLÜC 82] M. Glück, E. Hoffmann and E. Reya, Z. Phys. C13 (1982) 119.
- [GOTT 83] T.D. Gottschalk, CALT-68-1059 (1983).
- [GOTT 85] T.D. Gottschalk, CALT-68-1299 (1985).
- [GOUL 83] K. Goulianos, Phys. Rep. 101 (1983) 169.
- [GROS 73] D.J. Gross and F. Wilczek, Phys. Rev. D8 (1973) 3633.
- [GROZ 87] A.G. Grozin, Z. Phys. C34 (1987) 531.
- [GUNI 72] J.F. Gunion, S.J. Brodsky and R. Blankenbecler, Phys. Rev. D6 (1972) 2652.
- [HALZ 83] F. Halzen and P. Hoyer Phys. Lett. 130B (1983) 326.
- [HANN 81] C. Hanna et al., Phys. Rev. Lett. 46 (1981) 398.
- [HANS 82] G. Hanson et al., Phys. Rev. D26 (1982) 991.
- [HERR 85] W. Herr, Ph. D. Thesis, University of Heidelberg (1985).
- [HORG 80] R. Horgan and M. Jacob, Nucl. Phys. B179 (1980) 441.
- [INGE 85] G. Ingelman and R. Schlein, Phys. Lett. 152B (1985) 256.
- [INGE 86] G. Ingelman, DESY, 86-131.
- [ISR 71] The ISR Staff, Phys. Lett. 34B (1971) 425.
- [JACO 76] M. Jacob and P.V. Landshoff, Nucl. Phys. B113 (1976) 395.
- [JACO 78] M. Jacob and P.V. Landshoff, Phys. Rep. 48 (1978) 285.
- [JACO 87] M. Jacob, CERN TH 4813/87.
- [JAFF 89] D.E. Jaffe et al., Phys. Rev. D40 (1989) 2777.
- [JENG 73] R. Jengo, A. Krzywicki and B. Petersson, Nucl. Phys. B65 (1973) 319.

REFERENCES (Cont'd)

- [JONE 85] G.T. Jones et al., Z. Phys. C27 (1985) 43.
- [KAJA 80] K. Kajantie and E. Pietarinen, Phys. Lett. 93B (1980) 269.
- [KENY 82] I.R. Kenyon, Rep. Progr. Phys. 45 (1982) 1069.
- [KIM 88] V.T. Kim, Mod. Phys. Lett. A3 (1988) 909.
- [KOKK 69] J.J. Kokkedee, The Quark Model, W.A. Benjamin, 1969.
- [KOLL 79] K. Koller and H. Krasemann, Phys. Lett. 88B (1979) 119.
- [KONI 78] K. Konishi, A. Ukawa and G. Veneziano, Phys. Lett. 78B (1978) 243.
- [KOUR 79] R. Kourkouvelis et al., Phys. Lett. 84B (1979) 277.
- [KOUR 80] R. Kourkouvelis et al., Phys. Rev. Lett. 45 (1980) 966.
- [KOUR 82] R. Kourkouvelis et al., Z. Phys. C16 (1982) 101.
- [KRIP 77] J. Kripfganz and J. Ranft, Nucl. Phys. B124 (1977) 353.
- [KROL 88] P. Kroll, CERN-TH 4983/88.
- [KRZY 72] A. Krzywicki and B. Peterson, Phys. Rev. D6 (1972) 924.
- [KRZY 78] A. Krzywicki, Ecole d'été de physique des particules, 23 (1978).
- [LAND 77] P.V. Landshoff, Phys. Lett. 66B (1977) 452.
- [LAPE 82] L.V. Laperashvili, Sov. J. Nucl. Phys. 35 (1982) 431.
- [LARS 84] T.I. Larson, Phys. Rev. D29 (1984) 1013.
- [LEV 83] M. Lev and B. Petersson, Z. Phys. C21 (1983) 155.
- [LLOY 80] D. Lloyd Owen et al., Phys. Rev. Lett. 45 (1980) 89.
- [LOHS 82] T. Lohse, Diploma Thesis, University of Dortmund (1982) and refs therein.
- [LOHS 85] T. Lohse, Ph. D. Thesis, Univ. of Dortmund (1985).
- [LOHS 86] T. Lohse, private communication (1986).
- [MALH 83] P.K. Malhotra and R. Orava, Z. Phys. C17 (1983) 85.
- [MANK 88] R. Mankel, Ph.D. Thesis, Univ. of Dortmund (1988).
- [MARC 84] G. Marchesini and B.R. Webber, Nucl. Phys. B238 (1984) 1.
- [MATT 83] G. Matthiae, Proc. of the Int. Europhysics Conf. on High Energy Phys., Brighton 1983.
- [MCCU 81] N.A. McCubbin, Rep. Progr. Phys. 44 (1981) 1027.

REFERENCES (Cont'd)

- [MOLL 84] R. Møller, Proc. of the XV Int. Symp. on Multiparticle Dynamics, Lund (1984) 314.
- [NAUD 86] C. Naudet et al., Phys. Rev. Lett. 56 (1986) 808.
- [NICZ 81a] B. Niczyporuk et al., Z. Phys. C9 (1981) 1.
- [NICZ 81b] B. Niczyporuk et al., Phys. Lett. 100B (1981) 95.
- [OLSE 88] S.L. Olsen, Proc. of the XXIV Int. Conf. on High-Energy Physics, Munich (1988), eds R. Kotthaus and J. Kuhn.
- [OWEN 78] J.F. Owens, E. Reya and M. Glück, Phys. Rev. D18 (1978) 1501.
- [OWEN 84] J.F. Owens, Phys. Rev. D30 (1984) 943.
- [PANO 68] W.K.H. Panofsky, Proc. 14th Int. Conf. on High Energy Phys., Vienna 1968, p. 23.
- [POVL 83] J. Povlis et al., Phys. Rev. Lett. 51 (1983) 967.
- [SAXO 85] D. Saxon, Proc. Int. Europhysics Conf. on High Energy Physics, Bari 1985, p. 899.
- [SCHM 87] M. Schmelling, Ph. D. Thesis, Univ. of Dortmund (1987).
- [SCHW 63] A. Schwarzschild and C. Zupancic, Phys. Rev. 129 (1963) 854.
- [SIVE 76] D. Sivers, S. Brodsky, and R. Blankenbecler, Phys. Rep. 23C (1977) 753.
- [SJOS 86] T. Sjöstrand, Comp. Phys. Comm. 39 (1986) 374.
- [SMIT 85] A.M. Smith et al., Phys. Lett. 163B (1985) 267.
- [SMIT 86] A.M. Smith et al., Phys. Lett. 167B (1986) 248.
- [SMIT 87] A.M. Smith et al., Phys. Lett. 184B (1987) 293.
- [SOSN 83] R. Sosnowski, Proceedings of the Int. Europhysics, Conf. on High Energy Physics, Brighton 1983, p. 628.
- [STEV 86] P. Stevenson and H. Politzer, Nucl. Phys. B277 (1986) 758.
- [SUDA 56] V. Sudakov, Sov. Phys. JEPT 3 (1956) 65.
- [SZCZ 89] M. Szczekowski, Int. J. Mod. Phys. A4 (1989) 3985.
- [TANN 89] M.J. Tannenbaum, Int. J. Mod. Phys. A4 (1989) 3377.
- [VOSS 88] R. Voss, Nucl. Phys. B (Proc. Suppl.) 3 (1988) 581.
- [WALS 80] T.F. Walsh and P.M. Zerwas, Phys. Lett. 93B (1980) 53.
- [WEBB 84] B.R. Webber, Nucl. Phys. B238 (1984) 492.

REFERENCES (Cont'd)

- [WEBB 86] B.R. Webber, Ann. Rev. Nucl. Part. Sci. 36 (1986) 253.
- [WROB 73] A. Wroblewski, Acta Phys. Pol. 43 (1973) 857.
- [YEUN 87] R. Yeung, Ph.D. thesis, UCLA, 1987 UCLA-87-008 and private communication.