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EXPERIMENTS WITH ALPHA PARTICLES AT THE CERN INTERSECTING STORAGE RINGS

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ABSTRACT

The successful injection of alpha particles into the CERN Intersecting Storage Rings (ISR) allowed $\alpha\alpha$ and αp interactions to be studied at centre-of-mass (c.m.) energies $\sqrt{s} = 125$ GeV and 88 GeV, respectively, opening up a new energy regime for the investigation of nuclear collisions. Five experimental groups took data simultaneously, using detectors which had been developed during 10 years of pp experiments at the ISR.

This review begins with an outline of the technical, historical, and physics developments leading to these exceptional runs, followed by a short description of the five detectors which were employed. The results are presented in the order of increasing final-state complexity: elastic scattering, quasi-elastic interactions (nuclear break-up without particle production), inelastic processes leading to particle production, and hard (high- p_T) interactions.

The data are compared with results from pp interactions at corresponding c.m. energies, with extrapolations from nuclear interactions measured at lower energy, and with theoretical expectations and calculations. Whilst it seems that no new phenomena have been uncovered so far, these measurements provide important constraints and guidelines to the search for the predicted new state of matter — the quark-gluon plasma.

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1. INTRODUCTION

In 1980, two exceptional runs took place at the CERN Intersecting Storage Rings (ISR). For the first time, beams of alpha particles were accelerated in the Proton Synchrotron (PS) and transferred to the ISR. First, two alpha beams were stored and collided at a centre-of-mass (c.m.) energy of 125 GeV in the intersections of the rings. Then an alpha beam in one ring was brought to collision with a proton beam in the other ring at a c.m. energy of 88 GeV. New world records were established for c.m. energies of subatomic collisions at accelerators — they were only broken when the proton–antiproton collider at CERN started operating a year later.

The two runs lasted only 60 hours each (fig. 1), but the beam luminosities were high; in fact they met the most optimistic expectations of the experimental physicists who had proposed the runs. The successful operation of the machine complex was the fruit of several months of preparation before the runs and of experience accumulated during many years. Despite the fact that light-ion acceleration never did represent more than a side activity at CERN, this activity has been alive since 1964.

Owing to the high luminosities achieved, the planned survey experiment could be performed and completed within the two short runs. Five experimental teams took data simultaneously in different intersection regions. A large amount of data was collected. If the number of publications per beam hour is taken as a measure, the week with alpha particles in the ISR ranks as one of the most productive weeks in the CERN history [1–18].

The experiments with alpha particles have pioneered a new field: nuclear interactions at the highest possible energies, using a colliding beam machine and with sophisticated, large solid-angle detectors which had been developed during 10 years of pp experiments at the ISR.

The physics goals were born out of particle physics. The role of the nucleus consists in providing the microlaboratory for investigating various aspects of the strong interaction. The mean free path of hadrons in nuclear matter is comparable to or even smaller than the nuclear dimensions; therefore multiple interactions of the nucleons or partons (quarks, gluons) take place in the nucleus. The space–time intervals between these interactions are comparable to the range or characteristic time of the strong interaction itself. Thus the nuclear microlaboratory offers the unique possibility to study the strong interaction on a space and time scale matching the one of the interaction proper. In the detector laboratory, where only the late remnants of the products of the interactions arrive, we are faced with the difficult task of decoding the messages from the microlaboratory; the obvious way to solve this problem is to compare messages from different microlaboratories (i.e. to compare pp with ap and aa interactions).

No doubt the use of nuclear targets or projectiles increases the complexity of the already complex strong interaction. One of the potential rewards has just been indicated; another one was pointed out many years ago by Cocconi [19]. New complex states may be created in nucleus–nucleus collisions at very high energy — states which cannot be reached in pp collisions but which require the interaction of many nucleons under conditions which resemble those presumably prevailing in supernovae or neutron stars or at an early stage of the Universe: high temperature, pressure, and baryon density. What Cocconi visualized as new states of higher complexity may turn out to be the doorway to states of a new simplicity. Lattice quantum chromodynamics (QCD) predicts [20–22] that under the cited conditions a deconfinement phase transition will take place — nuclear matter composed of separate nucleons will melt into a quark–gluon plasma. The experimental proof for such a phase transition is considered to be an important test of QCD, almost as important as the finding of the W or Z boson was for the theory of the electroweak interaction.

The burning question is whether the condition necessary for the transition (an energy density around 1.5 GeV/fm^3) can be reached in nucleus–nucleus collisions at high energy. According to

theoretical estimates, the incoming kinetic energies per nucleon at which collisions of heavy ions have been studied so far (up to 2.1 GeV/N at Berkeley and 3.7 GeV/N at Dubna) are not sufficient. ISR energies corresponding to 500 GeV/N on a fixed target are high enough, but alpha particles are perhaps not heavy enough.

However, the moderate step towards higher complexity in going from pp over ap to aa collisions promises safe advancement. The alpha particle is the smallest “real”, compact nucleus — as opposed to the loosely bound deuteron. Thus on the one hand the number of nucleons is small enough to allow a microscopic description of the collision processes in terms of single and multiple nucleon–nucleon interactions. This will be most transparently illustrated by the case of elastic scattering, where the search for the new phenomenon requires the understanding, microscopically, of the elastic subprocesses which take place. On the other hand, the mean number $\langle \nu \rangle$ of nucleon–nucleon collisions per ap or aa collision is already appreciable [$\langle \nu(ap) \rangle = 4\sigma(pp)/\sigma(ap) \simeq 1.2$; $\langle \nu(aa) \rangle = 16\sigma(pp)/\sigma(aa) \simeq 1.8$].

Observations of the features of ap and aa events and comparison with the single nucleon–nucleon interaction will thus provide a guideline for estimating the conditions in collisions of heavier nuclei at ISR or higher energies — apart from being interesting in their own right. In this sense, the experiments with light ions at the ISR may establish a bridge between particle physics and the field of high-energy nucleus–nucleus collisions yet to be explored.

2. ACCELERATION AND STORAGE OF ALPHA PARTICLES

2.1 The history

“Meanwhile it is suggested that the possibilities might be borne in mind in planning, so that the options can be kept open for the future.”

F. Farley, CERN, 1970

The interest in light-ion beams at CERN has a long history; it remained dormant for many years, but it never died. And it was carried on by two communities, in parallel and sometimes independently: the accelerator and the particle physics communities.

Already in 1964 a study was performed at the 50 MeV Linac injector to the PS, which compared proton and deuteron acceleration and demonstrated the technical feasibility of ion acceleration [23]. An early message advertising the physics interest and feasibility of nucleus–nucleus collisions in the ISR came from Farley in 1970 [24]. Starting from Cocconi’s speculation [19] that in high-energy collisions of heavy nuclei new complex states of nuclear matter may be created, he investigated the possibility of using the ISR as a nucleus–nucleus collider and emphasized the importance of weakly or electromagnetically interacting particles (muons, muon pairs, or photons) as probes in such experiments.

In 1974 the construction of a new Linac at the PS was started. The projected availability of the old Linac for lengthy machine studies and the possibility of alternate pulse injection from the two Linacs into the PS stimulated ideas of using the old Linac for a light ion or polarized beam program in parallel with the main (proton) program. Therefore, around 1975 the studies at the accelerator were intensified [25], culminating in four short test runs in 1976 when deuterons were accelerated in the PS and stacked in the ISR [26]. Several experiments at the ISR took data. In the same year an ISR workshop was held which underlined the theoretical relevance of pd and dd interactions [27]. Topics considered to be of interest were: the intermediate inelastic states in elastic and quasi-elastic interactions, multiparticle production in soft inelastic interactions, neutron diffraction, and comparisons of the high- p_T hadron yield in dd, dp, and pp interactions.

The physics output from the short runs in 1976 was large. One experiment measured the yield of high- p_T π^0 at 90° [28] — and found that it increases by a factor consistent with 2 for pd and consistent with 4 for dd, as compared with the yield in pp interactions at the same $\sqrt{s}$ per nucleon–nucleon collision. The measurement of the differential elastic cross-sections produced evidence for the occurrence of intermediate inelastic states [29]. A large number of papers were published on the diffraction dissociation of neutrons and protons [30].

One may speculate whether these activities would have played a greater role at CERN during the following years if two important events had not occurred at the same time. First, after the discovery of the J/ψ at the end of 1974, the interest of the international particle physics community was largely directed towards the new quantum number charm; secondly, the possibility of $\bar{p}$ cooling was followed up very seriously. This opened the way for a sufficient number of $\bar{p}$ to be accelerated in the PS and stored in the SPS (and also in the ISR) together with protons, using both machines as $p\bar{p}$ colliders; indeed, it became one of the main CERN programmes (fig. 2).

Nevertheless the light ions were not forgotten [31]. In 1977, tests went on at the PS to find out which alpha beam intensity could be achieved at the exit of the Linac [32]. Already in 1976, after one of the test runs with deuterons, a first and seemingly successful attempt had been made to accelerate alpha particles in the PS. But it was afterwards found that the alpha beam was heavily contaminated by deuterons which had been absorbed in the walls of the duoplasmatron source. These deuterons can efficiently compete with He ions even after days of running the source on He gas. The tests in 1977 showed that the true alpha current at the Linac exit was only about 0.5 mA, instead of the 2 mA one believed to have accelerated in 1976.

At the beginning of 1979 the CERN–Heidelberg–Lund Collaboration submitted a proposal to study ap and aa collisions at the ISR using the Split-Field-Magnet (SFM) detector [33]. The main physics goals of the proposed survey experiment were: i) to measure elastic scattering in order to learn more about the intermediate inelastic states; ii) to study the multiparticle production in normal inelastic aa and ap interactions in order to test various soft-parton models which, at that time, were gaining in popularity — these investigations were a logical continuation of experiments on nuclear targets done at FNAL and the CERN SPS with similar goals in mind [34, 35]; iii) to search for the mechanism causing the “anomalous nuclear enhancement” of the yield of hadrons with high transverse momenta produced from nuclear targets, which had been observed at FNAL some years before [36, 37].

When further tests, performed in 1979 and 1980 at the Linac, had shown that alpha beams with sufficient intensity could be produced, the experiment was approved in April 1980 and took place in August of that year.

2.2 The story

“This is more exciting than watching the Olympic games.”

T.J.M. Symons, Communication from I4 to the ISR control room, 28 July 1980

2.2.1 Technical problems

At the time when alpha beams were requested for experiments at the ISR, it was uncertain whether they could be accelerated in the PS since the expected beam intensity was at the lower limit of the PS control system [38]. Setting up the machine with a deuteron beam prior to the alpha run had to be avoided because of the contamination problem mentioned before. The break-through was achieved when a nitrogen gas stripper was introduced after the preacceleration to 516 keV [39]. The stripper had to be pulsed to keep the gas load on the preinjector vacuum below the acceptable limit.

The standard duoplasmatron source was modified and optimized for He^+ production; 30% of the He^+ ions are ionized to He^{++} by the stripping gas.

As for deuterons, acceleration in the (Alvarez-type) Linac is done in the 4π or $2\beta\lambda$ mode. Protons, which are accelerated in the 2π or $\beta\lambda$ mode, move during one cycle of the applied RF field from one gap to the next gap. However, ions with a charge over mass ratio Z/m , equal to one half of the one for protons, would need twice the accelerating field in order for them to be accelerated to the same speed as protons; they would then have twice the proton momentum, thus requiring higher magnetic focusing fields. Both modifications are impracticable at the CERN Linac. Therefore the ions are accelerated with half of the proton velocity, so that the time elapsing whilst they move from one gap to the next is twice the RF cycle (fig. 3).

At injection into the PS the speed of the ions is $\beta = 0.15$ (half that of the protons). Thus a frequency variation by a factor of 6 is needed for acceleration up to $\beta \simeq 1$. However, the frequency of the RF cavities can only be varied by a factor of 3 (from 3 to 9.5 MHz) by actual tuning. This problem is solved by harmonic number switching. The injected ions are trapped at the same RF frequency as the protons, but the harmonic number h is equal to 40 (instead of $h = 20$ for protons), i.e. 40 buckets of particles are stacked around the circumference of the PS. By accelerating them to $\beta = 0.5$, one arrives at the upper frequency limit of 9.55 MHz. At this point the beam is adiabatically debunched, reducing the RF voltage. Then the frequency is retuned to 4.77 MHz while the beam is left coasting at an intermediate magnetic flat top. Thereafter the beam is adiabatically trapped at $h = 20$ for further acceleration.

After having been accelerated in the PS, alpha particles (or deuterons) have the same rigidity as protons, $p/Z = 26.05 \text{ GeV/c}$, corresponding to a momentum of 13 GeV/c per nucleon. Therefore the trajectories to and in the ISR (transfer trajectories and closed orbits) are identical. However, the higher rest mass of the alpha particles leads to slightly different relativistic parameters and hence to small changes in the beam dynamics parameters, such as the RF frequency and space-charge tune shift, and in the transverse stability parameters. Thus the change of the particle rest mass has repercussions on a large number of modules of the ISR control system. Nevertheless the conversion of the control system from protons to alphas takes less than a minute, since in the ISR control system the particle rest mass is a variable to be fetched from a data bank. Difficulties during the stacking of alpha particles into the ISR were caused by unexpected strong longitudinal instabilities, which could be cured by increasing the bunch length with an additional cavity.

2.2.2 The runs

In 1980 the ISR received altogether seven stacks of alpha beams. The first two alpha beams were stored during a machine development session on 21 May 1980. The currents achieved were around 700 mA, giving a luminosity of $3 \times 10^{27} \text{ cm}^{-2} \text{ s}^{-1}$. The c.m. energy of 105.2 GeV represented a new world record for subatomic collisions; this record was broken the same night when the beams were accelerated in the ISR to 62 GeV, allowing a few hours of data-taking at a c.m. energy of 124.8 GeV.

The second two alpha beams, of 3 A each, were stored on 24 July. After accelerating to 62 GeV the luminosity was $4 \times 10^{28} \text{ cm}^{-2} \text{ s}^{-1}$. The beams circulated for about 30 hours; they stubbornly went on circulating when a summer thunderstorm raged over Geneva. But when the thunderstorm had moved away towards the Alps of the Valais, and the man on shift in the ISR control room had written into the logbook “24⁰⁰ thunderstorm stopped — beams quiet”, a minute later a spike on the 18 kV Swiss power-line made the beams suddenly disappear.

The fifth stack was accomplished on 28 July after a 20 A proton beam was already stored in the other ring. This stack was the most successful, reaching 4.3 A. The luminosity for ap collisions was $8 \times 10^{29} \text{ cm}^{-2} \text{ s}^{-1}$ at 88 GeV (c.m. energy).

On 31 July, once more two alpha beams were stored to compensate for the interruption of the first aa run caused by nature.

Table 1 lists the best alpha beam intensities at various stages of the accelerating system [39]. The integrated luminosities, $6 \times 10^{33} \text{ cm}^{-2}$ for aa and $2 \times 10^{35} \text{ cm}^{-2}$ for ap , were as high as expected in the proposal [33], thanks to the careful preparations at the Linac and PS and to the competence and excellent co-operation of the teams operating the Linac, PS, and ISR, and thanks also to adequate beam-time allocation by the ISRC and Research Board.

2.3 Continuation and future projects

2.3.1 More light ions at the ISR

The alpha story had one direct continuation in August 1983, when two more weeks were allocated for experiments with alpha particles and deuterons in the ISR. It was planned to have four runs (aa , ap , dd , and dp), all at equal c.m. energy per nucleon–nucleon collision (31 GeV) in order to facilitate direct comparisons. These runs had been requested by several experimental groups [40, 41] in order to extend the previous measurements and to confirm the most surprising results from the first run.

On the accelerator side, further improvements were achieved. Instead of injecting the ions directly from the Linac into the PS, the PS booster (PSB) was included in the acceleration—as is normally done for proton acceleration. This results in an increase of the beam intensity for the following reasons [42]. The injection efficiency from the Linac to the PS (bypassing the PSB) is only 50%, and only $53 \mu\text{s}$ out of the $100 \mu\text{s}$ long Linac pulse can be used, injecting four turns. (The revolution time, i.e. one turn, in the PS is $13 \mu\text{s}$ for ions with $\beta = 0.15$ at injection.) The injection efficiency from the Linac to the four-ring PSB is higher (75%), and one uses the full Linac pulse. Thus the detour via the PSB results in a potential increase of the intensity by a factor of 3. The increase can be even higher if the beam intensity from the Linac is high. With 4×10^{11} alpha particles in the PS after RF capture, the space-charge limit was in fact already reached during the alpha run in 1980. The space-charge limit of the PSB is 11 times higher. Therefore a maximum number of 45×10^{12} alphas could be trapped in the PSB and injected from there into the PS. Provided the Linac can produce a high enough alpha current, the beam intensity from the PS can be even a factor of 11 higher than for direct injection into the PS. The higher space-charge limit is of use for the deuteron beams where high currents can be produced in the Linac.

Another advantage of including the PSB is that it takes some of the burden from the PS. The booster accelerates deuterons and alphas to a kinetic energy of 500 MeV per charge (800 MeV for protons). No change of the harmonic number in the PS is necessary anymore. However, it has to be changed in the PSB; the initial harmonic number $h = 10$ is changed to $h = 5$ at an intermediate magnetic flat top during PSB acceleration—the procedure is quite analogous to the one described above for the PS.

The light-ion run which took place in 1983 was another great success on the accelerator side. The maximum beam intensities again fulfilled the most optimistic expectations (see table 1). The integrated luminosity for aa interactions was a factor of 15 higher than in 1980 and the one for dd was about $4 \times 10^{35} \text{ cm}^{-2}$. However, nature retaliated by producing a series of local thunderstorms. Thus the pd run was completely lost and the integrated luminosity for the ap run was only 1/4 of the one in 1980.

2.3.2 Other projects

After the successful alpha runs at the ISR in 1980, the possibility of accelerating and storing heavier ions was discussed very intensively. There were two reasons. First, the technical feasibility of ion beams had been demonstrated once more. Secondly, just in the years 1979 to 1981,

high-energy nucleus–nucleus collisions became a subject of fundamental interest as a result of new developments in theoretical physics. Remarkably, they even became of common interest to various separate physics communities: nuclear physicists, astrophysicists, and particle physicists. Part of the motivations were not new: the hope was to simulate in the laboratory the high-density and high-temperature state of nuclear matter supposedly prevalent in neutron stars and supernovae. The idea that at sufficiently high temperature or density the state of nuclear matter would be a new phase — a quark–gluon plasma — was old too, almost as old as the idea of the quark structure of hadrons. The really new incentive came from QCD lattice calculations, which predict that a deconfinement phase transition will take place at around Hagedorn's limiting or critical temperature. Since the literature on this subject has grown enormously during the past years, the reader is referred to the articles by Shuryak [43] and Satz [44], and to the proceedings of three topical meetings [20–22] and references therein.

The employment of the ISR for heavy-ion collisions was forcefully advocated by Willis [45], T.D. Lee [46], and others. However, the European high-energy physics community had decided to close the ISR in 1984 — a sacrifice to the large electron–positron (LEP) accelerator; it soon became clear that this decision was irreversible.

The ISR is not the only machine at CERN capable of producing high-energy ion beams. Once the ions are accelerated in the PS they can just as well be injected into the Super Proton Synchrotron (SPS) and be accelerated there [47]. Thus from a technical point of view, alpha particle beams are already feasible at the SPS. A proposal by a GSI–LBL–Heidelberg–Marburg–Warsaw Collaboration [48] to study ^{16}O collisions with target nuclei ranging from ^{40}Ca to ^{206}Pb at PS energies has recently been approved. The Collaboration will provide the ion source. Therefore, starting from about 1986, even ^{16}O beams will become an option at the CERN accelerator complex.

At the second Bielefeld Workshop [21] experimentalists concentrated on the (hypothetical) use of the SPS for investigations relevant to the experimental verification of the deconfinement phase transition. A rich program was drawn up. At present it is not clear how much of it will be realized, since such a program will have to compete with the very tight future SPS schedule at CERN (fixed-target program with p beams, $p\bar{p}$ collisions, and acceleration of e^- and e^+ for LEP).

Outside CERN, proposals for dedicated high-energy heavy-ion accelerators have been made or are in the process of being made — such as the Venus and Tevalac proposals at Berkeley, the SIS18 and SIS100 proposals at Darmstadt, or the colliding nuclear beam accelerator at Brookhaven. However, none of these projects has been given the green light. The future of high-energy nucleus–nucleus collisions is uncertain beyond 1983.

3. STATUS OF DETECTORS AT THE ISR IN AUGUST 1980

During the alpha week in 1980, five teams took data at four different intersections. In the following the status of the detectors, as of August 1980, will be summarized together with the main physics goals of the different teams. As reference for the experimental set-ups, the official CERN experiment numbers are used. For instance, “R210” signifies the tenth approved experiment in intersection region I2 at the ISR.

3.1 Experiment R110

The detector of the CERN–Oxford–Rockefeller (COR) Collaboration consists of a superconducting solenoid magnet equipped with cylindrical drift chamber modules (DCM), two barrel hodoscopes (A,B) of scintillation counters, and four lead/scintillator sandwich calorimeters (fig. 4). Two lead-glass counter arrays with strip chambers in front of them are localized outside the magnet. This experiment determined the inclusive π^0 yield at large transverse momenta. For the

trigger and the data analysis, only the two lead-glass arrays of 336 individual counters (each 21 radiation lengths thick) and the two scintillation counter hodoscopes A and B were used [5]. The trigger required at least two coincident hits in the A counters to ensure a beam-beam interaction, and an energy deposit in either one of the lead-glass arrays of at least 3 (4) GeV for aa (ap) interactions. The solid angle coverage after the fiducial cuts is

$$\begin{aligned} |y| < 0.4, \quad |\phi| < 20^\circ & \text{ for the inside array, and} \\ |y| < 0.55, \quad |\phi| < 28^\circ & \text{ for the outside array,} \end{aligned}$$

where y is the rapidity and ϕ the azimuthal angle with respect to the horizontal plane in the c.m. system for symmetrical beams (for ap , the y window is shifted owing to asymmetric beam momenta).

Intersection 1 was operating with “low β ” quadrupoles which reduce the effective height of the beams at the intersection, so that the luminosity was about a factor of 2 higher than for all other intersections. This, combined with the large solid angle for π^0 detection, enabled the experiment to extend the measurement of the inclusive π^0 cross-section up to transverse momenta p_T of 8 GeV/c, whereas the three other experiments which measured large- p_T hadron yields reached about 5 GeV/c.

3.2 Experiment R210

The CERN–Napoli–Pisa–Stony Brook Collaboration used a non-magnetic detector to measure the differential elastic aa and ap cross-sections for scattering angles ranging from 3.5 to 16 mrad [4]. Their apparatus (fig. 5a) consisted of two symmetrical sets of hodoscopes located downstream of the interaction region. The circular hodoscopes H3 and H4, which covered the angular region $25 < \theta < 130$ mrad, served as luminosity monitor and as veto counters to reject inelastic events with secondaries in this angular region. Each of the small-angle hodoscopes TB was comprised of two planes (TB ϕ) of scintillation counters (fig. 5b); these provided the trigger and, for the off-line analysis, a measurement of the charge Z of the track, since the integrated pulse recorded is proportional to the energy loss and hence to Z^2 . An array of vertical finger counters (TB y), each counter 25 mm wide, measured the horizontal position. Two planes of drift tubes (TB z), 10 mm in diameter, measured the vertical position.

3.3. Experiment R418

The CERN–Heidelberg–Lund Collaboration performed a survey experiment using the SFM detector. Various triggers were applied to measure elastic scattering, quasi-elastic nuclear break-up, and soft (normal) and hard (high- p_T) inelastic interactions. The main magnet of the detector is a “split-field” magnet — two dipoles with opposite fields joined by common flux return yokes (fig. 6a). The maximum field strength is 1 T, and the airgap is 1 m high with a volume of 30 m³. The gap is equipped with planar multiwire proportional chambers (MWPCs) arranged in such a way as to optimally cover the full solid angle. Two compensator magnets downstream of the interaction are primarily needed to correct the trajectories of the non-interacting beams but, equipped with MWPCs, they also serve as part of small-angle spectrometers, in particular for elastically scattered tracks (fig. 6b). The total number of sense wires is around 70,000. Except for the holes caused by the elliptical beam pipes (< 7 mrad in the vertical coordinate), the MWPCs cover almost the full solid angle [49]. Two gas Cherenkov counters (Č1 and Č2) distinguish between pions and heavy particles (K or p) in the two high- p_T trigger regions around 45° with respect to the beams. A time-of-flight (TOF) scintillation counter system, covering azimuthal angles of about $\pm 13^\circ$ with respect to the horizontal plane, is indicated as well; it was used in connection with large- p_T and elastic triggers.

In general, the apparatus is triggered by means of wire information, and the decision is made in one to three consecutive steps depending on the desired selectivity. The first decision requires a certain number of track candidates in given regions of the detector; a track candidate is defined by coincidences of whole wire planes. The second decision uses signals from groups of wires—a group consists of, for example, 16 adjacent wires of one plane, corresponding to a 64 mm wide slab if the wire distance was 4 mm. The trigger logic selects certain coincidence patterns between such groups from different planes (“roads”), thus imposing limits on the stiffness of the track in the bending (horizontal) plane or linearity in the non-bending (vertical) plane. The third decision uses single-wire-hit information processed by an on-line computer to select, for instance, tracks with momenta above a given value [50].

3.4 Experiment R806

The apparatus of the Athens–Brookhaven–CERN Collaboration consisted of two identical fine-grained lead/liquid-argon calorimeters and of scintillation counter arrays surrounding the vacuum tube at the intersect (fig. 7). An inelastic interaction was indicated by a coincidence between two forward scintillator hodoscopes or by at least one hit in the barrel hodoscope around the intersection; in addition, the trigger required minimum energy deposits in the calorimeter. Both calorimeters were placed at a distance of 2.15 m from the interaction vertex and at a polar angle ($\theta = 90^\circ \pm 15$, i.e. at a rapidity $y = 0 \pm 0.25$ for symmetric beams) and covered an azimuthal angle of 15° . Each calorimeter contains $(14 + 14 + 44)$ gaps of liquid argon separated by thin lead plates. The gaps which house the electrode plane are grouped into three sections: the ϕ section, $2.5 \lambda_0$ ($=$ radiation lengths) deep; the u - v section ($2.5 \lambda_0$); and the θ section ($12 \lambda_0$). The electrode planes are subdivided transversely into strips oriented in four different directions providing the corresponding coordinates: ϕ (azimuthal angle); u and v (inclined); and θ (polar angle) [51]. The spatial resolution of the calorimeters was thus sufficient to separate the two electromagnetic showers from the $\pi^0 \rightarrow \gamma\gamma$ decay and to identify π^0 s with an experimental mass resolution $\sigma = 15$ MeV. The inclusive cross-section for π^0 production was determined in the p_T range from 2 to 5 GeV/c.

3.5 Experiment R807

The Axial Field Spectrometer (AFS), operated by a Brookhaven–CERN–Copenhagen (NBI)–Lund–Pennsylvania–Rutherford–Tel Aviv Collaboration, consists of a conventional C-type magnet. The magnetic field volume (0.5 T on axis) is equipped with a 42-layer cylindrical drift chamber (DC) surrounding the vacuum chamber; between the vacuum chamber and the inner radius of the DC the cylindrical (“barrel”) hodoscope (IH) of scintillation counters is located; this is identical to the one used by R806.

The vertex detector was implemented by a sequence of Cherenkov counters (aerogel, high-pressure, and atmospheric) and of planar MWPCs (PC1, PC2) at 90° with respect to the beams ($45^\circ < \text{polar angle } \theta < 135^\circ$) and covering an azimuthal wedge $\Delta\phi \simeq 45^\circ$ (fig. 8a) [8].

Two copper calorimeters were positioned 4.5 m downstream of the intersection and covered polar angles θ from 1.2° to 6° (fig. 8b) [2]. A scintillation counter hodoscope, consisting of four quadrants, was situated in front of each calorimeter. These calorimeters were used for studies of correlations between deflected energy and charged particle multiplicity in the central region. The spectrometer arm at 90° was used for a measurement of the inclusive single-particle cross-section as a function of p_T .

Various triggers were applied: a “minimum-bias” trigger, high-multiplicity triggers, and a single-particle large- p_T trigger. The minimum-bias trigger required either a coincidence between the two hodoscopes in front of the calorimeters or at least one hit in IH. The large- p_T trigger was

accomplished by a three-stage decision: The fast pretrigger selected “roads” defined by a threefold coincidence between one of the IH counters and corresponding wire groups of PC1 and PC2. The second stage, using a random access memory (RAM), required individual wires of PC1 and PC2 to be aligned with the interaction diamond and an adjacent pair of azimuthal sectors of the DC. The third stage requires a track in the central chamber on the basis of drift-time information and applies a sagitta cut — this stage uses a microprocessor [52].

4. ELASTIC SCATTERING AND TOTAL CROSS-SECTIONS

4.1 Differential elastic cross-section

Elastic scattering is the simplest phenomenon of a two-body interaction, from a phenomenological point of view. Nevertheless, the differential elastic scattering cross-section reflects much of the complexity of the interaction — and of the objects if one or both particles are composite — as, for instance, protons or, on a higher level, nuclei. Thus the elastic cross-section is a fundamental piece of information and is often one of the first things to be measured when a new energy regime is opened up.

The measurement of the elastic cross-section is relatively straightforward at the ISR. A symmetrical two-arm spectrometer capable of particle identification and momentum analysis is required. A veto on low-momentum particles is useful for fast rejection of inelastic events. The test of whether the two observed particles are collinear in the c.m. system is crucial and may even render a measurement of the absolute momenta superfluous.

Two properties of ap and aa interactions render a measurement in this case more difficult than a measurement of pp elastic scattering. First, the elastic ap and aa cross-sections $d\sigma/dt$ are concentrated at much smaller values of $t = -(p_{in} - p_{out})^2$, the invariant four-momentum transfer squared. The slope B of the diffraction cone [$d\sigma/dt \propto \exp(Bt)$] for t above the Coulomb scattering region is proportional to the effective interaction radius squared; hence it is considerably steeper for ap and aa than for pp . Moreover, for the same t value the scattering angle was a factor of 4 smaller for a scattered alpha particle than for a scattered proton owing to the twice higher incoming momentum (62 GeV/c instead of 31 GeV/c). Secondly, charged nuclear fragments, p , d , ^3He , and t , emerge at very small angles, too, and provide a background if particle identification is not complete.

Two experiments, R210 [4] and R418 [6] have measured elastic scattering in partly complementary, partly overlapping t regions. Their results are compiled in table 2 and shown in fig. 9. The discrepancy between the two experiments in their overlap region is apparent. This discrepancy has not yet been explained. For the time being one can only indicate the “weakest points” of the two experiments from which systematic errors could have originated. In the SFM detector (R418) the corrections for absorption in the vacuum chamber and the detector material (chamber frames etc.) were large at the smallest angles. They were tested by measuring the pp cross-section in the same angular range and by comparing it with the world’s best pp data. However, because of the higher absorption cross-section of alpha particles and the complicated geometry and nuclear composition of the absorbing material, the probability that there is a systematic error is not completely negligible. The non-magnetic detector (R210) did not distinguish between ^3He and ^4He (the charge was measured but not the momentum). Background from ^3He was only rejected by the collinearity requirement. The ^3He content of the data was estimated to be $< 5\%$, but since the t distribution of ^3He falls less rapidly than the elastic aa t distribution, this background supposedly increases with t and may have been underestimated at larger t .

The t range covered by the two experiments contains the first diffraction minimum for ap interactions and the first and second minima for aa interactions. This range is very interesting from a theoretical point of view.

4.2 Theoretical interpretation

4.2.1 The effective interaction radius

Elastic hadron scattering at very high energy, such as ISR energies, satisfies two conditions: i) The size R of the scattered particles is large compared with their wavelength ($kR \gg 1$); ii) the distance D between detector and interaction vertex is large, so that $kR^2/D \ll 1$. These are analogous to the conditions for Fraunhofer diffraction of light.

Based on the optical analogy (and ignoring spin-flip contributions) the scattering amplitude can be expressed as the Fourier–Bessel transform of the absorber profile

$$f(t) = ik \int b db \Gamma(b) J_0(b\sqrt{-t}). \quad (4.1)$$

Here $\Gamma(b)$ is the nuclear matter profile as a function of the impact parameter b . For an absorptive profile of the form

$$\Gamma(b) \propto \exp(-b^2/R^2) \quad (4.2)$$

the cross-section has the form

$$d\sigma/dt \propto |f(t)|^2 \propto \exp(-R^2 t/2) \quad [53, 54]. \quad (4.3)$$

Thus the logarithmic slope $B = R^2/2$ measured by R210 below the first diffraction minimum can be used to determine the effective interaction radius R

$$\begin{aligned} B(ap) &= 41 \pm 3 \text{ GeV}^{-2} \rightarrow R(ap) = 1.81 \pm 0.05 \text{ fm}, \\ B(aa) &= 100 \pm 10 \text{ GeV}^{-2} \rightarrow R(aa) = 2.8 \pm 0.2 \text{ fm}. \end{aligned} \quad (4.4)$$

4.2.2 Multiple scattering expansion

The explanation of the diffraction pattern at higher t values is more complicated. Glauber [55] has shown that the minima and maxima arise from the interference between single and multiple nucleon–nucleon (NN) scattering amplitudes. In fig. 10 the four terms of the multiple scattering expansion for the case of $p^4\text{He}$ scattering are shown separately [56]. The NN amplitude was assumed to be purely imaginary and Gaussian in $q^2 (= -t)$:

$$f(q) = (\sigma_{\text{tot}}/4\pi)(i+\rho) \exp(-Bq^2/2) \quad (4.5)$$

with $\rho = 0$. The ^4He form factor was assumed to be Gaussian in q^2 :

$$S(q) \propto \exp(-Cq^2/2). \quad (4.6)$$

Under these assumptions the n -fold convolution of the amplitude is easily performed, and one obtains

$$F_n \propto (-1)^{n-1} \exp[-(B+C)q^2/(2n)], \quad n = 1, \dots, 4, \quad (4.7)$$

with a slope proportional to the inverse number of scatterings n . The essential property of the multiple scattering expansion is that the successive terms alternate in sign. As a consequence, successive terms interfere destructively; for example, the first minimum is the place of maximum

destructive interference between single and double scattering. Thus the characteristic diffraction pattern is explained microscopically.

For quantitative calculations in the framework of the Glauber model, one needs as input more realistic a NN amplitude and a wave function of the target nucleus. The only nucleus whose wave function is accurately known is the deuteron. Usually, for alpha particles simple wave functions are chosen, fitting the measured ^4He charge form factor but ignoring the nuclear properties such as binding energy and S'- or D-state contributions. Thus for αp elastic scattering there are not only uncertainties associated with the first input quantity of the model, the elementary NN amplitude, but also appreciable uncertainties associated with the second input quantity, the nuclear wave function [57]. In fig. 9 two calculations based on the Glauber model are shown as dashed lines. For αp they were done by Proriot et al. [58] and for $\alpha\alpha$ by Alberi et al. [59]. The first authors have chosen a NN amplitude (4.5) with $\sigma_{\text{tot}} = 41.8$ mb, $\rho = 0.062$, $B = 12.92$ (GeV/c) $^{-2}$ and for the ^4He nuclear wave function a Fourier transform [60] fitted to the measured charge form factor. The second authors used $\sigma_{\text{tot}} = 40.2$ mb, $\rho = 0$, and $B = 12.2$ (GeV) $^{-2}$, and a Gaussian density distribution with $\langle R^2 \rangle^{1/2} = 1.66$ fm for ^4He . The agreement between the calculation and the αp data is reasonable apart from a few data points at t around 0.2 (GeV/c) 2 , and apart from the fact that no second minimum is seen by the experiment at $t \simeq 0.7$ (GeV/c) 2 . The pure Glauber calculation lies between the two inconsistent $\alpha\alpha$ data sets at $t > 0.4$ (GeV/c) 2 . The discrepancy between experiment and calculation around the minima is not serious — there are several theoretical refinements, all of which tend to fill up the minima, as for instance by using the correct real to imaginary ratio ρ of the NN amplitude, which is not equal to zero.

4.2.3 Intermediate inelastic states

Leaving aside the uncertainty related to the alpha wave function and the experimental discrepancy for the $\alpha\alpha$ cross-section at large t , we now turn to the perhaps most interesting aspect of nuclear elastic scattering at very high energy. It has been argued that at very high energy the Glauber multiple scattering expansion has to be modified by the inclusion of terms arising from the coherent excitation of intermediate inelastic states (IIS) [61]. This process is illustrated in fig. 11 for a double scattering event. At the first collision the projectile turns into an excited state of mass M , and at the second collision it returns to the ground state. In order not to destroy the nucleus (or the coherence of the outgoing wave) the coherence condition has to be fulfilled:

$$q\lambda < \hbar, \quad (4.8)$$

where λ is the distance between the two nucleons, and the minimum momentum transfer q to excite the intermediate mass M is

$$q = (M^2 - m^2)/(2p_{\text{in}}). \quad (4.9)$$

With increasing incoming momentum p_{in} , higher intermediate masses can be coherently excited. Thus an energy dependence is introduced to the nuclear scattering amplitude (in addition to the one due to the energy dependence of the elementary amplitude).

The intermediate states are quantum fluctuations of the incoming nucleon [54]. They coincide with the states identified as diffractive excitations. Here, as IIS, they become intermediately real as long as the nucleus provides the necessary momentum transfer q to raise them on the mass shell. Also, they are prepared in a very specific way. Thus they provide an independent test for models of diffractive dissociation.

The general experimental situation concerning IIS is not clear. Experiments at the CERN ISR [29] (pd and dd) and at FNAL [62] (ap) claim to have found evidence for IIS. More recent experiments (pd, ap) disagree, at least on the size of the effect [63, 64]. In fig. 9 two calculations are shown (continuous lines) for ap and aa which take into account the IIS correction [58, 59] to the Glauber formulae. The difference is small ($< 20\%$) and is below the experimental precision (which is of the order of 10% owing to the over-all normalization uncertainty) in the case of ap. However, for aa the correction is appreciable in the t region above the second diffraction minimum, and the authors of ref. 6 (lower data points) concluded that in order to describe their data the Glauber model needs a correction because of IIS.

4.3 Total cross-sections

A precise determination of the total cross-section was not possible because the usual methods are either not viable or are very difficult in the case of aa and ap collisions at the ISR. The traditional transmission method is not accessible at the ISR for obvious reasons. The methods of i) measuring the total interaction rate normalized to the luminosity, or ii) extrapolating the differential elastic cross-section down to $t = 0$ and applying the optical theorem, or iii) measuring the total interaction rate and the differential elastic cross-section simultaneously, are severely hindered by the presence of the beam pipes and the special properties of aa and ap interactions already discussed above. The elastic cross-section is concentrated at much smaller angles and there are more inelastic channels which can escape detection than in the case of pp interactions, namely those where the alpha particle breaks up into fragments which are either neutral or have angles close to 0° and rigidities close to the beam rigidity.

Nonetheless the measurements allow a rough estimate of the total elastic and inelastic cross-sections. Extrapolating the differential elastic cross-section below the first diffraction minimum to $t = 0$ and using the optical theorem (neglecting the real part of the amplitude) [53], one determines the total cross-section

$$(\frac{d\sigma}{dt})_{t=0} [\text{mb/GeV}^2] = \sigma_{\text{tot}}^2 / (\hbar^2 16\pi) = 5.1 \times 10^{-2} (\sigma_{\text{tot}} [\text{mb}])^2. \quad (4.10)$$

Experiment R210 [4] thus obtains

$$\sigma_{\text{tot}}(\text{ap}) = 30 \pm 20 \text{ mb} \quad \text{and} \quad \sigma_{\text{tot}}(\text{aa}) < 500 \text{ mb}, \quad (4.11)$$

where the cross-section for aa is regarded as an upper limit since the slope was only determined close to the diffraction minimum. Based on a more sophisticated five-parameter fit to the aa differential cross-section, R210 obtained

$$\sigma_{\text{tot}}(\text{aa}) = 250 \pm 50 \text{ mb}. \quad (4.12)$$

Another, independent determination of the cross-section based on the luminosity-monitor counters led to the estimate

$$\sigma_{\text{tot}}(\text{aa}) = 295 \pm 40 \text{ mb}. \quad (4.13)$$

The total production cross-section σ_{prod} was determined by experiment R418 [11]. The production cross-section does not contain the elastic (σ_{el}) and quasi-elastic ($\sigma_{\text{q.el.}}$) channels, where “quasi-elastic” refers to those inelastic channels where the nucleus breaks up but no particles are created. After correction for events with undetected or neutral pions, R418 obtained the result

$$\sigma_{\text{prod}}(ap) = 101 \pm 10 \text{ mb} \quad \text{and} \quad \sigma_{\text{prod}}(aa) = 265 \pm 26 \text{ mb}. \quad (4.14)$$

Since $\sigma_{\text{q.el.}}$ is estimated theoretically [65] to be small—of the order of 5% of the total cross-section—the values of σ_{prod} are consistent with the inelastic ap and aa cross-sections measured at the next lower energies (see table 3). The ap cross-sections measured by R210 also agree well with the measurements at FNAL and the CERN SPS. Their total and total elastic aa cross-sections are significantly lower than the LBL results [66]. The elastic pp cross-section decreases from about 20 mb (at the LBL Bevatron energy) to 8 mb at $\sqrt{s} = 31$ GeV. This decrease is proportional to the decrease observed for aa interactions.

5. QUASI-ELASTIC INTERACTIONS

5.1 General considerations

High-energy inelastic nuclear interactions contain a special class of events which are characterized by the break-up of the nucleus without accompanying produced particles. (At high energy this class represents a minority!) Such processes may be caused by direct quasi-elastic interactions between the nuclear constituents—nucleons or nucleon clusters. Thus the observed nuclear fragments are either “spectators” or “participants”. In the latter case they receive a momentum kick typical of an elastic interaction. Alternatively, less direct mechanisms may be responsible for the break-up: for instance, the coherent excitation of the nucleus to a higher energy state followed by the decay into various fragment combinations; or the elastic interaction of two nucleons—one out of each projectile—where one of them picks up (recombines with) a third nucleon thus leading to an emitted deuteron.

The highly relativistic kinematics, combined with the properties of the high-energy NN elastic scattering amplitude, have some interesting consequences.

i) The four-momentum transfer t is mainly transverse to the direction of the beams owing to the small scattering angle θ :

$$-t = (p_{\text{in}} - p_{\text{out}})^2 \approx (p\theta)^2 \approx p_T^2. \quad (5.1)$$

Thus quasi-elastic scattering is mainly reflected in the transverse momentum distribution of the fragments.

ii) Correspondingly, the longitudinal momentum transfer Δp_L is small:

$$\Delta p_L = p_{\text{in}}(1 - \cos \theta) \approx p_{\text{in}}\theta^2/2 \approx t/(2p_{\text{in}}). \quad (5.2)$$

Therefore the p_L distribution of the nuclear fragments is governed by the internal motion, independently of whether the fragment was a spectator or a participant.

iii) The internal motion is amplified in the c.m. system or the ISR laboratory system owing to the Lorentz boost. If p_L is the longitudinal momentum of a fragment f with mass m_f in the rest system of the nucleus A , then

$$p_L^{\text{cm}} = \gamma p_L + \eta E \approx \gamma(p_L + m_f), \quad (5.3)$$

where $\gamma = E_{\text{beam}}/(m \cdot A)$ and $\eta = \gamma \cdot \beta$. Thus for an alpha beam energy of 62 GeV, a longitudinal momentum of $p_L = 50$ MeV/c in the alpha rest system becomes amplified to a relative momentum of $\Delta p_L^{\text{cm}} \approx 820$ MeV/c with respect to the “nominal” fragment momentum, which is equal to $\gamma m_f/c$ (neglecting effects resulting from the binding energy and non-zero p_T).

iv) In order to transfer a fixed amount of energy onto the nucleus A, i.e. to raise its mass from m to M , the minimum momentum transfer needed decreases with increasing energy:

$$|t_{\min}| \approx (M^2 - m^2)^2/s, \quad (5.4)$$

where $t = -(PA_{\text{in}} - PA_{\text{out}})^2$ and $s = (PA_{\text{in}} + PB_{\text{in}})^2$; PA_{in} , PA_{out} are the incoming and outgoing four-momenta of the nucleus A; and PB_{in} is the incoming four-momentum of nucleus B. Expression (5.4) is the Lorentz-invariant form of the formula, which is well known, e.g. from the context of diffractive dissociation into high-mass states [cf. expression (4.9)].

In close analogy to the important role of diffraction dissociation in high-energy pp interactions, one may expect a similar phenomenon on the nuclear level in high-energy nuclear interactions. Therefore the nuclear break-up channels may partly be due to diffractive excitation of the nucleus, and the channels without pions represent the low mass spectrum of these processes.

5.2 Results

The quasi-elastic nuclear break-up was studied with the SFM detector (R418) [13]. In this detector the charged fragments p , d , t , and ${}^3\text{He}$ could be distinguished (see fig. 12) by measuring the energy loss ($\rightarrow$ charge) in a scintillation counter hodoscope, and the curvature ($\rightarrow$ momentum) in the magnetic field.

The trigger required one or two high-momentum charged tracks on either side of the intersection and vetoed on low-momentum charged particles. Such a trigger has a serious bias since at least one nuclear fragment is required on either side; moreover, most of the accepted events are incomplete owing to the acceptance hole caused by the vacuum tubes (< 7 mrad) and because neutrons were not detected. Nevertheless the inclusive distributions of the observed charged fragments contain valuable information.

5.2.1 Transverse-momentum distributions

In fig. 13 the p_T^2 distributions of protons, deuterons, and tritons are shown for $\alpha\alpha$ interactions. These distributions are averaged over p_L . (The p_L distributions will be shown below.) One observes that the average logarithmic slope

$$b = -p_T^{-2} [\ln(dN/dp_T^2)_{p_T=0} - \ln(dN/dp_T^2)_{p_T}] \quad (5.5)$$

of the p_T distribution is much lower than expected if the observed fragments were all spectators; the expected slopes for pure spectators are indicated (dotted line).

In evaluating this observation, we have to keep in mind that it applies only to the fragments observed in the SFM detector. The transverse momentum below which the fragments escape detection (because of the beam pipe) is 0.2, 0.3, 0.3, and 0.4 for d , t , ${}^3\text{He}$, and ${}^4\text{He}$, respectively. Thus the spectator fragments d , t , ${}^3\text{He}$, and α are most likely not detected at all. Because of these track losses one loses also many break-up events because at least one track is required on both sides. The event losses in turn cause a bias of the distribution of the proton fragments although the detection efficiency is high (between 40 and 70% at any p_T) for protons. For instance, channels with an unfragmented alpha on one side have almost no chance of being seen. The visible inclusive cross-section for the main break-up channel $\alpha \rightarrow p + X$ (one alpha fragmenting into a proton, without any produced particle but requiring one or two nuclear fragments on the other side), was about 0.5 mb in $\alpha\alpha$ interactions; the fully inclusive cross-section (without any requirements on the other side) is about 5 mb, as estimated from data samples where a minimum-bias trigger was applied.

Bearing in mind the trigger bias discussed above, one may ask what processes are responsible for the emission of the observed fragments. If a single exponential function is fitted to the distribution of protons, one obtains (dash-dotted line) a slope which is close to the slope for elastic NN scattering. From this we infer that quasi-elastic scattering is likely to be the dominant process, except at the lowest p_T where the distribution is somewhat steeper, presumably owing to the presence of true spectator protons.

The slopes for deuterons and tritons are significantly lower than expected for quasi-elastic pd or pt scattering (figs. 13b,c). The Gaussian widths σ_T of the distributions which are related to the slope

$$2\sigma_T^2 = b^{-1} \quad (5.6)$$

are given in table 4. It can be noticed that the widths for p, d, and t are all very similar. This leads to the conclusion that direct quasi-elastic collisions are not the only processes by which d and t are emitted. One may think of other mechanisms; for instance, a quasi-elastically scattered “participant” nucleon picks up one or two “spectator” nucleons in the same nucleus [67].

5.2.2 Longitudinal-momentum distributions

We next consider the longitudinal-momentum (p_L) distribution of protons in the rest system of the alpha particle (fig. 14a). The distribution is slightly asymmetric and the maximum is shifted to negative values. There are two possible causes for these two features. First, it can be shown that for quasi-elastic NN scattering the p_L spectrum is a convolution of an exponential function $\exp(2Bmp_L)$ and a Gaussian distribution

$$d\sigma/dp_L \propto \exp \{-(p_L - \langle p_L \rangle)^2/2\sigma_L^2\} \quad (5.7)$$

for the internal motion, where B is the logarithmic slope for NN elastic scattering and p_L' is the longitudinal component of the quasi-elastic momentum transfer ($p_L' < 0$); the shift of the mean value $\langle p_L \rangle$ is a consequence of binding energy and finite p_T . Secondly, at negative p_L the spectrum contains a background from inelastic channels with π^0 production (estimated to be smaller than 20%). The region $p_L > 0$ is virtually free of this background.

Therefore $d\sigma/dp_L$ was parametrized as in formula (5.7), and the fit to determine σ_L was done only for $p_L > 0$. Experimental values for σ_L and those corrected for the experimental momentum resolution are shown in table 4. The same fit was also done for various p_T bins (see table 5), and it can be seen that, up to $p_T = 400$ MeV/c, σ_L does not depend on p_T .

Comparing the longitudinal width σ_L with the transverse width σ_T , we notice that σ_L is a factor of 3 smaller than σ_T . The widths σ_T and σ_L are defined such that $\sigma_T = \sigma_L$ for pure spectator protons. The asymmetry between σ_L and σ_T results from the contribution of quasi-elastic scattering processes in the observed events. Let us recall that, as a consequence of relativistic kinematics, quasi-elastic scattering shows up in the p_T distributions but perturbs the p_L distributions only slightly. The weak dependence of σ_L on p_T in the spectator and quasi-elastic scattering region ($p_T < 400$ MeV/c) is perhaps the best illustration of this kinematical effect. A similar asymmetry has already been observed in ^4He fragmentation at lower energy [68].

Thus it can be concluded that σ_L is closely related to the width of the intrinsic momentum distribution of the proton before the break-up of the alpha. The width $\sigma_L = 70 \pm 3$ MeV/c obtained for protons ($0 < p_T < 450$ MeV/c) is in good agreement with the width predicted by the statistical model [69, 70],

$$\sigma_0 = \sqrt{(\langle p_L^2 \rangle - \langle p_L \rangle^2)} = p_F/\sqrt{5} = 72 \text{ MeV/c}, \quad (5.8)$$

if one uses for the Fermi momentum p_F the value $p_F = 160 \text{ MeV/c}$ as determined in electron scattering on ${}^6\text{Li}$ [71]. However, this value is higher than the one ($\sigma_L = 55 \text{ MeV/c}$) reported from a measurement of alpha fragmentation at LBL [72].

For fragments with mass K out of a nucleus A , Goldhaber's model [70] predicts for the width:

$$\sigma_L = \sigma_0 \sqrt{[K(A-K)/(A-1)]} = \sigma_0(2/\sqrt{3}) = 83 \text{ MeV/c}. \quad (5.9)$$

The width obtained for the deuterons, $\sigma_L = 122 \pm 8 \text{ MeV/c}$, is higher than this prediction.

Our understanding of the mechanisms involved in the nuclear break-up can be furthered by the study of correlations between fragments on the same side and on opposite sides. Naturally, exclusive measurements, where all fragments are detected simultaneously, are of the greatest value for this kind of investigation (e.g. $ap \rightarrow ppt$ or $ap \rightarrow pdd$). Such studies are still in progress [73].

6. PARTICLE PRODUCTION

6.1 Theoretical background

“This parameter alpha tells you really very little because alpha is the logarithm of something and you are just in the range where you see nothing by looking at alpha, dear friend.”

A. Biaś, Volendam, 1982

6.1.1 Models

The dominant mode of hadron-hadron interactions at high energy is inelastic and leads to multiple production of hadrons, mostly pions. Soon after these processes were discovered many years ago in cosmic-ray events it was realized that conventional (perturbative) approaches to describe them are ruled out owing to the strength of the interaction [74]. Many alternative approaches have been tried in the past 30 years but with only partial success.

The present, widely accepted theory of the strong interaction, QCD, deals mainly with short-distance (high- q^2 or high- p_T) phenomena—since here perturbative methods can be applied. However, the most common event in strong interactions, the production of many hadrons with low p_T , is still outside the domain of QCD since it is intimately connected with the unsolved problem of quark confinement.

Thus the normal (“soft” or low- p_T) inelastic hadron interactions are the realm of several competing models which are all QCD-inspired in the sense that the basic interaction is thought to be caused by colour forces between the hadron constituents. Of course (as a result of natural selection) all these models are able to reproduce many known properties of hadron-nucleon interactions—such as the slow increase in the multiplicity of produced particles with increasing energy or the limited transverse momenta of the bulk of produced particles. However, applied to hadron-nucleus or nucleus-nucleus interactions they lead to qualitatively different predictions. Therefore measurements of particle production from nuclei are considered as a test for these models.

The discussion of results in the following subsections will refer to some of these models. Their basic contents or assumptions are described below.

WOUNDED NUCLEON MODEL (WNM) [75]: Every “wounded” (= inelastically interacting) nucleon out of the nucleus A contributes, on the average, half of the multiplicity produced by a free NN

collision; it contributes only once, even if it has interacted several times with different nucleons in B. Thus the total average multiplicity produced in a collision of nuclei A and B is

$$\langle n(AB) \rangle = 0.5 (WN_A + WN_B) \langle n(NN) \rangle, \quad (6.1)$$

where WN_A and WN_B are the numbers of wounded nucleons in A and B, respectively.

The basic assumption — that a wounded nucleon contributes only once — was derived from a formation time argument: most of the hadrons emitted from the wounded nucleon of A are only formed after A has passed the nucleus B [76].

The original model [75] deals only with the total multiplicity distribution and does not specify the momentum distribution of the hadrons emitted from a wounded nucleon. However, the assumption of equal contribution from every wounded nucleon requires that the rapidity distribution resulting from a wounded target-nucleon is mirror-symmetric to the one from a wounded projectile-nucleon (in the c.m. system). This is illustrated in figs. 15a,b, which show two possible schematic decompositions of the rapidity distribution for a p-nucleus collision with two interacting target-nucleons (three wounded nucleons altogether). The case of fig. 15a corresponds to the old two-fireball model [77]; that of fig. 15b corresponds to the two-phase model by Fishbane and Trefil [78]. In figs. 15c,d also the predictions of two models are shown which are not consistent with mirror symmetry: the Energy Flux Cascade Model (EFCM) [79] and the Multiperipheral Parton Model (MPM), also called the Fan Diagram Dominance Model (FDDM) [80, 81].

ADDITIVE QUARK MODEL (AQM) [65, 81–83]: The wounded nucleons are replaced by the wounded quarks. The quarks involved in the interaction are hadron-like constituents, something like dressed valence quarks [84] or “valons” [85]. A meson contains two of them, nucleons contain three. Thus a meson or a nucleon can be wounded up to two or three times, respectively, traversing a thick nucleus. Every wounded quark contributes only once to the total multiplicity. It is assumed that only one quark of each nucleon is wounded in a NN collision. The idea of interacting constituent quarks goes back to the old additive quark model for hadron–proton interactions [86]. The hypothesis that a wounded quark contributes multiplicity only once is derived from the same formation time argument as that applied for the WNM.

Since, in addition to the wounded (participant) quarks, also the non-interacting (spectator) quarks fragment into hadrons, the prediction for the total multiplicity of produced hadrons and for their distribution in phase space implies more than the simple wounded quark counting. However, for the particle density in the central rapidity region ($y_{cm} \approx 0$) the prediction is simple: each pair of quarks (one out of the projectile, one out of the target), both of which have been wounded for the first time, spans a “colour string”; the string decays into particles giving rise to the same multiplicity around $y_{cm} = 0$ as a NN collision.

To illustrate the main difference between the WNM and the AQM let us consider a proton interaction with a nucleus. Assume that WQ_p quarks in the proton interact with WQ_A quarks in the nucleus A. Since in the model it is assumed that in a NN collision only one quark out of each N interacts, $WQ_A = WN_A$. According to the AQM the central particle density would be WQ_p times higher ($1 < WQ_p < 3$) than for a NN collision:

$$R_{cent}(pA/pp) = \langle n_{cent}(pA) \rangle / \langle n_{cent}(pp) \rangle = WQ_p. \quad (6.2)$$

The WQ_A wounded quarks in A contribute multiplicity only near the target fragmentation region of A. According to the WNM the ratio would be like the one of the total multiplicities (6.1):

$$R_{\text{cent}}(\text{pA/pp}) = (1 + WN_A)/2 = (1 + WQ_A)/2 \quad (6.1')$$

(see fig. 15).

DUAL PARTON MODEL (DPM) [87]: The nucleon wave function is a superposition of Fock states, the simplest one consisting of a (slow) valence quark and a (fast) valence diquark. Higher states contain, in addition, an increasing number of quark–antiquark pairs from the sea. In a NN collision the dominant process is one where the slow valence quarks of target and projectile are interchanged and, owing to colour forces between the separating quarks and diquarks, two decay chains are formed between the diquark of one nucleon and the slow valence quark of the other.

In multiple collisions of a nucleon in a nucleus the sea quark pairs become involved, and chains are formed between them and the valence quark or diquark of the interacting target nucleons. Whereas in the WNM the projectile nucleon could be wounded only once and in the AQM it could be wounded up to three times, in the DPM it can, in principle, be wounded an infinite number of times.

As in the AQM, the multiplicity of produced hadrons depends on the rapidity difference between the two constituents which span the chain or string. Whilst at finite energy the chains resulting from the interaction of a projectile sea quark with a target valence quark are shorter than chains between valence quarks and valence diquarks, at asymptotically high energy the rapidity lengths of the chains become equal. More precisely, the differences between the lengths of the chains is energy-independent since a fixed momentum fraction of the sea quarks corresponds to an energy-independent rapidity difference, but the length of all chains increases like the logarithm of the incoming energy. Therefore, for very high energy the model predicts the multiplicity in the central rapidity region to be proportional to the number ν of NN interactions,

$$R_{\text{cent}}(\text{AB/pp}) = \nu. \quad (6.3)$$

COLOUR NEUTRALIZATION MODEL (CNM) [88]: The nucleon is a superposition of Fock states, as in the DPM. Only partons which are close in rapidity interact or are interchanged. If we define the rapidities of the beam proton y_p and of the target nucleon y_A such that $y_p > y_A$, and y_i is the rapidity where the interaction takes place, then the interacting beam parton materializes into hadrons distributed from y_i to y_p , and the interacting target parton materializes into hadrons distributed from y_A to y_i . This corresponds to the formation of two chains (or strings) in a NN collision — as for all the previously sketched models. For any additional interaction in pA collisions, one more chain is produced extending from y_A to y_{i+1} if y_{i+1} is the rapidity location of the additional interaction. Moreover, if y_{i+1} is lower than any previous y_i , the beam fragment chain is extended from y_{i+1} to y_p . Thus the model is very similar to the DPM. There is, however, a subtle difference: since the interacting partons are assumed to be distributed homogeneously over the central rapidity region, the additional chains, ranging from y_A to y_{i+1} , occupy only half of the available rapidity range (y_A to y_p), on the average — in contrast to the DPM where they are only shorter by a constant energy-independent amount than $(y_p - y_A)$. At asymptotically high energy, the central multiplicity for a collision of a proton with ν target nucleons is

$$R_{\text{cent}}(\text{AB/pp}) = [\nu/2 + \nu/(\nu+1)]. \quad (6.4)$$

6.1.2 The average event

In the above section on models it was assumed that we know the number ν of elementary NN collisions which took place in an individual event and also the number of wounded nucleons.

Although special experiments have tried to find a way of measuring the number ν on an event-by-event basis [35, 89], this is generally not feasible. Therefore in many cases the average number $\langle \nu \rangle$ [90] is used:

$$\langle \nu(AB) \rangle = AB\sigma(NN)/\sigma(AB), \quad (6.5)$$

where $\sigma(NN)$ and $\sigma(AB)$ are the integrated cross-sections for a NN collision and one of nucleus A with nucleus B. Depending on whether we are interested in the average number of collisions leading to particle production, or in all inelastic collisions, or in any collisions, we may use σ_{prod} , or σ_{inel} , or σ_{tot} . For proton–nucleus (pA) collisions, the number of these is equal to the number of wounded nucleons in A:

$$\langle WN_A(pA) \rangle = \langle \nu(pA) \rangle. \quad (6.6)$$

For nucleus–nucleus collisions

$$\langle WN_A(AB) \rangle = A\sigma(pB)/\sigma(AB) \quad \text{and} \quad \langle WN_B(AB) \rangle = B\sigma(pA)/\sigma(AB) \quad (6.7)$$

and the total number of NN collisions $\langle \nu(AB) \rangle$ is equal to the average number of collisions one nucleon out of B (or A) experiences times the average number of wounded nucleons in B (or A):

$$\langle \nu(AB) \rangle = [A\sigma(NN)/\sigma(NA)] \cdot [B\sigma(NA)/\sigma(AB)] = AB\sigma(NN)/\sigma(AB). \quad (6.8)$$

Analogous “counting rules” can be written down for the number of wounded quarks and quark–quark interactions by replacing the nucleon cross-sections with the corresponding quark cross-sections and the numbers of constituent nucleons with the numbers of constituent quarks [83]. A few examples of $\langle \nu \rangle$ are given in table 6. Observables such as the final-state multiplicity of tracks are frequently shown as a function of $\langle \nu \rangle$.

Alternatively, the data are often plotted as a function of A, and it turns out that a parametrization such as A^α works surprisingly well in many cases. (This has led to the name “ A^α physics” for experiments studying A-dependences.) For instance, the total cross-sections are well described, at 280 GeV/c incoming momentum and for A larger than 6, by [91]:

$$\begin{aligned} \sigma(pA)_{\text{prod}} &= 1.14 \sigma(pp)_{\text{prod}} A^{0.718}, \\ \sigma(\pi A)_{\text{prod}} &= 1.27 \sigma(\pi p)_{\text{prod}} A^{0.753}, \\ \sigma(\bar{p}A)_{\text{prod}} &= 1.17 \sigma(p\bar{p})_{\text{prod}} A^{0.713}. \end{aligned} \quad (6.9)$$

It can be noticed that even for antiprotons the black disk limit $A^{2/3}$ is not quite reached. Equations (6.9) are convenient and very precise tools for calculating the average numbers of collisions $\langle \nu(pA) \rangle$ according to eq. (6.5) in hadron–nucleus collisions. A general parametrization of this form for an inclusive single-particle cross-section would then be

$$d^3\sigma/dp^3(AB \rightarrow hX) = c A^\alpha B^\beta d^3\sigma/dp^3(NN \rightarrow hX), \quad (6.10)$$

where α and β depend on the particle species and its momentum.

The same parametrization has even been used to describe the A-dependence of the distribution of global event parameters such as total transverse energy E_T per event (summed over all produced hadrons) [92].

Other possible parametrizations appear even more natural in the framework of certain models and fit the data equally well [see, for instance, eq. (7.2) in the next section]; thus there are practical but probably no deep theoretical reasons why those of the form (6.9) or (6.10) are preferred.

One warning at the end of this introduction to theoretical models is appropriate. When talking about multiparticle production in general, one usually refers to the global or average characteristics of these processes, e.g. to the multiplicity distribution or to the inclusive distribution of a certain particle species h as a function of one variable (transverse momentum p_T or rapidity y). In inclusive measurements $A+B \rightarrow h+X$, one averages over all unobserved particles and variables (X) characterizing the event. The situation is even worse for nuclei, since in addition one averages over various numbers of collisions.

Often the inclusive distributions are interpreted as representative of the “average event”. This is a mistake. Such an “average event” is perhaps one which never occurs. Consider, for instance, the average rapidity distribution of produced particles in a pp or AB collision and assume for the sake of the argument that the dominant mechanism would be single diffraction, i.e. the projectile gets excited, the target remains in the ground state and vice versa; in contrast to the assumption, the average event would suggest a double-diffraction mechanism.

Thus we should keep in mind that the information from inclusive measurements is limited and often leaves too much freedom to the theoretical interpretation. On the other hand, theoretical models which deal only with the “average event”, run the risk of simplifying the problem too much. Yet, at present, most data are inclusive and we are therefore forced to deal with the average event.

6.2 Multiplicity distributions

6.2.1 Total multiplicity of charged hadrons

The topological cross-sections $\sigma(n)$ for having n charged secondary hadrons in the final state were measured by R418 [11]. The applied “minimum-bias” trigger accepted all events with at least one charged track anywhere in the detector. Elastic and purely quasi-elastic events were eliminated in the off-line analysis. Therefore the event sample contains only truly inelastic events where particles are produced, diffractive inelastic interactions included. The corresponding cross-section is σ_{prod} , see subsection 4.3. Because of the full solid-angle coverage of the SFM detector, 95% of σ_{prod} are seen. The average track detection efficiency for negative tracks is between 80% and 85% depending on the magnetic field setting, and it is relatively homogeneous over the full phase space. In contrast, the acceptance for positive tracks with rigidities close to the beam rigidity vanishes at small angles owing to the beam pipe. It is therefore advantageous to consider only the multiplicity n_- of negative tracks; n_- is a direct measure for all produced charges (n_{ch}) since

$$n_{\text{ch}} = 2n_- + Z_1 + Z_2, \quad (6.11)$$

where Z_i are the charges of the incoming particles or nuclei.

The results for the multiplicity distribution, corrected for track detection efficiency, event losses, secondary interactions, and decays of K^0 and Λ , are listed in table 7. In fig. 16 the distributions are shown and compared with the multiplicity distributions for pp interactions at corresponding $\sqrt{s}$, which were measured and analysed in exactly the same way. In fig. 16a,b the previous data for pp multiplicity distributions are also shown [93]; the new pp data have significantly smaller errors than the older streamer chamber data.

The lines in the figures represent theoretical calculations made by Chao and Pirner [94]. According to their model, which is based on the DPM, two chains (or “superclusters”) are formed in each interaction. Furthermore, they assumed a stochastic and independent decay of each supercluster into hadrons. The model has three free parameters relating the final-state multiplicity

of hadrons to the number of superclusters. These parameters were determined by a fit to the energy dependence of the mean charged multiplicity in hadron-proton interactions,

$$\langle n(\text{hp}) \rangle = 0.88 + 0.44 \ln(s) + 0.118 \ln^2(s), \quad (6.12)$$

and the empirical relation between dispersion and mean multiplicity,

$$d(\text{hp}) = \sqrt{[\langle n^2(\text{hp}) \rangle - \langle n(\text{hp}) \rangle^2]} = 0.55 \langle n(\text{hp}) \rangle - 0.5. \quad (6.13)$$

The lines in figs. 16a,b represent the predicted multiplicity distribution after the three parameters have been fixed. The distribution of the number of superclusters was calculated in a geometrical “row on row” model [95]. There are $2^{16} - 1 = 65535$ different combinations of nucleons in aa collisions. Only 316 groups of combinations (diagrams) are physically distinguishable. (See reference [96] for a list of these diagrams together with their partial inelastic cross-sections.) These diagrams are composed of the following basic types of interactions (fig. 17):

- i) one nucleon of A interacts with one or more nucleons in B;
- ii) two or more parallel interactions take place;
- iii) one row of nucleons out of A collides with one row out of B at the same impact parameter.

In figs. 16c,d the predictions of the model are shown for ap and aa interactions; in the latter case the contributions from the three different types of interactions (i) to (iii) to the total multiplicity distribution are drawn separately. The agreement between the data and this calculation is very good apart from some deviation at high multiplicity for ap and at low multiplicity for aa . Such deviations are perhaps not surprising since the model is only relevant for non-diffractive processes, but the multiplicity distributions include both diffractive and non-diffractive events. Diffractive inelastic scattering amounts to almost 20% of the total NN cross-section [97], and coherent diffraction on the ^4He nucleus occurs on the level of 10% of the total cross-section according to Bialas et al. [65]. By fitting the three free parameters to the experimental multiplicity distributions, the possible errors due to this simplification are probably minimized, but not fully excluded.

In a recent paper, Takagi [98] calculates the multiplicity distribution of negative particles in aa collisions in the framework of the “Collective Quark Tube model”. The agreement between his results and the data is very good. The average quark tube size, i.e. the number of quarks in a row out of one alpha particle which collides with a quark or row of quarks out of the other alpha particle, is only 1.06. This number measures the collectivity (or number of multiple sequential collisions) on the quark level. The small deviation of this number from unity indicates that collective quark interactions are rare in aa interactions.

In order to make a comparison with several other models (most of which deal only with the “average event”), the measured and predicted ratios of the average total multiplicities of negative tracks,

$$R_{(aa/pp)} = \langle n_{(aa)} \rangle / \langle n_{(pp)} \rangle \quad (6.14)$$

$$R_{(ap/pp)} = \langle n_{(ap)} \rangle / \langle n_{(pp)} \rangle,$$

are given in table 8. All models agree with the data if the theoretical uncertainties of the order of 10%, which are mainly due to the uncertainty of the measured cross-sections, are taken into account. The cross-sections are needed in order to calculate the average number of inelastic collisions which determine the multiplicity according to the models. These numbers are

$$\begin{aligned}\langle \nu(ap) \rangle &= 4\sigma_{\text{inel}}(NN)/\sigma_{\text{prod}}(ap) = 1.31 \pm 0.13, \\ \langle \nu(aa) \rangle &= 16\sigma_{\text{inel}}(NN)/\sigma_{\text{prod}}(aa) = 2.0 \pm 0.2.\end{aligned}\tag{6.15}$$

(They are not identical to the numbers in table 6 which included also elastic collisions.) Clearly, more precise measurements of the cross-sections are needed for a sharper test of the models.

In fig. 18 the dispersions d_- of the multiplicity distributions are plotted versus $\langle n_- \rangle$ and compared with pp data [93, 99] at various energies. All pp data agree with a linear relationship between d_- and $\langle n_- \rangle$; the line represents the original fit by Wroblewski [100]. The aa and ap data points do not fall on the same line. This indicates that—as expected in multiple interaction models—nuclear collisions do not obey the same KNO scaling as pp collisions. In these models the apparent increase of the ratio $d_-/\langle n_- \rangle$ results from two effects working in opposite directions: the ratio is lowered (as $1/\sqrt{\nu}$) for a higher number ν of collisions, but the fluctuation of ν raises it again. The values for some higher moments of the multiplicity distributions have been given in table 7.

6.2.2 Charged particle multiplicity in the central rapidity region

Experiment R807 has studied multiplicity distributions of charged particles produced in the central rapidity region [7]. The trigger required at least one hit (for some data sets at least two hits) in the barrel counter hodoscope surrounding the intersection, or a coincidence between the scintillator hodoscopes in front of the two calorimeters at forward angles (see fig. 8). This trigger is not sensitive to inelastic diffractive events. In fig. 19 the distribution of tracks found in the cylindrical drift chamber within a pseudorapidity interval $|\eta| = |\ln \tan(\theta/2)| < 0.8$ and over nearly the full azimuthal angle, apart from two gaps of width $\Delta\phi = 16^\circ$, are shown for aa , ap , and pp interactions at corresponding $\sqrt{s_{NN}}$.

These distributions were first used to test whether the same KNO scaling of the form

$$F(z) = (a + bz + cz^2) \exp(-dz),\tag{6.16}$$

with $z = n/\langle n \rangle$, describes pp, ap , and aa data. The constants b , c , and d were determined from the pp data. Only the mean $\langle n \rangle$ and the over-all normalization a were fitted. The fit to the ap data is rather good, with a $\chi^2/\text{d.f.} = 1.1$, but the aa data show a systematic departure from the pp scaling function, the minimum $\chi^2/\text{d.f.}$ being 3.8. This is illustrated in fig. 20a, which shows the normalized squared deviation multiplied by the sign of the deviation as a function of the multiplicity. This result agrees qualitatively with the one obtained by R418 for the total multiplicity distribution.

The second investigation confronted the measured multiplicity distribution with the WNM and the AQM. It was assumed that the basic assumption (6.1) of the WNM is also valid for the multiplicity in the central region,

$$R_{\text{cent}}(AB/NN) = W/2 = (WN_A + WN_B)/2\tag{6.17}$$

and that the multiplicity distribution for W wounded nucleons is the W -fold convolution A_W of the multiplicity “amplitude” $A(n)$, which has the property that a twofold convolution

$$A_2(n) = \sum A(j)A(n-j)\tag{6.18}$$

yields the observed pp multiplicity distribution. The probabilities a_W for having W wounded nucleons were determined using as the nucleon density distribution for ^4He :

$$\rho(r)/\rho(0) = (1 + ar^2/c^2)/[1 + \exp(r/f - c/f)], \quad (6.19)$$

with $a = 0.445$, $c = 1.008$ fm, $f = 0.327$ fm [101], and using a geometrical model (no deflection of the trajectories in the nucleus). The probability a_w depends on the inelastic NN cross-section σ^* in the nucleus, which was treated as the only free parameter. The resulting expression for the multiplicity distribution

$$P(n) = \sum a_w(\sigma^*) A_w(n) \quad (6.20)$$

was fitted to the experimental distributions. The values for σ^* obtained from the fit, 25 ± 4 mb for ap and 20 ± 5 mb for aa , are in good agreement with the non-diffractive free pp cross-section $\sigma(pp) = 26$ mb [97]. The corresponding values for a_w are given in table 9. These values are compared in case of aa collisions with the corresponding numbers calculated by Takagi [98], by Białas and Kolawa [102], and by Stenlund [103] in eikonal approximations; the agreement between the fits and the theoretical calculations is fairly good. In fig. 20b the normalized squared deviation (times the sign) is shown as a function of the multiplicity. The average value $\chi^2/\text{d.f.}$ is good (≈ 1) for ap and acceptable for aa ($\chi^2/\text{d.f.} = 2.8$).

The same procedure was then used to test the AQM by calculating the number of strings between wounded quarks. A corresponding fit to the experimental data yields a NN cross-section σ^* which is too small (14 mb for ap and 11 mb for aa) and, in particular for aa , the $\chi^2/\text{d.f.} = 16$ signifies a bad fit. This is demonstrated in fig. 20b. From the small fitted value of σ^* it can be inferred that the AQM would give too high an average central multiplicity $\langle n_{\text{cent}} \rangle$ for aa collisions if one would fix σ^* at a reasonable value. The experimental ratio

$$R_{\text{cent}}(aa/pp) = \langle n_{\text{cent}}(aa) \rangle / \langle n_{\text{cent}}(pp) \rangle$$

equals 1.5 if one uses $\langle n_{\text{cent}}(pp) \rangle = 2.66$ as obtained by the best fit of the KNO scaling function to the pp data, and the ratio $R_{\text{cent}}(ap/pp)$ equals 1.15. Unfortunately, a direct comparison of the multiplicities and multiplicity ratios from this experiment with the results of R418 and the “standard” theoretical predictions (see table 8) is not possible because, firstly, the multiplicity is limited to a part of the central region. Secondly, the experiment has a trigger bias: the trigger of R807 is insensitive to inelastic diffraction of nucleons. Thus the multiplicity distribution for pp interactions is biased in favour of non-diffractive events. But events where diffractive and non-diffractive NN interactions occur either simultaneously or successively in an ap or aa collision will be accepted by the trigger. The fact that the trigger biases for aa and ap interactions are different from those for pp interactions tends to lower the multiplicity ratios $R_{\text{cent}}(aa/pp)$ and $R_{\text{cent}}(ap/pp)$.

It is difficult to judge whether the systematic uncertainties introduced by the not completely known trigger bias are smaller than the small differences between the WNM and the AQM. These uncertainties could, in principle, be eliminated by measuring the topological cross-sections or the total visible cross-section [$\sigma(n) = \sigma_{\text{vis}} P(n)$].

In a recent calculation Białas and Kolawa compare the multiplicity distributions in the wounded nucleon and the colour string model [102]. They conclude that: a) it is difficult to distinguish between the two models by investigations of the multiplicity distributions; b) the predictions for the multiplicity distributions in ap and aa interactions are to a large extent model-independent.

6.3 Transverse momenta and transverse energy

6.3.1 Transverse momenta and hadronic temperature

The inclusive invariant cross-sections $E(d\sigma^3/dp^3)$ for $aa \rightarrow h^-$ [3, 8], $aa \rightarrow \pi^0$ [1], and $pp \rightarrow h^-$ [163, 8] are shown in fig. 21 as a function of p_T at fixed rapidities (y_{cm} near 0). In the following, these cross-sections will be discussed with the main emphasis on the low- p_T part and on the connection with the concept of hadronic temperature; in the next section on hard interactions they will be discussed once more with the emphasis on the high- p_T part.

Let us first recall the natural separation between high p_T and low p_T : high p_T is defined in such a way that the corresponding dimension is much smaller than the nucleon diameter R , say by a factor of 10. Therefore $p_T > 10 \hbar/R \approx 2 \text{ GeV}/c$ can be called high p_T .

For p_T up to about $1 \text{ GeV}/c$ the invariant cross-section decreases rapidly with increasing p_T ; the dependence is approximately exponential, with the well-known slope $B \approx -6 (\text{GeV}/c)^{-1}$ for aa , as for pp (and ap) interactions. From pp interactions it is known that the slope does not depend on the incoming energy ($10 < E_{inc} < 2000 \text{ GeV}$, on a fixed target). This feature of “limited transverse momenta” has been explained by the existence of a maximum hadronic temperature [104, 105]:

$$T_{max} \approx B^{-1} \approx m_\pi, \quad (6.21)$$

The relation (6.21) between T_{max} and the inverse slope is only an approximate equality since the correct thermal p_T distribution is

$$d^2\sigma/dp_T^2 \propto m_T \sum_{n=0}^{\infty} (\mp)^n K_1[(n+1)m_T/T], \quad (6.22)$$

with $m_T = \sqrt{p_T^2 + m^2}$ and the $+$ sign for bosons, the $-$ sign for fermions; K_1 is the modified Bessel function which approaches $\exp(-p_T/T)/\sqrt{p_T}$ for $p_T \gg m$ [105].

At low p_T ($p_T < 0.2 \text{ GeV}/c$), which is often below the range covered by experiments, the function (6.22) deviates significantly from an exponential $\exp(-Bp_T)$.

In order to determine the temperature, several functions were fitted to the p_T distribution of negative particles in the p_T range from 0.1 to $2.0 \text{ GeV}/c$: the correct thermal distribution (6.22), and other, more common—although less justified—functions [16]. Using only negative particles has the advantage that the assumption of a pion mass for all of them is a good approximation, since over 80% are known to be pions. The results of these fits are shown in table 10 together with the assumed functional forms of the distribution. In addition, $\langle p_T \rangle$ is also given, which is related to T by the formulae given also in the table. It is seen that T and $\langle p_T \rangle$ vary significantly for the different fits. However, the values for aa are always larger than those for pp interactions by a few percent.

Above $p_T \approx 1 \text{ GeV}/c$ the decrease of the cross-sections (fig. 21) becomes less rapid; at $3 \text{ GeV}/c$ the pion yield is already higher by more than a factor of 10 than expected by extrapolation from the low- p_T data (see dashed line in fig. 21). This departure from $\exp(-Bp_T)$ depends on the incoming energy; it becomes more pronounced at higher energies. The discovery of this phenomenon, first made at the ISR [106], supported the evidence for a substructure of the nucleons—in close analogy to Rutherford’s conclusion on a point-like nucleus when he observed momentum transfers on alpha particles which were larger than expected. It was the first evidence for point-like constituents in strong interactions.

Some authors [43, 107] consider the decreasing exponential slope as evidence of temperatures higher than T_{max} . They assume that the hadrons with “medium-high p_T ” ($2\text{--}5 \text{ GeV}/c$) result from the fragmentation of partons (quarks or gluons) emitted directly from a hot plasma which was formed in the initial state of the collision, whereas the low- p_T hadrons are emitted after the

plasma expanded and cooled down to the temperature $T_{\max}$. In these new or extended thermodynamic models $T_{\max}$ plays the role of the critical temperature at which partons condense into hadrons. These models are mentioned here in the context of low- p_T interactions, and not in the next section on hard interactions, because there are no quantitative predictions (known to the author) from any of them which go beyond the explanation of the inclusive distributions. According to some of the authors, a quark-gluon plasma is already created in pp collisions at ISR energies. Obviously, the chances of creating such a plasma in a nucleus-nucleus collision at similar energy must be greater [108] and therefore one expects the component with the flatter slope in the p_T spectrum to increase. As we shall see in the next section, this expectation is qualitatively met by the experimental findings; however, alternative explanations are possible.

6.3.2 Multiplicity dependence of p_T distributions

If we relate the slopes of the inclusive p_T distribution to a temperature (or several temperatures, since the slope changes as a function of p_T), we have to keep in mind that these temperatures represent an average over many different events. We may ask whether the temperature changes as a function of some quantities characterizing the event, such as multiplicity or E_T of the produced particles. Here we study the dependence of the p_T distributions on the multiplicity and compare it for aa , ap , pp , and $p\bar{p}$ interactions.

A first investigation of the multiplicity dependence of the p_T distribution led to the result that, within statistical errors, the p_T distributions in the range $0.2 < p_T < 2$ GeV/c scale with the total multiplicity for aa interactions at $\sqrt{s} = 125$ GeV in the same way as for pp interactions at $\sqrt{s} = 63$ GeV [109]. In contrast with this observation, a surprising result was obtained at the CERN $p\bar{p}$ Collider at $\sqrt{s} = 540$ GeV [110]: the average $\langle p_T \rangle$ increases significantly with increasing multiplicity (fig. 22). For determining $\langle p_T \rangle$ we have to know the complete p_T distribution or to extrapolate it into the unknown regions (see previous subsection). Alternatively we may use a substitute for $\langle p_T \rangle$, e.g. the inverse logarithmic slope of the distribution in the known p_T region. If $d\sigma/dp_T^2$ would be proportional to $\exp(-Bp_T)$ for p_T from 0 to infinity, then

$$\langle p_T \rangle = 2/B. \quad (6.23)$$

This procedure was applied to the aa , ap , and pp data [111]. It was then found that $2/B$ does not differ significantly between pp and aa and is almost constant (or at most weakly increasing) as a function of multiplicity in the central region—in contrast with the $p\bar{p}$ data at $\sqrt{s} = 540$ GeV (fig. 22). The same result was obtained by the AFS Collaboration, calculating $\langle p_T \rangle$ for $0.08 < p_T < 2$ GeV/c (see fig. 5 in ref. [7]): A recent study with pp data samples of much higher statistics has led to the result [112] that for the highest ISR energy, $\sqrt{s} = 63$ GeV, the $\langle p_T \rangle$ does indeed rise, but very weakly, with n , whereas at $\sqrt{s} = 31$ GeV no such rise is seen. The rise of $\langle p_T \rangle$ as a function of n observed at the $p\bar{p}$ Collider has stimulated several interesting thoughts [113–115], but the weak n dependence of $\langle p_T \rangle$ at ISR energies cannot easily be fitted into the same picture.

A better insight into the multiplicity dependence of p_T distributions can be obtained if one studies the change of the whole distribution instead of the change of only the average p_T . This is done in fig. 23 where the ratios

$$R_{p_T}(n) = [(n\sigma_n)^{-1}d\sigma_n/dp_T^2] / [\langle n \rangle \sigma^{-1}d\sigma/dp_T^2] \quad (6.24)$$

are plotted as a function of p_T . The denominator is the fully inclusive (average) p_T distribution normalized by the corresponding $\langle n \rangle$; σ_n in the nominator signifies that the topological cross-section has n produced charged particles in the final state. This ratio has the additional

advantage that the acceptance cancels as long as it is multiplicity-independent. It can be seen that the distributions change in a potentially interesting way. If, for instance, the slope of the p_T distribution would simply decrease for increasing n , we would expect to see a continuous decrease of the ratio for increasing p_T , at low n , and a continuous increase of the ratio at high n . What is seen instead is a shallow maximum or minimum in the p_T region from 0.3 to 0.6 GeV/c. It looks as if the p_T distribution $\sigma_n^{-1} d\sigma_n/dp_T$ becomes broader for increasing n .

6.3.3 Distribution of total transverse energy

The transverse momenta of charged particles were determined by experiment R807 in the central rapidity region ($|y| < 0.8$) and, assuming their mass to be equal to the pion mass, the transverse energies $m_T = \sqrt{p_T^2 + m^2}$ of all tracks in an event were summed to obtain E_T . The resulting inclusive cross-sections $d\sigma/dE_T$ for aa and ap are shown in fig. 24 together with the corresponding cross-sections for pp interactions at the same $\sqrt{s}$ [18]. The pp data were multiplied by a factor of 2.5 and of 5 for the comparison with ap and aa , respectively, in order to normalize the curves at $E_T = 0.5$ GeV. It can be seen that the E_T distributions are wider for ap and aa than for pp interactions. The probability of having an event with $E_T = 10$ GeV is greater by a factor of 100 in aa than in pp interactions. The question is whether this increased probability is due to the more frequent occurrence of hard interactions, i.e. interactions where fast quarks or gluons get deflected to large angles (jet production), or whether it is due to soft, high-multiplicity processes where many hadrons with small p_T add up to a large E_T .

A study of the average multiplicity of tracks $\langle n(E_T) \rangle$ and of their average transverse momentum as a function of E_T , $\langle p_T(E_T) \rangle$, led to the conclusion that the high- E_T events are soft — at least up to the highest E_T reached in this experiment. Note that $\langle n(E_T) \rangle$ and $\langle p_T(E_T) \rangle$ are related by

$$E_T = \langle n(E_T) \rangle \langle p_T(E_T) \rangle. \quad (6.25)$$

The evidence is shown in fig. 25. With increasing E_T , $\langle n(E_T) \rangle$ grows linearly but then seems to reach a plateau. Correspondingly, $\langle p_T(E_T) \rangle$ grows slowly from 300 MeV/c to about 600 MeV/c, except at the highest E_T where it rises more quickly in correspondence with the saturation of $\langle n(E_T) \rangle$. Assuming that the observed multiplicity distribution $P(n)$ in the central y region and the single-particle inclusive cross-section $d\sigma/dp_T^2$ are uncorrelated, a Monte Carlo simulation led to a result which agrees remarkably well with the experimental observation (see fig. 25). This shows that the observed rise of $\langle p_T(E_T) \rangle$ or the flattening of $\langle n(E_T) \rangle$ at high E_T is a consequence of the choice that nature has to make given a rapidly falling multiplicity distribution and an also rapidly falling p_T distribution in order to reach high E_T . In other words, the information obtained from $d\sigma/dE_T$ is already contained in $d\sigma/dp_T$ and $P(n)$.

Therefore, from the R807 data it was concluded [18] that there is no need for a new mechanism (jet production), even at the highest values of E_T studied. However, it should be added that the data do not preclude the onset of a new mechanism; it is likely that at the highest E_T where $\langle p_T(E_T) \rangle = 0.8$ GeV/c, the events contain a non-negligible amount of hard interactions.

Likewise, the observed slow rise of $\langle p_T(E_T) \rangle$ from 0.3 to 0.6 GeV/c does not necessarily signify an increase of the “temperature” T . Even given a thermal distribution for one single T , a qualitatively similar rise would be enforced by the high- E_T trigger. But if the inclusive single-particle p_T distribution already displays more than one T — a possibility which was discussed before — the rise of $\langle p_T(E_T) \rangle$ can indeed indicate a rising temperature, i.e. events with higher local T in the central region are favoured with increasing E_T .

6.3.4 Event temperature fluctuations

In the previous subsections first the average (fully inclusive) temperature was discussed, and then temperatures were compared for event samples having different multiplicities or E_T . We can go a step further and define a temperature for each event by calculating the average transverse momentum of all charged particles in the final state,

$$\bar{p}_T = n^{-1} \sum_i (p_T)_i, \quad (6.26)$$

or, even preferably, of all charged particles contained within a certain region of longitudinal phase space. Since temperature is defined only for a “co-moving” region of longitudinal phase space [115], an event can have several temperatures; if we sum over tracks that are widely separated in rapidity, we risk averaging over different temperatures within one event.

It has been suggested that the thermal emission of hadrons in hadron-hadron collisions at high energy is the result of large transverse acceleration imparted to the hadron constituents [116]. For uniform transverse acceleration a_T the temperature is given by

$$kT = a_T/2\pi, \quad (6.27)$$

with k being the Boltzmann constant. Relating the transverse acceleration a_T to the closest approach of the hadrons, i.e. to the impact parameter and to the range of the force D and the constituent mass m_q in a geometrical picture,

$$\langle a_T \rangle = (\langle b \rangle D m_q / 2)^{-1}, \quad (6.28)$$

we arrive at the right order of magnitude for the temperature of the emitted particles. Expression (6.28) implies, in addition, that significant temperature fluctuations occur, depending on the impact parameter.

An investigation was carried out at the SFM detector, searching for temperature fluctuations in pp , aa , and ap interactions [10]. In trying to extract such fluctuations — which are of a dynamical origin — from real events, one has first to eliminate trivial, statistical fluctuations which result from finite numbers of particles. The latter fluctuations vanish with an increasing number of particles. More quantitatively: the dispersion of the $\bar{p}_T$ distribution for a fixed number of particles in the central region is defined as

$$d_n(\bar{p}_T) = \sqrt{\langle \bar{p}_T^2 \rangle_n - \langle \bar{p}_T \rangle_n^2}, \quad (6.29)$$

where the brackets “ $\langle \rangle_n$ ” indicate averaging over all events with the same n . The statistical contribution to $d_n(\bar{p}_T)$ must decrease as $1/\sqrt{n}$. Hence, assuming the non-trivial fluctuations do not depend on n , we may write:

$$d_n^2(\bar{p}_T) / \langle \bar{p}_T \rangle^2 = A + B/n. \quad (6.30)$$

This normalized squared dispersion is plotted in fig. 26 together with linear fits of the form (6.30) for pp ($\sqrt{s} = 63$ GeV), aa , and ap interactions. The fits are very good, and it is found that in all cases the intercept A is larger than 0. Figure 27 shows $\sqrt{A} = d_\infty(\bar{p}_T) / \langle \bar{p}_T \rangle$, as a function of $\sqrt{s}$, for all reactions studied. For pp interactions the observed effect is apparently energy-independent from $\sqrt{s} = 31$ to 63 GeV. The fluctuation $\sqrt{A}$ is lower by about 30% for ap and aa interactions than for pp interactions.

The interpretation of these observations is not straightforward. Only the energy-independence is predicted by the model [116]. The fact that the fluctuation decreases for ap and aa interactions may be attributed to multiple interactions; if in one and the same event more than one interaction occurs, the temperature fluctuation must decrease as a consequence of averaging. This effect was simulated by mixing two independent pp events and calculating $\bar{p}_T$ for the double event. The resulting intercept is as low as for ap and aa interactions. This may explain the decrease of A for aa collisions where, on the average, close to two parallel NN collisions take place, but for ap interactions the effect is larger than expected from multiple interactions.

The procedure applied to extract the non-trivial fluctuation has one (unavoidable) weak point, which becomes more obvious in the case of aa and ap interactions but is also present in pp interactions: for increasing multiplicity, the number of multiple collisions increases in aa and ap interactions, hence the non-trivial fluctuations decrease owing to averaging within one event. High multiplicity in pp interactions is usually associated with smaller impact parameters — for instance, the rise of $\langle p_T \rangle_n$ for increasing n was considered to be a consequence of decreasing impact parameters [114] — hence high multiplicity in pp interactions corresponds to higher temperatures. Therefore, to assume that A is independent of n , as was done in order to arrive at eq. (6.30), is not well justified for theoretical reasons. The experimental fact that the form (6.30) fits the data so well is a surprise. Possibly, the intercept A does not represent the full unbiased fluctuation of temperature.

6.4 Rapidity distributions

6.4.1 Fragmentation region and cosmic μ^+/μ^- ratio

In order to measure the distribution of charged particles over a wide rapidity range, we need either a movable spectrometer or a large-solid-angle magnetic detector. Except for the SFM detector, no other detector fulfils one of these requirements. The rapidity variable

$$y = 1/2 \ln [(E + p_L)/(E - p_L)] = \ln [(E + p_L)/m_T] \quad (6.31)$$

depends on the rest mass m_0 because m_T and E do. Thus, in principle, we need to know also the mass of the particle. Negative particles are dominantly pions, so we can safely assume that $m_0 = m_\pi$. Positive particles are mainly protons at high momenta and mainly pions at low momenta; it makes a big difference to the rapidity whether the proton or pion mass is used, therefore positive particles will be treated differently.

In fig. 28 the inclusive rapidity distributions

$$dn/dy = \sigma_{\text{prod}}^{-1} (d\sigma/dy) \quad (6.32)$$

of negative hadrons are shown for $aa \rightarrow h^-X$ and $ap \rightarrow h^-X$ and compared with $pp \rightarrow h^-X$ at corresponding $\sqrt{s}$. These distributions, measured by R418 [16], are corrected for acceptance, secondary interactions, and decays. All distributions have a round maximum at $y = 0$, and those for aa and ap are broader than the ones for pp collisions. The solid line represents a calculation by Capella et al. [117] in the framework of the DPM. This calculation agrees very well with the aa data; it is somewhat lower in case of the pp data.

For a better comparison with the pp data the ratios of the rapidity distributions

$$R_y(aa/pp) = [dn(aa \rightarrow h^-)/dy]/[dn(pp \rightarrow h^-)/dy], \quad (6.33)$$

and $R_y(ap/pp)$ are shown in fig. 29. The ratio rises very sharply at the highest rapidities, on both sides for aa and on the a side for ap collisions. In the central region the ratio is flat over almost four units for aa , and it rises slowly from the p side to the a side for ap .

Let us first consider the sharp rise in the a fragmentation regions. It is known that the π^+/π^- ratio in pp interactions rises from a value 1 at $x_F = p_L/p_{\max} \simeq 0$ to a value 5 at $x_F > 0.8$ [118]. In "soft" quark models this observation is connected to the fact that the p is a (uud) state and hence favours $\pi^+ (= u\bar{d})$ over $\pi^- (= d\bar{u})$ [119]; the π^+/π^- ratio at high x_F is higher than the naïvely expected factor of 2 because the u quarks carry more than twice the momentum of the d quark. The alpha particle has as many neutrons as protons. The ratio π^-/π^+ from an interacting neutron ($= ddu$) is expected to be the same as the ratio π^+/π^- from an interacting proton owing to isospin symmetry. Therefore the yields of π^- and π^+ should be equal in aa collisions at any y . In ap collisions the two yields would then be equal on the a side if the p and a fragmentation regions are well separated. Consequently, one expects the ratio to rise at high y on the a side. A crude estimate shows that probably most of the sharp rise can be explained by this isospin effect. Assume that the yield of particles produced in the a fragmentation region is proportional to the number of wounded nucleons WN_A . The yield $\sigma(N \rightarrow \pi^-)$ of negative pions produced from an average nucleon ($N = 0.5n + 0.5p$) is related to the π^+ and π^- yield from a proton:

$$\sigma(N \rightarrow \pi^-) = 0.5 [\sigma(p \rightarrow \pi^-) + \sigma(p \rightarrow \pi^+)] = 0.5\sigma(p \rightarrow \pi^-)[1 + \sigma(p \rightarrow \pi^+)/\sigma(p \rightarrow \pi^-)]. \quad (6.34)$$

Substituting y for the Feynman x_F ,

$$y \simeq \ln(x_F \sqrt{s}/m_T) = \ln x_F + \ln(\sqrt{s}/m_T), \quad (6.35)$$

using $\sqrt{s} = 44$ GeV and $m_T = 0.4$ GeV, and taking the ratio $[d\sigma(p \rightarrow \pi^+)/dx_F]/[d\sigma(p \rightarrow \pi^-)/dx_F]$ from reference [118], we obtain for the ratio

$$\begin{aligned} [d\sigma(a \rightarrow \pi^-)/dy]/[d\sigma(p \rightarrow \pi^-)/dy] &= [WN_A d\sigma(N \rightarrow \pi^-)/dy]/[d\sigma(p \rightarrow \pi^-)/dy] \\ &= 0.5 WN_A \{[1 + d\sigma(p \rightarrow \pi^+)/dy]/[d\sigma(p \rightarrow \pi^-)/dy]\}, \end{aligned} \quad (6.36)$$

the curves which are indicated in fig. 29 using $WN_A(ap) = 1.31$ and $WN_A(aa) = 4\sigma(ap)/\sigma(aa) = 1.52$.

The first observation of the large ratio between the π^- yield in the a fragmentation region and the π^- yield in the p fragmentation region has helped to answer a long-standing riddle in cosmic-ray physics. The measured cosmic μ^+/μ^- ratio at sea level is 1.27 ± 0.02 for energies from 20 to 1000 GeV. The excess is, of course, related to the dominance of protons in the primary cosmic flux, leading to more π^+ than π^- which then decay into muons. Calculations of air showers used inclusive cross-sections $pp \rightarrow \pi^\pm$ measured at the Brookhaven AGS (30 GeV) and relied on Feynman scaling for the extrapolation to higher energies. These calculations arrived at an expected μ^+/μ^- ratio of 1.57, considerably higher than the measured one. This discrepancy was only resolved after the large π^- yield at high y in aa collisions was reported at an ISR Workshop [120]. Using the known cosmic component of alpha particles and of heavier nuclei as the first new input, and the inclusive cross-sections for $pp \rightarrow \pi$ measured at the ISR as second new input, Jacob and Muraki [121] obtained a ratio $\mu^+/\mu^- = 1.44$. This ratio can be further reduced to 1.3–1.35 if we take into account the high-energy neutrons produced by the cascade of primary protons through the atmosphere. Thus the obvious discrepancies between the observed μ^+/μ^- ratio and the cascade calculations based on scaling in hadron production have been eliminated.

6.4.2 Central region

Consider now the particle ratios at $y_{cm} = 0$. The measured ratios of negative particles have the values

$$R_{y=0}(ap/pp) = 1.10 \pm 0.05 \quad \text{and} \quad R_{y=0}(aa/pp) = 1.67 \pm 0.05. \quad (6.37)$$

These values are very close to the ratios of the total average negative multiplicities given in table 8. This is not too surprising if one notices (compare fig. 29 with fig. 28) that, over the whole central rapidity range, $R_y(aa/pp)$ is almost constant, and $R_y(ap/pp)$ rises linearly; it is only at the edges that the ratio changes drastically, but the edges contribute very little to the total multiplicity. Thus the difference between the ratio of the total produced negative multiplicities $\langle n_- \rangle$ and the ratio of the negative charged-particle densities at $y_{cm} = 0$ is small. However, the predictions of most models are particularly precise and transparent for the central ratio. The general formulae for $R_{y=0}(AB/pp)$ are listed below.

WOUNDED NUCLEON MODEL:

$$R_{y=0}(AB/pp) = [\langle WN_A \rangle + \langle WN_B \rangle] / 2 = [A\sigma(pB) + B\sigma(pB)] / 2\sigma(AB). \quad (6.38)$$

This ratio is identical to the ratio of the total charged multiplicity for mirror-symmetric rapidity distributions resulting from wounded target and projectile nucleons [see eq. (6.1)]; the number of wounded nucleons, $\langle WN_A \rangle$ and $\langle WN_B \rangle$, is defined by eq. (6.7).

ADDITIVE (OR WOUNDED) QUARK MODEL:

$$R_{y=0}(AB/pp) = \langle WQ_A(AB) \rangle / \langle \nu(qA) \rangle = \sigma(qA)\sigma(qB) / \sigma(qq)\sigma(AB) \quad (6.39)$$

where $\langle WQ_A(AB) \rangle = 3A\sigma(qB)/\sigma(AB)$ is the average number of wounded quarks in A, and $\langle \nu(qA) \rangle = 3A\sigma(qq)/\sigma(qA)$ is the average number of quarks wounded in A by one single (B) quark [122].

COLOUR NEUTRALIZATION MODEL:

$$R_{y=0}(AB/pp) = \langle WN_A \rangle [1 - 0.5^{\nu(pB)}] + \langle WN_B \rangle [1 - 0.5^{\nu(pA)}]. \quad (6.40)$$

This ratio is slightly different from the one integrated over the central region, eq. (6.4) [88].

DUAL PARTON MODEL [87]:

$$R_{y=0}(AB/pp) = \langle \nu(AB) \rangle = AB\sigma(pp) / \sigma(AB). \quad (6.41)$$

This expression is only valid for asymptotically high energy; see eqs. (6.3) and (6.5).

The special formulae and numerical values for ap and aa interactions are given in table 11 and compared with the data. It is at present impossible to draw any firm conclusion from the comparison with the theoretical predictions. First, it should be recalled that all predictions suffer from an approximately 10% uncertainty due to the uncertainties of the ap and aa total cross-sections. Given this uncertainty, all predictions are consistent with the data. Secondly, the experimental value of $R_{y=0}(ap/pp)$ is believed to be somewhat low; comparing the measured average total and central multiplicity in pp interactions at $\sqrt{s} = 31, 44$, and 62 GeV, it seems that the value

for 44 GeV which forms the denominator of $R_{y=0}(ap/pp)$ is “on the high side” (see fig. 3 of reference [123] and the curve by Capella et al. [117] in fig. 28). If, however, we rather take an interpolation between the pp data at $\sqrt{s} = 31$ and 62 GeV, $R_{y=0}(ap/pp)$ would turn out to be 1.18 [124].

Even if the experimental uncertainties of the total cross-sections should one day be reduced, it will be very difficult to distinguish between the AQM and the CNM or DPM at low values of $\langle v(pA) \rangle$. This is demonstrated in fig. 30, which shows the central ratio $R_{y=0}(pA/pp)$ as a function of $\langle v(pA) \rangle$ for various data [125, 126], and which shows also the predictions of the different models discussed or mentioned above (solid and dashed lines).

6.4.3 Leading protons

It was stated at the beginning of this section that there is a difficulty in obtaining the rapidity distribution of the positively charged hadrons if the masses are not identified. However, in aa collisions this difficulty can be overcome in the following way. If we assign a proton mass to all particles and subtract the rapidity distribution of the negative particles from that of positive particles, we obtain, to a good approximation, the rapidity distribution of protons:

$$dn(aa \rightarrow p)/dy = [dn(aa \rightarrow h^+)/dy] - [dn(aa \rightarrow h^-)/dy] \quad (6.42)$$

assuming only that the yield of positive and negative mesons are equal everywhere and hence cancel in the difference. The resulting distribution is shown in fig. 31. The humps seen on both sides are composed of a narrow peak of spectator protons, expected to be at $y = 3.5$, and a wider peak of leading protons. The width ($\Delta \approx 1.5$ units) of the nucleon fragmentation region can be extracted from this figure.

It is also instructive to consider the difference between positive and negative particles as a function of p_L or Feynman $x_F = 2p_L/\sqrt{s}$ (fig. 32). Here we see spectator peaks at the edge of a very flat distribution; the experimental uncertainties in the central region ($p_L < 3$ GeV/c) are large because the difference $\Delta dn/dp_L$ is very small (of the order of 1.5%) compared with the sum $(dn^+ + dn^-)/dy$, and is indeed comparable to or even smaller than the systematic uncertainties of the acceptance in the central region. However, it is already known for pp interactions that the distribution $dn(pp \rightarrow p)/dp_L$ (after subtraction of $p\bar{p}$ pairs) is approximately constant down to low p_L [127]; this indicates that even at high energies protons are “stopped” in the c.m. system.

6.5 Correlations

6.5.1 Energy flow at small angles and central multiplicity

The two downstream calorimeters of R807 allowed measurement of the energy flow into an angular region $19 < \theta < 90$ mrad, and comparison of the energy deflected into the calorimeters from protons interacting with protons or with alpha particles, or of the energy deflected from alpha particles interacting with protons or with alpha particles [9]. A simultaneous measurement of the multiplicity in the central barrel counter hodoscope made it possible to explore the correlations between central multiplicity and forward-going energy [2].

The calorimeter has a good chance of detecting interacting nucleons out of the alpha particle; these are distributed approximately homogeneously in p_L from 0 to 15 GeV/c (see previous subsection). The detection efficiency is low for non-interacting (spectator) fragments with mass number f which carry a momentum $p_L \approx 15 \cdot f$ GeV/c; the lower p_T cut for these nuclear fragments is about $0.3 \cdot f$ GeV/c and their p_T distribution is steeper than the one of non-interacting nucleons (see Section 5). Most of the heavier fragments (d, t, ^3He) will pass through the hole in the calorimeter, but a crude estimate based on the available information given in Section 5 shows that up to 50% of the nucleons close to the spectator peak may hit the calorimeter — these consist of quasi-elastically

scattered nucleons, of spectator nucleons in the true sense, and of nucleons diffractively excited into a low-mass state. In addition to these nucleons, the calorimeters detect also mesons produced at high x_F . For a $p_T = 0.3$ GeV/c, the range in x_F covered by the calorimeter is $0.2 < x_F < 1$.

In the case of αp interactions, a different portion of the proton debris is detected by the calorimeter downstream of the incoming proton, which has a momentum of 32 GeV/c, than by the calorimeter downstream of the alpha particle with 15 GeV/c momentum per nucleon. The acceptance is worse for the leading nucleons, which are distributed approximately flat in p_T from 0 to 32 GeV/c, because their angles are, on the average, smaller by a factor of 2. For a $p_T = 0.3$ GeV/c the range in x_F covered by the calorimeter is about $0.1 < x_F < 0.5$ in this case.

Figure 33 compares the distribution of detected energy on the proton side for a $p\alpha$ interaction ($31.5 + 63$ GeV/c incoming momenta) with a pp interaction ($31.5 + 31.5$ GeV/c), and the distribution of energy on the alpha side for $p\alpha$ interactions with the one for $\alpha\alpha$ interactions ($63 + 63$ GeV/c).

The energy deposition in the proton-side calorimeter is lower for $p\alpha$ interactions than for pp interactions. This decrease must be due to a reduction either of the number of particles hitting the calorimeter or of their energy. With the given granularity of the calorimeters it was not possible to decide between these alternatives. But it was possible to examine whether the angular distribution of the forward-going particles changes. An increased lateral spread of the energy flow would result in a reduced number of particles intercepted by the calorimeter. Indeed, it was found that on the average the angular distribution is wider for $p\alpha$ than for pp . This effect is connected with the higher frequency of high-multiplicity events in $p\alpha$ interactions as compared to pp interactions. It was found that for samples of events with a central multiplicity above a certain high value, the differences between the lateral spreads for $p\alpha$ and $\alpha\alpha$ become more pronounced, whereas for low-multiplicity events the differences are smaller than on the average. Hence it was concluded that the larger lateral spread can be explained by the multiple interactions of the proton in the alpha which at the same time produce the higher multiplicity of particles in the central region, and that there is no evidence that the hadronization products interact strongly with other nucleons in the alpha particle. This conclusion agrees with the accepted view that the fast particles produced (relative to the nucleus) hadronize far outside the nucleus owing to the relativistic dilation of the final-state evolution time.

The similarity between the pulse-height distribution in the alpha-side calorimeter for αp and $\alpha\alpha$ interactions is even more striking (fig. 33b). However, the deflected energy is much larger in both cases than the one for pp interactions. Since some of the spectator fragments are detected in the calorimeters as well, the measured energy cannot be directly related to the number of interacting nucleons or to details of their hadronization. It is worth noting, however, that not only the pulse-height distribution but also the lateral distribution and the mean calorimeter module multiplicity as a function of the pulse height are very similar on the alpha side for both cases, i.e. αp and $\alpha\alpha$ interactions.

The correlation between energy deposit in one of the two calorimeters and the number of hits in the barrel counter hodoscope is displayed in fig. 34 [2]. For αp interactions, only the energy deposit on the alpha side is shown. The number of hits in the barrel counter hodoscope has not been corrected for efficiency, multiple hits, and background, so it is not a true multiplicity, but it is strongly correlated with the true multiplicity of charged tracks in the central region.

The observed correlation between energy deposit and central multiplicity can be understood qualitatively as follows. As the multiplicity increases, the energy deposit first rises because more nucleons participate. Then it decreases because the angular spread of the deflected energy becomes wider; since very many particles are detected centrally, some energy must be taken away from the forward region. As discussed earlier, the calorimeter sees, unfortunately, some of the spectator

nucleons; were it not the case, then the qualitatively expected behaviour of the correlation function just described would presumably be more pronounced.

In fig. 34a the correlation function for pp interactions ($\sqrt{s} = 31$ GeV) is shown and, moreover, the calculated convolution of two, three, and four independent pp interactions. At very large barrel multiplicity (> 24) the aa data and the fourfold pp convolution merge. This suggests that aa interactions with very high central multiplicity are very much like four independent pp interactions. It was also shown by the authors that other features—such as lateral spread and pulse-height distribution—are very much alike for aa and fourfold pp collisions at high barrel multiplicity.

A similar study was carried out for pa interactions (fig. 34b). Again, at very high multiplicity, the energy deposit on the alpha side is four times higher than for pp interactions ($\sqrt{s} = 44$ GeV). Here the pp data were scaled down by the ratio 15/22 in order to account for the fact that the nucleons out of the alpha carry only 15 GeV/c; this does not account for the different calorimeter acceptance for debris from 22 GeV and 15 GeV protons. At very large multiplicity the energy deposit is larger than for aa interactions (240 units in pa versus 180 units in aa). It is not clear whether this difference is due to the different dynamics of pa and aa events (in the first case all four nucleons are hit by one proton; in the second case the four nucleons of one alpha interact presumably with several different nucleons out of the other alpha particle); or whether it is due to a non-correspondence of the multiplicity scale. The authors of ref. 2 argue that very high multiplicity events of the pa type, where the proton interacts with all four alpha nucleons, do not contribute much to the high-multiplicity events in aa interactions because the number of these events (with barrel multiplicity > 24) is much smaller in pa than in aa interactions.

The conclusions from this study are as follows: First, very high multiplicity aa and ap events involve all four nucleons of the alpha particle. Secondly, the characteristics of the energy flow into the angular region covered by the calorimeters are very similar for aa and ap events and for four independent pp interactions. There is no evidence of any new behaviour or of the presence of collective effects. The latter statement, however, has to be weighted by the sensitivity to new phenomena of this kind of experiment which measures the energy flow into a relatively small angular region with a coarse-grained apparatus.

6.5.2 Two-particle correlations in rapidity

Many of the final-state hadrons (pions, kaons, etc.) in a strong interaction are not “directly” produced but are the decay products of resonances (or of clusters or fireballs). This experimental fact was established by detailed studies of resonance production cross-sections and of particle correlations [128]. All present models of multiparticle production incorporate this property in a natural way. The processes leading to a multiparticle final state are conceptually divided into two steps. In the first one, clusters are produced by the primary mechanism. For instance, a colour string decays into colour-neutral clusters. In the second step, these clusters decay, largely independently of each other, into the final-state hadrons.

It is natural to wonder whether in interactions of nuclei the clusters produced in the first step are any different from those produced in pp interactions. The hopes placed upon high-energy nuclear collisions is linked to the expectation that in these collisions larger volumes of nuclear matter become thermalized, i.e. larger “fireballs” are generated than in pp collisions. If the energy densities in these fireballs are large enough, the deconfinement phase transition takes place according to QCD lattice calculations.

The incentive for the careful study [12] of inclusive two-particle correlations at the SFM detector was to search for possible differences between the clusters produced in aa , ap , and pp interactions. The problem of how to prove experimentally whether a deconfinement phase transition occurred is not the subject of this report. Neither will we discuss the sensitivity of this investigation

for recognizing massive fireballs — for this undertaking more information on the decay properties of such fireballs would be needed (e.g. does the hot fireball cool down by emitting high-momentum partons or by first expanding and then emitting soft particles?). In the following, only the experimental observations will be reported, followed by a few remarks on how one could proceed in order to eventually improve the sensitivity of this search for large clusters.

It is common to study two-particle correlations as a function of the rapidity y . The normalized correlation function is defined as

$$R(y_1, y_2) = \rho^{(2)}(y_1, y_2) / [\rho^{(1)}(y_1)\rho^{(1)}(y_2)] - 1, \quad (6.43)$$

where $\rho^{(1)}(y) = \sigma_{\text{prod}}^{-1} d\sigma/dy$ and $\rho^{(2)}(y_1, y_2) = \sigma_{\text{prod}}^{-1} d^2\sigma/dy_1 dy_2$ are the charged-particle densities. The longitudinal phase-space variable y is used because the clusters produced are spread over a wide range in longitudinal phase space, whereas the occupied transverse phase space is “limited”. The short-range correlation width of roughly half a unit in rapidity corresponds to a longitudinal momentum difference $\Delta p \simeq 200 \text{ MeV}/c$ in the rest system of a cluster. This is comparable to one component of the average transverse momentum and thus reflects the same characteristic dimension (1 fm) of strong interactions.

The normalized correlation function has the advantage that the experimental efficiency cancels in the ratio. It was observed [12] that the fully inclusive $R(y_1, y_2)$ scale well for pp interactions from $\sqrt{s} = 31$ to $\sqrt{s} = 62 \text{ GeV}$ as a function of the normalized rapidities

$$\psi_i = y_i / y_{\text{inc}}, \quad (6.44)$$

where y_{inc} is the rapidity of the incoming beams in the c.m. system. Thus this variable is suitable for comparisons at different $\sqrt{s_{\text{NN}}}$. In fig. 35 the ratio $R(\psi_1, \psi_2)$ is compared for aa , ap , and pp interactions for two values of ψ_1 ($\psi_1 = 0$ and $\psi_1 = -0.5$). Pronounced differences are seen. Towards high absolute ψ_2 the correlation function does not decrease below 0 — on the alpha side for ap interactions and on both sides for aa interactions, in contrast with the pp interactions (fig. 35a). The proton side in ap interactions follows the pp correlation function. It has long been known that the so-called long-range correlations displayed by the fully inclusive two-particle correlation function can be largely explained by the fact that the sample of all events is a mixture of events with different structure, such as diffractive and non-diffractive events. If one particle is selected in a certain region of phase space, then one of the two classes of events is favoured, and the second particle is more likely to be distributed according to the event structure of this class — this causes a change of $\rho^{(2)}(y_1, y_2)$ with respect to the uncorrelated product $\rho^{(1)}(y_1)\rho^{(1)}(y_2)$ of the average particle densities. In the paper [12] it is quantitatively shown in a very general way how the “long-range correlations” result from mixing events of different structure — this was done by decomposing the total sample of events into subsamples with fixed charged-particle multiplicity.

If, however, we consider semi-inclusive correlation functions $R_n(y_1, y_2)$ for subsamples with a fixed total multiplicity n , a very pronounced narrow structure can be seen (fig. 36). This structure, known for pp interactions since a long time [129], is associated with the short-range correlations in rapidity; the width (FWHM) is around 1 unit. In order to be able to compare the width and height of the correlation peak systematically as a function of n for aa , ap , and pp interactions, a parametrization of the peak and of the background level is needed. Such a parametrization was derived from a simple cluster model. The model assumes that clusters of similar properties are produced all over the central rapidity region; each cluster decays into K charged particles ($K = 1, 2, \dots$) at rapidities y_i ($i = 1, \dots, K$), so that the y_i are uncorrelated and Gaussian-distributed with

width δ around their mean value y . The formula derived for the normalized correlation function is similar to the one found in the literature [129] but is somewhat more general:

$$R_n(y_1, y_2) = F_n \rho_n^{(1)}(y_a) / [\rho_n^{(1)}(y_1) \rho_n^{(1)}(y_2)] \times (2\sqrt{\pi}\delta_n)^{-1} \exp[-(\Delta y)^2/4\delta_n] - (1+F_n)/n, \quad (6.45)$$

where F_n is the strength defined as $F_n = [\langle K(K-1) \rangle] / \langle K \rangle^2$, with K being the charged-particle multiplicity of the cluster decay, n the total charged-particle multiplicity in the event, δ_n the width, $\Delta y = y_1 - y_2$, and $y_a = (y_1 + y_2)/2$.

The fits were done for the central rapidity region $|y_1|, |y_2| < 2$. The background level $(1+F_n)/2$ and the height of the correlation peak are correlated; an independent fit of both yielded values of F_n which agreed within 3%.

The results of this analysis are shown in figs. 37 and 38 as a function of n_- , the multiplicity of negative particles. Both δ_n and F_n first decrease and then level off for increasing multiplicity n_- . The dependence of δ_n on n_- is energy-independent for the three pp energies, but F_n shows some energy dependence. There are no visible differences for δ_n and its dependence on n_- when comparing aa and ap with pp interactions; the data points for F_n coincide, within the error bars, for ap and pp interactions, and also for aa and pp interactions except in the region $n_- \simeq 8$ in the latter comparison.

Under simplifying assumptions the cluster decay multiplicity $\langle K \rangle_n$ can be obtained from the strength F_n . The result is shown in figs. 38c,d. The average cluster decay multiplicity decreases from about 2 at low total multiplicity to 1.5 at higher multiplicity.

Two conclusions can be drawn from the results of this analysis of two-particle correlations in aa , ap , and pp interactions. First, the decrease of the width and cluster strength or cluster multiplicity with increasing multiplicity n_- has not been observed before; rather, an increase of $\langle K \rangle$ with multiplicity was once reported [129]. An increase of $\langle K \rangle$ is expected if the cluster production mechanism is n -independent, since in this case a bias in favour of low cluster multiplicity is present in low-multiplicity events and one in favour of high cluster multiplicity in high-multiplicity events. However, one can imagine mechanisms which introduce a different n -dependence: for instance, a colour string could break up in many pieces, yielding many low-mass clusters and high total multiplicity; alternatively, if it decays into less pieces these may form fewer and more massive clusters. This qualitative argument could perhaps explain the observed tendency: larger cluster width δ_n and strength F_n at lower multiplicity. A different explanation may be provided by associating low-multiplicity events with inelastic diffractive interactions; but then the interesting question is whether clusters differ in diffractive and non-diffractive interactions. One may also wonder whether the increase of cluster width and strength is connected with the small increase of $\langle p_T \rangle_n$ for decreasing n at small values of n (see ref. [7] and fig. 24). Clearly, a quantitative comparison of theoretical predictions with these recent experimental observations, which are common to aa , ap , and pp interactions, is needed.

The second conclusion concerns the comparison between pp, ap , and aa interactions. The analysis has not shown any serious, systematic difference between the three cases, and one has to conclude that cluster production mechanisms are the same. However, it should again be pointed out that a study of average features such as $\langle K \rangle$ and δ_n may be quite insensitive to new phenomena. These are more likely to show up in the tails of the distributions. For instance, it is possible that although the probability for producing very high mass clusters [130] is low in the case of pp interactions, this would significantly increase for aa interactions but still without affecting the average $\langle K \rangle$ in any significant way. A study of the distribution of K would require a different technique, e.g. "minimum spanning tree" or "nearest neighbour in y " [131]. Such a study has not yet been done on the alpha data.

6.5.3 Second-order interference of identical hadrons

Elastic scattering, which was discussed at the beginning of this report, is the classic example of first-order interference in high-energy physics. We will now deal with another interference phenomenon which is referred to under various names in high-energy physics (Goldhaber–Goldhaber–Lee and Pais effect, Hanbury-Brown and Twiss effect, Kopylov–Podgoretski effect, Bose–Einstein correlations, pion interferometry). This phenomenon [132–134] is believed to be the analogue of the second-order interference of photons [135] which was discovered by Hanbury-Brown and Twiss [136] and was used by them to measure the angular size of distant stars. First-order interference refers to the interference between phase-correlated wave function amplitudes which leads to characteristic patterns of the intensity (square of the amplitudes) in momentum space—for instance, the one of the differential elastic cross-section. Second-order interference refers to the interference between intensities when the amplitudes have random phases with respect to one another.

Assume that there are two sources (S_1 and S_2), separated by the microscopic transverse distance r , which radiate particles (photons, pions, or protons) (fig. 39); the phases of the radiation from S_1 and S_2 are uncorrelated. If the particles emitted from S_1 and S_2 are identical and coincident in time, and the two observers O_1 and O_2 cannot distinguish the sources (by the uncertainty principle), then the amplitudes must be summed. In the notation of fig. 39, only the phase due to the paths D_1 and D_2 is explicitly shown; all the rest is contained in A and B . The expectation value of the intensity observed by either one of the observers is

$$\begin{aligned} \langle I \rangle = \langle |A \exp(ikD_1) + B \exp(ikD_2)|^2 \rangle &= \langle |A|^2 + |B|^2 \rangle \\ &+ \langle AB^* \exp[ik(D_1 - D_2)] + A^*B \exp[ik(D_1 - D_2)] \rangle. \end{aligned} \quad (6.46)$$

The interference terms cancel since $\langle AB^* \rangle = \langle A^*B \rangle = 0$ because of the assumed random phase between the two sources. But the product of the two intensities has an interference term which does not vanish:

$$\langle I_1 I_2 \rangle = \langle (|A|^2 + |B|^2)^2 \rangle + 2\langle |A|^2 |B|^2 \cos[2k(D_1 - D_2)] \rangle. \quad (6.47)$$

The argument of the cosine can be expressed by the transverse separation r between S_1 and S_2 and by the angle θ between O_1 and O_2 :

$$2k(D_1 - D_2) = k\theta r. \quad (6.48)$$

One says there is second-order interference if

$$\langle I_1 I_2 \rangle \neq \langle I_1 \rangle \langle I_2 \rangle. \quad (6.49)$$

The condition for second-order interference is that the sources radiate incoherently. For sources with a fixed relative phase, one would have first-order but no second-order interference.

In fig. 39 a special and particularly simple geometry was shown — this was Cocconi's example [133]. More detailed and general expressions have been derived by Kopylov and Podgoretski and others [134]. This example explains very well the essential application of the second-order interference: from a measurement of intensity correlations of identical particles in phase space (here, $k\theta$) the space-time distance between the sources (here, r) can be determined. Whilst for photons the effect is well understood [135], there are theoretical uncertainties associated with it

when the observed particles are strongly interacting, mainly because of possible final-state interactions [137].

In many experiments at high energies the correlations between identical pions, kaons, or protons were measured to determine the dimensions of the source emitting these particles, i.e. the interaction volume. It has also been suggested that a measurement of the strength of the correlation can be used to determine the chaoticness, i.e. the amount of incoherence, of the pion-emitting source [138].

The main experimental problem is to determine the coincidence rate in the absence of second-order interference, i.e. the product $\langle I_1 \rangle \langle I_2 \rangle$ against which one compares the signal $\langle I_1 I_2 \rangle$. The problem arises because there are the other, dynamical (short range in y) correlations, already discussed, which cannot easily be separated from the second-order interference. Basically, three different approaches have been used in past experiments to determine $\langle I_1 \rangle \langle I_2 \rangle$:

- i) take the coincidence rate of non-identical particles;
- ii) reshuffle tracks in an event such that the correlations are destroyed;
- iii) combine tracks from different events.

The disadvantage of (ii) and (iii) is that also the dynamical correlations are destroyed. This is not the case for (i), but since the magnitude of the dynamical correlations is not identical for opposite and equally charged particles, even (i) does not provide the ideal background, which should contain the dynamical correlations but not the second-order interference. In all three approaches it is very important to understand the detection efficiency and momentum resolution as a function of the relative momentum of the tracks. This is the reason why measurements of this effect are very delicate.

In the publication by R807 [17] the following parametrization was used to describe the effect:

$$R = \langle I_1 I_2 \rangle / \langle I_1 \rangle \langle I_2 \rangle = 1 + \lambda [J_1(q_T r) / (q_T r)]^2 [1 + (q_L \tau c)]^{-2}, \quad (6.50)$$

where J_1 is the first-order Bessel function, q_T and q_L are the transverse and longitudinal components of $p_1 - p_2$ with respect to $p_1 + p_2$, respectively, and p_i are the three-momentum vectors of the particles. This expression has been derived [134] assuming that point-like sources are uniformly distributed over the surface of a sphere with radius r and with oscillator time τ ; an alternative interpretation of τc as the depth of the pion-emitting region was proposed by Cocconi [133]. The parameter λ measures the degree of chaoticness or incoherence of sources; if the relative phases are all random, then $\lambda = 1$.

This ratio was studied by R807 for like-sign pairs of identified pions produced in pp , $p\bar{p}$, and aa interactions. The uncorrelated background was calculated by method (iii). The effect is clearly seen in both variables (see fig. 40). The fitted parameters r^{++} , r^{--} , λ^{++} , λ^{--} , $(\tau c)^{++}$ and $(\tau c)^{--}$ are shown in fig. 41 as a function of the charged-particle multiplicity n_{ch} in the central rapidity region. It can be seen that r and τc increase for increasing multiplicity, whereas the incoherence parameter λ seems to decrease. The values obtained for positive and negative pion pairs do not differ in a statistically significant way. Moreover, the fitted parameters do not show any systematic dependence on the incident beam particles. In particular, for fixed n_{ch} no A dependence was observed.

Barshay explains the increase of r and τc for increasing multiplicity by relating increasing multiplicity with decreasing impact parameter [114]; as a consequence of decreasing impact parameter the original overlap volume increases. The fact that τc is smaller than r supports the interpretation of τc as the depth of the pion-emitting region. Knowing that most pions are not directly produced but are decay products of resonances, it may be concluded that the pion-emitting volume is not identical to the original interaction volume but a shell of depth τc and radius r which contains the decay points of the resonances.

7. HARD INTERACTIONS

7.1 The anomalous nuclear enhancement

Some of the first fixed-target experiments studying high- p_T hadron production were done with nuclear targets. In order to obtain the yield corresponding to a proton target, one (incorrectly) assumed the target mass (A) dependence to be the same as that of the total hadronic cross-section ($\propto A^{2/3}$) [36]. When the Chicago–Princeton (CP) Collaboration at FNAL performed the first systematic measurement of the A dependence of inclusive cross-sections $p + A \rightarrow h^\pm + X$ (with $h^\pm$ being an identified charged hadron $\pi^\pm$, $K^\pm$, p , or $\bar{p}$), and when they parametrized the cross-section as

$$d\sigma/dp_T \propto A^{\alpha(p_T)}, \quad (7.1)$$

they found that the power $\alpha(p_T)$ increases with p_T from about 0.8 to 1.1 (for pions) or to 1.3 (for K^- , p , or $\bar{p}$) [37]. Similar observations were reported shortly afterwards by other experimentalists [139].

Considerable theoretical activity was stimulated by these experimental findings. The rise of $\alpha(p_T)$ to a value 1 could be readily explained in the framework of parton models [140–142]. Here the basic subprocess leading to large- p_T hadron production is supposed to be large-angle elastic scattering of “hard” partons (those carrying appreciable momentum fractions of the incoming hadrons). Since deflection to large angles is rare, these partons are almost not attenuated when passing through nuclei. Thus the cross-section is expected to approach an A^1 dependence for increasing p_T .

The rise of $\alpha(p_T)$ above the value 1 for $p_T > 2$ GeV/c was harder to understand. Quite a number of possible mechanisms were proposed to explain this “anomalous nuclear enhancement”, as it was called. The proposed mechanisms fall into three categories depending on which of the three ingredients contributing to high- p_T particle production is modified by replacing the proton target by a nucleus:

- i) the structure function of the target nucleon;
- ii) the parton scattering process; and
- iii) the fragmentation or recombination of the scattered parton(s) into hadrons.

In models of category (i) the structure function of a nucleus is not simply $ZF(p) + (A-Z)F(n)$ where $F(p)$ and $F(n)$ are the structure functions of a free proton or neutron, respectively. The idea of the nucleus behaving like one big bag of quarks and gluons [143], of having collectively acting nucleon tubes [144], of density fluctuations [145], and the addition of the Fermi motion of nucleons within the nucleus to the Fermi motion of partons within the nucleon [146, 147] — all these ideas imply that the nucleus structure function differs from a simple superposition of nucleon structure functions.

Models of category (ii) explain the anomalous yield at high p_T by multiple parton scattering from different target nucleons (within one nucleus). Either the beam parton which leads to the trigger jet or trigger hadron undergoes several hard deflections [147–149]; or several beam partons are scattered simultaneously, leading to several jets [150–152]; or several partons, scattered simultaneously, recombine to the trigger hadron. This latter process may be particularly relevant for high- p_T baryon production [142]. The probability for each scattering is proportional to the thickness of the target, hence $\propto A^{1/3}$. If we multiply by the geometrical cross-section ($\propto A^{2/3}$) we obtain the multiple scattering expansion of the cross-section (up to triple scattering):

$$\sigma(\text{parton} + A \rightarrow \text{parton} + X) = C_1 A + C_2 A^{4/3} + C_3 A^{5/3} + \dots \quad (7.2)$$

Multiple hard-scattering processes can compete with the — already rare — single hard-scattering in certain p_T regions because the convolution of elementary parton–nucleon scattering probabilities, each of which decreases rapidly with increasing p_T , leads to a less rapidly falling function of p_T .

A third category of mechanisms [category (iii)] is conceivable where the fragmentation of hard-scattered partons into hadrons is modified by the presence of the nuclear medium. Hadronization of fast partons takes place mainly outside the nucleus owing to the long characteristic times involved; thus in general little A dependence is expected [153]. However, one can think of specific mechanisms which should depend on A , such as the forming of a large- p_T hadron by combining the scattered beam parton with one or more partons out of the target.

There are various possible experimental tests by which the relevance of the above listed mechanisms can be studied. Obviously the cleanest way of learning whether the nuclear structure function is responsible, is to measure the structure function using lepton probes. When recent experiments investigating the A dependence of Drell–Yan lepton pair production found consistency with a linear A dependence for a certain range of pair-masses and p_T [154], mechanisms of category (i) were usually abandoned, although the kinematic range and the precision of these experiments were limited. This was somewhat premature, as we know today, after the European Muon Collaboration published their very precise comparison of the structure functions (per nucleon) for iron and deuterium targets as a function of the Bjorken $x = Q^2/2M\nu$. The ratio of the structure functions normalized to one nucleon,

$$R = F_2(\text{Fe})/F_2(\text{D}), \quad (7.3)$$

was found to decrease from about 1.15 at low x to 0.9 at x around 0.6 [155]. An analysis of old SLAC data confirmed these results and extended the range to higher x [156]. In the light of these measurements some of the old speculations of type (i) [143] have to be reconsidered [157].

To test mechanisms of type (iii), deep-inelastic lepton scattering experiments again provide the cleanest method, which consists of comparing the fragmentation of the struck quark for targets with different mass number A . Such experiments have been performed [158]. A -dependent (attenuation) effects have been seen, but they cannot explain the anomalous nuclear enhancement.

There is no better testing ground for mechanisms of type (ii) than the one provided by the hard hadron–nucleus and nucleus–nucleus interactions themselves (although the same mechanisms enter also, as higher-order corrections, into deep-inelastic lepton interactions and Drell–Yan pair production [159]). One promising way of obtaining more information on these mechanisms is to examine details of the final-state structure, i.e. distributions of and correlations between the particles produced together with a high- p_T trigger hadron or jet. The first investigations of this kind were done at FNAL [160], comparing targets of different A .

At the CERN ISR, experiments investigating the event structure in high- p_T pp interactions have a long tradition [161, 162]. When alpha particles were stored for the first time in 1980, the three large-solid-angle detectors (experiments R110, R418, and R807) and the small-solid-angle detector (R806) were ready to take data, to measure inclusive production cross-sections, and to study the event structure for hard ap and aa collisions.

7.2 The trigger hadron

7.2.1 Inclusive cross-section as a function of p_T

All four experiments applied single hadron triggers (as opposed to jet triggers) in the central rapidity (y) region. The y -windows covered by the different experiments were: $|y| < 0.5$ (R110); $y = 0.7 \pm 0.3$ and $y = -0.7 \pm 0.3$ (R418); $|y| < 0.25$ (R806); and $|y| < 0.6$ (R807), for aa or pp collisions where the beam energies were equal in both rings. For ap collisions the rigidities p/Z of

both beams were equal, but not the momenta. The alpha beam had a momentum of 62 GeV/c (or 15.5 GeV/c per nucleon), the proton beam a momentum of 31 GeV/c. Therefore the nucleon-proton c.m. moves with a rapidity of 0.35 along the p direction in the ISR lab. system. Consequently, for ap interactions, the y-windows are shifted in all experiments by +0.35 if the rapidity of the incoming a is positive.

Experiments R418 [3] and R807 [8] measured charged hadrons; R110 [5] and R806 [1] detected electromagnetically decaying neutral hadrons π^0 and η . The measured invariant cross-sections from three of the experiments (R110 did not publish absolute cross-sections) are compiled in table 12. The cross-sections for negative particles (R418) are shown in fig. 21 and in fig. 42 together with the π^0 cross-section (R806) and the charged hadron cross-section (from R807), the latter multiplied by the ratio $h^-(h^+ + h^-)$ as measured by R418. Also shown are the cross-sections for pp collisions at the corresponding c.m. energies (per NN collision) [8, 163]; the charged hadron cross-section published by R807 for pp interactions was multiplied by the ratio $h^-(h^+ + h^-)$ as measured by the CP Collaboration [37].

The cross-sections fall over 9 orders of magnitude within the p_T range covered by the experiments—these plots are perhaps the best illustration for the large integrated luminosity obtained (and used) in the short aa and ap runs in 1980.

7.2.2 Rapidity dependence

For ap interactions, R418 observes a more rapid decrease of the cross-section on the alpha-side ($\langle y \rangle \simeq 1.00$) than on the proton side ($\langle y \rangle \simeq -0.3$) (fig. 42a). It is important to note this y dependence, not only because it plays a role in our comparison with pp data but also because it may signify some interesting physics. The y dependence becomes very significant at large p_T —in contrast with the weak y dependence observed by the British-Scandinavian (BS) Collaboration [163] for pp interactions at $\sqrt{s} = 44$ GeV. However, the observed y dependence in ap collisions is in fair agreement with the hypothesis of radial scaling [164, 165], which has recently been supported again by a measurement of the inclusive $pp \rightarrow \pi^0$ cross-section [166]. Parametrizing the y dependence with a form suggested by radial scaling $(1 - x_R)^n$, one obtains from the ratio of the ap cross-sections at $\langle y \rangle = 1$ to the one at $\langle y \rangle = -0.3$, values of $n = 4.9 \pm 0.2$ for positive hadrons (h^+) and $n = 4.0 \pm 0.4$ for h^- , approximately independent of p_T in the range 2 to 4.5 GeV/c (with $x_R = 2 p_T \cosh(y)/\sqrt{s_{NN}}$). These values are comparable to those found in pp interactions [164–166]. The question is whether the old BS Collaboration data [163] are incorrect or misleading, and the observed y dependence of the inclusive cross-sections $ap \rightarrow h$ reflects nothing but radial scaling for NN interactions; or whether the BS data are correct and the observed y dependence is a nuclear effect—perhaps the jet emitted from the alpha particle is more strongly attenuated than the one from the proton.

7.2.3 Particle composition at high p_T

The dependence of the power $a(p_T)$ (7.1) on the particle species found by the CP Collaboration [37] is equivalent to a change of the particle composition going from light to heavier nuclear targets. This is an interesting experimental effect. It has been explained by Krzywicki et al. as a result of the higher gluon scattering cross-section as compared to the quark scattering cross-section [149]. According to their QCD parton model, the nucleus will act as a gluon filter if multiple hard scattering plays a dominant role. Since there are perhaps other explanations like one which derives from differences in the final-state interactions of the produced hadrons (e.g. the kaons have smaller reabsorption cross-sections than the pions), it is desirable to extend the range of experimental data on the particle composition as a function of p_T , incoming energy, and target or projectile mass number.

Charged pions with laboratory momenta larger than 3.8 GeV/c were identified by R418 using gas Cherenkov threshold counters [3, 167]. Reconstructing the invariant mass of the two resolved decay photons from the $\eta \rightarrow 2\gamma$ and $\pi^0 \rightarrow 2\gamma$ decay, R806 measured the η/π^0 ratio in the range $2.5 < p_T < 3$ GeV/c [1]. The measured ratios of $\pi^\pm/h^\pm$ and η/π^0 are given in table 13 and compared with pp data. Most ratios in *ap* and *aa* collisions are close to the corresponding ratios in pp and pd collisions; however, for *aa*, the ratio π^+/h^+ is lower, i.e. the fraction of heavy particles, K^+ and p , is larger than for pp and pd collisions at the same $\sqrt{s_{NN}}$. A similar decrease of π^+/h^+ can be seen on the alpha-side for *ap* interactions, whereas π^-/h^- increases.

7.2.4 Comparison of cross-sections with pp data

In order to compare the yields of charged hadrons, we consider the inclusive cross-section ratios

$$R_{p_T}(aa/pp) = [d\sigma(aa)/dp_T^2]/[d\sigma(pp)/dp_T^2] \quad (7.4)$$

and $R_{p_T}(ap/pp)$ as a function of p_T in fig. 43. Cross-section ratios for π^0 production were also published by R110, which had the largest solid angle of all four detectors; in addition, since this intersection was equipped with low- β quadrupoles, the luminosity was higher by a factor of 2 than in the other intersections. Thus R110 could cover a p_T range from 4 to 8 GeV/c.

A horizontal line and a curve are drawn on each plot. The lines at $R_{p_T}(ap/pp) = 4$ and $R_{p_T}(aa/pp) = 16$ are the values that can be expected in the absence of any “shadowing” effects or of any anomalous nuclear enhancement.

The curve in fig. 43a was calculated as:

$$R_{p_T}(ap/pp) = \left[\sum_i 2^{a_i(p_T)} d\sigma(pd \rightarrow h_i)/dp_T^2 \right] / \left[\sum_i d\sigma(pp \rightarrow h_i)/dp_T^2 \right], \quad (7.5)$$

using the inclusive cross-section $d\sigma/dp_T^2$ for $pd \rightarrow h_i$ and $pp \rightarrow h_i$, and $a_i(p_T)$, the power to the mass number A , as measured by the CP Collaboration [37] ($h_i = \pi^+, \pi^-, K^+, K^-, p, \bar{p}$). This extrapolation from the A dependence measured at FNAL to the ratio $R_{p_T}(ap/pp)$ at the ISR is meaningful only if the A dependence depends but weakly on the incoming energy ($\sqrt{s_{NN}} = 44$ GeV for *ap* versus $\sqrt{s_{NN}} = 20$ and 27 GeV at FNAL). The energy dependence was measured by the CP Collaboration for the limited range of 200 to 400 GeV/c incoming momenta; $a(p_T)$ clearly depends on the energy for low p_T ($p_T < 1$ GeV/c) but for high p_T ($p_T > 3$ GeV/c) the energy dependence seems weak except for $\bar{p}$ (see fig. 19 of reference [37]). Isospin effects show up when comparing positive and negative hadron yields from H_2 and nuclear targets separately. However, these effects largely cancel if one sums over all charges, and they are absent for π^0 . Thus this extrapolation is not unreasonable.

The ratios $R_{p_T}(ap/pp)$ measured by R418 and R806 up to $p_T \simeq 5$ GeV/c are, within the given statistical and systematic errors, compatible with a factor of 4 but could be also consistent with the extrapolated curve (in the case of R418 a 10% p_T -independent cross-section normalization uncertainty is not included in the errors bars). However, the data of R110, starting from 5 GeV/c up, are significantly lower than 4 — an unexpected result. At this point it should be recalled that, owing to the motion of the *ap* c.m. none of the experiments compare *ap* with pp cross-sections at equal rapidity. The *ap* cross-section measured by R418 is compared with pp data taken by the BS Collaboration ($\sqrt{s} = 44$ GeV and $y = 0$ or $y = 0.6$ [163]); only the *ap* data at small rapidities (proton side) are used. If the *ap* cross-section at the alpha side ($y \simeq +1$) is compared with the pp data at $y = 1$, the ratio decreases to about 3 in the range $2 < p_T < 3.5$ GeV/c [3] because the alpha side cross-section decreases more quickly than the pp cross-section at $y = 1$, as discussed above.

In the case of R110 and R806, the rapidity window for pp collisions is centred around $y = 0$, but for ap it is around $y = 0.35$; yet this small difference in y can presumably not explain a decrease from an expected value of around 4.5 down to 3. Thus in conclusion, the situation for ap collisions is at present still confusing.

The ratio $R_{p_T}(aa/pp)$ is shown in fig. 43b. All four experiments have published measurements of this ratio; R110 in the range of p_T above 4 GeV/c, the other three up to about 5 GeV/c. Put together, these data are consistent with a continuous rise from about 13 at low p_T to 50 at the highest p_T . For R418, again there is an open question connected with the rapidity dependence. For the data points shown, pp data [163] at $y = 0$ ($\sqrt{s} = 31$ GeV) are used in the denominator of the ratio, whilst the aa data are taken at $y \approx 0.7$. Since for pp interactions at $\sqrt{s} = 31$ GeV the BS Collaboration found a significant decrease (by about a factor of 2 at $p_T = 2.6$ GeV/c) of the cross-section going from $y = 0$ to $y = 0.6$, the ratio of aa to pp data at equal y would show a dramatic increase above 16 already at $p_T = 2.6$ GeV/c. Are the old BS data at $y = 0.6$ too low, as suggested by R807 [8]? Or is the strong rise of the ratio at $y = 0.6$ real? The other three experiments do not have a y -problem in this case, since they compare data at equal y .

We now turn to the theoretical understanding of the combined data: a rise of $R_{p_T}(aa/pp)$ up to a value of 22 at $p_T \approx 4$ GeV/c could be explained by a parton multiple scattering model [168] (see lower curve in fig. 43b) and more recently by a QCD parton model [169], but the increase beyond this value was, at first sight, a surprise, and led to hopes that aa collisions would reveal some new phenomena. However, in a recent paper Staszek and Wilk [170] explain even the continuation of the rise between 4 and 7 GeV/c by taking into account the Fermi motion of nucleons (see upper curve in fig. 43b). Even so it is desirable to have confirmation of this result, which was perhaps the most surprising one so far obtained from the alpha runs at the ISR.

7.3 The associated secondaries

The subject of this subsection is the (charged) secondaries produced together with the high- p_T trigger hadron, their phase-space distributions, and correlations between them and the trigger particle.

The final-state structure of high- p_T pp interactions is well known, in particular in the central rapidity region, from many experimental investigations with large-solid-angle detectors at the ISR, and with the FNAL accelerator and the CERN Super Proton Synchrotron [162]. The gross features have been understood in the frame of the hard parton-parton scattering model [171] which predicts two jets of hadrons due to the fragmentation of the two scattered partons: the trigger jet of which the trigger hadron is the "leading particle", and the away jet. In addition, two jets result from the fragmentation of the remaining, non-interacting ("spectator") partons out of the incoming beams. Today, QCD provides even more quantitative descriptions of the very complicated dynamics of these interactions but, experimentally and theoretically, a number of problems still persist, requiring experiments with higher statistics, higher p_T , and better detectors.

The goal of studying high- p_T ap and aa interactions was to reveal the specific mechanism(s) responsible for the anomalous nuclear enhancement. In order to achieve this goal, one relies on the comparison with high- p_T pp interactions. The interest is in ratios or differences of distributions rather than in the distributions themselves. Thus the task of interpreting the measurements is largely simplified, compared with recent studies of the high- p_T pp interaction where the aim is to reveal the complete dynamics, i.e. the sum of all mechanisms.

The first sub-subsection will deal with the dependence of the total multiplicity distribution on the $p_T(\text{trig})$ of the trigger hadron; the second one will concentrate on the central rapidity region — the region of the trigger and away jets; the third one will show correlations between the a -fragmentation region, more specifically the spectator protons (protons! — not partons) and the charge of the trigger hadron.

7.3.1 Total multiplicity of associated particles

The distribution of the total multiplicity of negative tracks per event is shown in fig. 44a for high- p_T aa interactions with $p_T(\text{trig}) > 3$ GeV/c, and compared with minimum-bias aa interactions. These distributions were measured by R418 [172]. They are not corrected for acceptance.

There is a significant increase in the average multiplicity of large- p_T events compared with that of minimum-bias events, but the dispersions $d_- = \sqrt{\langle n_-^2 \rangle - \langle n_- \rangle^2}$ are about equal. The question is: To what degree does the distribution for large- p_T events reflect the different dynamics, for instance having four or more jets instead of two jets in minimum-bias events, or having more central collisions (smaller impact parameter) or multiple interactions? Before we examine the multiplicity distribution as a function of $p_T(\text{trig})$, a comment on the specific trigger bias of multiplicity distributions for an inclusive single-particle trigger may be appropriate.

It is by now well known that large-transverse-energy triggers with calorimeters introduce a strong bias on the multiplicity distribution [173], but less well known is the fact that also a single-particle inclusive trigger biases the multiplicity distribution. If the slope of the p_T distribution and the shape of the y distribution of particles in the trigger region does not depend on the total multiplicity N , then each event to be accepted by such a trigger acquires a weight proportional to the multiplicity of tracks in the event. For this case it can easily be shown that the multiplicity distribution for an inclusive trigger is related to $P(n)$ for minimum-bias events by

$$P_{\text{incl}}(n) = (n/\langle n \rangle)P(n), \quad (7.6)$$

independent of p_T . Indeed, the slope of the p_T distribution does not depend strongly on n at least up to $p_T \simeq 2$ GeV/c (see fig. 23), the y distribution depends only weakly on n in the central region (if we disregard the lowest multiplicity channels), and the multiplicity distributions for $P_{\text{incl}}(n)$ resemble very much the distribution of $(n/\langle n \rangle)P(n)$ although they are not quite independent of $p_T(\text{trig})$.

In fig. 44b the dependence of the average negative multiplicity $\langle n_- \rangle$ and the dispersion d_- on $p_T(\text{trig})$ is shown for aa , ap , and pp [174] interactions. The average negative multiplicity rises to about $p_T(\text{trig}) = 2$ GeV/c, but the differences between aa , ap , and pp multiplicity are weakly p_T -dependent and the dispersion remains roughly constant. The rise must be indicative of the different dynamics of large- p_T processes.

A precise study of differences of the average multiplicity of charged tracks in the central region ($|y| < 0.8$) as a function of p_T was made by R807 [18]. They find that the difference $\Delta(aa-pp) = \langle n(aa) \rangle - \langle n(pp) \rangle$ rises with p_T , and the difference $\Delta(ap-pp) = \langle n(ap) \rangle - \langle n(pp) \rangle$ is constant — compared with pp data at the same $\sqrt{s_{NN}}$. The rise and constancy correspond to the rise and constancy of the inclusive cross-section ratios $R_{p_T}(aa/pp)$ and $R_{p_T}(ap/pp)$ (see fig. 45). This correspondence is in qualitative agreement with expectations from multiple scattering models: if multiple scattering causes the rise of the cross-section, then at the same time the associated multiplicity must also increase.

7.3.2 Central rapidity region

Here we will look at several projections of the momentum distributions of associated secondaries in the central rapidity region: the distribution in azimuthal angle around the beam axis, relative to the trigger hadron, $\Delta\phi = |\phi - \phi(\text{trig})|$; the rapidity distributions of particles near the plane defined by the incoming beams and the trigger hadron; and the p_T distributions near this plane.

The $\Delta\phi$ distribution in aa collisions has been shown and compared with those in pp collisions by R807 [8] for $|y| < 0.8$ and for three p_T intervals of the secondaries (fig. 46). The particle density has maxima at 180° and 0° , which are more pronounced the higher the p_T of the selected

secondaries. These maxima are attributed to the away jet and trigger jet, respectively. The minimum around 90° is at a higher level for aa collisions (compared with pp), but the rise from the minimum to the maxima is similar. In other words, the difference between the aa and pp $\Delta\phi$ -distributions is rather uniformly distributed over $\Delta\phi$ —except for the highest p_T window where the away-side enhancement seems somewhat broader for aa interactions (see also fig. 7 of reference [8]). Such a broadening would be expected in parton rescattering models.

For the following pictures from ref. [14] cuts in $\Delta\phi$ have been applied to select the region where the away jet and trigger jet dominate—these regions coincide with the region where the momentum resolution and acceptance of the SFM detector is best. The cuts are:

$$\begin{aligned} \Delta\phi &= |\phi - \phi(\text{trig})| < 35^\circ \text{ ("towards side")}, \\ \Delta\phi &> 145^\circ \text{ ("away side")}. \end{aligned} \quad (7.7)$$

The average rapidity distributions dn/dy for secondaries with $p_T > 1 \text{ GeV}/c$ are shown in fig. 47 for ap and aa collisions. Positive values of dn/dy represent particles on the towards side, negative values those on the away side. The corresponding distributions for minimum-bias events are added for comparison. The y positions of the trigger particle which is not contained in this plot or in the following ones, are indicated by arrows. In this projection, the maximum of the $\Delta\phi$ distribution (previous figure) around 0° can be recognized as a narrow enhancement around the trigger rapidity on the towards side, and the maximum around 180° , on the away side, as a wide enhancement. The reason for this striking difference between the towards and the away side is that the rapidity of the trigger jet almost coincides with the rapidity of its leading hadron—the trigger particle whose y is fixed; however, the away jets in different events are distributed over a wide rapidity range depending on the c.m. motion of the two scattered partons relative to the c.m. of the incoming hadrons. In case of ap interactions, the towards-side enhancement is very similar to the one observed in corresponding pp interactions at $\sqrt{s} = 44 \text{ GeV}$.

For aa interactions the enhancement is lower and broader in y ; however, the statistical errors of the data points are too large to allow a more detailed investigation of this possibly interesting change of the towards-side structure; part of this change can presumably be explained by the expected higher trigger bias [175]. The trigger bias implies that the fraction of the trigger-jet momentum carried by the leading (= trigger) hadron is larger than for an unbiased jet, and that the fraction carried by the other jet-associated particles is lower. Since $\sqrt{s_{NN}}$ is lower for aa than for ap , but $\langle p_T(\text{trig}) \rangle$ is kept constant, the variable $x_T = 2p_T(\text{trig})/\sqrt{s_{NN}}$, which determines the size of the trigger bias, is larger.

After having localized the regions in y and $\Delta\phi$ where the trigger and away jets are found, we now turn to the momentum distribution of particles in this region. The momentum distribution of jet-associated particles is directly related to the parton fragmentation function which approximately scales in the variable:

$$z = \vec{p} \cdot \vec{p}(\text{jet}) / |\vec{p}(\text{jet})|^2. \quad (7.8)$$

In order to be able to use this variable, it is necessary to determine the jet axis and the parton-parton centre of mass.

In purely hadronic interactions the task of determining the away-jet axis is very difficult, in particular for the $p_T(\text{trig})$ range considered here, and therefore it has become common to use

$$x_E = -\vec{p}_T \cdot \vec{p}_T(\text{tj}) / |\vec{p}_T(\text{tj})|^2, \quad (7.9)$$

where $\vec{p}_T(tj)$ is the projection of the trigger-jet momentum onto a plane orthogonal to the beams [sometimes $p_T(\text{trig})$ of the trigger hadron is used instead of $p_T(tj)$]. This is a good approximation to z as long as the trigger jets and away jets are at large angles with respect to the beams. For the trigger jet one could do better; here the axis is known but it is only known in the laboratory, not in the parton-parton c.m. system. Therefore the same variables, p_T and x_E , were used also for the trigger jet.

By comparing the p_T and x_E distributions for aa and ap interactions with those for pp , one can test, in principle, whether there is any parton rescattering. If the parton leading to the trigger hadron rescattered once or more, there should be two or more away jets. Consequently the p_T or x_E distribution should “soften”, i.e. should fall more rapidly with increasing p_T or x_E , as compared with the case where only one away jet recoils against the trigger jet.

The p_T distributions dn/dp_T in aa and ap interactions with a $p_T(\text{trig}) > 3$ GeV/c are shown in fig. 48 for particles with $|y| < 2$ and $\Delta\phi < 35^\circ$ (negative p_T) or $\Delta\phi > 145^\circ$ (positive p_T). The p_T distribution in minimum-bias aa events is indicated as well. The distributions for the three large- p_T samples are rather similar; they fall off less rapidly than the one for minimum-bias events, on the away side. On the towards side the distributions are steeper, and those for large- p_T events are closer to the aa minimum-bias distribution in the same $\Delta\phi$ and y region.

For the purpose of comparison between aa , ap , and pp , R418 has subtracted the p_T distribution for minimum-bias events from those for high- p_T events in order to isolate the particles associated with the away jet and trigger jet. They rely on the hypothesis that the particles produced centrally in minimum-bias events are a good approximation to the “background” of particles in high- p_T events which is associated with the spectator parton fragments. It can be seen from fig. 47 that (for $p_T > 1$ GeV/c) this background is almost negligible on the away side, and it corresponds to the level of dn/dy outside the trigger-jet enhancement on the towards-side. In fig. 49b,c the ratio

$$R_{x_E}(ap/pp) = [dn(ap)/dx_E]/[dn(pp)/dx_E] \quad (7.10)$$

is shown for ap interactions with $p_T(\text{trig}) > 3$ GeV/c. The pp data sample [$\sqrt{s} = 44$ GeV, same p_T range, $\langle y(\text{trig}) \rangle = 0.7$], used for the denominator was taken with the same apparatus and the same trigger by the CERN–Dortmund–Heidelberg–Warsaw (CDHW) Collaboration (experiments R416 and R419). The ratios for the two trigger sides, $\langle y(\text{trig}) \rangle = 1.1$ and $\langle y(\text{trig}) \rangle = -0.4$, are shown separately.

In fig. 49a the x_E distributions for aa and pp ($\sqrt{s} = 31$ GeV) interactions measured by R807 are shown [with $p_T(\text{trig}) > 2.5$ GeV/c, $\langle y(\text{trig}) \rangle = 0$, and x_E defined using $p_T(\text{trig})$ instead of $p_T(tj)$ in (7.9)]. (Also shown is a pp interaction result from R413 [176].) No background was subtracted from these distributions; but by subtracting the p_x distributions with

$$p_x = x_E |\vec{p}_T(\text{trig})| \quad (7.11)$$

for high- p_T pp from high- p_T aa interactions, R807 finds that the difference

$$\Delta(dn/dp_x)_{p_T} = [dn(aa)/dp_x]_{p_T} - [dn(pp)/dp_x]_{p_T} \quad (7.12)$$

has the same shape as $[dn/dp_x(pp)]_{\text{mibi}}$ in pp minimum-bias events. Likewise the difference

$$\Delta(dn/dp_x)_{\text{mibi}} = [dn(aa)/dp_x]_{\text{mibi}} - [dn(pp)/dp_x]_{\text{mibi}} \quad (7.12')$$

has the same shape as $(dn/dp_x)_{\text{mibi}}$. However, the difference is somewhat larger for high- p_T than for minimum-bias events:

$$\Delta(dn/dp_x)_{\text{high-}p_T} = (1.5 \pm 0.3) \Delta(dn/dp_x)_{\text{mibi}}. \quad (7.13)$$

From fig. 49 it can be inferred that the momentum distributions of particles in the trigger-jet and away-jet regions are very similar in ap , aa , and pp interactions once the soft minimum-bias-like background has been subtracted. This can be used to put an upper limit on hard parton rescattering.

Let us assume that the trigger parton has rescattered once and acquired about equal p_T at each scatter. Thus two away jets are produced, each balancing about one half of the trigger-jet momentum $p_T(tj)$. Since the fragmentation function for each away jet scales in

$$x_A = \vec{p}_T \vec{p}_T(aj) / |p_T(aj)|^2, \quad (7.14)$$

where the away-jet momentum $|p_T(aj)|$ equals, by assumption, about $0.5 |p_T(tj)|$, the x_E distributions for such events fall more rapidly than if only a single away jet is produced:

$$[dn(x_E)/x_E]_{\text{resc}} = 2[dn(2x_E)/dx_E]_{\text{single}}. \quad (7.15)$$

Thus the decrease of the ratio $R_{x_E}(ap/pp)$ and $R_{x_E}(aa/pp)$ below 1 at high x_E (say, $x_E > 0.4$) measures directly the amount of rescattering of this kind. Unfortunately, reality is more complicated since, in addition to double hard scattering, also combinations of hard and soft scattering may contribute. Only a proper convolution of the parton scattering probabilities can predict the quantitative softening of the x_E distributions. This has not yet been done.

It can be seen that the sensitivity of this test for multiple scattering is of about the same size as the expected effect. The expected effect is about 20% at $p_T > 3$ GeV/c by extrapolation from the FNAL data, and less than this relying on the measured ap inclusive cross-sections; the error of the ratio $R_{x_E}(ap/pp)$ for $x_E > 0.4$ is about 20%.

7.3.3 Nuclear fragmentation region

Studies of correlations between the high- p_T trigger particle in the central rapidity region and the nuclear fragments were the privilege of the SFM detector [15] owing to its solid-angle coverage extending down to angles of 7 mrad with respect to the beams. The acceptance for spectator (non-interacting) protons out of the alpha particle is around 50%. Since they have half the rigidity of the beams, the magnetic field bends them out of the beam pipe (see Section 5).

Correlations between the trigger track and the spectator protons have been studied primarily to see whether the Fermi motion of nucleons helps in building up the p_T of the trigger hadron. This hypothesis was not supported by the data, although here again the statistics of the data limit the strength of the conclusion. However, an unsuspected, interesting effect can be seen in fig. 50, which shows the longitudinal momentum flow of the positive excess particles in hard ap interactions for four different trigger conditions. For comparison, the momentum flow in minimum-bias events is shown in each figure as a smooth curve. For a positive trigger particle on the alpha side (positive p_L) the peak of spectator protons at 15 GeV/c shrinks; for a negative trigger it increases. For a positive or negative trigger on the proton side (negative p_L) the spectator peak is unchanged relative to the one in minimum-bias interactions. A plausible interpretation of these observations is as follows. Suppose that, for a fraction of the events, the trigger particle on the alpha side contains valence quarks of the interacting nucleon. In this case a positive trigger particle is more likely to result from

an interacting proton than from a neutron; therefore there are less spectator protons. For a negative trigger the interacting nucleon is, for the same reasons, more likely to be a neutron; therefore there are more spectator protons. On the other hand, valence quarks of the alpha do not, or only rarely do, constitute the trigger particle in the proton hemisphere; therefore there is no correlation with the spectator protons for such triggers.

Here we have a way of studying the relative amount of valence-quark versus gluon involvement in large- p_T particle production, because neutron projectiles are compared with proton projectiles by counting the number of spectator protons. Of course this can be done just as well with deuteron beams in the ISR—perhaps even better. The track-finding efficiency for spectator protons has been improved by inserting a three-plane multiwire proportional chamber in each compensator magnet of the SFM detector. In addition, calorimeters for detecting spectator neutrons were installed for the run in 1983.

8. CONCLUSIONS AND OUTLOOK

The two short runs with alpha particles in the CERN ISR have provided a wealth of new data on nuclear interactions at the highest c.m. energy available at accelerators. The integrated luminosity limited the investigations to purely hadronic physics—other interesting processes such as (di)lepton or direct photon production were beyond the reach of these runs.

It is difficult to summarize in a few words what has been learnt from the measurements. It seems as if more new questions have been raised than old questions answered. This may, in part, be attributed to the exploratory nature of the runs and to a certain lack of precise theoretical predictions, in part also to the shortness and uniqueness of the runs.

In the case of elastic scattering the unresolved discrepancy between two experiments prevents any firm conclusions being arrived at concerning the size of the inelastic shadow effect.

For purely quasi-elastic interactions, the momentum distribution of protons and of deuterons was determined in the rest system of the alpha particle. The p_L distribution contains clean information on the internal momentum distribution owing to the high-energy kinematics.

Many aspects of inelastic interactions leading to particle production were studied. The most striking differences between these and pp interactions at the same c.m. energy, as seen for instance in the multiplicity or E_T distributions, can be accounted for by multiple nucleon or quark interactions. Other striking differences were qualitatively explained by the different quark composition of alphas and protons. These observations called for direct comparisons between $\alpha\alpha$ and dd interactions and between αp and dp interactions, which were consequently proposed for the second light-ion run in 1983. The test of different models was limited mainly by the uncertainties of the measured αp and $\alpha\alpha$ total cross-sections, which blurred the theoretical predictions. The differences found in pp, αp , and $\alpha\alpha$ collisions for quantities which are not “additive” in the number of nucleon or quark collisions were either insignificant (two-particle correlations) or small ($\langle p_T \rangle$). Such comparisons have demonstrated that traces of new phenomena (such as those mentioned in the Introduction) have to be searched for in fluctuations rather than in average quantities.

Perhaps the most surprising result was the measured yield of hadrons at high p_T ($p_T > 4$ GeV/c). The inclusive cross-section for $\alpha p \rightarrow \pi^0$ was found to be less than four times the one for $pp \rightarrow \pi^0$, and the one for $\alpha\alpha \rightarrow \pi^0$ is higher than 40 times the inclusive cross-section $pp \rightarrow \pi^0$ at $p_T \approx 7$ GeV/c. Both results differ from extrapolations based on the results of the FNAL experiment which discovered the anomalous nuclear enhancement. The data at lower p_T (< 5 GeV/c) lie in between the extrapolation from the FNAL experiment and the simple factor of 4 (16) for αp ($\alpha\alpha$) interactions compared with pp interactions. Studies of the properties of associated particles have so far not revealed any significant evidence for the existence of a different mechanism in high- p_T

particle production from nuclei, as compared to high- p_T particle production from nucleons. The sensitivity of these studies was just about of the same size as the expected effects; thus higher precision is needed, i.e. more events at large p_T .

The direct comparison between collisions of two compact nuclei (aa) and two loosely bound nuclei (dd), having the same c.m. energy per nucleon and the same nucleon composition, now accessible owing to the 1983 runs, should be very useful for singling out those effects which are due to multiple simultaneous and near-by nucleon interactions. The corresponding comparison between ap and dp interactions will be lacking because, unfortunately, there was no dp run (mainly as a consequence of thunderstorms).

Almost certainly the results from the second (and last) run in 1983 will answer some of the questions. The total aa , ap , and dd cross-sections will be determined in intersection I2 (R210) with an estimated precision of 5% — a considerable improvement compared to the first measurements. With a 10 times higher integrated luminosity for aa interactions and increased solid angle for single-hadron high- p_T and jet triggers in I1 (R110) and I8 (R807), the number of events at high p_T will be increased by two orders of magnitude. New triggers have been applied: a hadron pair trigger in I1, electron and photon triggers in I8, total transverse energy triggers in I1 and I8, single-hadron high- p_T triggers at small angles (10° and 20°) in I4, and diffractive triggers in I6.

In 1983 the Nuclear Science Advisory Committee in the USA recommended the construction of a colliding beam accelerator for nuclei — from protons up to uranium. The projected maximum energy to be attained is higher than that of the existing ISR: 50 to 100 GeV c.m. energy per NN collision. This recommendation was inspired by the hope that in collisions of heavy and very energetic nuclei, the QCD-predicted deconfinement transition may become detectable.

The alpha runs in the ISR in 1980 have allowed a first glimpse into this new field of interest. It was fortunate that a second run with light ions could take place before the final shutdown of the ISR.

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Table 1

Alpha and deuteron beam intensities (best values achieved in 1980 and 1983)
at various stages of the accelerating system

	α (1980)	α (1983)	d (1983)
Linac			
After 516 kV pre-injector (at gas stripper)	350 mA		
Entrance tank I (He^+ , He^{++})	250 mA		
Entrance tank II (He^{++})	12 mA		
Exit Linac (100 μs pulse)	10 mA	10 mA	13 mA
Corresponding number of ions per pulse	3×10^{12}	3×10^{12}	9×10^{12}
PS Booster			
Accepted		8 mA	10 mA
Corresponding number of ions per pulse		2.5×10^{12}	7.0×10^{12}
Accelerated		1.1×10^{12}	3.3×10^{12}
Before injection to PS		0.8×10^{12}	2.4×10^{12}
PS			
Accelerated to 26 GeV/c	30 mA	105 mA	165 mA
Corresponding number of ions per pulse	2×10^{11}	7×10^{11}	22×10^{11}
ISR			
Number of PS pulses stacked	≈ 300		
Current before acceleration	4.3 A	13.0 A	25 A
After acceleration to 31.5 GeV/c	3.8 A	11.5 A	22 A
Number of ions in one ring	4×10^{13}	12×10^{13}	46×10^{13}

Table 2

Differential elastic cross-sections

a) Results from R210 [M. Ambrosio et al., Phys. Lett. 113B(1982) 347.

$aa \rightarrow aa$ at $\sqrt{s} = 126$ GeV				$ap \rightarrow ap$ at $\sqrt{s} = 89$ GeV	
$-t$ (GeV ²)	$d\sigma/dt$ (mb/GeV ²)	$-t$ (GeV ²)	$d\sigma/dt$ (mb/GeV ²)	$-t$ (GeV ²)	$d\sigma/dt$ (mb/GeV ²)
0.056	42.4 $\pm$ 0.4	0.232	1.40 $\pm$ 0.10	0.052	97.5 $\pm$ 2.0
0.059	28.0 $\pm$ 0.4	0.239	1.20 $\pm$ 0.08	0.056	86.2 $\pm$ 2.0
0.060	23.0 $\pm$ 0.3	0.246	1.10 $\pm$ 0.06	0.060	81.1 $\pm$ 1.8
0.063	19.4 $\pm$ 0.3	0.254	0.92 $\pm$ 0.04	0.063	69.8 $\pm$ 1.7
0.069	14.0 $\pm$ 0.3	0.264	0.80 $\pm$ 0.04	0.067	60.5 $\pm$ 1.6
0.073	10.8 $\pm$ 0.3	0.271	0.64 $\pm$ 0.04	0.071	46.2 $\pm$ 1.4
0.081	8.5 $\pm$ 0.2	0.280	0.55 $\pm$ 0.03	0.075	41.2 $\pm$ 1.3
0.085	7.8 $\pm$ 0.2	0.286	0.42 $\pm$ 0.03	0.079	31.2 $\pm$ 1.1
0.090	7.4 $\pm$ 0.2	0.298	0.34 $\pm$ 0.03	0.084	24.8 $\pm$ 1.0
0.095	6.8 $\pm$ 0.2	0.310	0.28 $\pm$ 0.03	0.088	20.4 $\pm$ 0.9
0.103	6.9 $\pm$ 0.2	0.322	0.18 $\pm$ 0.02	0.099	13.8 $\pm$ 0.4
0.107	6.1 $\pm$ 0.2	0.340	0.144 $\pm$ 0.020	0.112	10.7 $\pm$ 0.7
0.111	8.1 $\pm$ 0.2	0.360	0.104 $\pm$ 0.016	0.117	8.2 $\pm$ 0.6
0.116	8.5 $\pm$ 0.2	0.382	0.050 $\pm$ 0.010	0.122	6.1 $\pm$ 0.5
0.122	8.4 $\pm$ 0.2	0.410	0.058 $\pm$ 0.010	0.127	3.3 $\pm$ 0.4
0.127	9.2 $\pm$ 0.2	0.440	0.068 $\pm$ 0.012	0.133	3.9 $\pm$ 0.4
0.132	8.7 $\pm$ 0.2	0.480	0.040 $\pm$ 0.008	0.144	1.7 $\pm$ 0.3
0.139	8.0 $\pm$ 0.2	0.514	0.028 $\pm$ 0.006	0.149	2.0 $\pm$ 0.3
0.145	7.1 $\pm$ 0.2	0.548	0.028 $\pm$ 0.008	0.155	1.3 $\pm$ 0.2
0.156	6.0 $\pm$ 0.2	0.583	0.024 $\pm$ 0.007	0.161	1.2 $\pm$ 0.2
0.166	5.5 $\pm$ 0.2	0.624	0.018 $\pm$ 0.006	0.167	1.1 $\pm$ 0.2
0.175	4.5 $\pm$ 0.2	0.676	0.018 $\pm$ 0.006	0.174	0.96 $\pm$ 0.20
0.182	4.0 $\pm$ 0.2	0.720	0.009 $\pm$ 0.003	0.180	0.46 $\pm$ 0.12
0.190	3.6 $\pm$ 0.2			0.186	0.38 $\pm$ 0.12
0.196	3.20 $\pm$ 0.14			0.193	0.25 $\pm$ 0.16
0.206	2.60 $\pm$ 0.14			0.210	0.48 $\pm$ 0.12
0.214	2.16 $\pm$ 0.10			0.220	0.80 $\pm$ 0.18
0.223	1.72 $\pm$ 0.10			0.228	1.08 $\pm$ 0.20

Table 2 (cont.)

b) Results from R418 [W. Bell et al., Phys. Lett. **117B** (1982) 131].

$-t$ (GeV ²)	$d\sigma/dt$ (10 ⁻³ mb/GeV ²)	
	$aa \rightarrow aa, \sqrt{s} = 125 \text{ GeV}$	$ap \rightarrow ap, \sqrt{s} = 88 \text{ GeV}$
0.2375		50.21 ± 9.43
0.2625	277.29 ± 31.94	112.30 ± 8.69
0.2875	160.49 ± 19.76	160.79 ± 9.38
0.3125	71.27 ± 11.45	191.91 ± 9.16
0.3375	30.04 ± 6.34	181.67 ± 8.43
0.3625	15.06 ± 3.56	139.46 ± 7.07
0.3875	4.77 ± 1.50	120.41 ± 7.08
0.4125	5.60 ± 1.52	79.94 ± 5.67
0.4375	4.83 ± 1.21	62.46 ± 4.82
0.4625	5.71 ± 1.23	38.88 ± 3.19
0.4875	7.57 ± 1.42	32.38 ± 2.87
0.5125	5.10 ± 1.07	24.18 ± 2.42
0.5375	5.95 ± 1.12	17.97 ± 1.91
0.5625	3.81 ± 0.87	9.43 ± 1.36
0.5875	3.72 ± 0.85	6.48 ± 1.17
0.6125	3.00 ± 0.73	5.24 ± 1.02
0.6375	2.53 ± 0.65	5.14 ± 1.15
0.6625	0.50 ± 0.28	2.99 ± 0.90
0.6875	1.25 ± 0.46	1.96 ± 0.70
0.7125	0.50 ± 0.28	3.39 ± 1.39
0.7375	0.76 ± 0.36	
0.7625	0.20 ± 0.18	

Table 3

Integrated cross-sections in millibarns for ap and aa interactions σ_{el} = elastic cross-section σ_{prod} = cross-section for particle production $\sigma_{q.el.}$ = quasi-elastic cross-section $\sigma_{in} = \sigma_{prod} + \sigma_{q.el.}$ $\sigma_{tot} = \sigma_{in} + \sigma_{el.}$

	$\sqrt{s_{NN}}$ (GeV)	σ_{tot} (mb)	σ_{el} (mb)	σ_{in} (mb)	σ_{prod} (mb)	$\sigma_{q.el.}$ (mb)	Ref.
ap	44	130 $\pm$ 20	20 $\pm$ 4	110 $\pm$ 20			R210 ^{a)}
	44				101 $\pm$ 10		R418 ^{b)}
	27.4	129 $\pm$ 3.1	22.9 $\pm$ 1.3	103 $\pm$ 3			FNAL ^{c)}
	23.8	131.2 $\pm$ 0.8					NA8 ^{d)}
aa	31.2	250 $\pm$ 50	45 $\pm$ 15				R210 ^{a)}
		295 $\pm$ 40					
	31.2				265 $\pm$ 26		R418 ^{b)}
	2.7	408 $\pm$ 6			262 $\pm$ 19		LBL ^{e)}
	31.2	385.4	112.9	272.5	256.9	15.6	Cracow ^{f)}

a) M. Ambrosio et al., Phys. Lett. **113B** (1982) 347.b) W. Bell et al., Phys. Lett. **117B** (1982) 131.c) A. Bujak et al., Phys. Rev. D **23** (1981) 1895.d) J.P. Burq et al., Nucl. Phys. **B187** (1981) 205.e) J. Jaros et al., Phys. Rev. **C18** (1978) 2273.f) A. Białas et al., Z. Phys. **C13** (1982) 147.

Table 4

The widths σ_T and σ_L in MeV/c obtained from Gaussian fits to the p_T and p_L distributions of nuclear fragments in quasi-elastic $\alpha\alpha$ interactions. The widths are corrected for detector resolution.
[W. Bell et al., submitted to Nucl. Phys.]

	p	d	t	^3He
σ_T	217 ± 3	236 ± 3	205 ± 4	259 ± 7
σ_L	70 ± 3	122 ± 8		

Table 5

The widths σ_L (in MeV/c) for protons in different p_T slices (in MeV/c) for quasi-elastic interactions.
[W. Bell et al., submitted to Nucl. Phys.]

	$0 < p_T < 150$	$150 < p_T < 300$	$300 < p_T < 450$	$450 < p_T < 600$
σ_L	67 ± 2	64 ± 2	81 ± 3	97 ± 3

Table 6

Cross-sections, cross-section defects, and mean numbers of collisions for various proton–nucleus and nucleus–nucleus interactions

AB	σ_{AB} (mb)	$\sigma_x = AB\sigma_{pp} - \sigma_{AB}$	$\langle \nu_{AB} \rangle = (AB\sigma_{pp}/\sigma_{AB})$
pd	74.8 ^{a)}	5.2	1.07
dd	158 ^{b)}	34	1.21
pa	129 ^{c)}	31	1.24
aa	408 ^{d)}	360	1.83
pPb	3300 ^{e)}	4780	2.45
¹⁶ O ¹⁶ O	2100 ^{f)}	10190	5.85
PbPb	13770 ^{g)}	2×10^6	150

a) G. Alberi and G. Goggi, Phys. Rep. **74**(1)(1981) 1.

b) J. Jaros et al., Phys. Rev. **C18** (1978) 2273.

c) A. Bujak et al., Phys. Rev. **D23** (1981) 1895.

J.P. Burq et al., Nucl. Phys. **B187** (1981) 205.

d) G. Bellettini et al., Nucl. Phys. **79** (1966) 609.

e) $\sigma_{tot} = 2\pi r^2(2A^{1/3} - b)^2$ with $r = 1.33$ fm and $b = 0.7$.

NB: The pp cross-sections were taken at corresponding c.m. energy.

Table 7

Topological cross-sections $\sigma(n_-) \equiv \sigma_{n_-}$ (in millibarns)
of the negative secondaries produced in ap and aa collisions,

total production cross-section ($\sigma_{\text{prod}} = \sum \sigma_{n_-}$),

and moments of the σ_{n_-} distributions

$$\langle n_- \rangle = \sum n_- \sigma_{n_-} / \sigma_{\text{prod}}, \quad C_-^k = \langle n_-^k \rangle / \langle n_- \rangle^k, \quad D_-^k = \left[\sum (n_- - \langle n_- \rangle)^k \sigma_{n_-} / \sigma_{\text{prod}} \right]^{1/k}$$

[From W. Bell et al., Phys. Lett. 128B (1983) 349.]

n_-	$ap \sqrt{s_{NN}} = 44 \text{ GeV}$	$aa \sqrt{s_{NN}} = 26.3 \text{ GeV}$	$aa \sqrt{s_{NN}} = 31.2 \text{ GeV}$
0	5.1 $\pm$ 1.1	12.9 $\pm$ 2.8	11.9 $\pm$ 2.6
1	10.0 $\pm$ 1.6	22.5 $\pm$ 3.6	22.2 $\pm$ 3.9
2	10.2 $\pm$ 1.5	25.8 $\pm$ 3.8	23.3 $\pm$ 4.0
3	10.9 $\pm$ 1.4	29.4 $\pm$ 3.9	23.9 $\pm$ 3.7
4	10.9 $\pm$ 1.4	27.2 $\pm$ 3.6	24.4 $\pm$ 3.4
5	10.8 $\pm$ 1.1	26.1 $\pm$ 3.5	24.3 $\pm$ 3.3
6	9.6 $\pm$ 1.0	22.1 $\pm$ 2.8	21.7 $\pm$ 2.4
7	7.89 $\pm$ 0.87	20.8 $\pm$ 2.6	20.0 $\pm$ 2.2
8	6.42 $\pm$ 0.75	16.5 $\pm$ 2.0	16.5 $\pm$ 1.7
9	5.59 $\pm$ 0.67	13.4 $\pm$ 1.5	14.1 $\pm$ 1.4
10	3.97 $\pm$ 0.52	11.1 $\pm$ 1.1	12.8 $\pm$ 1.2
11	3.18 $\pm$ 0.40	9.78 $\pm$ 0.94	10.4 $\pm$ 1.0
12	2.35 $\pm$ 0.32	7.71 $\pm$ 0.85	8.93 $\pm$ 0.83
13	1.51 $\pm$ 0.22	6.05 $\pm$ 0.77	7.15 $\pm$ 0.83
14	1.19 $\pm$ 0.22	3.53 $\pm$ 0.72	6.15 $\pm$ 0.97
15	0.62 $\pm$ 0.21	2.67 $\pm$ 0.68	4.55 $\pm$ 0.90
16	0.39 $\pm$ 0.15	2.02 $\pm$ 0.51	3.57 $\pm$ 0.69
17	0.30 $\pm$ 0.17	2.22 $\pm$ 0.55	2.80 $\pm$ 0.76
18	0.09 $\pm$ 0.07	1.41 $\pm$ 0.68	2.32 $\pm$ 0.62
19	0.08 $\pm$ 0.08	0.95 $\pm$ 0.60	1.35 $\pm$ 0.49
20		0.54 $\pm$ 0.43	1.27 $\pm$ 0.42
21		0.36 $\pm$ 0.38	0.82 $\pm$ 0.49
22			0.33 $\pm$ 0.21
23			0.25 $\pm$ 0.25
24			0.04 $\pm$ 0.07
σ_{prod}	101 $\pm$ 10 mb	265 $\pm$ 26 mb	265 $\pm$ 26 mb
$\langle n_- \rangle$	5.33 $\pm$ 0.13	5.82 $\pm$ 0.14	6.47 $\pm$ 0.16
C_-^2	1.46 $\pm$ 0.03	1.50 $\pm$ 0.03	1.51 $\pm$ 0.03
C_-^3	2.60 $\pm$ 0.11	2.81 $\pm$ 0.12	2.81 $\pm$ 0.12
C_-^4	5.29 $\pm$ 0.36	6.11 $\pm$ 0.44	6.05 $\pm$ 0.41
D_-^2	3.61 $\pm$ 0.07	4.12 $\pm$ 0.11	4.60 $\pm$ 0.10
D_-^3	3.22 $\pm$ 0.12	3.94 $\pm$ 0.17	4.31 $\pm$ 0.14
D_-^4	4.80 $\pm$ 0.12	5.63 $\pm$ 0.18	6.18 $\pm$ 0.14

Table 8

The ratios of average multiplicities of all negative particles produced in ap and aa to those in pp collisions at corresponding c.m. energy, and predictions of different theoretical models. The uncertainties of all predicted ratios are of the order of 10% owing to the uncertainties of the ap and aa cross-sections.

	$R_{-}(ap/pp)$	$R_{-}(aa/pp)$
Experiment R418 ^{a)}	1.20 ± 0.04	1.74 ± 0.06
WNM ^{b)}	1.18	1.62
AQM ^{c)}	1.25	1.76
DPM ^{d)}	1.30	2.06
CNM ^{e)}	1.22	1.72

a) W. Bell et al., Phys. Lett. **128B** (1983) 349.

b) A. Białas et al., Nucl. Phys. **B111** (1976) 461.

c) A. Białas et al., Z. Phys. **C13** (1982) 147.

d) A. Capella and J. Tran Thanh Van, Phys. Lett. **93B** (1980) 146.

e) S.J. Brodsky et al., Phys. Rev. Lett. **39** (1977) 1120.

Table 9

The relative numbers a_i for i wounded nucleons in ap and aa collisions. These were obtained from the best fits to the multiplicity distributions ^{a)}. The following lines show the corresponding calculated numbers.

	a_2	a_3	a_4	a_5	a_6	a_7	a_8
ap ^{a)}	1.0	0.31 ± 0.03	0.10 ± 0.02	0.013 ± 0.01			
aa ^{a)}	1.0	0.42 ± 0.06	0.32 ± 0.05	0.21 ± 0.05	0.13 ± 0.05	0.06 ± 0.03	0.012 ± 0.01
aa ^{b)}	1.0	0.304	0.310	0.160	0.070	0.010	5×10^{-4}
aa ^{c)}	1.0	0.344	0.320	0.184	0.10	0.04	5×10^{-3}
aa ^{d)}	1.0	0.420	0.359	0.239	0.145	0.065	0.017
ap ^{d)}	1.0	0.26	0.039	0.0023			

a) T. Åkesson et al., Phys. Lett. **119B** (1982) 464.

b) F. Takagi, Tohoku University preprint TU/83/264 (1983).

c) A. Białas and A. Kolawa, Z. Phys. **C22** (1984) 231.

d) E. Stenlund, private communication.

Table 10

Temperature T and average transverse momentum $\langle p_T \rangle$ obtained from fitting two different parametrizations to the experimental p_T distributions of charged particles produced in the central region ($|y| < 1$) in aa , ap , and pp collisions^{a)}.

[W. Bell et al., submitted to Z. Phys. C.]

		T (MeV)		$\langle p_T \rangle$ (MeV/c)	
		a)	b)	a)	b)
aa	$\sqrt{s_{NN}} = 31 \text{ GeV}$	155 ± 7	178 ± 7	365 ± 14	383 ± 14
pp		150 ± 6	171 ± 7	356 ± 13	369 ± 14
ap	$\sqrt{s_{NN}} = 44 \text{ GeV}$	157 ± 7	179 ± 7	368 ± 14	386 ± 14
pp		153 ± 6	175 ± 7	357 ± 12	378 ± 14

a) $dn/dp_T^2 \propto \sqrt{(p_T^2 + m^2)} \sum_{n=1}^{\infty} K_1 [n\sqrt{(p_T^2 + m^2)}/T]$ and

$$\langle p_T \rangle = \sqrt{(\pi m T/2)} \sum_{n=1}^{\infty} K_{3/2} [(nm)/T] / \sum_{n=1}^{\infty} K_2 [(nm)/T]$$

b) $dn/dp_T^2 \propto \exp(-p_T/T)$ and

$$\langle p_T \rangle = 2 T$$

Table 11

Measured and predicted ratios of charged particle densities at $y_{\text{cm}} = 0$

Experiment	$R_{y=0}(ap/pp)$	$R_{y=0}(aa/pp)$
R418 ^{a)}	1.095 ± 0.05	1.67 ± 0.05
WNM ^{b)}	$1/2 [1 + (4\sigma_{pp}/\sigma_{ap})] = 1.15$	$4\sigma_{ap}/\sigma_{aa} = 1.52$
AQM ^{c)}	$3(\sigma_{aq}/\sigma_{ap} = 1.21$	$\sigma_{aq}^2/\sigma_{qq}\sigma_{aa} = 1.71$
DPM ^{d)}	$4\sigma_{pp}/\sigma_{ap} = 1.31$	$16\sigma_{pp}/\sigma_{aa} = 1.99$
CNM ^{e)}	$1 + 2\sigma_{pp}/\sigma_{ap} - (1/2)^4\sigma_{pp}/\sigma_{ap} = 1.25$	$8\sigma_{ap}/\sigma_{aa} [1 - (1/2)^4\sigma_{pp}/\sigma_{ap}] = 1.82$

a) W. Bell et al. (CERN-Heidelberg-Lund Coll.), "Charged particle production in aa and ap collisions at the CERN ISR", preliminary data, to be submitted to Z. Phys. C.

R. Szwed, "Particle production in aa and ap collisions at the CERN ISR", invited talk, Proc. Int. Conf. on High-Energy Nuclear Physics, Balatonfüred (Hungary), 1983 (ed. J. Erö), p. 197.

b) A. Białas et al., Nucl. Phys. **B111** (1976) 461.

c) A. Białas et al., Z. Phys. **C13** (1982) 147.

d) A. Capella and J. Tran Thanh Van, Phys. Lett. **93B** (1980) 146.

e) S.J. Brodsky et al., Phys. Rev. Lett. **39** (1977) 1120.

NB: Production cross-sections used:

$\sigma_{pp} = 33 \pm 0.14$ mb [U. Amaldi et al., Nucl. Phys. **B166** (1980) 301];

$\sigma_{qq} = (1/9) \sigma_{pp}$;

$\sigma_{ap} = 101 \pm 10$ mb;

$\sigma_{aa} = 265 \pm 26$ mb [W. Bell et al., Phys. Lett. **128B** (1983) 349];

$\sigma_{aq} = 40.8$ mb [see c) above].

Table 12

Inclusive invariant cross-section $E(d^3\sigma/dp^3)$ [mb/GeV²] for hadron production in aa collisions at $\sqrt{s_{NN}} = 31.2$ GeV and ap collisions at $\sqrt{s_{NN}} = 44$ GeV

a) π^0 production at $y \approx 0$ [A. Karabarounis et al., Phys. Lett. **B104** (1981) 75].

p_T (GeV/c)	$aa \rightarrow \pi^0$	$ap \rightarrow \pi^0$
2.15	$(2 \pm 2) \times 10^{-3}$	$(9 \pm 1) \times 10^{-3}$
2.66	$(69 \pm 9) \times 10^{-4}$	$(21 \pm 4) \times 10^{-4}$
3.18	$(13 \pm 2) \times 10^{-4}$	$(45 \pm 6) \times 10^{-5}$
3.69	$(25 \pm 4) \times 10^{-5}$	$(12 \pm 2) \times 10^{-5}$
4.19	$(56 \pm 13) \times 10^{-6}$	$(25 \pm 5) \times 10^{-6}$
4.7		$(5 \pm 1) \times 10^{-6}$
5.2		$(23 \pm 6) \times 10^{-7}$

b) Charged hadron production at $y \approx 0$ [T. Åkesson et al., Nucl. Phys. **B209** (1982) 309.]

p_T (GeV/c)	$aa \rightarrow h^\pm$
2.25	$(63 \pm 2) \times 10^{-3}$
2.38	$(417 \pm 11) \times 10^{-4}$
2.55	$(232 \pm 7) \times 10^{-4}$
2.72	$(136 \pm 5) \times 10^{-4}$
2.96	$(56 \pm 2) \times 10^{-4}$
3.30	$(18 \pm 1) \times 10^{-4}$
3.64	$(62 \pm 6) \times 10^{-5}$
4.05	$(20 \pm 3) \times 10^{-5}$
4.65	$(40 \pm 8) \times 10^{-6}$
5.81	$(38 \pm 7) \times 10^{-7}$

Table 12(cont.)

c) Charged hadron production at $y \neq 0$ [W. Bell et al., Phys. Lett. 112B (1982) 271].

p_T (GeV/c)	$aa \rightarrow h^+$ $\langle y \rangle = 0.83$		$aa \rightarrow h^-$ $\langle y \rangle = 0.63$		$ap \rightarrow h^+$ $\langle y \rangle = 1.18$		$ap \rightarrow h^-$ $\langle y \rangle = 0.98$		$ap \rightarrow h^+$ $\langle y \rangle = -0.48$		$ap \rightarrow h^-$ $\langle y \rangle = -0.28$	
0.3	360 ± 4		330 ± 4		100 ± 3		100 ± 3		110 ± 3		100 ± 3	
0.5	120 ± 2		110 ± 2		35 ± 1		35 ± 1		39 ± 1		36 ± 1	
0.7	40 ± 1		35 ± 1		11 ± 0.5		10 ± 0.5		12 ± 0.5		10 ± 0.5	
0.9	15 ± 0.5		12 ± 0.4		4.1 ± 0.3		3.7 ± 0.3		4.6 ± 0.3		3.7 ± 0.2	
} 10^0												
1.1	590 ± 26		480 ± 23		150 ± 13		120 ± 12		190 ± 15		140 ± 13	
1.3	260 ± 16		190 ± 14		65 ± 8		62 ± 8		72 ± 9		53 ± 7	
1.5	110 ± 10		81 ± 8		31 ± 5		27 ± 5		34 ± 5		21 ± 4	
1.7	56 ± 6		36 ± 5		21 ± 4		12 ± 3		14 ± 3		7.6 ± 3	
1.9	21 ± 4		17 ± 3		7 ± 2		41 ± 1		7 ± 2		4.2 ± 2	
} 10^{-2}												
2.1	950 ± 40		600 ± 16		260 ± 10		160 ± 4		350 ± 10		190 ± 4	
2.3	470 ± 18		280 ± 8		120 ± 5		78 ± 2		180 ± 5		93 ± 2	
2.5	230 ± 10		140 ± 5		59 ± 3		38 ± 1.2		90 ± 2.5		47 ± 1	
2.7	110 ± 6		64 ± 3		28 ± 1.2		19 ± 0.7		48 ± 1.5		24 ± 0.7	
2.9	57 ± 3		33 ± 2		15 ± 0.8		10 ± 0.4		27 ± 1		14 ± 0.5	
} 10^{-4}												
3.1	2300 ± 180		1500 ± 93		690 ± 44		440 ± 23		1200 ± 55		680 ± 27	
3.3	1200 ± 110		640 ± 56		320 ± 27		250 ± 17		630 ± 36		330 ± 16	
3.5	700 ± 81		360 ± 40		180 ± 14		99 ± 7		320 ± 17		220 ± 10	
3.7	290 ± 48		180 ± 27		82 ± 9		59 ± 5		190 ± 12		110 ± 7	
3.9	120 ± 29		120 ± 21		37 ± 6		34 ± 4		100 ± 8		77 ± 6	
} 10^{-6}												
4.1	1000 ± 250		440 ± 120		340 ± 49		190 ± 28		580 ± 59		340 ± 34	
4.3	610 ± 190		390 ± 110		130 ± 29		120 ± 21		400 ± 46		220 ± 26	
4.5	190 ± 94		230 ± 80		97 ± 24		86 ± 17		200 ± 31		110 ± 18	
4.7	85 ± 60		100 ± 59		37 ± 14		29 ± 10		120 ± 21		61 ± 14	
4.9					15 ± 8		28 ± 10		41 ± 12		48 ± 12	
} 10^{-7}												
5.1							7 ± 5		67 ± 15		26 ± 9	
5.3							9 ± 5		30 ± 10		24 ± 8	
5.5									25 ± 9		18 ± 7	
5.7									17 ± 8		11 ± 6	
} 10^{-7}												

Table 13

Particle composition at high p_T

	aa $\sqrt{s_{NN}} = 31 \text{ GeV}$	ap $\sqrt{s_{NN}} = 44 \text{ GeV}$	pp $\sqrt{s} = 27.4 \text{ GeV}$	pd $\sqrt{s_{NN}} = 27.4 \text{ GeV}$	pp $\sqrt{s} = 44 \text{ GeV}$
π^+/h^+	$0.43 \pm 0.03^a)$	$0.54 \pm 0.02^b)$ $0.59 \pm 0.02^c)$	$0.53 \pm 0.02^d)$	$0.52 \pm 0.02^d)$	$0.60 \pm 0.03^e)$
π^-/h^-	$0.77 \pm 0.04^a)$	$0.83 \pm 0.02^b)$ $0.72 \pm 0.02^c)$	$0.75 \pm 0.02^d)$	$0.76 \pm 0.02^d)$	$0.72 \pm 0.03^e)$
η/π^0	$1.0 \pm 0.3^f)$	$0.51 \pm 0.15^f)$	$0.57 \pm 0.08^g)$ $(31 < \sqrt{s_{pp}} < 62 \text{ GeV})$		

a) W. Bell et al., Phys. Lett. **112B** (1982) 271.b) W. Bell et al., Phys. Lett. **112B** (1982) 271.c) W. Bell et al., Phys. Lett. **112B** (1982) 271.d) D. Antreasyan et al., Phys. Rev. D **19** (1979) 764.e) B. Alper et al., Nucl. Phys. **B100** (1975) 237.f) A. Karabarounis et al., Phys. Lett. **104B** (1981) 75.g) A. Karabarounis et al., Phys. Lett. **104B** (1981) 75. $3 < p_T < 4.5 \text{ GeV/c}; \langle y \rangle \approx 0.7.$ $3 < p_T < 4.5 \text{ GeV/c}; \langle y \rangle \approx 1.1.$ $3 < p_T < 4.5 \text{ GeV/c}; \langle y \rangle \approx -0.4.$ $p_T \approx 3.8 \text{ GeV/c}; \langle y \rangle \approx 0.$ $p_T = 3.5 \text{ GeV/c}; \langle y \rangle \approx 0.$ $2.5 < p_T < 3 \text{ GeV/c}; \langle y \rangle \approx 0.$ $3 < p_T < 7 \text{ GeV/c}; \langle y \rangle \approx 0.$

Figure captions

- Fig. 1 : Distribution of ISR time devoted to physics in 1980.
- Fig. 2 : The CERN particle accelerator network with the Proton Synchrotron (PS), PS Booster (PSB), Intersecting Storage Rings (ISR), and a section of the Super Proton Synchrotron (SPS).
- Fig. 3 : Acceleration in the $\beta\lambda$ and $2\beta\lambda$ mode with an Alvarez structure.
- Fig. 4 : Experiment R110. Superconducting solenoid detector, vertical cut. A and B are scintillation counter hodoscopes, DCM1–DCM3, are drift chamber modules.
- Fig. 5 : Experiment R210. a) Plan view of intersection with scintillation counter hodoscopes H_1 and small-angle hodoscopes TB. b) Exploded view of TB hodoscopes.
- Fig. 6 : Experiment R418. a) Artist's view of the Split-Field Magnet (SFM) detector with the upper iron yoke and part of the multi-wire proportional chambers (MWPC) removed. b) Horizontal cut through the detector. CM1 and CM2 are the downstream compensator dipole magnets, Č1, Č2 the threshold gas Cherenkov counters, TOF, and T1 and T2 the time-of-flight scintillation counter hodoscopes.
- Fig. 7 : Experiment R806. Liquid-argon/lead-plate sandwich calorimeters, vertical cut. SC_1 , SC_2 are scintillation counter hodoscopes.
- Fig. 8 : Experiment R807: a) Axial Field Spectrometer (AFS), vertical cut. b) One downstream calorimeter with barrel counter hodoscope surrounding the beam intersect.
- Fig. 9 : Differential elastic ap (a) and aa (b) cross-sections. Data from experiments R210 [4] and R418 [6] and calculations by Proriot et al. [58] and Alberi et al. [59].
- Fig. 10 : Contributions of the four multiple scattering terms $F_n(-t)$ to the pa elastic scattering amplitude [56]. The vertical coordinate is normalized to the pp amplitude $F_{pp}(t=0)$. Arrows indicate places of maximum destructive interference (minima).
- Fig. 11 : a) single scattering, b) double scattering, and c) double scattering with intermediate inelastic state in elastic pd scattering.
- Fig. 12 : Number of tracks (vertical coordinate) as a function of the rigidity P/Z and energy loss $\Delta E \propto Z^2$ in a scintillation counter hodoscope surrounding the beam pipe ($7 \text{ mrad} < \theta < 50 \text{ mrad}$).
- Fig. 13 : Transverse momentum distribution:
- a) Protons. The dashed line represents a Gaussian fit to the data with $\sigma_T = 225 \text{ MeV/c}$. The dash-dotted line represents the distribution expected for quasi-elastic scattering folded with Fermi motion ($\sigma = 76 \text{ MeV/c}$) and momentum resolution ($\sigma = 42 \text{ MeV/c}$); total width $\sigma_T = \sqrt{(220^2 + 76^2 + 42^2)} = 237 \text{ MeV/c}$. The dotted line corresponds to the distribution expected for spectator protons, characterized by the convolution of Fermi motion and momentum resolution $\sigma_T = \sqrt{(76^2 + 42^2)} = 87 \text{ MeV/c}$. Different curves are normalized to the same content.
 - b) Deuterons. As in (a) the lines represent a Gaussian fit ($\sigma_T = 245 \text{ MeV/c}$), pd quasi-elastic scattering ($\sigma_T = \sqrt{(147^2 + 110^2 + 78^2)} = 199 \text{ MeV/c}$), and spectator deuterons ($\sigma_T = \sqrt{(110^2 + 78^2)} = 135 \text{ MeV/c}$).
 - c) Tritons with Gaussian fit ($\sigma_T = 224 \text{ MeV/c}$) [13].
- Fig. 14 : Longitudinal momentum distribution in the rest frame of the fragmenting alpha nucleus for:
- a) Protons. The dashed line is a Gaussian fit (5.7) with width $\sigma_L = 95 \text{ MeV/c}$, average peak position $\Delta p_L = 80 \text{ MeV/c}$. The dash-dotted line is a convolution of a Gaussian with an exponential distribution (see text) yielding $\langle p_L \rangle = -36 \text{ MeV/c}$.
 - b) Deuterons. The dashed line is a Gaussian fit to the data with $\sigma_L = 179 \text{ MeV/c}$ [13].

- Fig. 15 : Schematic decomposition of normalized rapidity distributions of produced particles for pA collisions, where exactly two NN interactions took place ($\nu = \text{WNA} = 2$). Fixed-target reference system with $y_A = 0$, $y_p = Y$. The critical rapidities y_c below which “cascading” occurs, i.e. particle densities increase, are indicated.
- a) Two-fireball model [77]
 - b) Two-phase model [78]
 - c) Energy-flux cascade model [79]
 - d) Additive quark model [82]
 - f) Colour neutralization model [88].
- Fig. 16 : Multiplicity distribution of negative particles $P(n_-) = \sigma(n_-)/\sigma_{\text{prod}}$:
- a) and b) pp interactions at $\sqrt{s} = 44$ and 30 GeV; the data are from experiment R418 [11] and streamer chamber [93] using the relation $n_- = n_{\text{ch}} - 2$; the solid line represents a calculation by Chao and Pirner [94].
 - c) and d) ap interactions ($\sqrt{s_{\text{nn}}} = 44$ GeV) and aa interactions ($\sqrt{s_{\text{nn}}} = 31$ GeV); data from R418 [11], lines are predictions [94]; the curves marked (i) to (iii) represent contributions from different types of interactions, see text and fig. 17.
- Fig. 17 : Different types of multiple nucleon-nucleon (NN) interactions in nucleus-nucleus collisions:
- i) nucleon-row collision (“proton-nucleus type”);
 - ii) two or more parallel NN interactions;
 - iii) row on row collision.
- Fig. 18 : Dispersion $d_- = \sqrt{\langle n_-^2 \rangle - \langle n_- \rangle^2}$ of the multiplicity distribution versus average multiplicity $\langle n_- \rangle$ for pp interactions below ISR energy [99], at ISR energies (streamer chamber data [93] and R418 [11]) and for ap and aa interactions [11]. The line indicates the original Wroblewski fit [100] adopted to negative particles.
- Fig. 19 : Multiplicity distribution of charged particles in the central pseudo-rapidity region $|\eta| < 0.8$ for pp, ap, and aa interactions [7].
- Fig. 20 : Comparison of fits of several models to the multiplicity distributions shown in the previous figure. The vertical axis is equal to $\chi^2 = |P(n)_{\text{fit}} - P(n)_{\text{data}}|^2 / \sigma_{\text{data}}^2$ multiplied by the sign of the difference between data and fit.
- a) Fit of a KNO scaling function (6.16) to the aa and ap data.
 - b) Fit of the wounded nucleon and gluon string model to the aa data (see text) [7].
- Fig. 21 : Inclusive invariant cross-section for $aa \rightarrow h^-$ measured by R418 [3] and R807 [8], and for $aa \rightarrow \pi^0$ measured by R806 [1] compared with $pp \rightarrow h^-$ [8, 163] and $pp \rightarrow \pi^0$ [1] as a function of p_T .
- Fig. 22 : Average transverse momentum $\langle p_T \rangle$ of charged particles in the central rapidity region as a function of particle density $\Delta n / \Delta y$. Comparison of aa data at $\sqrt{s} = 125$ GeV and pp data at $\sqrt{s} = 63$ GeV (data by R418 [16, 111]) with $p\bar{p}$ data at $\sqrt{s} = 540$ GeV [110].
- Fig. 23 : Normalized transverse momentum distribution (6.24) for different multiplicity bins as a function of p_T for aa (full circles) and pp (open circles) data at $\sqrt{s} = 125$ and 63 GeV, respectively [16].
- Fig. 24 : Distribution of total transverse energy E_T of charged particles in the central rapidity region $|y| < 0.8$ for aa, pa, and pp interactions. The pp cross-sections at $\sqrt{s} = 31$ and 44 GeV were scaled up by factors of 5 and 2.5, respectively [18].
- Fig. 25 : Average multiplicity $\langle n \rangle$ and average single-particle transverse momentum $\langle p_T \rangle$ of charged tracks versus E_T for
- a) pp interactions at $\sqrt{s} = 31.5$ GeV; data compared with Monte Carlo calculation.
 - b) pa and aa interactions at $\sqrt{s} = 88$ GeV and 125 GeV, respectively [18].

- Fig. 26 : Normalized squared dispersion of the $\bar{p}_T$ distribution for n charged particles in the central region ($|x_F| < 0.3$) as a function of $1/n$.
a) pp interactions, $\sqrt{s} = 63$ GeV; b) aa interactions, $\sqrt{s_{nn}} = 31.2$ GeV; c) ap interactions, $\sqrt{s_{nn}} = 44$ GeV. The straight lines are fits according to (6.30) [10].
- Fig. 27 : Square root of the intercept A obtained from the fits in the previous figure as a function of $\sqrt{s_{nn}}$ for aa , ap , and pp interactions and for double pp events [10].
- Fig. 28 : Comparison of the rapidity distributions for negative hadrons produced in ap , aa , and pp interactions [16]:
a) and b) ap compared with pp ; c) and d) aa compared with pp .
a) and c) linear, b) and d) logarithmic vertical scale. The solid lines correspond to predictions by Capella et al. [117].
- Fig. 29 : Ratio of the rapidity distributions of negative hadrons [eq. (6.33)] produced in aa/pp and in ap/pp interactions [16]. The curves are obtained using eq. (6.36).
- Fig. 30 : Ratio of particle densities in the central rapidity region $R_{y=0}$ as a function of the mean number of inelastic collisions $\langle \nu \rangle$ defined in eq. (6.15) compared with theoretical predictions. Data from R418 [16], NA5 [125], and a compilation by Shabelsky and Shekhter [126]. Theoretical predictions: MPM [80], EFCM [79], AQM [81–83], WNM [75], and CNM [88].
- Fig. 31 : Rapidity distribution of positive particles and of the positive excess (= leading and spectator protons) in aa interactions at $\sqrt{s} = 125$ GeV [16].
- Fig. 32 : Longitudinal momentum distribution of protons in aa interactions [15].
- Fig. 33 : Distribution of energy deposited in downstream calorimeters a) on the proton side for pa interactions (31 + 62 GeV incoming energy) compared with pp interactions (31 + 31 GeV), and b) on the alpha side for ap interactions (62 + 31 GeV) compared with aa interactions and with pp interactions (15 + 15 GeV, solid line) [9].
- Fig. 34 : Mean calorimeter pulse height versus barrel counter multiplicity: a) for aa interactions ($\sqrt{s_{nn}} = 31$ GeV) compared with pp interactions ($\sqrt{s} = 31$ GeV) and with the twofold, threefold, and fourfold convolutions of pp pulse-height distributions; b) for pa interactions ($\sqrt{s_{nn}} = 44$ GeV) on the alpha side compared with pp interactions ($\sqrt{s} = 44$ GeV); the pulse heights were scaled down by a factor of 31/44 in the case of pp interactions [7].
- Fig. 35 : Comparison of the correlation function $R(\psi_1, \psi_2)$ defined by Eqs. (6.43) and (6.44) for aa and ap interactions with pp interactions (solid line) at two values of the normalized rapidity: a) $\psi_1 = 0$ and b) $\psi_1 = -0.5$ [12].
- Fig. 36 : The correlation function $R_n(y_1, y_2)$ at fixed observed negative multiplicity $n_- = 5$ for aa , ap , and pp interactions: a) $y_1 = 0$, b) $y_1 = -1.5$ [12].
- Fig. 37 : Cluster width δ_n as a function of the corrected negative multiplicity n_- for a) pp interactions at $\sqrt{s} = 30.4, 44$, and 62.2 GeV; b) ap and aa compared with pp interactions (shaded area) [12].
- Fig. 38 : Cluster strength F_n versus n_- (corrected) for: a) ap and pp interactions at $\sqrt{s_{NN}} = 44$ GeV; b) aa and pp interactions at $\sqrt{s_{NN}} = 31$ GeV. Cluster multiplicity $\langle K \rangle_n$ versus n_- : a) ap and pp at $\sqrt{s_{NN}} = 44$ GeV; b) aa and pp at $\sqrt{s_{NN}} = 31$ GeV [12].
- Fig. 39 : Experimental arrangement for second-order interference: two sources S_1 and S_2 , separated by the transverse distance r , two observers O_1 and O_2 at distances D_1 and $D_2 \gg r$, and separated by the angle θ with respect to the sources. The sources emit waves with amplitudes indicated in the sketch. Relative phases between radiation from S_1 and S_2 are contained in A and B and assumed to be random distributed.

- Fig. 40 : Ratio of the coincidence rate between pairs of positive pions from the same event to pairs from different events: a) as a function of q_T with $q_L < 0.15$ GeV, and b) as a function of q_L with $q_T < 0.15$ GeV for pp interactions. The fitted function (6.50) is shown. Fitted parameters were $r^{++} = 1.1 \pm 0.1$ fm, $(\tau c)^{++} = 0.9 \pm 0.1$ fm, and $\lambda^{++} = 0.49 \pm 0.06$ [17].
- Fig. 41 : Fit parameters obtained with parametrization (6.50) for aa , pp, and $p\bar{p}$ event subsamples with selected charged-particle multiplicity n_{ch} as a function of $\langle n_{ch} \rangle$: a) r^{++} , b) $(\tau c)^{++}$, c) r^{--} , d) $(\tau c)^{--}$, e) λ^{++} , f) λ^{--} [17].
- Fig. 42 : Inclusive invariant cross-section for a) $ap \rightarrow h^-$ measured by R418 [3] and $ap \rightarrow \pi^0$ measured by R806 [1]; b) $aa \rightarrow h^-$ [3, 8] and $aa \rightarrow \pi^0$ [1] compared with $pp \rightarrow h^-$ at the same $\sqrt{s_{NN}}$ [8, 163].
- Fig. 43 : Ratio of the cross-sections shown in the previous figure with additional data points for $p_T > 4$ GeV/c from R110 [5]. a) $R_{p_T}(ap/pp)$: the dashed line is an extrapolation from the measurement by the CP Collaboration [37]; b) $R_{p_T}(aa/pp)$: the dashed line is a prediction by Sukhatme and Wilk [168], the dotted line is one by Staszal and Wilk [170].
- Fig. 44 : a) Observed multiplicity distribution of negative tracks (uncorrected for acceptance) for aa minimum-bias interactions and for events with a high- p_T trigger hadron ($p_T(\text{trig}) > 3$ GeV/c). b) Dependence of $\langle n_- \rangle$ and d_- on $p_T(\text{trig})$ (data from R418 [172]).
- Fig. 45 : Comparison of associated multiplicity differences and cross section ratios as a function of p_T : a) $R_{p_T}(ap/pp)$; b) $\Delta(pa - pp)$; c) $R_{p_T}(aa/pp)$; d) $\Delta(aa - pp)$ [18].
- Fig. 46 : The $\Delta\phi$ distribution of associated secondaries within $|y| < 0.8$ and in three p_T intervals for pp and aa interactions with $p_T(\text{trig}) > 2.5$ GeV/c [8].
- Fig. 47 : Rapidity distributions of associated particles with $p_T > 1$ GeV/c in two $\Delta\phi$ wedges defined by (7.7) for events with $p_T(\text{trig}) > 3$ GeV/c: a) ap , $y(\text{trig}) = 1.1$; b) ap , $y(\text{trig}) = -0.3$; c) aa , $y(\text{trig}) = 0.7$. Particles on the toward side are plotted with positive values of dN/dy , those on the away side with negative values. The rapidity distribution was shifted for each event in such a way as to have $y(\text{trig})$ at an exact position corresponding to $\langle y(\text{trig}) \rangle$. The location of the trigger hadron is indicated by an arrow. The histogram represents the distribution for minimum-bias events [14].
- Fig. 48 : Distribution in p_T of the associated secondaries in the central region ($|y| < 2$) and in the two $\Delta\phi$ wedges (7.7). Negative (positive) values of p_T correspond to tracks on the toward (away) side. The small data points represent the p_T distribution for aa minimum-bias events [14].
- Fig. 49 : a) The x_E distributions on the away side for particles within $|y| < 0.8$ and $p_{out} < 0.5$ GeV/c for aa and pp interactions with $p_T(\text{trig}) > 2.2$ GeV/c (data from R807 [8]); the dashed line represents the result for pp interactions from R413 [176].
b) Ratio $R_{x_E}(ap/pp)$ of x_E distributions on the away side (positive x_E) and the towards side (negative x_E) in the central rapidity region ($|y| < 2$) for events with $p_T(\text{trig}) > 3$ GeV/c after subtraction of minimum-bias background; ap interactions with $y(\text{trig}) \simeq 1.1$.
c) The same as (b) but for ap interactions with $y(\text{trig}) \simeq -0.4$; for the pp event sample used in the denominator $y(\text{trig}) = 0.7$ [14].

Fig. 50 : Longitudinal momentum flow of positive excess particles $p_L \cdot (dn_+/dp_L - dn_-/dp_L)$ for ap collisions with a high- p_T trigger [$p_T(\text{trig}) > 3 \text{ GeV}/c$] compared with ap minimum-bias interactions (solid line):
a) positively charged trigger hadron with $y(\text{trig}) > 0$; b) $y(\text{trig}) < 0$;
c) negatively charged trigger hadron with $y(\text{trig}) > 0$; d) $y(\text{trig}) < 0$. Preliminary data from R418 [15].

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