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PARTICLE PRODUCTION IN HIGH ENERGY COLLISIONS AND THE NON-RELATIVISTIC QUARK MODEL

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ABSTRACT

"Here is a subject full of mystery, lying far from the high energy frontier. The mystery is why the non-relativistic quark model works so well, and sometimes, why it does not" [6].

The present review deals with multiparticle production processes at high energies using ideas which originate in the non-relativistic quark model. Consequences of the approach are considered and they are compared with experiment.

АННОТАЦИЯ

Настоящий обзор посвящен рассмотрению процессов множественного рождения при высоких энергиях с помощью идей берущих начало в нерелятивистской кварковой модели. Обсуждаются следствия такого подхода и их сравнение с экспериментом.

KIVONAT

Nagyenergiás sokrészecske-keltéses folyamatokat vizsgálunk olyan meggondolások segitségével, amelyek a nem-relativisztikus kvark-modellen alapulnak. Egybevetjük e közelités következményeit a kisérleti eredményekkel.

Introduction

The aim of the present paper is to investigate the multiparticle production processes in the framework of quark combinatorics. We want to discuss here the results and the difficulties of this approach as well as its future possibilities.

The quark combinatiorial calculus, which was first proposed in [1,2], is based on the hypothesis of the existence of dressed constituent quarks. The notion of quark [3] appeared in the early sixties as a mathematical expression of the SU(3) symmetry properties of the hadrons. Since then it had gone through a long way of evolution, being now the starting point of any serious attempt to create the theory of strong interactions.

Quarks as objects existing inside the hadrons were considered first in the constituent quark models /see [4] and several papers following it/. Further, it turned out, that in the framework of the quark model not only the hadron spectra can be obtained, but - using the impulse approximation - hadron-hadron collision processes at high energies can also be handled. The investigation of hard processes /such as deep inelastic scatterings of electrons, muons and neutrino on nucleons, $\mu^+ \mu^-$ production with large effective masses in hadron collisions, e^+e^- annihilations into hadrons/ led also to the quark structure of hadrons. The quantitative description of these processes on the basis of the parton hypothesis [5] required the introduction of point-like objects, the symmetry properties of which coincided with those of the constituent quarks.

Nowadays physical theories in general, and the theory of strong interactions /quantum chromodynamics/ in particular are constructed on the basis of non-abelian gauge theories. In QCD the phenomenon of asymptotic freedom provides the field theoretical explanation of the parton model, and, on the other hand, gives a possibility to calculate the deviations from the latter one. There is a serious hope that in the framework of QCD one will be able to observe the confinement of quarks. At the same time the increase of the effective charge in the infrared region means, that at large distances one has to deal with all the problems connected with strong interactions. That's why we think that it is reasonable to describe soft processes, i.e. processes at large distances, in a different, semi-phenomenological way, which gives a good agreement with the experimental data and is based, to a certain extent, on the features of the exact theory. Such a point of view is quite frequent and at present there exist several attempts to realize it in fields where the perturbational approach of QCD is not applicable. First of all this concerns hadron spectroscopy, where the introduction of dressed /constituent/ quarks is a great success /see, e.q. [6-11]/. The consideration of relativistic or non-relativistic constituent quark models with different types of potentials in here equally possible; the common feature of these models is the introduction of sufficiently massive quarks /e.q. $M_u \approx M_d \approx 360$ MeV, $M_s/M_d \simeq 1,5$ [6]/. The same masses lead to very reasonable values for the baryon magnetic momenta [6,12,13] and even for the magnetic momenta of radiational vector meson decays $V \to P + \gamma$ [14-16]. Not everything fits well, however /e.g. extracting the mass values in the framework of the non-relativistic quark model from the measured magnetic momenta of Σ and Ξ one gets $M_s < M_n$/, but for the time being it is impossible to tell, if that is a failure of the constitutent quark model or if the discrepancies can be eliminated by some corrections which one can calculate remaining within the model /see [6,12,13,17]/.

The situation is not less mysterious in the field of hadron interactions at high energies [18-20]. The additive quark model gives, in good agreement with experiment, the ratio of the total cross section in NN and ΠN scatterings as $\sigma_{tot}(NN)/\sigma_{tot}(\Pi N)=3/2$. The measurements of the total cross sections of strange hyperons [21,22] lead to results fulfilling the predictions of the model $\sigma_{tot}(pp)-\sigma_{tot}(\Lambda p)\simeq\sigma_{tot}(pp)-\sigma_{tot}(\Sigma p)\simeq\frac{1}{2}(\sigma_{tot}(pp)-\sigma_{tot}(\Xi p))$ with the accuracy of 20%. Again, there are some confusing facts. For example, the experimental value of $\sigma_{tot}(NN)/\sigma_{tot}(\Pi N)$ differs from 3/2 by 10% only; it is, however, larger than the predicted value, and that - if we take all the predictions of the model literally - means, that the double scatterings of the constituent quarks lead not to shadowing but to antishadowing effects. Meanwhile, investigations of the elastic pp-scattering in the framework of the constituent quark model show that the observed minimum in $\frac{d\sigma}{dt}$ at $|t|\sim 1,4$ can be very well described by shadowing effects connected with the double scatterings of quarks [23-25].

The mentioned problems are of minor importance from the point of view of the present review. It is clear, that the additive quark model which is based on the introduction of constitutent quarks is able to catch as a whole the physics of the discussed phenomena in the description of both the static features of the hadrons and the hadron collisions at high energies. Note that the additivity of the model carries much more information about the hadron structure investigating high energy collision processes than considering static properties of hadrons. The additivity in the hadron scattering processes can be understood assuming a hadron picture due to which the constituent quarks are separated in space inside the hadrons. On the other hand, it is rather difficult to give an explanation for this additivity in the framework of bag models or models in which the hadron collision processes are connected with long-range color interactions.

The main question is, of course, whether additivity really can be observed in high energy hadron collisions, i.e. if in fact the impulse approximation holds. The most convincing arguments in favour of such a picture are given by the investigations of hadron-nucleus collisions at high energies [26-35]. The additive quark model enables us to calculate the ratios of the multiplicities of the secondary particles in the central region [26,27,29] and the ratios of inclusive cross-sections in the fragmentation region [30] in the πA and pA collisions without any parameters besides the quark-nucleon cross section $\sigma_{inel}(qN) \simeq \frac{1}{3}\,\sigma_{inel}(NN)$ and the density of nucleons in the nuclei. All the obtained relations are well satisfied by the experimental data. Note here, that the experimental ratio $\frac{pA\to baryons}{pp\to baryons}$ is in good agreement with the predicted value at $x = \frac{p}{p_{max}} = \frac{2}{3}$. This means, that in the collision process we in fact observe constituent quarks which carry the mass and the momentum of the proton.

Assuming such a nucleus-like structure of hadrons, several questions arise. Indeed, one has to see how the picture of hadrons formed by quark-partons and gluons can be consistent with the existence of spatially separated dressed quarks inside the hadrons. It has to be cleared if the introduction of relatively small quarks is not in contradiction with the quantum chromodynamical understanding of the nature of strong interactions. One has to see, finally, whether the picture of hadrons consisting of quasifree /i.e. almost real/ quarks is self-consistent at all knowing that real quarks do not exist as observable objects. For the time being our knowledge of quantum chromodynamical forces at large distances is not sufficient to give definite answers to these questions. However, the existence of spatially separated constituent quarks inside the hadrons can be reconciled with the parton picture assuming that a fast moving hadron is a system of three /or two/ spatially separated clouds of partons, each containing a valence quark, a sea of quark-antiquark pairs and gluons. Such a parton structure of hadrons was proposed in [36,37]; in [38] parton clouds corresponding to the dressed constituent quarks are named valons.

The existence of relatively small constituent quarks is connected with the question, if QCD can produce besides the radius of hadron /or the radius of quark confinement/ another, much smaller characteristic size, the radius of the constituent quark. This is, however, possible, if the characteristic size of the constituent quarks are determined mainly by gluon states in the t-channel [39]; $r_q \sim 1/M_g$ ($M_g \sim 2$-3 GeV[40]) is in this case much less than the hadron radius.

The introduction of almost real constituent quarks is by no means in contradiction with the lack of real, observable quarks. One can, for example, consider the amplitudes of quark processes as spectral diagrams over the quark masses [41].

Hence, having justified the picture of hadrons with spatially separated quarks inside them, we can present the approach of quark combinatiorial calculus which enables us to handle soft processes. In the framework of this approach two main assumptions are made. The first one concerns the spectator mechanism [18], which is based on the described hadron picture, and which is responsible for the fragmentational production of hadrons in hadron-hadron and hadron-nucleus collision processes. The spectator mechanism, however, can be proven directly only in experiments in which total cross sections are measured. To see all the consequences one has to translate the quark language into the hadron language, i.e. to determine the way hadrons are formed. The second assumption of our approach is connected with this question: we presume, that the production of secondary hadrons which is due to quarks joining each other obeys simple statistical rules [1,2]. The constituent quarks form hadrons independently of their spin and isospin states, of their flavours and of the fact if quarks or antiquarks are joining each other. It is convenient to obtain the different statistical relations with the help of a combinatorial calculus. The quark combinatorial calculus leads to several predictions, which, as it will be seen, are fulfilled well enough by experiment to consider the picture underlying our approach to be true as a whole. The deviations from the predictions will help us to understand the details of the picture.

II. Dressed quarks, quark structure of hadrons

Let us remind the well-known arguments supporting the impulse approximation in hadron collision processes at high energies. Comparing theoretical predictions with the experimental data, it turned out, that the processes

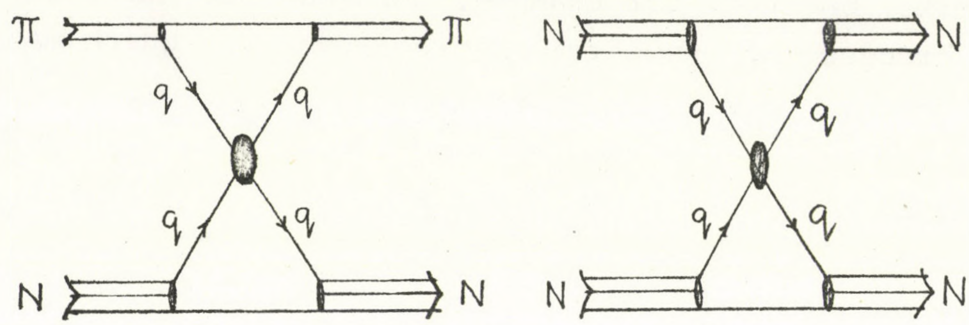


Fig. 1.

described sufficiently well the ratio of the total cross sections in NN and ΠN scattering [6-8].

$$\frac{\sigma_{tot}(NN)}{\sigma_{tot}(\Pi N)} = \frac{3}{2} \qquad (1)$$

as well as the decrease of the elastic pp-cross-section with the increase of the momentum transfer [7].

$$\frac{d\sigma_{pp\to pp}(t)}{dt} \sim F_p^4(t) \qquad (2)$$

where $F_p(t)$ is the proton form-factor.

Accepting the hadron picture with two radii, we assume, that hadrons are similar to light nuclei: the meson, consisting of a quark and an antiquark sufficiently far from each other reminds the deuteron while the baryon contains three constituent quarks in the same way as H_3 or He_3 is build up. The constituent quarks are surrounded by their "coat" of virtual particles. The radius of this "coat" is in fact the radius of the constituent quark. The mean distances between the constituent quarks determine the size of the hadron [24,36,42].

The radius of the constituent quark can be estimated from the total hadron-hadron cross-section, which, as it follows from Fig.1., can be expressed in terms of the total quark-quark cross-section. At moderately high energies $\sigma_{tot}(qq) \sim \frac{1}{9}\sigma_{tot}(qN) \simeq 4{,}5$ mb. Assuming, that the total quark-quark cross-section is determined by the geometrical sizes of the colliding quarks $\sigma_{tot}(qq) \simeq 2\Pi(2r_q)^2$ we obtain

$$r_q^2 \simeq 0{,}5 \ \mathrm{GeV}^{-2} .$$

There is another way of obtaining the radius of the constituent quark in the framework of the parton hypothesis. Without going into details, we give here only the results: due to the latest experimental result at Fermilab

$$r_q^2 \simeq 3\alpha_p' \simeq 0{,}45 \ \mathrm{GeV}^{-2}.$$

Hence, having $R_h \simeq 17 \ \mathrm{GeV}^{-2}$

$$r_q^2/_{R_h^2} \simeq 1/30$$

We consider here, naturally, coloured quarks. Since the quark confinement is due to the colour forces, we are bound to accept the following hadron picture. /In the following we consider a nucleon/. At large momenta /but $P < 10^8$ GeV/c/

the nucleon contains three clouds of quark-partons /Fig. 2a/.

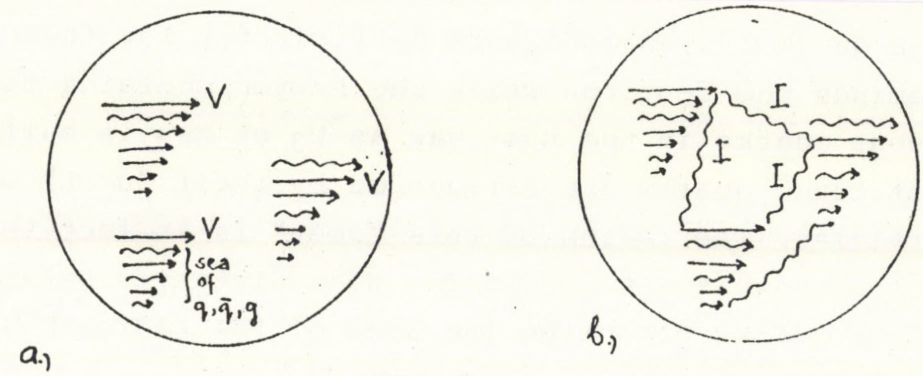


Fig.2.

Each of the clouds contains a coloured quark-parton which carries the quantum numbers of the constituent quark, and a sea of quark-antiquark pairs and gluons, which is colourless and has zero quantum numbers. The gluon interaction which keeps the constituent quarks inside the hadrons is taking place between the fast parton components /I/ [43]. The gluon exchange is inprobable between the partons carrying a relatively small fraction of the momentum /II/.

The transverse dimension of a cloud increases with the energy as $\sqrt{\alpha_p' \ln P/P_O}$, $P_O \sim 10$ GeV/c. Up to $P \ll 10^8$ GeV/c r_q remains essentially less than R_h, and, practically, the three /or in the case of a meson, two/ clouds do not overlap. When a fast hadron collides with the target, only one of the constituent quarks participates in the interaction; the other constituent quarks, or quark-parton clouds, remain spectators. The situation is different in the case of a hadron-nucleus interaction, i.e. when the target in large, and not only one, but two or three constituent quarks of the incident hadron can interact. We will come to this question later. As soon as $r_q^2 \ll R_h^2$, repeated collisions of the quarks are not probable. The interaction with the target is due to the slow components of the partons /a parton carrying energy E needs a time of the order of $\tau \sim \frac{E}{\mu^2}$ to interact/. The quark-parton cloud the slow component of which participated in the interaction breaks into partons. These partons then, interacting with each other, obtain their own "coats" and become constituent quarks, giving rise to the production of new particles /Fig.3./.

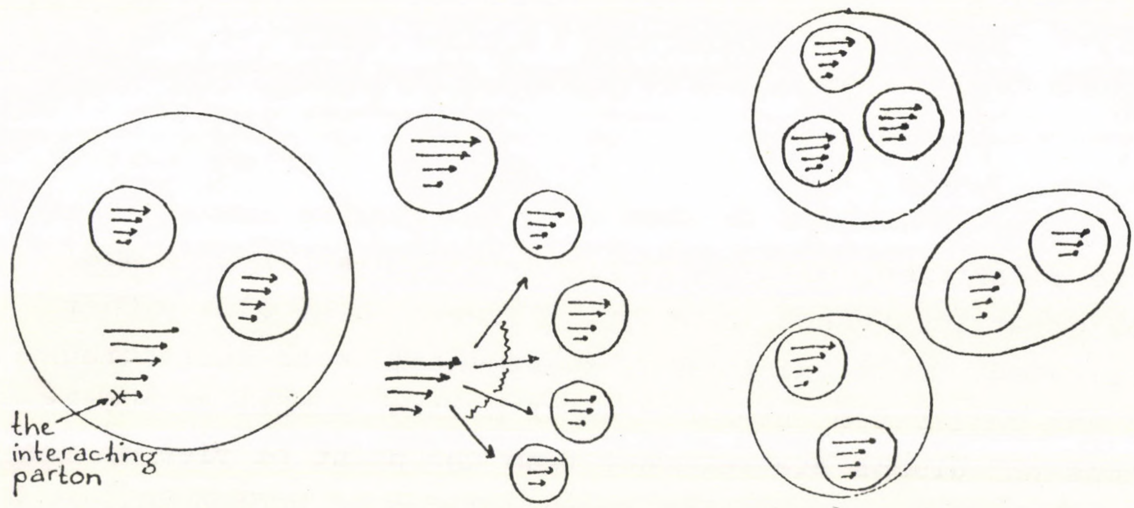


Fig.3.

The approach we are presenting deals in fact with the second and the third steps: the interactions of the partons and the gluons with each other which lead to the formation of constituent quarks, and the transition of theses constituents into hadrons /mesons, baryons, meson and baryon resonances/ in such a way that the set of hadron states corresponds to the states of the constituent quarks in the multiperipheral ladder. This approach is by no means the only possibility to handle the problem of the quark-hadron transfer. Very popular is recently the recombination model. Here the recombination of the quark-partons into the observable hadrons is investigated neglecting the intermediate states of this process like constituent quarks and resonances. This approach leads to an impressive agreement between the π^+/π^- ratio in proton collisions and the ratio of the proton structure functions $u(x)/d(x)$ measured in deep inelastic lepton-nucleon interactions. It is not clear, however, how the obtained pion spectra are connected with the spectra of resonances /ρ,ω etc./ the decays of which are relevant from the point of view of the spectra of long living particles.

In [38] the recombination of the quark-partons into hadrons is investigated introducing in the last step the "dressed" quarks /i.e., in our language, constituent quarks/. That means, an attempt is made to calculate the distribution of the constituent quarks. To find such a distribution would be of great importance; however, it seems to be also rather complicated. There exist some experimental facts indicating that the collective interactions of a large number of quarks-partons and gluons are relevant from the point of view of the formation of the constituent quark spectra in hadron collisions. In other words, the coherence of the initial state of partons and gluons plays an important role.

Until now we underlined the similarity between the structure of the systems of spatially separated quarks and the structure of light nuclei. There is, however, a very serious difference between them, which is connected with the fundamental property of the quark systems: the phenomenon of quark confinement. This can be seen very distinctly in the hadron diffraction dissociation processes. Let us compare in the following the diffraction dissociation of the deuteron and that of the meson. If a nucleon of the deuteron is elastically scattered on a target, three types of processes are possible /see Fig.4/. After the collision the nucleons might interact and form again a deuteron (a); they might interact and not form a deuteron (b) and finally, they might not interact after the collision (c) in the final state.

If one of the constituent quarks of a meson is scattered /Fig.5/, the quark-antiquark system can form in the final state either a meson analogous to the initial one, or an

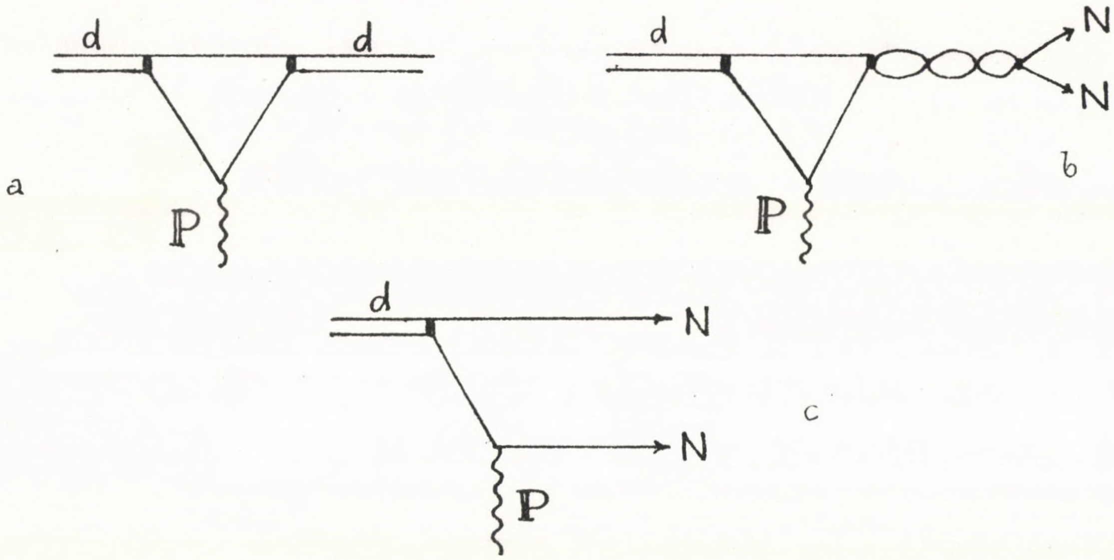


Fig. 4. The process of diffraction dissociation of the deuteron at high energies /P denotes the pomeron/.

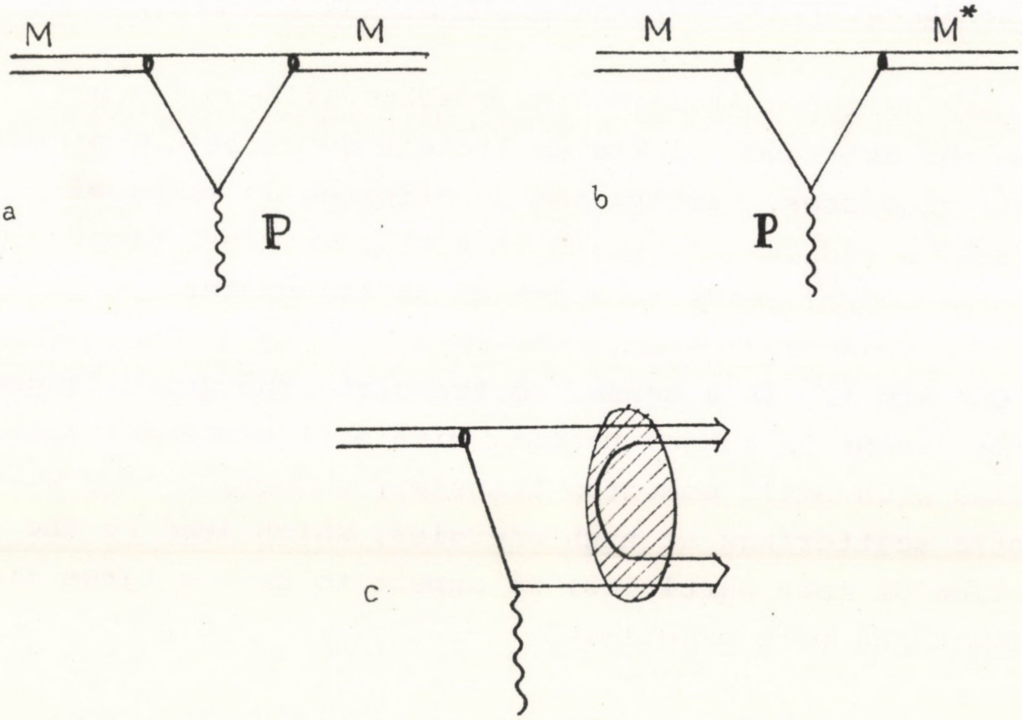


Fig. 5. The process of diffraction dissociation of the meson.

excited state M^* /or a set of excited states/. Quarks, however cannot dissociate into free particles. A new quark-antiquark pair has to be born and then the meson dissociates into two mesons.

The last process differs essentially from processes which are possible for nuclei. In the case of fast mesons each of their constituent quarks carries $x \sim \frac{1}{2}$ $(x = p/p_{max})$. After the dissociation of the meson the two new mesons /Fig. 5c/ will have also $x \sim \frac{1}{2}$, i.e. in such a production process of a quark-antiquark pair the initial quarks transmit a part of their momenta easily, "softly".

III. The structure of multiparticle production processes

Considering a picture with quark confinement, one assumes the existence of two equivalent descriptions of the physical processes, namely: the description in terms of quark states and that in terms of real particles, since each quark state corresponds to a set of hadron states.

Our aim is, in a sense, to translate the quark language into the hadron language. Dealing with soft processes /i.e. processes with small momentum transfer/ and especially with inelastic scatterings at high energies, which lead to the production of many particles, we expect to have a large field for comparison with experiment.

Let us see first of all, how to describe the process of hadron production in e^+e^- annihilation assuming the hadron structure with two characteristic sizes. The virtual γ-quantum produces a pair of point-like partons /a quark-antiquark pair which in the c.m. system scatters in different directions/. One assumes here, that the energy $\sqrt{s}$ of the e^+e^- pair is sufficiently large. At small s when $s \lesssim \frac{6}{r_q^2}$, the virtual γ produces with a high probability a pair of constituent quarks instead of a quark-parton pair. That's why the picture which we present is valid only at $s \gg \frac{6}{r_q^2} \simeq 6\ \mathrm{GeV}^2$.

The quark-antiquark pair produces new partons /quarks and gluons/, which at sufficiently large distances become dressed quarks /and dressed gluons/. The constituent quarks then, joining each other, form hadrons. A constituent quark and a constituent antiquark give a meson, three quarks - a baryon. If there exist heavy gluonic mesons, they can be formed by dressed gluons /see Fig.6/. The particles which are produced this way form two jets of hadrons flying in the opposite directions /if, of course, at the first stage hard gluons are not produced, which would then lead to a third jet /.

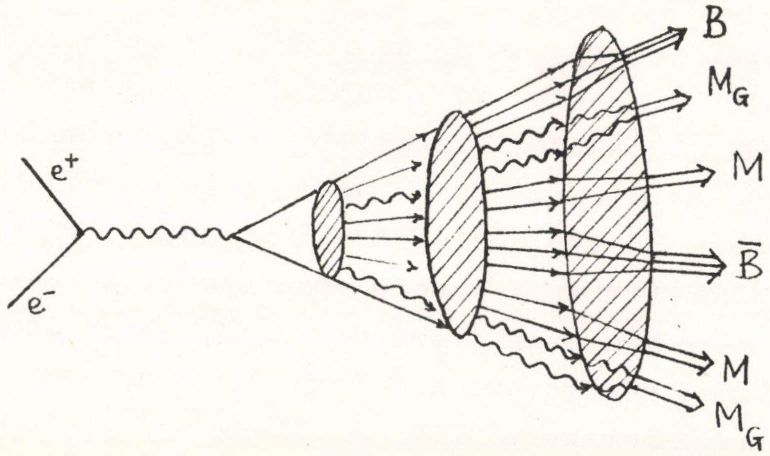


Fig.6.

In another hard process which is connected with multiparticle production of hadrons, namely the deep inelastic scattering of leptons on nucleous, the virtual γ^* interacts with one of the quark-partons of the nucleon. This quark-parton flying in the direction of the γ^* gives rise to the formation of new quark-partons and gluons analogously to the case of e^+e^- annihilation. Again, these particles transform into dressed quarks and heavy gluon formations. The difference is, that here in the beam going in the direction of the baryon remain two constituent quarks of the incident nucleon which did not take part in the interaction, i.e. quarks-spectators /Fig.7/

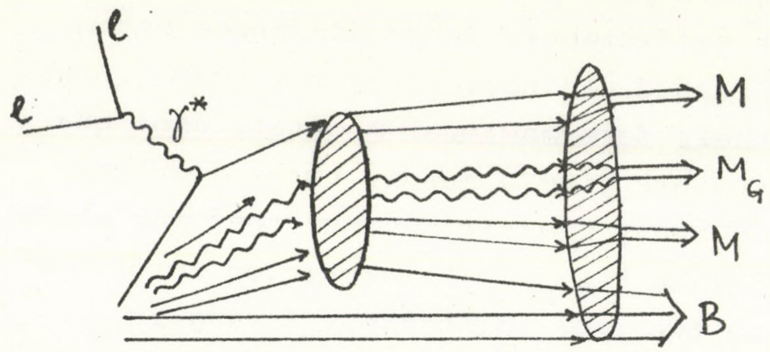


Fig.7.

In hadron-hadron collisions the hadron production is considered in a similar way. As we told already before, we assume the spectator mechanism as a natural consequence of the picture of spatially separated quarks. In inelastic hadron scattering at high energies in fact two dressed quarks collide: one constituent quark of the incident hadron and on of the target. The other constituents remain spectators. As a result of the collision many new quarks are produced, which afterwards join the quarks-spectators and form fast secondary hadrons, observable in experiment. Fig. 8 shows

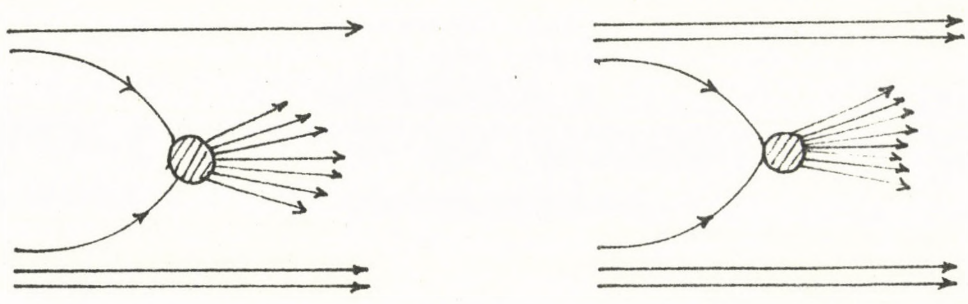


Fig. 8.

a picture of meson-baryon and baryon-baryon collisions of this kind.

If the hadron consists of discrete dressed quarks, then inside a fast baryon each of them has to carry about 1/3 of the total baryon momentum, while inside a meson - about half of the meson momentum. Consequently, multiparticle production processes in hadron-hadron collisions can be divided into two energetically different regions: the central and the fragmentation ones /I. and II. in Fig. 9 /

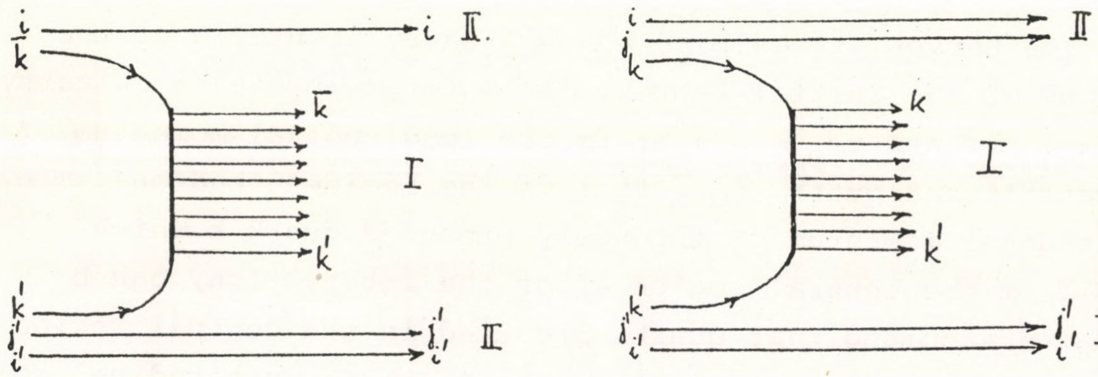


Fig. 9.

The quarks in the central region are seaquarks, carrying a small fraction of the incident momentum. Joining each other, they form the spectrum of slow hadrons.

The quark-spectators of the colliding particles (q_i, q_j and q_i', q_j' in Fig. 9) join quarks /or antiquarks/ of the sea forming the hadrons in the fragmentation region. The pair of quarks q_k and q_k' produced in the central region after the interaction "remember" their origin and have to be regarded as belonging to the fragmentation region.

Consider now what processes are possible in the fragmentation region. /For the sake of simplicity we consider a baryon fragmentation process/. The interacting quark q_k can join the spectators, forming a baryon state containing the same quarks as the incident one /Fig. 10.a./ if the collision of q_i, q_j and q_k is coherent, then the produced hadron B_{ijk} is analogous to the initial state /in the case of an incident proton that means $p \to p$ transition/. If the collision is not coherent, then the produced B^*_{ijk} state is some superposition of possible real hadrons /e.g. $p \to p$, $p \to \Delta^+$ etc./.

The spectators q_i, q_j can join a sea quark, in this case a baryon state B_{ij} is formed /Fig.10.b./. At the same time q_k together with a sea antiquark form a meson state M_k.

The baryon states B_{ijk} and B_{ij} carry about 2/3 of the momentum of the initial hadron. The interacting quark q_k carry away $x \sim \frac{1}{3}$ (where $x = \frac{p_L}{p_{max}}$; p_L is the longitudinal momentum of the constituent quark, p_{max} that of the incident hadron/. The longitudinal momentum of the newly produced quark k which comes from the central region after the interaction, can be estimated assuming that quarks produced in the central region distribute homogeneously in log x, i.e. their longitudinal momenta follow the geometrical progression law. This is the so called comb regime which leads to a Regge-pole exchange in elastic scattering. If so, the fastest produced quark has a momentum equal to a half of the incoming quarks momentum, the next one 1/4 of it etc. That means, that the meson state M_k is produced in the $x \lesssim 0{,}15$ region.

If one spectator joins two sea quarks, a baryon state B_i $(x \sim \frac{1}{3})$ is formed; the other spectator joining a sea antiquark form a meson state $M_j (x \sim \frac{1}{3})$. /Fig. 10.c./ There are also cases when only meson states are produced /Fig. 10.d,e/.

The meson fragmentation process can be considered in the same way. /Fig. 10.f,g,h/

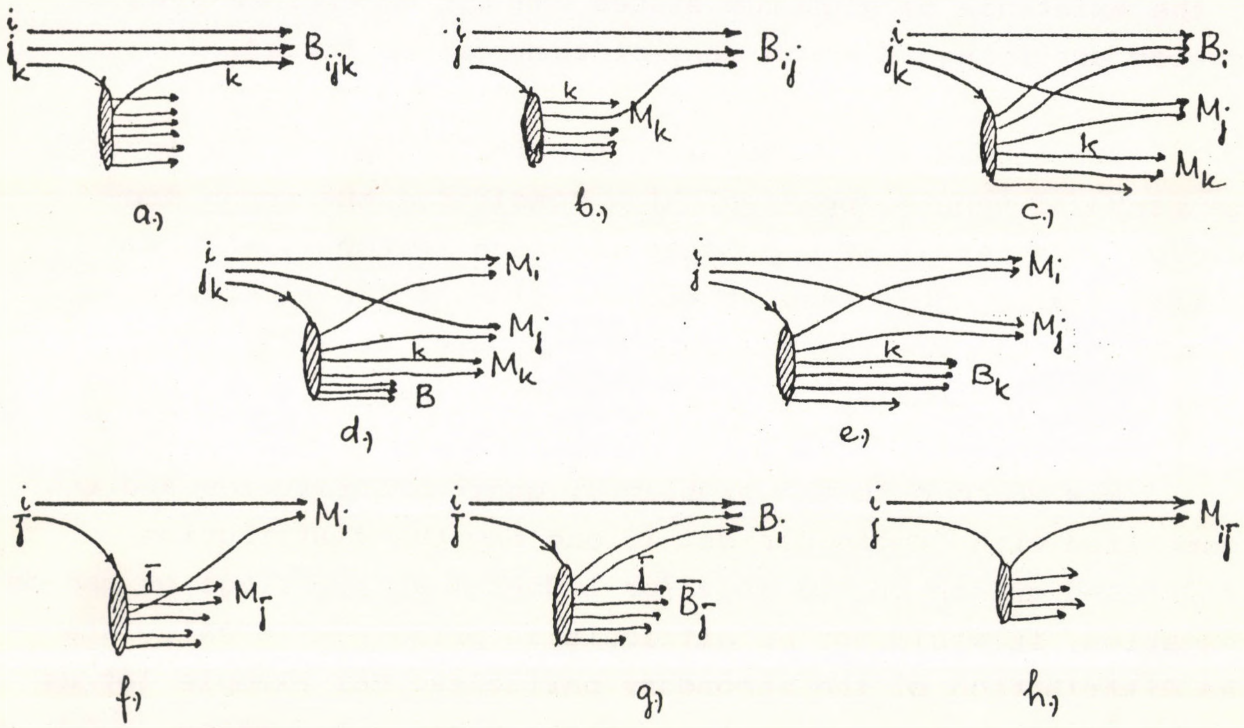


Fig. 10

IV. The calculation of the probabilities of hadron state production with the help of statistical rules - quark combinatorics

The second assumption which is made in the quark combinatorial calculus is connected with the newly produced particles. Our aim is to calculate the probabilities of particle production with the help of some statistical rules. To do so, we have to make first two remarks.

The quark model is SU(6) symmetric. The lowest hadron states are formed of constituent quarks. The existence of pure gluonic meson states i.e. gluonium states or glueballs remain, however, an open question. Several papers are considering this problem, coming to the conclusion, that, if gluonium exist at all it has to be relatively heavy - about 2-4 GeV. /However, in [44] arguments in favour of light glueballs are stated/. Since there are practically no informations available about them, we will consider only meson states formed by quarks. However, we think it is necessary to remember the possibility of the existence of gluonium states and not to exclude even a version according to which many of them can be formed.

The second remark concerns the following. In the last few years several papers appeared which considered the meson production as a result of a process in which partons /and not dressed quarks/ join each other. In spite of the similarities, this approach differs essentially from ours. We will return later to the comparison of the two approaches.

Now we come to the problem of quark distribution and its connection with combinatorics. If one new the distribution of dressed quarks in the jets /or ladders/ at the stage of hadron formation, it would not be difficult in principle to calculate the distribution of the secondary particles. For example let us see the inclusive cross section of the meson production $x \frac{d\sigma}{dx d\vec{k}_\perp}$ /x is the part of the momentum which is carried by the fast meson is the jet, $\vec{k}_\perp$ is the transversal component of its momentum/. The cross section is determined by the distribution of the dressed quark and the dressed antiquark $F(x_1,\vec{k}_{1\perp}; x_2,\vec{k}_{2\perp})$ and by the square of the wave function of the constituents in the meson $|\Psi(s_{12})|^2$ which depends on the energy of their relative motion $s_{12} = (K_1+K_2)^2 = \sum_{i=1,2} M^2_{i\perp} \cdot \frac{x}{x_i} - k^2_\perp$ (here $M^2_i = M^2 + k^2_{i\perp}$, M is the mass of the constituent quark: $M_u = M_d \simeq 330 \div 360$ MeV, $M_s/M_d \cong 1{,}5$):

$$x \frac{d\sigma}{dx d\vec{k}_\perp} = \int \prod_{i=1,2} \frac{dx_i}{x_i} d\vec{k}_{i\perp} \, \delta(x_1+x_2-x)\,\delta(\vec{k}_{1\perp}-k_{2\perp}-k_\perp) \,\cdot$$

$$\cdot F(x_1,\vec{k}_{1\perp}; x_2,\vec{k}_{2\perp}) \cdot |\Psi(\sum_{i=1,2} M_{i\perp}\frac{x}{x_i} - k^2_\perp)|^2 \qquad (3)$$

The behaviour of $|\Psi(s_{12})|^2$ can be sufficiently well estimated from the behaviour of the hadron /here-meson/ form factor at small momentum transfer. From the constituent quark models it is known, that $|\;(s_{12})|^2$ decreases quite quickly with the increase of s_{12}. This means, that joining each other are quarks with small relative rapidities i.e. usually neighbours on the rapidity axis. However, $F(x_1,\vec{k}_1\,;x_2,\vec{k}_2\,)$ is not known beforehand. If, nevertheless, there is a lot of information available about the inclusive spectra and about the relations between the newly produced particles this is connected with the smallness of the energy of the relative motion of the quarks joining each other. Indeed, at small energies for the interactions between quarks SU(6) symmetry holds, i.e. the interactions do not depend on quark quantum numbers like spin and isospin. Hence, if in the quark distributions there are no correlations with other quarks which are far on the rapidity axis and break this property, then the SU(6) symmetry remains valid for the production processes of secondary particles also. This means that in multiparticle production processes not only stable particles appear, but resonances also, independently of their quantum numbers and thus the production probabilities of all hadron states belonging to one SU(6) multiplet are equal. The probability of the hadron production within one SU(6) multiplet is proportional to the number of spin states of these hadrons, i.e. 2J+1.

Note, that the independence of the meson production probabilities on their quantum numbers can be obtained assuming the following factorization of the distribution function:

$$F(x_1,\vec{k}_{1\perp};\,x_2,\vec{k}_{2\perp}) = F(x_1,\vec{k}_{1\perp})F(x_2,\vec{k}_{2\perp}) \qquad (4)$$

Besides the absence of correlations between the quantum numbers of the quarks and antiquarks one can assume that no correlations exist between quark-antiquark pairs with conjugate quantum numbers-this means, that such pairs go apart far enough on the rapidity axis. Assuming the absence of colour correlations also /this is done in practically all papers considering the processes $q,g \to$ hadrons, although the opposite situation can also be taken into account [46,47]/one comes to simple rules for calculating relations between secondary hadrons - the rules of quark combinatorics. Using them, we first obtain expressions for some integral characteristics - the average multiplicities. In this case we can avoid the problem of not knowing the quark distribution functions in the jets.

Let us consider some, sufficiently large interval on the rapidity axis in the central region at the stage when dressed quarks and antiquarks are formed. /For the sake of simplicity we do not take into account gluons./ An arbitrarily chosen particle might be a quark or an antiquark with the same probability $\frac{1}{2}\,q + \frac{1}{2}\,\bar{q}$. We assume, that only the nearest neighbours are joining each other; such a restriction does not affect our considerations. The nearest neighbour is again either a quark, or an antiquark. The probability of the states qq, $\bar{q}\bar{q}$ is then

$$(\tfrac{1}{2}q + \tfrac{1}{2}\bar{q})(\tfrac{1}{2}q + \tfrac{1}{2}\bar{q}) \to \tfrac{1}{4}qq + \tfrac{1}{4}\bar{q}\bar{q} + \tfrac{1}{2}q\bar{q} \to \tfrac{1}{4}qq + \tfrac{1}{4}\bar{q}\bar{q} + \tfrac{1}{2}M \text{ where}$$

$M=q\bar{q}$ is a meson state. Taking into account a third possible quark or antiquark, one gets

$$(\tfrac{1}{4}qq + \tfrac{1}{4}\bar{q}\bar{q} + \tfrac{1}{2}M)(\tfrac{1}{2}q + \tfrac{1}{2}\bar{q}) \to \tfrac{1}{8}B + \tfrac{1}{8}\bar{B} + \tfrac{3}{4}M(\tfrac{1}{2}q + \tfrac{1}{2}\bar{q})$$

where $B=qqq$, $\bar{B}=\bar{q}\bar{q}\bar{q}$. Further iteration lead to the following multiplicity of particles produced in the central region:

$$(q,\bar{q} - \text{sea}) \to 6N\cdot M+N\cdot B+N\cdot\bar{B} \qquad (5)$$

The number N depends on the total energy of the colliding particles, and increases with the growth of s. Supposing that the multiplicity $N(s)$ is increasing logarithmically, it is convenient to write $N(s)=\ b\ \ln\frac{s}{s_o}$ at asymptotic energies. The parameters b and s_o can not be determined by quark combinatorics, but have to be the same for all processes. Hence, the relation between the produced mesons M, baryons B and antibaryons $\overline{B}$ is [1]:

$$M : B : \overline{B} = 6 : 1 : 1 \qquad (6)$$

In the same way one can get relations between baryons and mesons in the fragmentation region too [45]. In this case one considers an incident quark q_i, which, joining a quark or an antiquark of the sea, forms with the probability 2:1 mesons or baryons containing this quark:

$$(q_i + q,\bar{q} - \text{sea}) \rightarrow \frac{1}{3}B_i + \frac{2}{3}M_i + \frac{1}{3} M + N(s) \cdot (6M + B + \overline{B}) \qquad (7)$$

Here $B_i = q_i qq$, $M_i = q_i \bar{q}$; N(s) is a large number which is characterized by the number of quarks in the sea.

A similar relation is valid for the case when a pair of quarks q_i, q_j transforms into hadrons [45]:

$$(q_i q_j + q,\bar{q}\text{-sea}) \rightarrow \frac{1}{2} B_{ij} + \frac{1}{12}(B_i + B_j) + \frac{5}{12}(M_i + M_j) + \frac{1}{6} M +$$

$$+ N(s) \cdot (6M + B + \overline{B}) \qquad (8)$$

The baryon state B_{ij} contains both incident quarks: $B_{ij} = q_i q_j q$.

To complete the transition from quarks to hadrons, one has to solve a very important problem: to understand, what real hadrons correspond to the mesonic and baryonic states B_{ij}, B_i , M_i etc. Indeed, quark combinatorics, while operating with constituent quark states $q\bar{q}$ and qqq does not answer the question by what real particles they are saturated. In [1] the dominance of the lowest SU(6) multiplets was supposed, i.e. the meson 36-plet /$J^P = 0^{-1}, 1^{-1}$) for the $q\bar{q}$ states and the baryon 56-plet ($J^P = \frac{1}{2}^+, \frac{3}{2}^+$) for qqq, respectively. This is a rather rough approach, and, of course, a contribution of hadrons belonging to higher multiplets is quite natural.

The determination of hadrons which are saturating the meson and baryon states is in fact an experimental question, which, in a sense, characterizes the quark confinement. The analysis of experimental data shows, that the contribution of hadrons with L = 1 is quite significant: 20-30% of the produced particles. The share of L=2 multiplets seems to be about 10% /Fig. 11/.

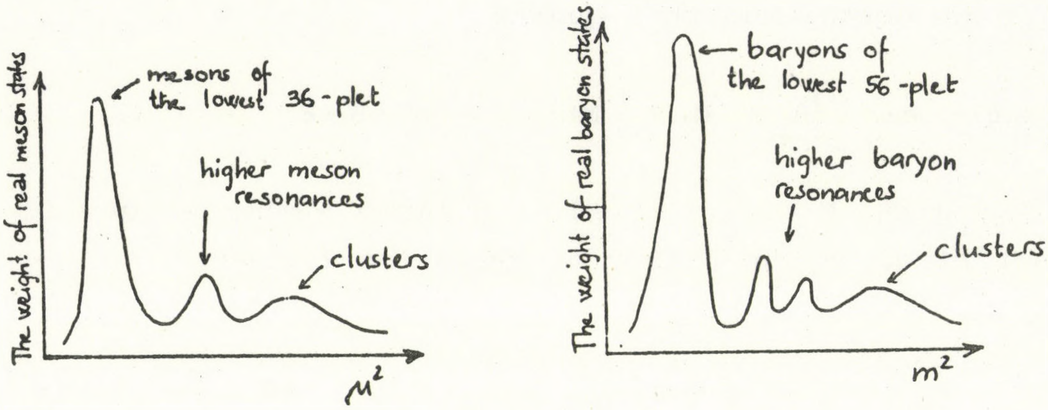


Fig. 11

Another problem which has to be cleared is the following. It is known, that the SU(6) symmetry is broken: this is connected mainly with the features of the s-quark which are different from those of the u and d quarks. In order to understand, how to introduce the SU(6)-breaking, let us consider the expression (3). The changes in $|\psi|^2$ will probably not give a serious contribution since their depend mainly on the changes of the binding energy and not of the masses of the constituents. Deviations in $F(x_1,x_2)$ can be quite noticeable These deviations can be taken into account if one considers a non-symmetrical quark sea with a relatively suppressed production of strange quarks. This suppression is characterized by a parameter $\lambda \le 1$; in the case of $\lambda = 1$ the symmetry between quarks u,d,s is restored. The values of λ might be different in the central and fragmentational regions, respectively [48]. This difference is due to the fact that the distributions of the produced strange and non-strange quarks can change with the increase in x in a different way. This means that $\lambda = \lambda(x)$, and $\lambda(0) = \lambda_{\text{central region}}$. However, since the present knowledge of particle production in the fragmentation region is not very accurate, one can use instead of $\lambda(x)$ the rougher description of an effective constant λ_f at not too small x.

Now we are in the position to express the states B_{ij}, B_i, M_i, B, M in the terms of real particles. For the meson states we consider the possibility of multiplets with L=0 and L=1. What concerns the baryons, the experimental evidence on baryon resonance production in multiparticle production processes is rather poor, therefore we restrict ourselves to the lowest L=0 multiplet.

The meson states M_i and M can be written as

$$\begin{aligned} &\sum_L \alpha_i(L) M_i(L) \\ &\sum_L \alpha(L) M(L) \end{aligned} \qquad (9)$$

The indices L=0,1 correspond to the s and p-wave states, respectively. The probabilities $\alpha_i(L)$ and $\alpha(L)$ are fixed by the conditions $\sum_L \alpha_i = 1$, $\sum_L \alpha(L) = 1$.

Denoting the real mesons belonging to the L=0 multiplet as $h_{M(0)}$ and those with L=1 as $h_{M(1)}$ we can write the decomposition of $M_i(L)$ and $M(L)$ into the real meson states in the form

$$\begin{aligned} M_i(L) &= \sum_h \mu_h^L(i) h_{M(L)} \\ M\,(L) &= \sum_h \mu_h^L\, h_{M(L)} \end{aligned} \qquad (10)$$

The coefficients $\mu_h^L(i)$ and μ_h^L /which are the probabilities of observing the meson $h_{M(L)}$ in the states $M_i(L)$ and $M(L)$ respectively/ are given in Table 1. The decay modes and their relative probabilities are taken from [49].

Similarly, the real hadron content of the states B, B_i and B_{ij} can be written as

$$B_i = \sum_L \tilde{\alpha}_L^{(1)} B_L(q_i) , \qquad B = \sum_L \tilde{\alpha}_L B_L \quad \text{etc.} \tag{11}$$

For the lowest 56-plet with L=O we have

$$\begin{aligned} B_{ij}(56;0) &= \sum_h \beta_h(ij)\, h_{B(56\ 0)} \\ B_i(56\ ;\ 0) &= \sum_h \beta_h(i)\, h_{B(56\ ;\ 0)} \\ B(56\ ;\ 0) &= \sum_h \beta_h \ h_{B(56\ ;\ 0)} \end{aligned} \tag{12}$$

The coefficients $\beta(i,j)$, $\beta(i)$ and β are presented in Table 2. Similary to α, the coefficients β and μ fulfill the normalization conditions $\sum_h \beta_h(i,j) = 1$, $\sum_h \mu_h(i) = 1$ etc.

V. Verification of the rules of quark statistics

As it was told in the Introduction, the hypothesis of the hadron structure with quasi-free dressed quarks leads to two serious consequences: to the spectator mechanism and to the statistical rules for the calculation of secondaries in multiparticle production processes. Let us begin with the latter one. We will consider two types of these statistical rules; those which appear in processes when quarks join each other independently of their spins and those which are connected with the situation that hadron states are produced independently of the fact if quarks join quarks or antiquarks . In the first case we get relations between secondary particles with different quark spins, in the second one - relations between the produced mesons and baryons. The experimental data give quite a definite argument in favour of the existence of statistical rules for particles with different quark spins. The situation with the relations between mesons and baryons is more ambiguous. We think that the best way to prove it is to investigate the particle production in e^+e^--annihilation - this will be done in the last part of the present chapter.

Considering Tables 1 and 2 it is easy to discover, that the production probabilities of directly produced particles /i.e. those which are not formed as results of decay processes/ with similar quark content obey some simple relations. For example: $\rho^+ : \pi^+ = 1:3$; $K^{*o}(890):K^o = 3:1$, $\Delta^+:p = 2:1$ etc.

These relations are consequences of the assumption that quarks join each other and form hadrons independently of their spins. They can be understood in the following way. We consider the dressed quarks formed in jets /or multiperipherial ladders/ as a gas of quarks and antiquarks with non-correlated spin projections [50]. In such a "gas" the number of $q\bar{q}$ pairs with definite total spin values $s_{q\bar{q}}$ is proportional to the statistical weight $2s_{q\bar{q}}+1$ of these states, i.e. the ratio of number of pairs with $s_{q\bar{q}}=1$ and of those with $s_{q\bar{q}}=0$ is 3:1. If the mesons are formed by quarks and antiquarks independently of their spin projections, then this ratio is true for the produced mesons too; the multiplicity of meson with $s_{q\bar{q}}=1$ is proportional to the multiplicity of $s_{q\bar{q}}=0$ states as 3:1. In hadron-hadron collisions this relations is true for both the fragmentational and the central regions. Examples for that can be the widely discussed ρ/π and $K^*(890)/K$ relations /see [51-54]/. There is a difficulty in the experimental proof of the relation 3:1. Name-

ly, one has to separate the directly produced mesons for which this relation is valid from those which appear as a decay product of some resonances.
The observed ρ/Π ratio, for example, might seriously change if there are Π mesons present produced by decays of unidentified resonances. That's why it is more convenient to prove 3:1 on secondary K-mesons which appear in decay processes to a much less extent. We expect that in strange particle production processes 75% of all particles have total quark spin $s_{q\bar{q}} = 1$ and only 25% is of $s_{q\bar{q}} = 0$. All particles with $s_{q\bar{q}} = 1$ are resonances and therefore /if we do not consider the decays of non-strange resonances into K-mesons/ about 75% of all observed K-mesons have to be decay products of resonances with $s_{q\bar{q}} = 1$.

Experimental data on the production of K-resonances in pp [55] and K^-p[56] collisions provide a possibility to test the quark statistical condition 3:1. In the mentioned works the inclusive cross sections of K, $K^*(890)$ and $K^*(1420)$ production were measured. Due to the SU(6) classification, the first two particles belong to the lowest 36-plet of mesons with L=O while the tensor resonance $K^*(1420)$ belongs to the L=1 multiplet. The results of the measurements are given in Table 1.

The meson multiplet with L=1 contains 4 SU(3) nonents $J^P = 0^+, 1^+, 1^+$ and 2^+. The statistical weight of each of these nonets is proportional to 2J + 1. That means, that 5/12 of the particles belonging to the multiplet with L=1 has to be tensor mesons, i.e. the total amount of particles produced in this multiplet is $12/5 \cdot T$. 75% of them have to be mesons with $s_{q\bar{q}} = 1$. Thus $V + 9/5 \cdot T$ is the contribution of resonances with $s_{q\bar{q}} = 1$ to the K-meson production. As it is seen from the data quoted in Table 3, the experimental value of the above mentioned quantity is in each case near to 75% of the total cross section of kaons. Such an agreement of experimental data with the theoretical predictions clearly speaks for the hypothesis of the "gas" of non-correlated quarks and antiquarks in multihadron production processes.

Let us mention here, that in our calculations we did not take into account those possible resonances, which belong to the L=2 SU(6) multiplet, and therefore all the kaons produced by decays of these resonances were added to the directly produced K-mesons with $s_{q\bar{q}} = 0$. The contribution of resonances with L=2 can be roughly estimated considering the cross section of the g-meson production. The contribution of these resonances turns out to be about 5-10%. It is interesting, that according to the most accurate measurement of the K^- and K^0 production [56] the contribution of $s_{q\bar{q}} = 1$ resonances is somewhat less that 75%. The addition of the resonances with L=2 will probably increase this value. In this case the agreement with the condition 3:1 will be much better than the accuracy of about 10%, which is usual in quark models.

Further, let us investigate the relations between the productions of mesons and baryons. These relations are, as we told before, consequences of quark statistics for quarks and antiquarks: they appear if the hadronization of the gas of constituent quarks is independent of the fact if quarks join quarks or antiquarks. To this mechanism correspond formulae (6)-(8). We consider in the following (7). This expression means, that as a consequence of statistical rules the baryon number of the quark q_i manifests itself as the probability of the production of the baryon states B_i by this q_i. The easiest way to prove $M_i : B_i = 2:1$ is to investigate reactions in which the multiparticle production is initiated by one quark - the current region in deep inelastic scatterings and the e^+e^- annihilation.

We have to consider the fragmentation region only, and forget about sea quarks - that means it is necessary to investigate the spectra of secondary hadrons. The spectrum of a secondary hadron is to be found from comparison with experiment. However, knowing the spectrum of one particle, the spectra of other particles of the same SU(6) -multiplet can in principle be obtained [57].

In the deep inelastic νN and $\bar{\nu} N$ collisions the current fragmentation is determined by quarks-partons knocked out from the nucleon. As it is known, the distributions of the quarks-partons depend on the value of x_B /where $x_B = \frac{Q^2}{2m(E-E')}$, Q is the momentum transferred to the lepton, E-E′ - the difference of the energies before and after the collision with the nucleon in the lab.system, m-the nucleon mass/. At $x_B > 0{,}1$ the structure functions of the nucleons are defined almost entirely by the valence quarks-partons; the contribution of the sea quarks-partons is here small. In this region of x_B the current fragmentation is determined mainly by valence u and and d quarks. The multiplicities of hadrons in jets generated by quarks can be written in the form

$$\frac{1}{\sigma} \int dx \frac{d\sigma}{dx} (\nu N \to u) = \frac{1}{3} B_u + \frac{2}{3} M_u + \frac{1}{3} M + N(6M+B+\bar{B})$$

$$\frac{1}{\sigma} \int dx \frac{d\sigma}{dx} (\bar{\nu} N \to d \cos\theta_c + s \sin\theta_c) = (\frac{1}{3} B_d + \frac{2}{3} M_d) \cos^2\theta_c + \qquad (13)$$

$$+ (\frac{1}{3} B_s + \frac{2}{3} M_s) \sin^2\theta_c + \frac{1}{3} M + N(6M+B+\bar{B})$$

where θ_c is the Cabibbo angle. The hadron states B_i, B, M_i and M can be expanded in terms of real hadrons, corresponding to different SU(6) multiplets. Hadrons belonging to one SU(6) multiplet have equal distributions in x. Hence the inclusive cross-sections of the hadrons in the jets of current fragmentation are given by universal functions:

$$\frac{1}{\sigma}\frac{d\sigma}{dx}(\nu N \to u) = Q_u(x)$$

$$\frac{1}{\sigma}\frac{d\sigma}{dx}(\bar{\nu}N \to d\cos\theta_c + s\sin\theta_c) = Q_d(x)\cos^2\theta_c + Q_s(x)\sin^2\theta_c \tag{14}$$

The introduced functions $Q_i(x)$ are

$$Q_i(x) = \frac{1}{3}\sum_L F_L(x)\tilde{\alpha}_L^{(1)} B_L(q_i) + \frac{2}{3}\sum_L \Phi_L(x)\alpha_L^{(1)} M_L(q_i) + \sum_L [\varphi_L(x)\alpha_L M_L + f_L(x)\tilde{\alpha}_L(B_L + \bar{B}_L)] \tag{15}$$

Here $\Phi_L(x)$ and $F(x)$ are distribution functions of fragmentational meson and baryon states belonging to the multiplet L. They fulfill the normalization conditions

$$\int_0^1 dx F_L(x) = 1, \qquad \int_0^1 dx \Phi_L(x) = 1 \tag{16}$$

$\varphi_L(x)$ and $f_L(x)$ are the distribution functions of meson and baryon states of the multiplet L in the central region. They are normalized to the number of states in the considered jet and therefore the normalization depends on the total energy of the jet $\sqrt{s}$:

$$\int_{\frac{m}{\sqrt{s}}}^1 dx\varphi_L(x) = 6N(s) + \frac{1}{3}, \qquad \int_{\frac{m}{\sqrt{s}}}^1 dx f_L(x) = N(s) \tag{17}$$

where m is a constant of the order of the mass of the particles. The behaviour of $\varphi_L(x)$ and $f_L(x)$ at small x defines the law due to which multiplicities increase in the jet; e.g. if $\varphi_L(x) \sim \frac{1}{x}$ and $f_L(x) \sim \frac{1}{x}$ at small x, this corresponds to the logarithmical increase: $N(s) \sim \ln s$. As it was told already, the details of (16) have to be chosen in order to fit the experiment. We will here consider only the lowest multiplets /35 and 56/ and take the following parametrization, which, as it will be seen, describe the experimental data quite well.

$$\varphi(x) = \frac{(1-x)^{n_M}}{x}\,(a_M+b_Mx+c_Mx^2)$$

$$f(x) = \frac{(1-x)^{n_B}}{x}\,(a_B+b_Bx+c_Bx^2) \qquad (18)$$

$$\Phi(x) = \frac{1-x}{\sqrt{x}}\,(A_M+B_Mx+C_Mx^2)$$

Let us explain these expressions. The parameters n_M and n_B have to be obtained from experiment; they determine the speed of the decrease of the contribution of sea hadrons at $x \to 1$. The same is true for a,A,b,B, etc. The presence of the factor 1-x leads for the meson formfactor at $x\to 1$ to a q^{-2} dependence /q is the momentum transfer/. Note, that the constituent quarks are not point-like objects and therefore the considered relations can not be understood literally as consequences of the Bloom-Gilman duality [58]. Nevertheless there exists a region $m_\rho^2 \ll q^2 \ll r_q^2$ /m_ρ is the characteristic hadron mass/ for the meson form factor where the quark structure is not important yet, but the form factor behaves already like q^{-2}. Apparently this corresponds to a relatively smooth transition from small to large q^2. The presence of $\sqrt{x}$ in the denominator of (18) is connected with the fact, that the probability to find in the meson a valence quark with small x is suppressed in comparison with the analogous probability for sea quarks. Indeed, a valence constituent quark will have a small x if its quantum numbers diffuse through the multiperipherial ladder. Such a transfer of quantum numbers can be described by the exchange of secondary reggeons with $\alpha_o(0) = \frac{1}{2}$ what leads to $x^{-1/2}$. This behaviour of the distribution function for fragmentation mesons is analogous to the behaviour of the distribution function for valence quark-partons in standart quark-parton models /see e.g. [59]/.

In the case of baryons the form-factor behaves like q^{-4} already at relatively small q^2 values and thus we assume that the asymptotical behaviour of the distribution function $F_o(x)$ is $(1-x)^3$ at $x \to 1$. In order to be able to compare the transition function of a constituent quark into a meson $\Phi_0(x)$ or into a baryon $F_0(x)$ we introduce the following parametrization:

$$F(x) = k \frac{\Phi(x)}{1+\frac{\varkappa^2}{(1-x)^2}} \qquad (19)$$

It is easy to see that $F(x)$ coincides with $\Phi(x)$ everywhere except the region $1-x \leq \varkappa$. As it will be seen from the results in the description of baryon spectra, the value $\varkappa$ turns out to be about 0,1. This is quite natural since the asymptotics $(1-x)^3$ become valid only near $x = 1$. The parameters $\varkappa$ u k are not independent because of the normalization condition (16) : k will be defined by the value of $\varkappa$.

The inclusive spectra in e^+e^- annihilation are determined by similar universal functions:

$$\frac{1}{2\sigma}\frac{d\sigma}{dx}(e^+e^- \to \text{hadrons}) = \frac{1}{3}Q_u(x) + \frac{1}{3}Q_{\bar{u}}(x) + \frac{1}{12}Q_d(x) + \frac{1}{12}Q_{\bar{d}}(x) + \frac{1}{12}Q_s(x) + \frac{1}{12}Q_{\bar{s}}(x) \qquad (20)$$

In the e^+e^- annihilation process there are two quark jets, that's why in the left hand side of (20) stands a factor 1/2. Besides, we so far do not consider the production of new heavy particles, and, correspondingly, we do not take into account the contribution of c and b-quarks. Hence, (20) can pretend only to the description of experimental data either lower than the threshold of the production of new particles, or, if over the threshold, then in the region of large x, in order to avoid the jet generated by the heavy quark.

There are not too many data on particle production in deep inelastic scattering. The situation is better for the particle spectra in e^+e^- annihilation. We will here take into account only the lowest multiplets to saturate B_i, B and M_i, M. The decays of resonances belonging to higher SU(6) multiplets imitate to a certain extent the direct production of hadrons belonging to lower multiplets. Hence, if the contribution of higher multiplets is about 20-30% the error in the spectra resulting from their omission will be less. In our case there is a supplementary reason for the decrease of the share of higher multiplets in the particle spectra at not very small x. Indeed, in the region $x > 0,2$ the spectra decrease quickly in x, and therefore the role of heavy resonances is relatively suppressed. The influence of heavy resonances can be essential only at small x. However, in the region of small x ($x < 0,1-0,2$) the experimental data are far from being definite, and give no possiblity to test our approach.

For the sake of simplicity we assume, that the x dependence of the fragmentation function $Q_i(x)$ equals for all kinds of quarks and antiquarks. This means that we ignore the mass difference between strange and not-strange quarks, which in fact has to give some observable effects in the fragmentation functions.

Figs. 12,13,14, show the experimental inclusive spectra $\frac{s}{\beta}\frac{d\sigma}{dx_E}$ of π,K and ρ^0 mesons and antiprotonos measured in e^+e^- annihilation at moderately high energies s = 16-25 GeV^2 [60]. Here β is the velocity of the secondary particle in the c.m. system of the colliding particles, $x_E = \frac{2E}{\sqrt{s}}$. At the considered energies scale invariant dependence on x_E is observed for these spectra /see [60,61]/. On the same Figs. the curves calculated for inclusive spectra with the help of (20) are also presented. Since quarks and antiquarks enter (20) in a symmetric way, different possibilities exist to describe the behaviour of the distribution function at $x \to 1$. The solid line

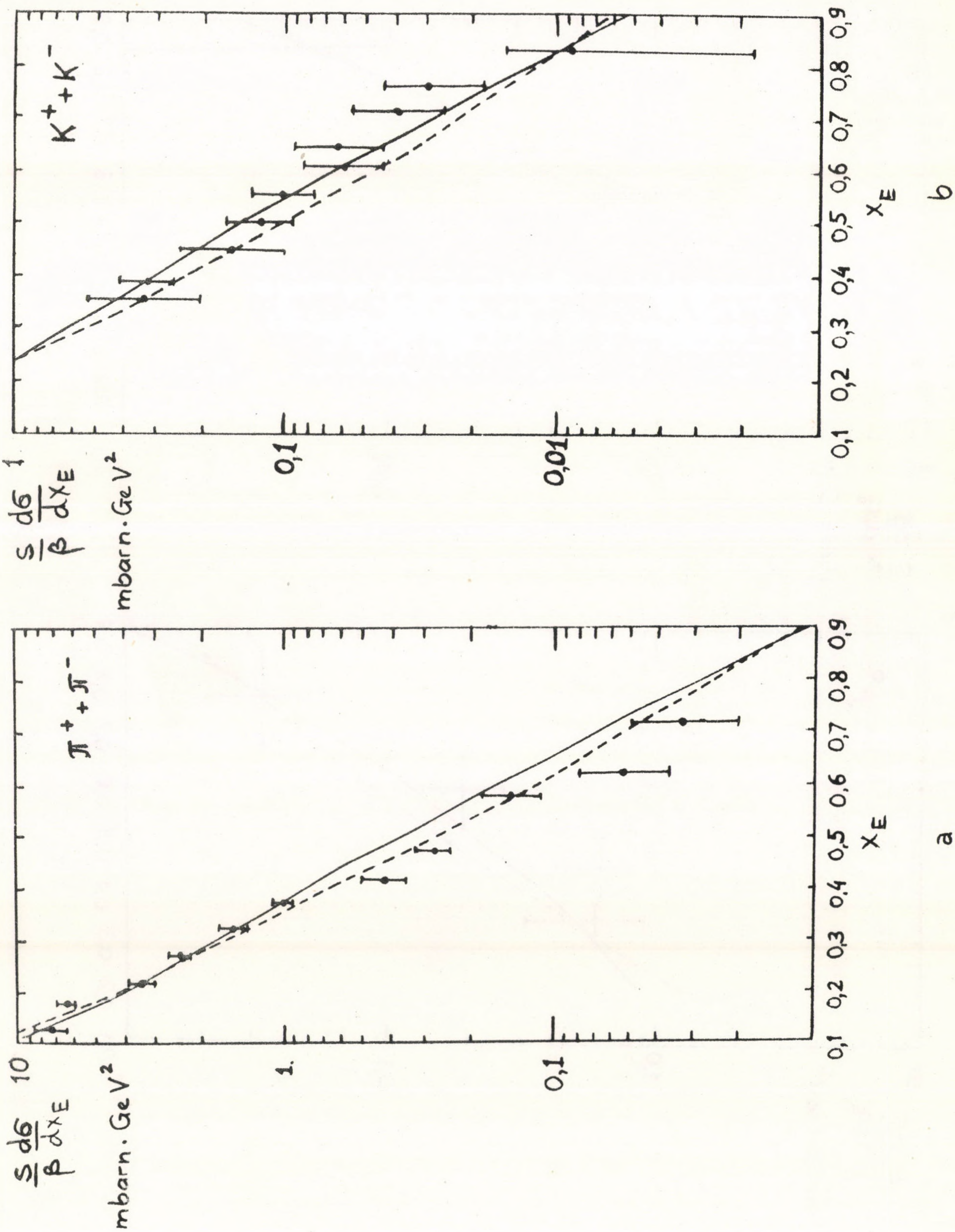


Fig. 12. Spectra of pions and Kaons

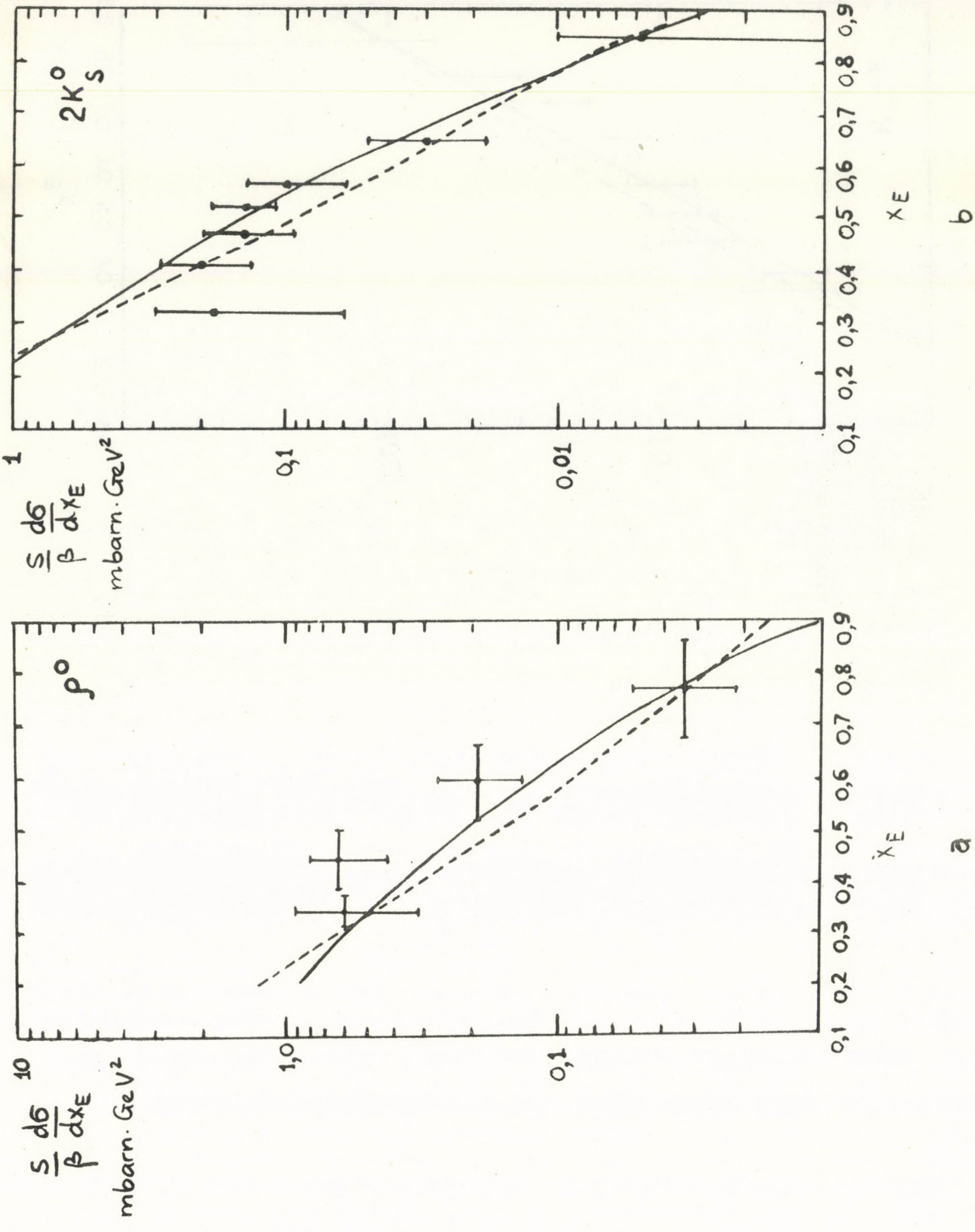


Fig. 13. Spectra of kaons and ρ^0-mesons
/ρ^0-preliminary date/

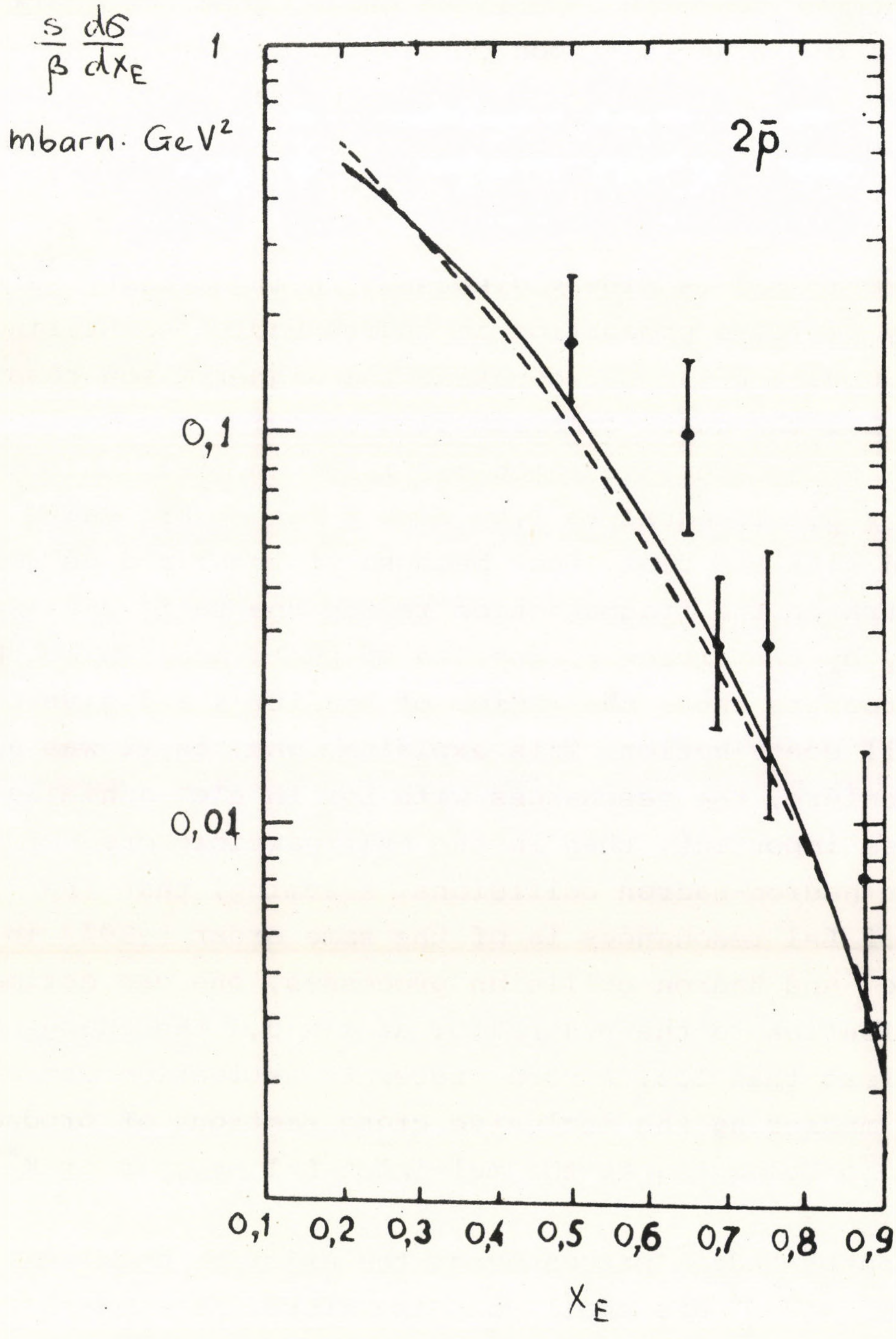


Fig. 14 Spectra of antiprotons

corresponds to the fastest decrease of $\varphi_0(x)$ and $f_0(x)$ at $x \to 1$, the dotted line to the slowest one; both of them fit the experimental data within the experimental errors.

The comparison with experiment shows, that the quark combinatorial calculus is a good candidate for a model to describe the spectra in the quark $\to$ hadron transitions. Especially impressive is the coincidence of experimental and theoretical results in the spectra of ρ_0 mesons.

The ratio ρ/Π is widely discussed in investigations of secondary particle production in hadron-hadron scatterings, where $\rho/\Pi \approx 0{,}09 - 0{,}13$. Not counting the higher meson resonances with L=1 the quark combinatorial calculus gives $\rho/\Pi \approx 0{,}17 - 0{,}21$. The calculated ρ/Π ratio in e^+e^- annihilation at $x > 0{,}4$ is equal to 1,5. Such a rather big value is connected with the fact, that because of the rapid decrease of the spectra in the fragmentation region the ratio ρ/Π is given mainly by the direct production of particles, while the resonance products enter the region of smaller x and give there a small contribution. This explains, why, as it was already told before, the resonances with L=1 in e^+e^- annihilation are less important, than in the multiparticle production processes in hadron-hadron collisions. Assuming, that the production of L=1 resonances is of the same order ($\sim 30\%$) in both the e^+e^- and hadron collision processes, one can estimate their contribution to the ρ/Π ratio: at $x > 0{,}4$ the change of it will be less than 10%. A more accurate estimation can be given after measuring the inclusive cross sections of production of mesons belonging to the multiplet L=1, e.g. f or K^{**}.

In lepton-hadron processes at the existing energies the values W^2 and Q^2 are small, and therefore the comparison of the predictions and the experimental data has to be performed only at large z /z is the momentum fraction of the fragmenting quark carried by the considered hadron:

$z = \frac{pk}{pQ} \frac{E^h_{lab}}{E_\nu - E_\mu}$, where p is the four-dimensional momentum of the proton of the target, Q - the four - momentum transferred from the leptons to the hadrons, and K - the four -momentum of the detected hadron/. The experiment [62] gives $\rho^0/\Pi^- = 0{,}7 \pm 0{,}4$ at $0{,}6 < z > 0{,}8$, and the statistics for this ratio are collected mainly at $z = 0{,}6$. The calculated spectra on Figs. 12,13 give $\rho^0/\Pi^- = 1{,}0$ at $x=z=0{,}6$ i.e. agree with experiment. This agreement is spoiled at smaller z values, but, as we told before, these values of z correspond to small Q^2 and W^2. It is, however, clear, that in this region the experimental evidence is not sufficient to prove the statistical rules.

The problem is, in fact, the following. In the general case one can write for the quark → hadron transition

$$q_i + (q,\bar{q}\text{ -sea}) \to \alpha B_i + (1-\alpha)\, M_i + \beta B + \\ + \left(\xi\frac{\alpha+\beta}{2} + \alpha - 1\right) M + N(\xi M + B + \overline{B}) \qquad (21)$$

where α, β and ξ are some free parameters, $0 \le \alpha \le 1, \xi > 0$. Since in the process of e^+e^- it is difficult to distinguish between the contributions of B_i and B, the coincidence of the experimental spectrum of the antiprotons with the results of the calculations indicates only, that the value of $\alpha + \beta$ is near to 1/3. That means, that the baryon number of the quark appears here as a production probability of baryon states: one can write $\beta = \frac{1}{3}$. In order to determine αB_i i.e. to devide the contributions of αB_i and βB in (21) it is necessary to measure secondary baryons with different charges, e.g. p and n or Σ^+ and Σ^-. It will be of great importance to measure baryon production in the current fragmentation region in deep inelastic νN and $\bar{\nu} N$ scatterings.

It seems to us, that a good possibility to measure α in (21) can be offered by the investigation of the production of charmed particles in e^+e^--annihilation at energies $\sqrt{s} \simeq 15 \div 25$ GeV. The $B_c = cqq$, $M_c = c\bar{q}$ and $M_{\bar{c}} = q\bar{c}$ states will be produced very probably at these energies in the fragmentation region, and if relation (7) is fulfilled /i.e. when $\alpha = \frac{1}{3}$ / we get for the total number of particles

$$B_c : \bar{B}_c : M_c : M_{\bar{c}} = 1 : 1 : 2 : 2.$$

VI. Multiplicities of the secondary particles in the fragmentation region and in the central region

If a quark q_k belonging to the baryon B_{ijk} hits the target, fast particles are produced with the following probabilities [63]:

$$B_{ijk} \to \Delta \cdot B_{ijk} + \Delta^* \cdot B^*_{ijk} + (1-\Delta-\Delta^*) \cdot [\frac{1}{2} B_{ij} +$$

$$+ \frac{1}{12}(B_i + B_j) + \frac{5}{12}(M_i + M_j) + \frac{1}{2}\,_k + \frac{2}{3} M_k + \frac{1}{2} M] + \ldots\ldots\ldots \qquad (22)$$

Here Δ and Δ^* are the probabilities of the coherent and incoherent transitions $B_{ijk} \to B_{ijk}$ and $B_{ijk} \to B^*_{ijk}$ respectively; they have to be determined from the experiment. In (22) the contribution of hadrons produced in the central region is not written down.

Analogously, the probability of production of fast hadrons after the collision of a meson $M_{i\bar{j}}$ with the target is

$$M_{i\bar{j}} \to \delta \cdot M_{i\bar{j}} + \delta^* \cdot M^*_{i\bar{j}} + (1-\delta-\delta^*) \cdot [\frac{1}{3} (B_i + \bar{B}_{\bar{j}}) + \frac{2}{3} (M_i + M_{\bar{j}}) + \frac{2}{3} M] \ldots \qquad (23)$$

Here the probabilities δ and δ^* of the processes $M_{ij} \to M_{i\bar{j}}$ and $M_{i\bar{j}} \to M^*_{i\bar{j}}$ cannot be defined in the framework of quark combinatorics. It can be shown, that in the quark model the probabilities Δ, Δ^* and δ, δ^* can depend on the initial hadron and on the type of the collision, thus is fact one has to write $\Delta_p(pp)$, $\Delta_p(kp)$. $\delta_k(kp)$ and so on. For the sake of simplicity, we will not take this into account.

The relations (22) and (23) and the expression of B_{ij}, M_i, B_i in terms of the real hadrons enable one to get easily the fragmentation multiplicities. For this purpose one has to take the wave function of the incident particle and to consider all the possible interactions of its constituent quarks.

As an example we consider in detail the fragmentation of the proton. We assume that the incident proton is completely polarized /this fact will be of no significance from the point of view of the result./ The proton wave function in this case is

$$\psi(p^{\uparrow}) = \frac{2}{3} \{ u^{\uparrow}u^{\uparrow}d^{\downarrow}\} - \frac{1}{3} \{ u^{\uparrow}u^{\downarrow}d^{\uparrow}\}$$

It is implied that the functions are symmetrized with respect to the SU(6) indices, e.g. $\{ u^{\uparrow}u^{\uparrow}d^{\downarrow}\} = \frac{1}{3} (u^{\uparrow}u^{\uparrow}d^{\downarrow} + u^{\uparrow}d^{\downarrow}u^{\uparrow} + d^{\downarrow}u^{\uparrow}u^{\uparrow})$. It can be seen immediately, that for the quarks-spectators the probability of being in a $\{ u^{\uparrow}u^{\uparrow} \}$ state is 2/9 /interacting is the quark $d^{\downarrow}$/, in $\{u^{\uparrow}u^{\downarrow}\}$ - 1/9, in $\{u^{\uparrow}u^{\uparrow}\}$ also 1/9 respectively. While the quark $u^{\uparrow}$ is interacting, the spectators are in a state which is described as
$\frac{1}{\sqrt{5}} (2\{u^{\uparrow}d^{\downarrow}\} - \{u^{\downarrow}d^{\uparrow}\}) = (ud)p.$

Thus, we have

$$B_{ij} = \frac{2}{9}B(u^{\uparrow}u^{\uparrow}) + \frac{1}{9}B(u^{\uparrow}u^{\downarrow}) + \frac{1}{9}B(u^{\uparrow}d^{\uparrow}) + \frac{5}{9} B_p(ud).$$

The decompositions of $B(u^{\uparrow}u^{\uparrow})$ and $B(u^{\uparrow}u^{\downarrow})$ into the real hadrons of the 56-plet lead to equal results, and therefore we write

$$\frac{1}{9} B(u^{\uparrow}u^{\downarrow}) + \frac{2}{3} B(u^{\uparrow}u^{\uparrow}) = \frac{1}{3} B(uu).$$

For the sake of simplicity we introduce the notation $B(u^{\uparrow}d^{\uparrow}) = B_1(ud)$. In the case of an incident proton the states B_i and M_i are equal to

$$B_i = \frac{2}{3} B(u) + \frac{1}{3} B(d) \qquad \text{and} \qquad M_i = \frac{2}{3} M(u) + \frac{1}{3} M(d)$$

respectively. As a result we can write

$$p \to \Delta_p \cdot p + \Delta_p^* \cdot B_p^* + (1-\Delta_p-\Delta_p^*) \{ \frac{1}{2}[\frac{5}{9}B_p(ud) + \frac{1}{3}B(uu) + \frac{1}{9}B_1(ud)] + \\ + \frac{1}{2} [\frac{2}{3}B(u) + \frac{1}{3}B(d)] + \frac{3}{2} [\frac{2}{3} M(u) + \frac{1}{3} M(d)] \tag{24}$$

Expanding the right-hand side in terms of the hadron states h /i.e. the meson states $h_{M(L)}$ and the baryon states h_B/ we finally, obtain

$$p \to \Delta_p \cdot p + \sum_h h \{\Delta_p^* \beta_h(p) + (1-\Delta_p-\Delta_p^*) [\frac{5}{18} \beta_h (ud_p) + \\ + \frac{1}{6} \beta_h(uu) + \frac{1}{18} \beta_h (ud_1) + \frac{1}{3} \beta_h (d)] \} + \\ + (1-\Delta_p-\Delta_p^*) \cdot \sum_{L=0,1} \sum_h h_{M(L)} \alpha_i(L)[\mu_h^L(u) + \frac{1}{2} \mu_h^L(d)] \tag{25}$$

Similarly to the incident proton case, the multiplicities of Λ and Σ^+ hyperons in the fragmentation region can be calculated:

$$\Lambda \to \Delta_\Lambda \cdot \Lambda + \sum_h h_B \{\Delta_\Lambda^* \beta_h(\Lambda) + (1-\Delta_\Lambda-\Delta_\Lambda^*) [\frac{\xi}{2(2+\xi)} \beta_h(ud_\Lambda) + \\ + \frac{1}{4(2+\xi)} (\beta_h(us_1) + \beta_h(ds_1) + \beta_h (us_o) + \beta_h(sd_o)) + \\ + \frac{2\xi + 1}{6(2+\xi)} \beta_h(s) + \frac{5 + \xi}{12(2+\xi)} (\beta_h(u) + \beta_h(d))]\} + \\ + (1-\Delta_\Lambda-\Delta_\Lambda^*) \sum_{L=0,1} \sum_h h_{M(L)} \alpha_i(L) [\frac{5+4\xi}{6(2+\xi)} \mu_h(s) + \\ + \frac{5\xi + 13}{12(2+\xi)} (\mu_h(u) + \mu_h (d))] \tag{26}$$

$$\Sigma^{+} \to \sum_h {}_h(\Sigma^+)\cdot h = \Delta_\Sigma \cdot \Sigma^+ + \sum_h h_B \{ \Delta^*_\Sigma \beta_h(\Sigma^+) +$$

$$+ (1-\Delta_\Sigma-\Delta^*_\Sigma)[\frac{\xi}{2(2+\xi)} \beta_h (uu) + \frac{5}{6(2+\xi)} \beta_h(us_\Sigma) +$$

$$+ \frac{1}{6(2+\xi)} \beta_h (us_1) + \frac{\xi + 5}{6(2+\xi)} \beta_h(u) + \frac{1 + 2\xi}{6(2+\xi)} \beta_h (s)]\} +$$

$$+ (1-\Delta_\Sigma-\Delta^*_\Sigma)\cdot \sum_{L=0,1} \sum_h h_{M(L)}\cdot\alpha_i(L)[\frac{5\xi+13}{6(2+\xi)} \mu_h (u) +$$

$$+ \frac{5+4\xi}{6(2+\xi)} \mu_h (s)] \qquad (27)$$

Differently form the proton case, in (26) and (27) it is taken into account, that the cross section of the interaction is less for the strange quark than for the non-strange one. Their ratio $\xi = \sigma_{inel}(sq)/\sigma_{inel}(qq)$ is near to 2/3.

Formula (23) enables us to calculate the fragmentation secondaries for incident mesons. In the cases of π^+ and K^+ we obtain the following

$$\pi^+ \to \sum_h F_h(\pi^+)\cdot h = \delta_\pi\cdot\pi^+ + \sum_{L=0,1} \sum_h h_{M(L)}\cdot\alpha_i(L)\cdot\{\delta^*_\pi \mu^L_h(\pi^+) +$$

$$+ (1-\delta_\pi-\delta^*_\pi)[\frac{2}{3} \mu^L_h(u) + \frac{2}{3} \mu^L_h(\bar{d})]\} +$$

$$+ \sum_h h_B (1-\delta_\pi-\delta_{\pi *})\cdot \frac{1}{3} \beta_h(u) + \sum_h h_{B^-} (1-\delta_\pi-\delta_{\pi *})\cdot\frac{1}{3} \beta_h(\bar{d})$$

$$(28)$$

$$K^+ \to \sum_h F_h(K^+)\cdot h = \delta_K\cdot K^+ + \sum_{L=0,1} \sum_h h_{M(L)}\cdot\alpha_i(L)\cdot\{\delta^*_K\mu^L_h(K^+) +$$

$$+ (1-\delta_K-\delta^*_K)\cdot[\frac{2}{3} \mu^L_h(u) + \frac{2}{3} \mu^L_h(\bar{s})]\} +$$

$$+ \sum_h h_B(1-\delta_K-\delta^*_K)\cdot \frac{1}{3} \beta_h(u) + \sum_h h_{\bar{B}}(1-\delta_K-\delta^*_K) \frac{1}{3} \beta_h(\bar{s})$$

$$(29)$$

In the central region the multiplicity of secondary particles is given by (5). Due to the additive quark model, the energy which is used for the production of new /sea/ quarks is determined by the energy of colliding quarks. In the pion-nucleon collision the square of this energy is about 1/6, in nucleon-nucleon collision about 1/9 of the total energy of hadrons. That means, that in the case of pion-nucleon collision we have

$$N_{\pi N}(s) = b \ln \frac{s}{6s_o} = b \ln \frac{s}{s^o_{\pi N}} \qquad (30)$$

while for the nucleon-nucleon case

$$N_{NN}(s) = b \ln \frac{2}{9s_o} = b \ln \frac{s}{s^o_{NN}} \qquad (31)$$

In the collision processes of strange particles on has to remember the difference between the cross-sections of the interaction of strange and non-strange quarks, and the fact that the heavier strange quark takes away a large part of the hadron momentum. Hence, for the kaon-nucleon collision one obtains

$$N_{KN}(s) = \frac{b\xi}{1+\xi} \ln \frac{s}{3(1+\mu)s_o} + \frac{b}{1+\xi} \ln \frac{s\mu}{3(1+\mu)s_o} = b\ln \frac{s}{s^o_{KN}} \qquad (32)$$

where $\mu = \frac{m_q}{m_s} \approx \frac{2}{3}$ is the ration of the strange and non-strange quarks. Finally,

$$N_{\Lambda N}(s) = N_{\Sigma N}(s) = \frac{2b}{2+\xi} \ln \frac{s\mu}{3(1+2\mu)s_o} + \frac{\xi b}{2+} \ln \frac{s}{3(1+2\mu)s_o} =$$

$$= b \ln \frac{s}{s^o_{\Lambda N}} \qquad (33)$$

The obtained expressions give a possibility to calculate the absolute values of average multiplicities of secondary particles in hadron-hadron collision. The parameters are fitted to the experimental data and according to them the coefficients in (30) - (33) are calculated. /For example the value of λ is selected to give the best agreement with the experimental K/π ratio in the central region and is found to be 0,3/. Supposing that the probabilities Δ and δ of the coherent processes $B_{ijk} \to B_{ijk}$ and $M_{\bar{i}\bar{j}} \to M_{\bar{i}\bar{j}}$ are mostly of diffractional origin, the value of these probabilities is estimated using the data on diffraction scattering. In the additive quark model the cross sections of diffraction processes in the meson-nucleon and baryon-nucleon scatterings are determined by the diagrams in Fig. 15.

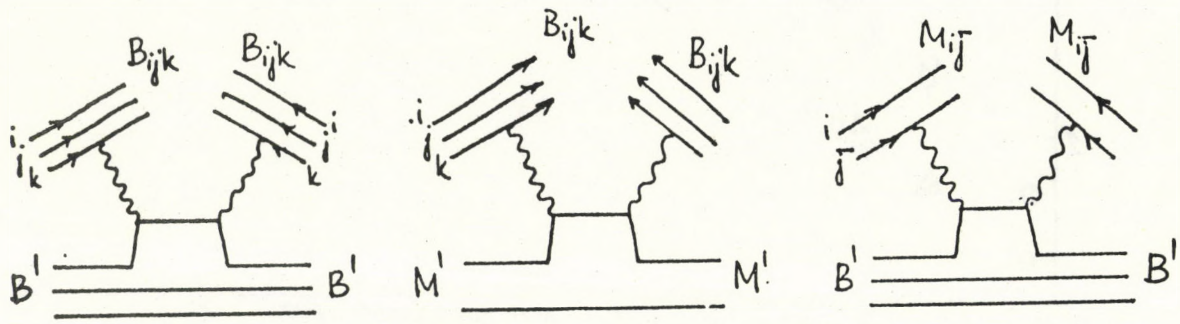


Fig. 15

/For the values of Δ and δ see [63] /.

The experimental data on average multiplicities of secondary hadrons in the pp, πp and Kp collisions permit us to prove the basic statements of quark combinatorics.

Consider first the meson production processes. In Fig. 16-18 the data on average multiplicities of secondary mesons in pp /Fig. 16/, $\pi^{\pm}p$ /Fig. 17/ and $K^{-}p$ /Fig. 18/ collisions at high energies are presented. The straight lines correspond to the predictions of the quark combinatorial calculus. In each case there is a satisfactory agreement of the theory and the experiment.

a./

b./

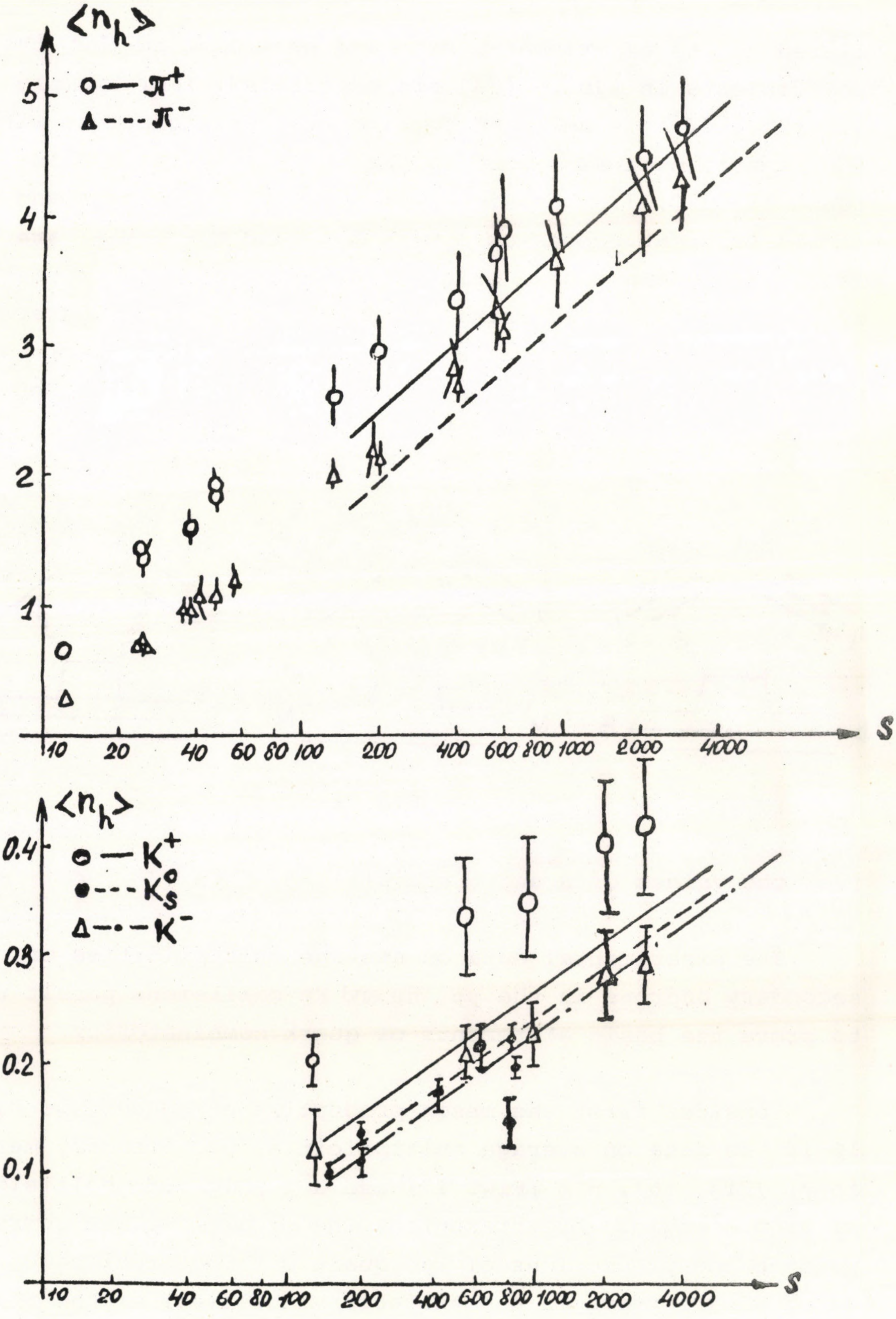
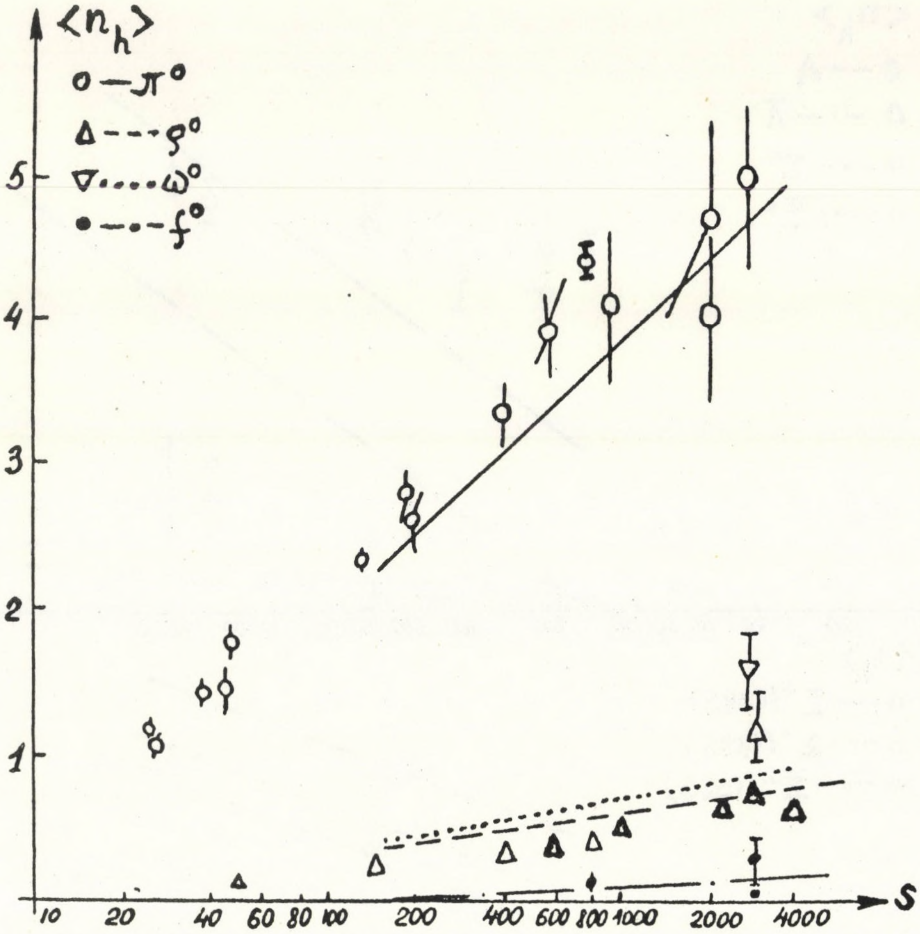


Fig. 16 Average multiplicities of secondary particles in pp collisions. The straight lines correspond to the predictions of the model.

c./


$\langle n_h \rangle$
o — π^0
△ - - - ρ^0
▽ ..... ω^0
● — • — f^0
5
4
3
2
1
10 20 40 60 80 100 200 400 600 800 1000 2000 4000
S

d./

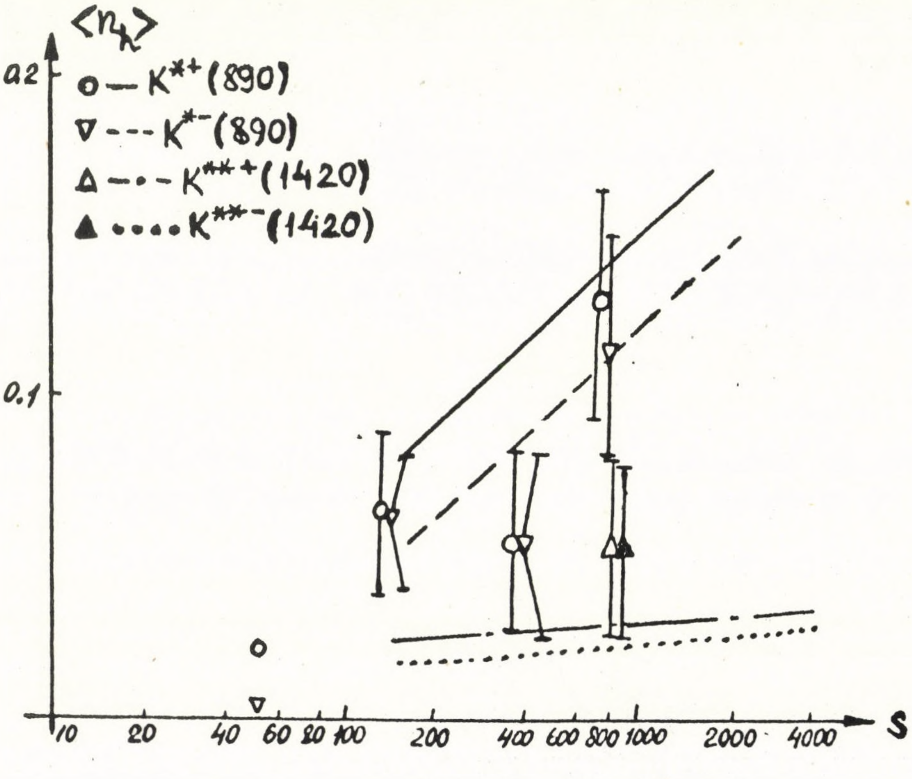

$\langle n_h \rangle$
0.2
o — $K^{*+}(890)$
▽ - - - $K^{*-}(890)$
△ — • — $K^{**+}(1420)$
▲ • • • • $K^{**-}(1420)$
0.1
10 20 40 60 80 100 200 400 600 800 1000 2000 4000
S

Fig. 16 c,d

e./

f./

g./

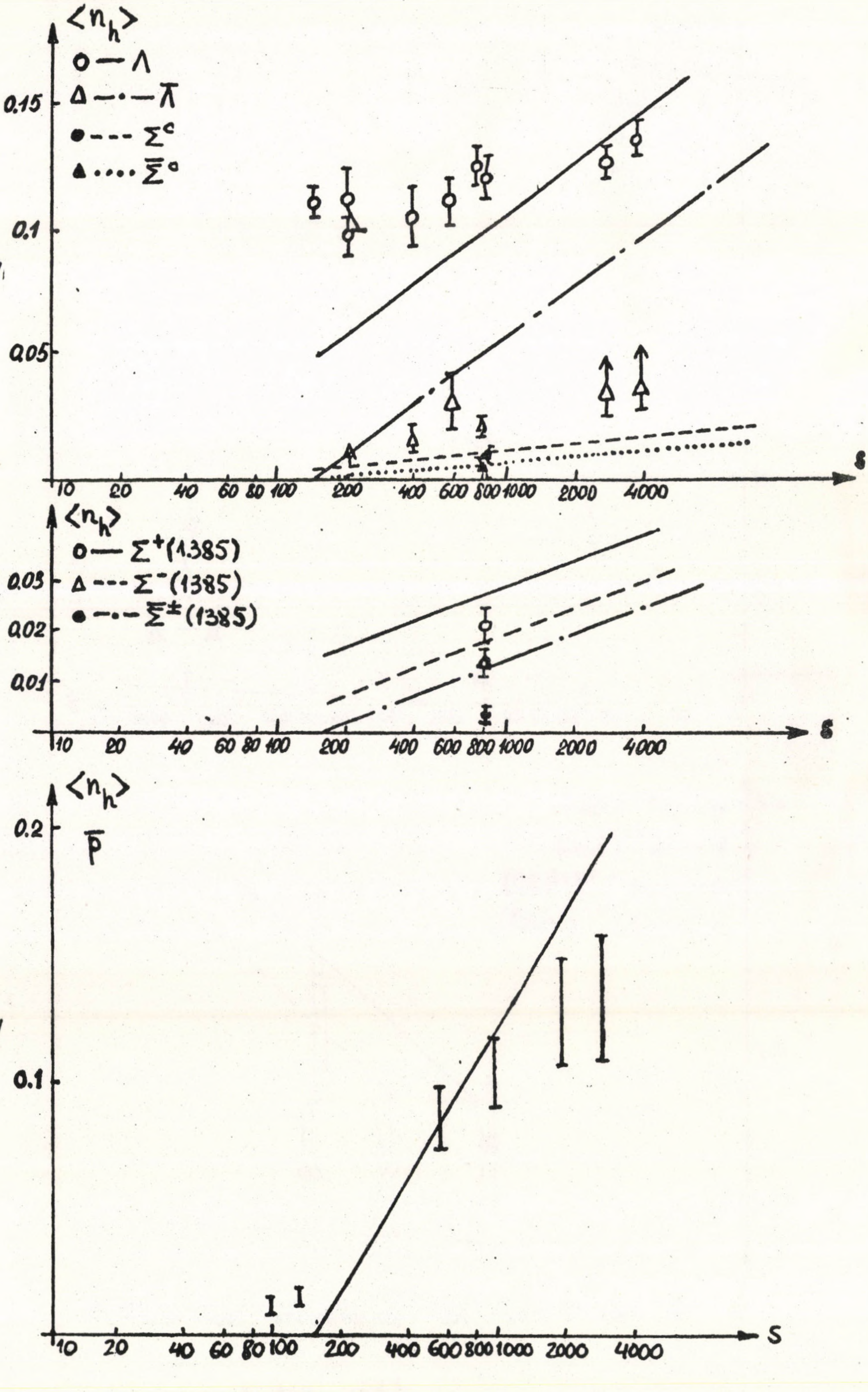


Fig. 16 e,f,g

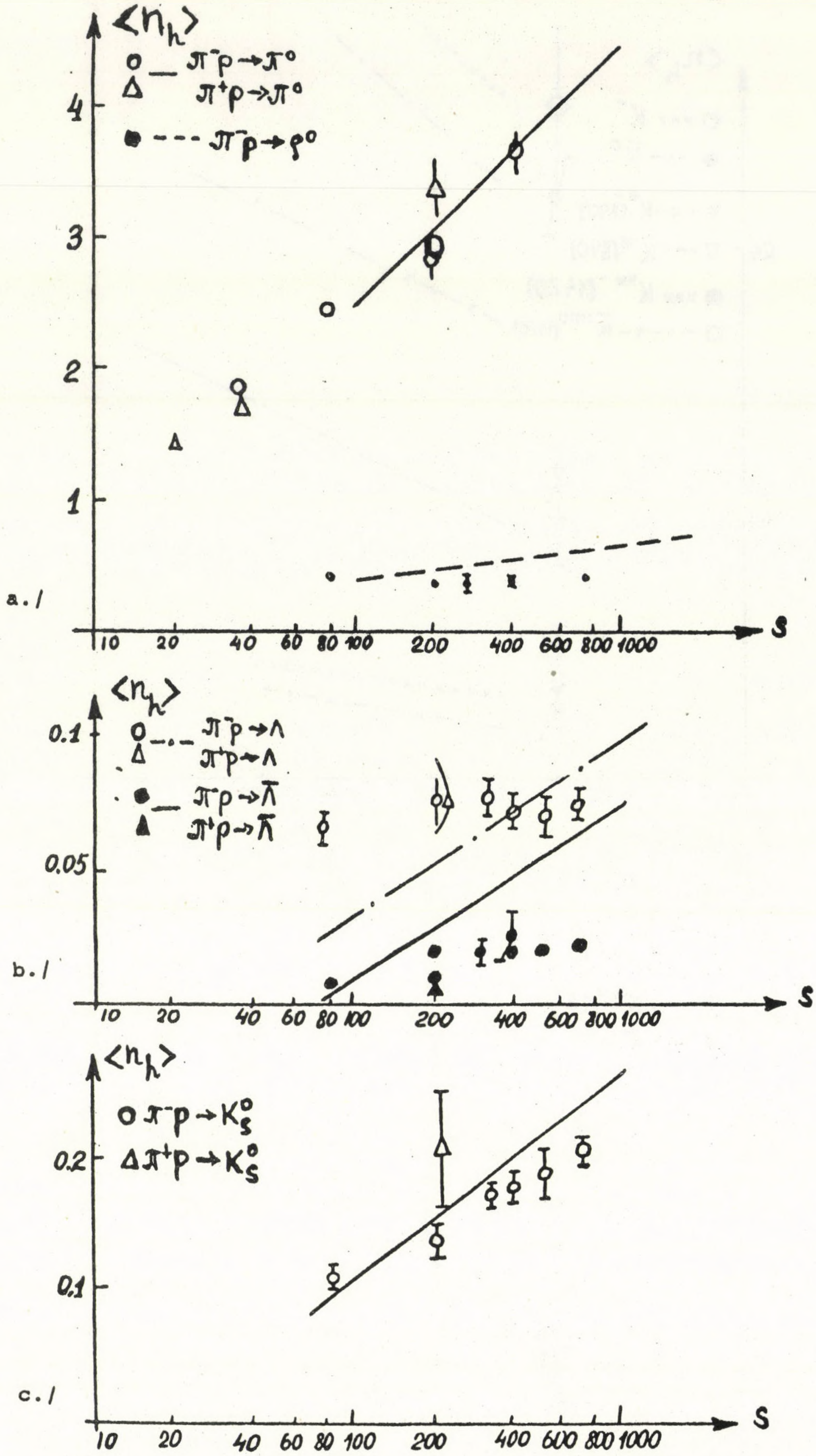


Fig. 17 Average multiplicities of secondary particles in $\pi^{\pm}p$ collisions

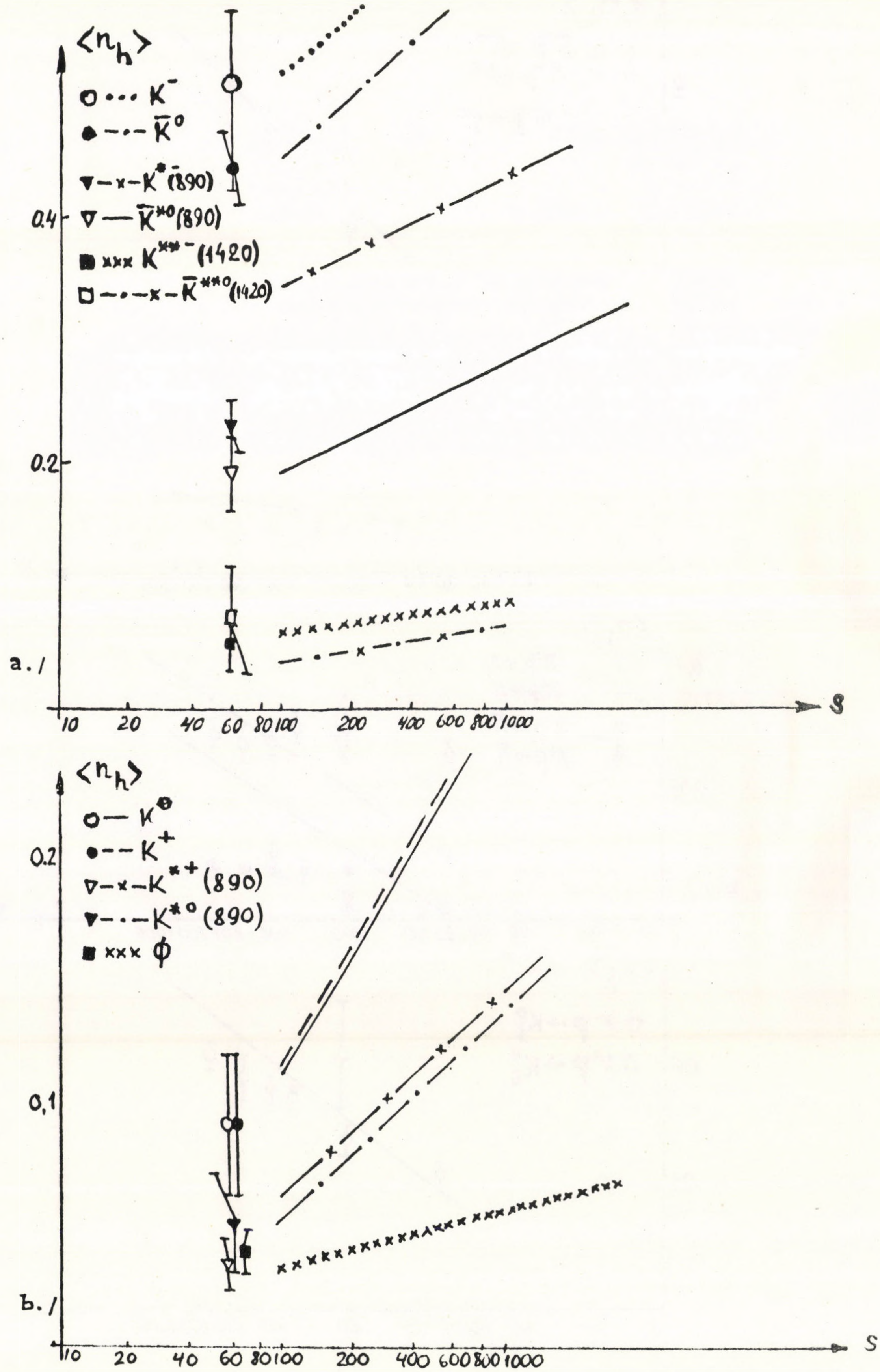
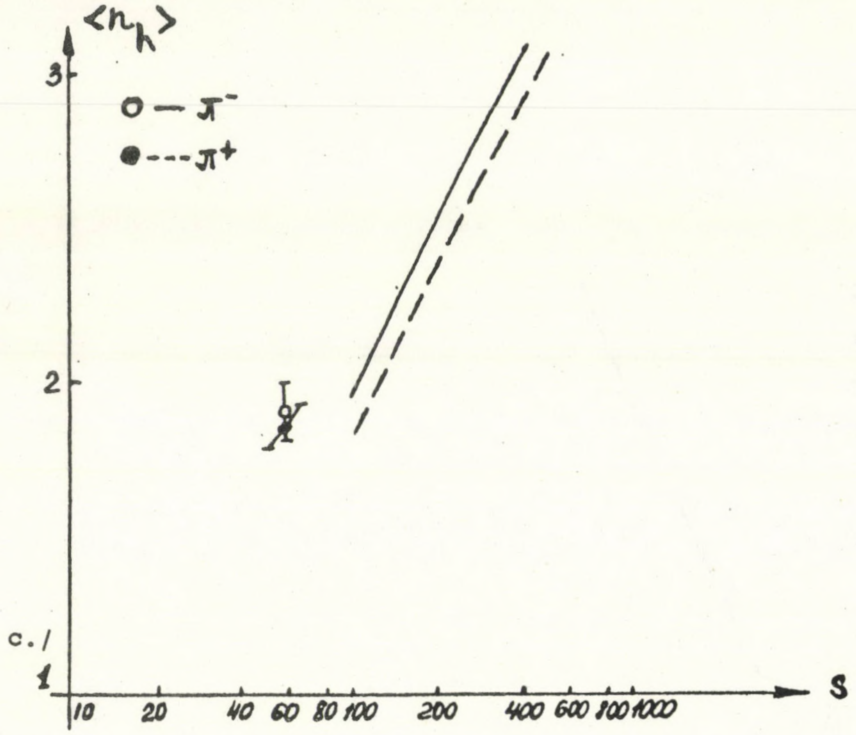
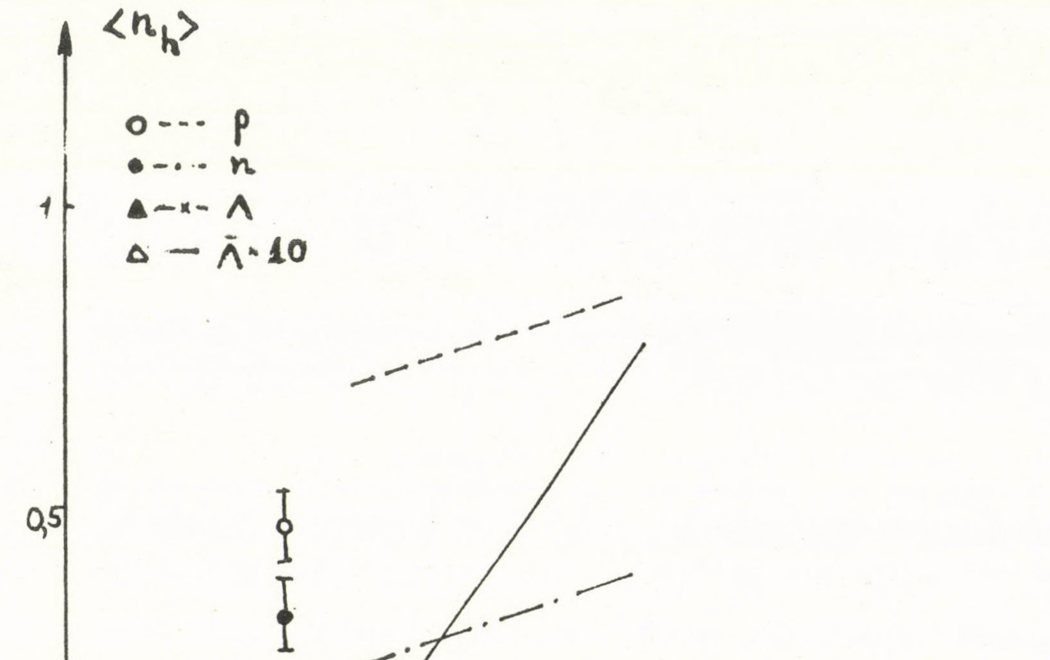


Fig. 18 a,b Average multiplicities of secondary particles in $K^{-}p$ colissions

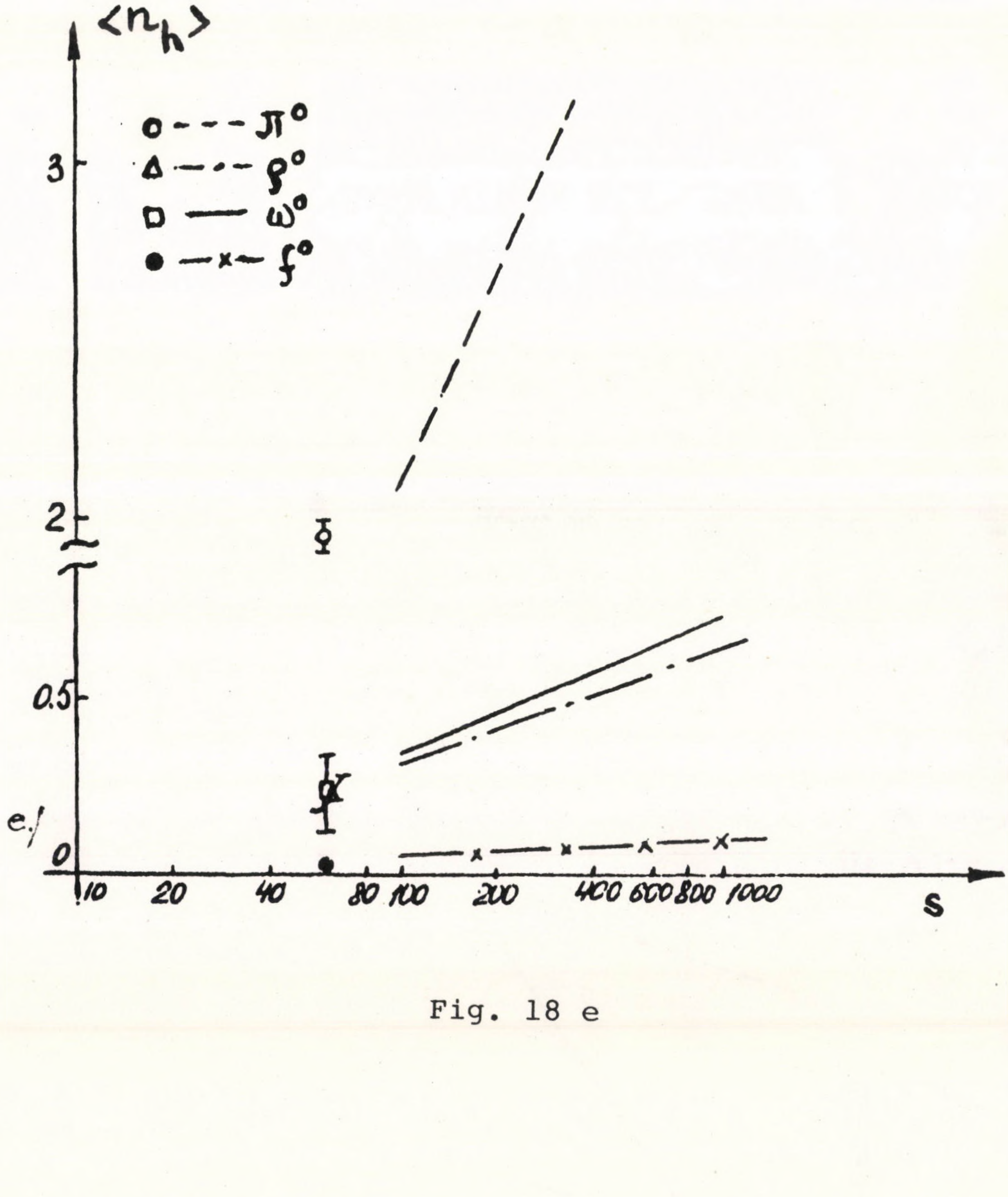


<n_h>
3
2
1
c./
o — π⁻
● --- π⁺
10
20
40
60
80
100
200
400
600
800
1000
s



<n_h>
o --- p
● -·-· n
▲ -x- Λ
△ — Λ̄·10
1
0,5
d./
10
20
40
60
80
100
200
400
600
800
1000
S

Fig. 18 c,d



$\langle n_h \rangle$
π^0
ρ^0
ω^0
f^0
3
2
0.5
e./
0
10
20
40
80
100
200
400
600
800
1000
s

Fig. 18 e

What concerns the baryons and the baryon resonances, the experimental data and the corresponding predictions of the quark model agree only roughly. For example, the ratios Λ/Σ^o, $\Sigma^+(1385)/\Sigma^o$ and $\Sigma^-(1385)/\Sigma^o$ satisfy the prediction quite well, what corresponds to the idea of baryons produced in SU(6) multiplets. The same ratios indicated that there might be a significant contribution of higher resonances. /For details see [63]/.

VII. Hadron-nucleus interactions

In the previous paragraph it was demonstrated that the investigation of multiparticle production processes provides a good possibility to prove the main assumption of the present approach, especially that concerning the extension of SU(6) symmetry. There are, however, processes, which allow to observe in a relatively pure way the consequences of the spectator mechanism, i.e. to prove the hypothesis which is teh crucial one from the point of wiev of the hadron structure. These processes are the hadron-nucleus collisions at high energies. They enable us to test the hadron structure because of the well-known fact that the fast secondary hadrons do not multiply by possible repeated collisions with the nuclear matter. This can be explained by the parton hypothesis: secondaries need time to be formed [5,43]. For fast particles this time increases with their momentum $p: \tau \sim \frac{p}{m^2}$. That means, that the constituents go through the nucleus before forming a secondary hadron, and, of course, they do not interact repeatedly with the nucleus.

As it was told already, in hadron-hadron collisions only one pair of consituent quarks takes part in the interaction /Fig. 9/. In a collision with a heavy nucleus, however, while going through the nuclear matter, the other constituents of the incident hadron can also interact. In the case of a superheavy nucleus all the constituents of the projectile would interact, so that all the three or two quark of an incident baryon or meson would break up. As a result, for example the multiplicity ratio of the secondaries in the central region for πA and pA interactions would be ~ 2/3. For real nuclei /even for A ~ 200/ a part of the constituent quarks still goes through a nucleus without interacting. The quarks which go through the nucleus without interaction determine the number of the fragmentational hadrons i.e. hadrons in the region of large x.

Hence, in baryon-nucleus collisions three different processes are possible: one quark is interacting, two go through the nucleus; two quarks are interacting, one goes through the nucleus; and finally, all three quarks interact. In meson-nucleus interactions one or two quarks of the incident meson can take part in the interaction. /These processes are shown in Fig. 19/.

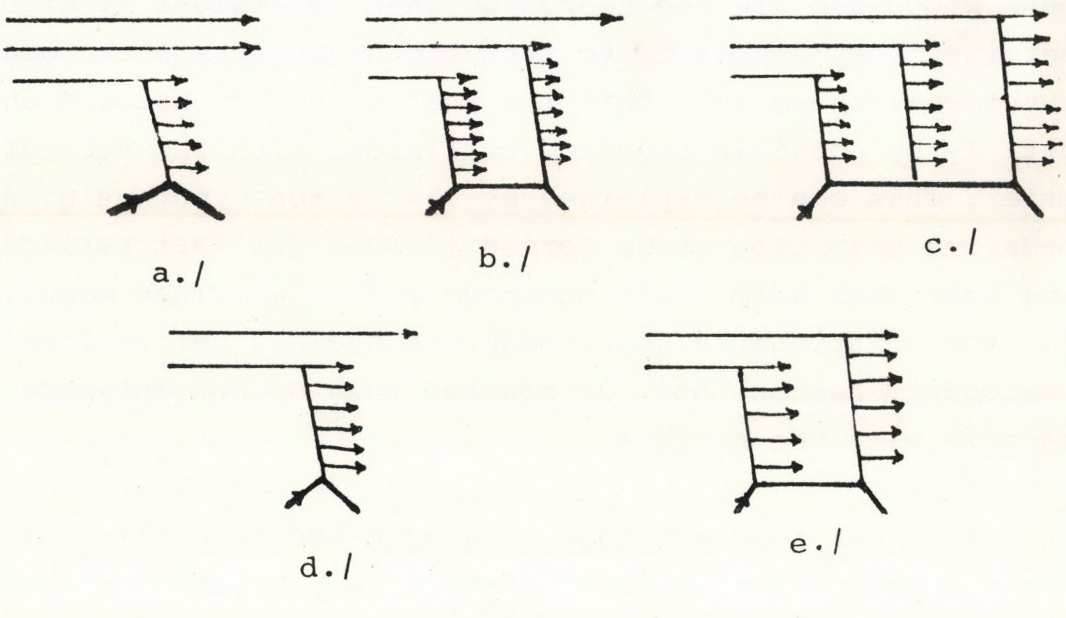


Fig. 19

Accepting the hadron picture with spatially separated quarks we assume that the constituent quarks interact with the nuclear matter in a independent way. The probability for a quark to interact is calculated as a function of the nuclear matter density and the quark-nucleon cross section

$$\sigma_{inel}(qN) \simeq \frac{1}{3}\sigma_{inel}(NN) \simeq \frac{1}{2}\sigma_{inel}(\pi N)$$

The probabilities of the processes can be written as

$$V_K = \frac{n!}{(n-k)!k!\sigma_{prod}} \int d^2b e^{-(n-k)\sigma_{inel}(qN)T(b)} [1-e^{-\sigma_{inel}(qN)T(b)}]^k \tag{34}$$

where k is the number of the interacting quarks, and h_n is the incident hadron consisting of n quarks [64].
The probability

$$\sigma_{prod} = \int d^2b \; [1-e^{-n\sigma_{inel}(qN)T(b)}] \tag{35}$$

has the meaning of the inelastic hadron-nucleus cross-section with the production of at least one secondary hadron and is obtained from the condition

$$\sum_{k=1}^{n} V_k^h = 1$$

The function T(b) is expressed in terms of the nuclear density

$$T(b) = A \int_{-\infty}^{\infty} dz \rho(b,z) \tag{36}$$

For $\rho(r = \sqrt{b^2+z^2})$, the Fermi parametrization,

$$\rho(r) = \frac{\rho_0}{1+\exp[(r-c_1)/c_2]} , \qquad 4\pi \int_0^\infty \rho(r) r^2 dr = 1 \tag{37}$$

is accepted. The c_1 and c_2 parameters are taken from the data on eA scattering [65, 66].

In the following we present the relative multiplicities of secondary particles in the central region. The multiplicities of secondaries n_{pA} and $n_{\pi A}$ in the pA and πA collisions can be expressed, using the formulae (34), in the form

$$R(\frac{pA}{qA}) = \frac{n_{pA}}{n_{qA}} = \sum_{k=1}^{3} kV_k^p = \frac{3}{\sigma_{prod}^{pA}} \int d^2b(1-e^{-\sigma_{inel}(qN)T(b)})$$

$$R(\frac{\pi A}{qA}) = \frac{n_{\pi A}}{n_{qA}} = \sum_{k=1}^{2} kV_k^\pi = \frac{2}{\sigma_{prod}^{\pi A}} \int d^2b(1-e^{-\sigma_{inel}(qN)T(b)}) \tag{38}$$

The ratio of the multiplicities in the meson-nucleus and nucleon-nucleus scatterings does not depend on n_{qA} and is equal to

$$R(\frac{\pi A}{pA}) = \frac{n_{\pi A}(y)}{n_{pA}(y)} = \frac{2\sigma_{prod}^{pA}}{3\sigma_{prod}^{\pi A}} = \frac{V_1^\pi(A) + 2V_2^\pi(A)}{V_1^p(A) + 2V_2^p(A) + 3V_3^p(A)} \tag{39}$$

The multiplicities $n_{\pi A}$ and n_{pA} might depend on the value of the rapidity of the corresponding secondaries. The comparison of the right-hand side of (39) with the experimental data is presented in Fig. 20.

The calculated value of $R(\frac{\pi A}{pA})$ is in agreement with experiment in the interval $1,5 < \eta < 3,5$ for the nuclei

C/A=12/ and Pb /A=207/ and for the fotoemulsion $\frac{1}{2}$ Ag + $\frac{1}{2}$ Br. The considered region for the values of the quasirapidity

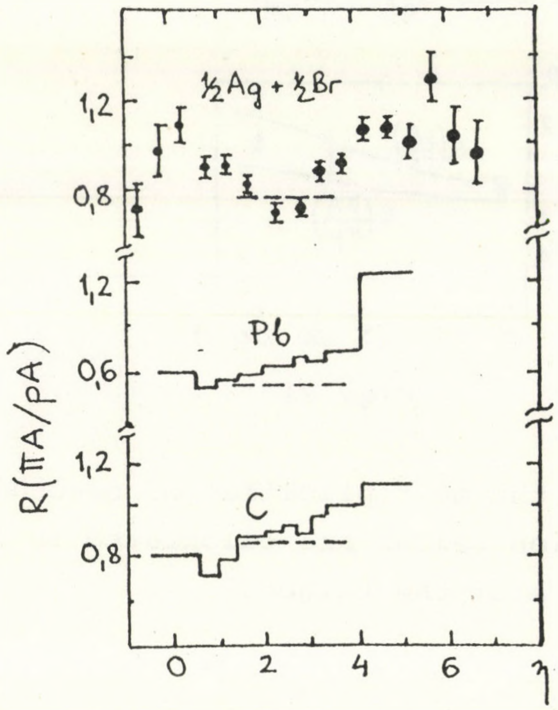


Fig.20

$\eta = -\ln\text{tg}\,\theta/2$ corresponds just to the central region of the collision processes.

The ratios of the secondaries in πA and πp scatterings and in pA and pp scatterings depend, due to (38), on the ratios n_{qA}/n_{qq} /where n_{qq} is the multiplicity in the quark-quark collision/:

$$R\left(\frac{\pi A}{\pi p}\right) = \left[V_1^{\pi}(A) + 2V_2^{\pi}(A)\right] \frac{n_{qA}(y)}{n_{qN}(y)} \tag{40}$$

$$R\left(\frac{pA}{pp}\right) = \left[V_1^{p}(A) + 2V_2^{p}(A) + 3V_3^{p}(A)\right] \frac{n_{qA}(y)}{n_{qN}(y)} \tag{41}$$

In Fig. 21 the experimental values averaged in the interval

$2,5 < \eta < 3,5$ are shown for $R(\frac{pA}{pp})$ (▲) and $R(\frac{\pi A}{\pi p})$ (▼) as functions of A. In this interval (39) is fulfilled for $R(\frac{\pi A}{pA})$, and one can take $n_{qA} \simeq n_{qN}$.

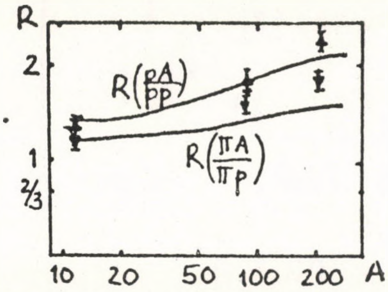


Fig. 21

Int the following the multiplicities of secondary hadrons in the fragmentation region are calculated as functions of the atomic number A of the target.

The $V_1^p(A)$, $V_2^p(A)$, $V_3^p(A)$ and σ_{prod}^{pA} values are shown in Figs. 22 and 23.

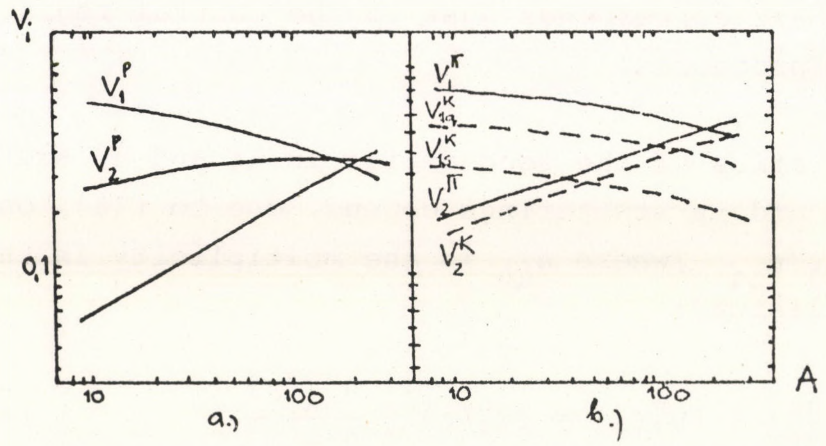


Fig. 22

The probability of absorbing a different number of incident quarks in hadron-nucleus interactions. The probabilities to absorb one, two or three quarks in a pA collision /a/. The quark absorption probabilities for pA /solid lines/ and KA /dashed lines/ interactions.

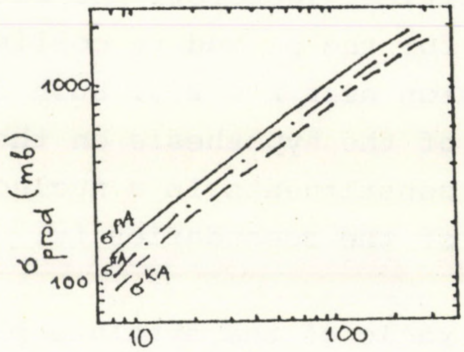


Fig. 23

The inelastic hadron-nucleus cross sections with the production of at least one secondary hadron as functions of A.

One sees that for light nuclei the most important is the process of Fig. 19a;however even for Be the probability of the process of Fig. 19b,with two interacting quarks, is not small ($\simeq$ 25 %). For $A>30$, the probability of the process of Fig.19a with two spectators decreases roughly as $A^{-1/3}$. For $A > 100$, the probabilities of all three processes are of the same order. As for the proton-nucleus cross section σ^{pA}_{prod} in Fig. 23, it increases as $A^{2/3}$ for $A > 30$, in full accordance with expectations.

As already said, the model with three spatially separated quarks enables one to express the multiplicity of a fast secondary baryon with $x \simeq \frac{2}{3}$ for proton-nucleus collisions, in terms of the similar quantity for pp interactions. Production of that fast baryon proceeds in both cases by picking up a newly made quark of the sea by the two non-interacting spectators. The upper vertices in Figs. 9b and 19a are the same, so they cancel in the ratio of the cross section or multiplicities. Therefore the ratio of the inclusive cross sections for the pA and pp collisions must not depend on x in a region near $x = 2/3$. Such independence of x represents a test of the hypothesis on the spatial separation of the three constituents in a nucleon, whatever the formation mechanism of the secondaries is.

The calculated ratio of the absolute proton yields, with $x \simeq 2/3$, from the nucleon and proton targets is

$$\frac{\frac{d^2\sigma}{dpd\Omega}(pA \to px)}{\frac{d^2\sigma}{dpd\Omega}(pp \to px)} = V_1^p(A)\,\frac{\sigma_{prod}^{pA}}{\sigma_{inel}^{pp}} \tag{42}$$

The results of our calculation are displayed in Fig. 24a for the Be, Al, Cu and Pb nuclei together with the data obtained at 19,2 GeV/c. Theory and experiment are consistent in the wide range $0.55 \leq x \leq 0.85$ where the experimental x-dependence of the ratio (42) is essentially flat. This indicates the absence of a substantial spread in momenta /with $\Delta x \gtrsim 1/6$/ of the constituents.

The experimental magnitudes of the V_1^p obtained from the data of ref. [67] by using eg. (42) are shown in Fig. 24b. to be consistent with our calculation.

The ratio of the meson yields near $x = \frac{1}{3}$ is obtained using the expression (7), (8):

$$\frac{\frac{1}{\sigma_{prod}^{pA}}\,\frac{d^2\sigma}{dpd\Omega}(pA \to Mx)}{\frac{1}{\sigma_{inel}^{pp}}\,\frac{d^2\sigma}{dpd\Omega}(pp \to Mx)} = V_1^p(A) + \frac{4}{5}V_2^p(A) \tag{43}$$

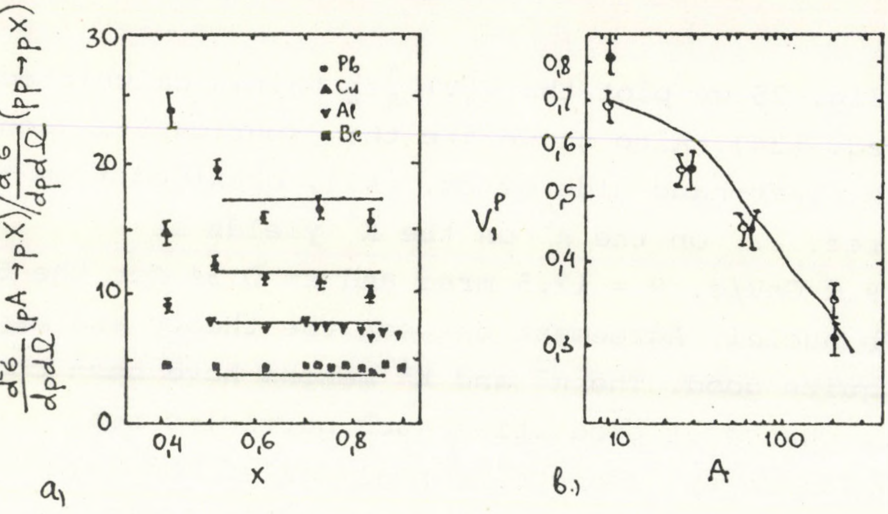


Fig. 24

a./ The cross-section ratios for nuclei and for hydrogen for p=19, 2GeV/c θ=12,5 mr as a function of x of the secondary . b./ Multiplicity of the secondary protons, averaged over the interval $0.52 \le x \le 0.85$ for p_o=19,2 GeV/c and θ= 12,5 mr /closed circles/, and for p_o=24 GeV/c and θ= 17 mr /open circles/, as a function of A.

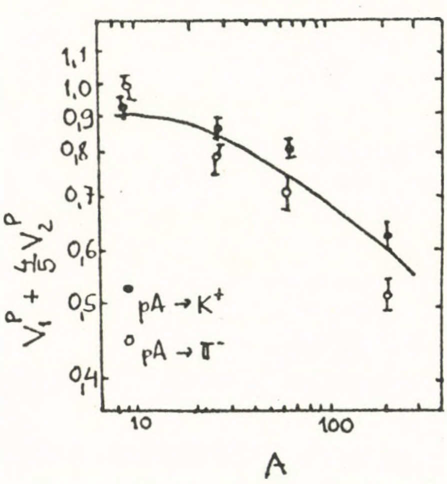


Fig. 25

Multiplicities of mesons in proton-nucleus interactions at p_o = 19,2 GeV/c and θ=12,5 mrad. The closed and open circles correspond to the production of K^+ and π^- respectively at x=0,34

In Fig. 25 we plot the $V_1^p V + \frac{4}{5} V_2^p$ values calculated according to eq. (24). Also shown are the experimental magnitudes of the left-hand side of eq. (43), obtained from the data of ref. on the π^- on the K^+ yields at $P_{lab} = 19.2$ GeV/c, $\theta = 12{,}5$ mrad and $x = 0.34$ for the Be, Al, Cu and Pb nuclei. Agreement between the theory and experiment is quite good. The π^- and K^+ mesons have been chosen since the chance of producting such particles near $x = \frac{1}{3}$ as resonance decay products is negligible. The opposite case of π^+ production at $x \simeq \frac{1}{3}$ is probably dominated just by the baryonic resonance decays, and therefore is not considered here.

When a pion strikes a nucleus or a proton, the ratio of inclusive spectra of the same fragments at $x \simeq \frac{1}{2}$ containing one of the pion quarks, must be

$$\frac{\dfrac{1}{\sigma^{\pi A}_{prod}} \dfrac{d^2\sigma}{dp d\Omega} (\pi^- A \rightarrow hx)}{\dfrac{1}{\sigma^{\pi p}_{inel}} \dfrac{d^2\sigma}{dp d\Omega} (\pi^- p \rightarrow hx)} \qquad h = \pi^-, \pi^o, p, n, \ldots \tag{44}$$

independently of the kind of the secondary. Therefore the single-hadron yield ratios, say π^-/K^-, π^-/p etc., must be the same /at $x \simeq \frac{1}{2}$/ for all nuclei in $\pi^- A$ interactions. The theoretical A dependence of V_1^π shown in Fig. 22b, can be approximated for $A > 60$, by

$$V_1^\pi(A) \simeq 1.75\ A^{-0.24}$$

If the incident particle is a kaon, the production of a fragment containing the strange quark is determined by the probability to absorb the non-strange quark, V_q^K. For instance, for the K^- beam the spectra of strange secondaries K^-, $\bar{K}^o$, Λ, Σ etc. must be in the ratio

$$\frac{\frac{1}{\sigma^{kA}_{prod}}\frac{d^2\sigma}{dpd\Omega}(K^-A \to h_s x)}{\frac{1}{\sigma^{kA}_{incl}}\frac{d^2\sigma}{dpd\Omega}(K^-p \to h_s x)} = V^K_q(A)(1+\frac{\sigma s}{\sigma_q}), \qquad h_s = {}^-,\Lambda,\Sigma,\ldots \tag{45}$$

According to Fig. 22b, $V^K_q(A) \simeq 0,82A^{-0.15}$ for $A > 30$.

On the other hand, the spectra ratio of the non-strange fragments like π^o, π^-, $\bar{N}$ etc., is determined by the probability to absorb the strange quark:

$$\frac{\frac{1}{\sigma^{kA}_{prod}}\frac{d^2\sigma}{dpd\Omega}(k^-A \to hx)}{\frac{1}{\sigma^{kp}_{incl}}\frac{d^2\sigma}{dpd\Omega}(k^-p \to hx)} = V^k_s(A)(1+\frac{\sigma_q}{\sigma_s}, \qquad h = \pi^-,\pi^o,\ \bar{p},\ \bar{n}, \tag{46}$$

For $A > 30$, $V^k_s(A) \simeq 0.6\ A^{-0.21}$. Therefore the ratios π^-/k^-, $\bar{p}/\Lambda$ etc., are predicted to decrease slightly in the K^-A collisions, as $V^k_s(A)/V^k_q(A) \sim A^{-0,06}$.
It means that, to some extent, a nucleus works like a filter detaining more non-trange quarks than the strange ones.

The comparison with the data given in [21] on the reactions $\pi^+A \to \pi^\pm, K^+$, p at $x \sim 1/2$, $pA \to p$ at $x \sim 2/3$ and $pA \to \pi^\pm$ at $x \sim 1/3$ with the calculated values [29] is presented on Figs. 26.,27. The experimental cross sections of $\pi^\pm$, $K^\pm$ and p production are given in [29] being parametrized as

$$\sigma(A) = \sigma_o \cdot A^\alpha \tag{47}$$

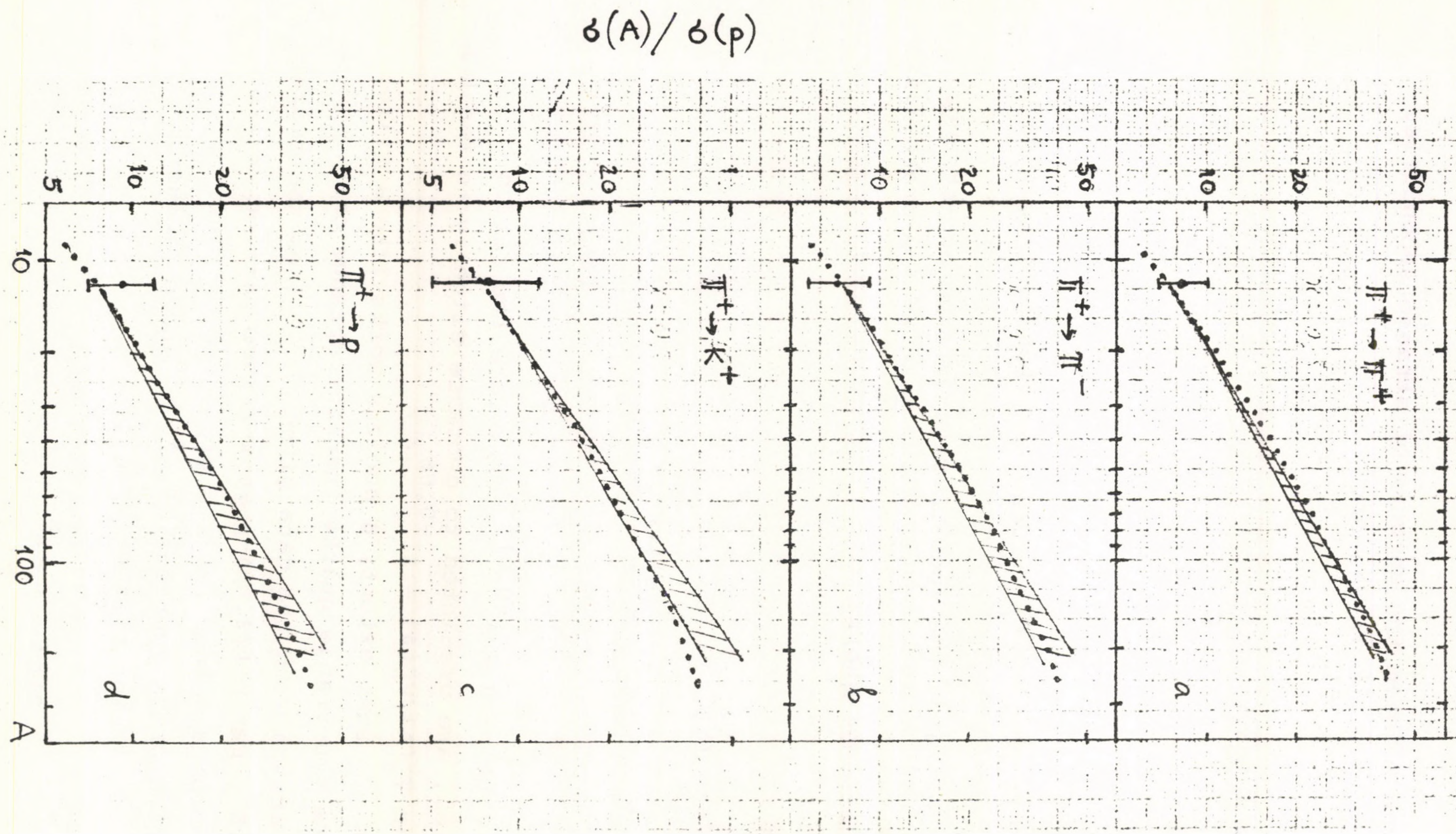


Fig. 26 Production cross section ratios of π^-(a), π^-(b),K^+(c) and p(d) in π^+A and π^+p collisions at 100 GeV/c as a function of A. Quark model predictions of ref.[2] are shown by dotted lines. Experimental point correspond to carbon to hydrogen cross section ratios, and the shaded areas correspond to experimental uncertainties of parameter /eq. 47/. The data were used at x= 0.5, p= 0.3 GeV/c.

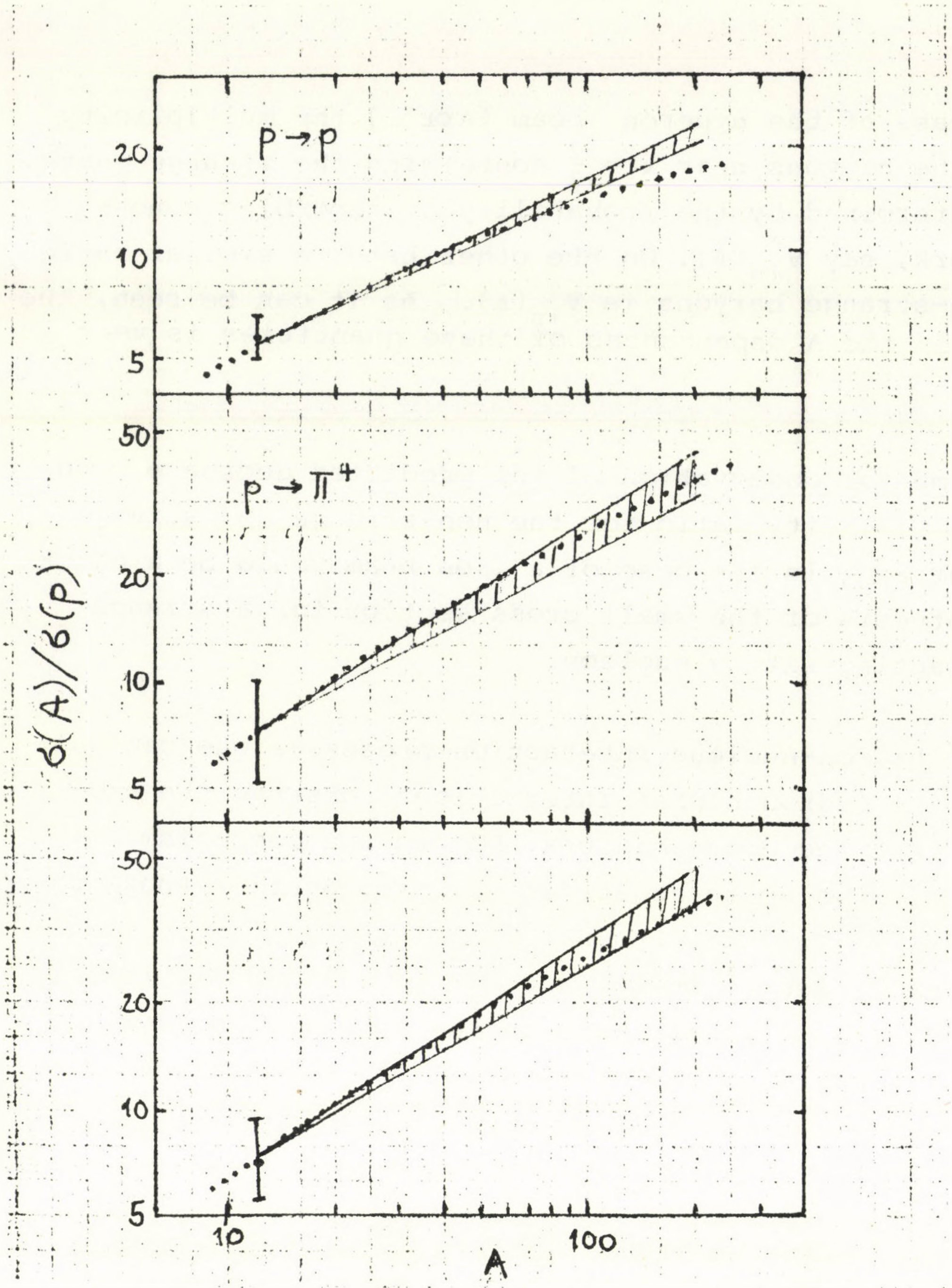


Fig. 27. Production cross section ratios of p(a), π^-(b) and π^-(c) in pA and pp collisions at 100 GeV/c as a function of A. All marks are the same as in fig 26. The data were used at x=0.7, p = 0.3 GeV/c for secondary protons and at x=0.3, p =0.3 GeV/c for assondary protons and at x=0.3, p = 0.3 GeV/c for pions.

In a case of the hyperon beam (Λ or Σ) the multiplicity ratio for the baryons near $x = \frac{2}{3}$ containing the strange quark, is again determined by the probability ov absorbing a non-strange quark, say $V^{\Lambda}_{1q}(A)$. On the other hand, a similar ratio for the non-strange baryons is $V^{\Lambda}_{1s}(A)$. As it can be seen, the difference in the A dependences of these quantities is very small.

Experimental observation of the predicted decrease with A of the multiplicity ratio for the non-strange and strange hadrons near $x = \frac{1}{2}$ in the case of a kaon beam would be a check of the hypothesis of the small cross section for a strange quark interacting with a nucleon.

In the hadron-nucleus interaction processes one can, similarly to the hadron-hadron interactions, observe the production of fast secondary hadrons. Due to the mechanism of the interaction we spoke about, we have to consider those cases, when one or two constituents of the incident baryon / $x \sim \frac{2}{3}$ and $x \sim \frac{1}{3}$, respectively/ and one constituent of the incident meson ($x \sim \frac{1}{2}$) participate in the interaction. For the baryon-nucleus collision we have, using the expressions (7) and (8):

$$\begin{aligned} & V_1^b(A)(q_i q_j + q,\bar{q}\text{-sea}) + V_2^b(A)(q_i + q,\bar{q}\text{-sea}) \rightarrow \\ & \rightarrow V_1^b\left(\frac{1}{2}B_{ij} + \frac{1}{12}(B_i + B_j) + \frac{5}{12}(M_i + M_j)\right) + \\ & + V_2^b\left(\frac{1}{3}B_i + \frac{2}{3}M_i\right) \end{aligned} \tag{48}$$

Besides, some distribution functions have to be introduced: $f_{ij}(x,p^2)$ for B_{ij}, $f_i(x,p^2)$ for B_i and $\varphi(x,p^2)$ for M_i. We consider $f_{uu} = f_{ud} = f_{dd}$; $\varphi_u = \varphi_d$, $f_u = f_d$.

Instead of (34) we have then

$$V_1^b(A)[\ \frac{1}{2}\ f_{ij}(x)B_{ij}\ +\ \frac{1}{12}\ (B_i f_i(x)\ +\ B_j f_j(x))\ +\ \frac{5}{12}(M_i\varphi_i(x)\ +$$

$$+\ M_j\varphi_j(x))]\ +\ V_2^b(A)(\frac{1}{3}\ f_i(x)B_i\ +\ \frac{2}{3}\ \varphi_i(x)M_i) \tag{49}$$

The meson-nucleus collision can be described as:

$$V_1^m\ (A)(q_i\ +\ q,\bar{q}\text{-sea})\ \rightarrow\ V_1^m(A)(\frac{1}{3}\ B_i\ \ +\ \frac{2}{3}M_i)\ \rightarrow$$

$$\rightarrow\ V_1^m(A)(\ \frac{1}{3}f_i(x)B_i\ +\ \frac{2}{3}\varphi_i(x)M_i) \tag{50}$$

Similarly to the hadron-hadron collision case, one can easily get the secondary particles produced in, pA, ΛA, ΣA, πA etc. processes. The results of the calculations and the comparison of the predictions with experiment will be given soon.

Concluding remarks

There are three groups of phenomena in which it is impossible to avoid the notion of constituent quarks. The fairly good agreement of the predictions of this approach and the experiment shows that the question of the existence of dressed quarks inside the hadrons has to be taken seriously. The spectator mechanism is proved by different theoretical and experimental results. The data on the production of secondary mesons support the assumption of quark combinatorics due to which secondary particles are produced in SU(6)-multiplets. It is quite natural, however, that the results of so simple calculations differ from the experiment by 10-15%, sometimes even more.

Further progress demand a better theoretical understanding and further detailed experimental proofs. We hope, that investigations in the not too far future will give the answers.

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$h_{M(0)}$		$L=0$						$h_{M(1)}$				$L=1$					
$J=0$	$J=1$	$\mu_h^0(u)$	$\mu_h^0(d)$	$\mu_h^0(s)$	μ_h^0	$\mu_h^0(\pi^+)$	$\mu_h^0(K^+)$	$J=0$	$J=1$		$J=2$	$\mu_h^1(u)$	$\mu_h^1(d)$	$\mu_h^1(s)$	μ_h^1	$\mu_h^1(\pi^+)$	$\mu_h^1(K^+)$
π^+	ρ^+	$\frac{n_f^J}{8}$	—	—	$\frac{n^J}{16}$	$\frac{2J+1}{4}$	—	δ^+	A_1^+	B^+	A_2^+	$\frac{n_f^J}{24}$	—	—	$\frac{n^J}{48}$	$\frac{2J+1}{12}$	—
π^0	ρ^0	$\frac{n_f^J}{16}$	$\frac{n_f^J}{16}$	—	$\frac{n^J}{16}$	—	—	δ^0	A_1^0	B^0	A_2^0	$\frac{n_f^J}{48}$	$\frac{n_f^J}{48}$	—	$\frac{n^J}{48}$	—	—
π^-	ρ^-	—	$\frac{n_f^J}{8}$	—	$\frac{n^J}{16}$	—	—	δ^-	A_1^-	B^-	A_2^-	—	$\frac{n_f^J}{24}$	—	$\frac{n^J}{48}$	—	—
K^+	K^{*+}	$\frac{n_f^J}{8}\lambda_f$	—	—	$\frac{n^J}{16}\lambda$	—	$\frac{2J+1}{4}$	$\varkappa^+$	Q_1^+	Q_2^+	K^{**+}	$\frac{n_f^J}{24}\lambda_f$	—	—	$\frac{n^J}{48}\lambda$	—	$\frac{2J+1}{12}$
K^0	K^{*0}	—	$\frac{n_f^J}{8}\lambda_f$	—	$\frac{n^J}{16}\lambda$	—	—	$\varkappa^0$	Q_1^0	Q_2^0	K^{**0}	—	$\frac{n_f^J}{24}\lambda_f$	—	$\frac{n^J}{48}\lambda$	—	—
$\bar{K}^0$	$\bar{K}^{*0}$	—	—	$\frac{n_f^J}{8}\lambda_f$	$\frac{n^J}{16}\lambda$	—	—	$\bar{\varkappa}^0$	$\bar{Q}_1^0$	$\bar{Q}_2^0$	$\bar{K}^{**0}$	—	—	$\frac{n_f^J}{24}\lambda_f$	$\frac{n^J}{48}\lambda$	—	—
K^-	K^{*-}	—	—	$\frac{n_f^J}{8}\lambda_f$	$\frac{n^J}{16}\lambda$	—	—	$\varkappa^-$	Q_1^-	Q_2^-	K^{**-}	—	—	$\frac{n_f^J}{24}\lambda_f$	$\frac{n^J}{48}\lambda$	—	—
	ω	$\frac{n_f^J}{16}$	$\frac{n_f^J}{16}$	—	$\frac{n^J}{16}$	—	—	ε	D	?	f	$\frac{n_f^J}{48}$	$\frac{n_f^J}{48}$	—	$\frac{n^J}{48}$	—	—
	ϕ	—	—	$\frac{n_f^J}{8}\lambda_f$	$\frac{n^J}{16}\lambda^2$	—	—	S	?	?	f'	—	—	$\frac{n_f^J}{48}\lambda_f$	$\frac{n^J}{48}\lambda^2$	—	—
η		$\frac{n_f^J}{48}$	$\frac{n_f^J}{48}$	$\frac{n_f^J}{24}\lambda_f$	$\frac{n^J}{48}(1+2\lambda^2)$	—	—										
η'		$\frac{n_f^J}{24}$	$\frac{n_f^J}{24}$	$\frac{n_f^J}{12}\lambda_f$	$\frac{n^J}{48}(2+\lambda^2)$	—	—										

Expansion of $M_{ij}(L)$, $M_i(L)$ and $M(L)$ in terms of real meson states. The normalisation coefficients n_f^J and n^J are: $n_f^J=(1+\lambda_f/2)^{-1}(2J+1)$, $n^J=(1+\lambda/2)^{-1}(2J+1)$. The mixing of singlet and octet states in η and η' is not taken into account $\eta=\eta_1$, $\eta'=\eta_8$. It is assumed, that the states ω, ε, D and f consist of only non-strange quarks while the states ϕ, s, f' only of strange ones.

Table 1

	B_{ij}								B_i			B	B^*_{ijk}			
h_B	$\beta_h(uu)$	$\beta_h(ud)_1$	$\beta_h(ud)_p$	$\beta_h(ud)_\Lambda$	$\beta_h(us)_1$	$\beta_h(us)_0$	$\beta_h(us)_\Sigma$	$\beta_h(ss)$	$\beta_h(u)$	$\beta_h(d)$	$\beta_h(s)$	β_h	$\beta_h(p)$	$\beta_h(\Lambda)$	$\beta_h(\Sigma^+)$	$\beta_h(\Xi^-)$
p	$\frac{1}{9}n_{qq}$	$\frac{1}{18}n_{qq}$	$\frac{41}{90}n_{qq}$	$\frac{1}{2}n_{qq}$					$\frac{2}{15}n_q$	$\frac{1}{15}n_q$		$\frac{1}{10}n_0$	$\frac{17}{27}$			
n		$\frac{1}{18}n_{qq}$	$\frac{41}{90}n_{qq}$	$\frac{1}{2}n_{qq}$					$\frac{1}{15}n_q$	$\frac{2}{15}n_q$		$\frac{1}{10}n_0$				
Λ			$\frac{3\lambda_f}{10}n_{qq}$	$\frac{\lambda_f}{3}n_{qq}$	$\frac{1}{10}n_{qs}$	$\frac{1}{10}n_{qs}$	$\frac{1}{10}n_{qs}$		$\frac{\lambda_f}{15}n_q$	$\frac{\lambda_f}{15}n_q$	$\frac{1}{10}n_s$	$\frac{\lambda}{10}n_0$		$\frac{2\xi+1}{2(2+\xi)}$		
Σ^+	$\frac{\lambda_f}{9}n_{qq}$				$\frac{1}{15}n_{qs}$	$\frac{1}{3}n_{qs}$	$\frac{41}{15}n_{qs}$		$\frac{2\lambda_f}{15}n_q$		$\frac{1}{10}n_s$	$\frac{\lambda}{10}n_0$			$\frac{42+9\xi}{27(2+\xi)}$	
Σ^0		$\frac{\lambda_f}{9}n_{qq}$	$\frac{\lambda_f}{90}n_{qq}$		$\frac{1}{30}n_{qs}$	$\frac{1}{6}n_{qs}$	$\frac{41}{150}n_{qs}$		$\frac{\lambda_f}{15}n_q$	$\frac{\lambda_f}{15}n_q$	$\frac{1}{10}n_s$	$\frac{\lambda}{10}n_0$		$\frac{1}{2(2+\xi)}$		
Σ^-										$\frac{2\lambda_f}{15}n_q$	$\frac{1}{10}n_s$	$\frac{\lambda}{10}n_0$				
Ξ^0					$\frac{\lambda_f}{15}n_{qs}$	$\frac{\lambda_f}{3}n_{qs}$	$\frac{41\lambda_f}{75}n_{qs}$	$\frac{1}{6}n_{ss}$	$\frac{1}{15}\lambda_f^2 n_q$		$\frac{\lambda_f}{5}n_s$	$\frac{\lambda^2}{10}n_0$				
Ξ^-								$\frac{1}{6}n_{ss}$		$\frac{\lambda_f^2}{15}n_q$	$\frac{\lambda_f}{3}n_s$	$\frac{\lambda^2}{10}n_0$				$\frac{9+42\xi}{27(2+\xi)}$
Δ^{++}	$\frac{2}{3}n_{qq}$								$\frac{2}{5}n_q$			$\frac{1}{5}n_0$				
Δ^+	$\frac{2}{9}n_{qq}$	$\frac{4}{9}n_{qq}$	$\frac{2}{45}n_{qq}$						$\frac{4}{15}n_q$	$\frac{2}{15}n_q$		$\frac{1}{5}n_0$	$\frac{10}{27}$			
Δ^0		$\frac{4}{9}n_{qq}$	$\frac{2}{45}n_{qq}$						$\frac{2}{15}n_q$	$\frac{4}{15}n_q$		$\frac{1}{5}n_0$				
Δ^-										$\frac{2}{5}n_q$		$\frac{1}{5}n_0$				
Σ^{*+}	$\frac{2\lambda_f}{9}n_{qq}$				$\frac{8}{15}n_{qs}$	$\frac{4}{15}n_{qs}$	$\frac{4}{75}n_{qs}$		$\frac{4\lambda_f}{15}n_q$		$\frac{1}{5}n_s$	$\frac{\lambda}{5}n_0$			$\frac{4+6\xi}{9(2+\xi)}$	
Σ^{*0}		$\frac{2\lambda_f}{9}n_{qq}$	$\frac{\lambda_f}{45}n_{qq}$		$\frac{4}{15}n_{qs}$	$\frac{2}{15}n_{qs}$	$\frac{2}{75}n_{qs}$		$\frac{2\lambda_f}{15}n_q$	$\frac{2\lambda_f}{15}n_q$	$\frac{1}{5}n_s$	$\frac{\lambda}{5}n_0$		$\frac{1}{2+\xi}$		
Σ^{*-}										$\frac{4\lambda_f}{15}n_q$	$\frac{1}{5}n_s$	$\frac{\lambda}{5}n_0$				
Ξ^{*0}					$\frac{8\lambda_f}{15}n_{qs}$	$\frac{4\lambda_f}{15}n_{qs}$	$\frac{4\lambda_f}{75}n_{qs}$	$\frac{1}{3}n_{ss}$	$\frac{2\lambda_f^2}{15}n_q$		$\frac{2\lambda_f}{5}n_s$	$\frac{\lambda^2}{5}n_0$				
Ξ^{*-}								$\frac{1}{3}n_{ss}$		$\frac{2\lambda_f^2}{15}n_q$	$\frac{2\lambda_f}{5}n_s$	$\frac{\lambda^2}{5}n_0$				$\frac{6+4\xi}{9(2+\xi)}$
Ω^-								$\lambda_f n_{ss}$			$\frac{3\lambda_f^2}{5}n_s$	$\frac{\lambda^3}{5}n_0$				

Expansions of $B^*_{ijk}(L)$, $B_{ij}(L)$ and $B(L)$ in terms of the real baryon states of the 56-plet.

Table 2

	$K^-p\to K^-$ /33/		$K^-p\to \bar{K}^0$ /33/		$K^-p\to K^+$ /33/		$K^-p\to K^0$ /33/		$pp\to K^0_s$ /32/	
	mbarn	%	mbarn	%	mbarn	%	mbarn	%	mbarn	%
$\sigma_{inclusive}$	8,3±1,5	100	8±0,5	100	1,6±0,5	100	1,6±0,5	100	7,4±0,5	100
vector mesons (V)	4±0,4	48	4,2±0,3	53	1,1±0,3	69	0,9±0,2	56	3,4±1	46
tensor mesons (T)	0,9±0,2	11	0,8±0,2	10	0,08±0,01	5	0,08±0,01	5	1,7±0,8	23
mesons with L=0 (4/3 V)	5,3±0,5	64	5,6±0,4	71	1,5±0,4	92	1,2±0,3	75	4,5±1,3	61
mesons with L=1 (12/5·T)	2,2±0,5	25	1,9±0,5	24	0,19±0,02	12	0,19±0,02	12	4,1±1,9	55
mesons with L=2 (estimation)		11		5		–		13		–
mesons with $s_{q\bar{q}}=1$ $(V+9/5\,T)$	5,6±0,4	68± ±5	5,6±0,3	71± ±4	1,24±0,3	72± ±19	1,04±0,2	65± ±13	6,5±1,8	87± ±24

Kaon production in K^-p collisions at 32 GeV/c [56] and in pp collisions at 405 GeV/c [55]. The inclusice cross-section is decreased in comparison with the data of [55] by the value of the cross-section of the diffraction dissociation.

Table 3.

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Kiadja a Központi Fizikai Kutató Intézet
Felelős kiadó: Szegő Károly
Szakmai lektor: Lovas István
Nyelvi lektor: Hraskó Péter
Példányszám: 460 Törzsszám: 81-431
Készült a KFKI sokszorositó üzemében
Felelős vezető: Nagy Károly
Budapest, 1981. julius hó