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Abstract

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A model for description of hadron elastic scattering and diffraction dissociation processes is proposed in this paper. The model is based on the solution of QFT single-time dynamic equations for the amplitudes of 2 → 2 and 2 → n transitions and also on the notions on the quark structure of hadrons.

Аннотация

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Модель процессов упругого рассеяния и дифракционной диссоциации адронов при высоких энергиях. Серпухов, 1983.

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В работе предложена модель для описания процессов упругого рассеяния и дифракционной диссоциации адронов. Модель основана на решении одновременных динамических уравнений в квантовой теории поля для амплитуд переходов 2 → 2 и 2 → n и представлениях о кварковой структуре адронов.

INTRODUCTION

In the present paper we consider a model of hadron elastic scattering and diffraction dissociation processes at high energies. The model is based on general properties of the scattering matrix such as unitarity and analyticity and physical features of composite hadron structure. Here we proceed from the three-dimensional dynamical equations in QFT for the amplitudes of the transitions $2\to 2$ and $2\to n$[1,2]. Hadron quark structure features are used for the construction of the kernels for these equations.

Hadron elastic scattering, as well as its diffractive dissociation processes including polarization in hyperon production processes are described. The model is appropriate for the description of the processes with any value of momentum transferred, and thus its validity range includes, e.g., from the QCD standpoint, both the momentum region, where perturbation theory over $\alpha_s(Q^2)$ can be justified and the momentum region where non-perturbative effects are essential ones, and QCD type calculations fail in selfconsistent description of these effects. Elastic scattering processes and diffraction dissociation-type processes have been analysed in various approaches, in which attempts were made for model account of the composite hadron structure in amplitude calculations, although limitations were imposed on the allowed region of momentum transferred or only elastic processes were considered. Such approaches may conventionally be divided into three groups. First of all, one should single out papers where quark counting rules are used to introduce and analyse cross-section dependence on the number of valence quarks in hadrons[3,6,25]. Another widespread class of models, including QCD type models also, brings hadron-hadron interaction to the level of consistuent interactions. Such models suppose introduction and parametrization of quark-quark interaction amplitudes and the composition law for hadron-hadron interaction[26]. Finally, there are models[27] related to introduction of several components usually two in hadron, each with a contribution to the total amplitude. For instance, when hadron is described as a two-component object with internal and peripheral parts, and the eikonal is the basic quantity of the model, one represents it as a sum of $\chi\to\chi_1+\chi_2$.

The model considered in the present paper should be related to the second one in the mentioned classes. In the first place the aim was to apply the model to consideration of the most interesting and actively studied processes.

1. THE BASIC EQUATIONS AND FORMULATION OF THE MODEL

The transition amplitudes of $p_1 + p_2 \to p_1' + p_2'$ and $p_1 + p_2 \to p_1' + k_2 + \ldots + k_n$, which we denote by $F(\vec{p}_1, \vec{p}_1')$ and $F(\vec{p}_1, \vec{p}_1', \vec{k}_2, \ldots, \vec{k}_n)$, respectively, are solutions of the following three-dimensional equations in quantum field theory:

$$F(\vec{p}_1,\vec{p}_1')=G(\vec{p}_1,\vec{p}_1') + \frac{i}{8\pi^2}\int\frac{d\vec{q}_1}{2q_1^o 2q_2^o}\delta(E_p-E_q)K(\vec{p}_1, \vec{q}_1)F(\vec{q}_1, \vec{p}_1'), \quad (1)$$

$$F(\vec{p}_1, \vec{p}_1', \vec{k}_2,\ldots, \vec{k}_n) = G(\vec{p}_1, \vec{p}_1', \vec{k}_2,\ldots, \vec{k}_n) +$$

$$+ \frac{i}{8\pi^2}\int\frac{d\vec{q}_1}{2q_1^o 2q_2^o}\delta(E_p-E_q)K(\vec{p}_1, \vec{q}_1)F(\vec{q}_1, \vec{p}_1', \vec{k}_2,\ldots, \vec{k}_n). \quad (2)$$

Here $E_p = \sqrt{\vec{p}_1^{\,2} + m_1^2} + \sqrt{\vec{p}_2^{\,2} + m_2^2}$ is the energy of a two-particle system, $K(\vec{p}_1, \vec{q}_1)$ is the real function which below the inelastic threshold coincides with the reaction matrix of quantum mechanics.

Though the set of the functions $\{G(\vec{p}_1,\vec{p}_1',\vec{k}_2,\ldots\vec{k}_n)\}$ is unknown as well as the set $\{F(\vec{p}_1,\vec{p}_1',\vec{k}_2,\ldots,\vec{k}_n)\}$, equations (1) and (2) follow from explicitly unitary representation of tne S-matrix:

$$S = (1 - \tfrac{i}{2}R)/(1 + \tfrac{i}{2}R), \quad (3)$$

and hence their use for calculation of the amplitudes of two-particle and multiparticle processes allows one to take into account non-trivial manifestations of the interaction dynamics related to the unitarity requirements.

The functions $G(\vec{p}_1, \vec{p}_1')$ and $G(\vec{p}_1,\vec{p}_1', \vec{k}_2,\ldots, \vec{k}_n)$ are defined by the equations: $(S = 1-iT)$

$$\langle p_1,p_2|R|p_1', p_2'\rangle + \frac{i}{2}\sum_{m>2}\langle p_1, p_2|R|m\rangle\langle m|T|p_1', p_2'\rangle =$$

$$= \frac{(2\pi)^4\,\delta^4(p_1+p_2-p_1'-p_2')}{\sqrt{2p_1^o 2p_2^o 2p_1'^o 2p_2'^o}}G(\vec{p}_1, \vec{p}_1'), \quad (4)$$

$$\langle p_1,p_2|R|p_1', k_2,\dots,k_n\rangle + \frac{i}{2}\sum_{m>2}\langle p_1,p_2|R|m\rangle\langle m|T|p_1',k_2,\dots,k_n\rangle =$$

$$= \frac{(2\pi)^4\delta^4(p_1+p_2-p_1'-\sum_{i=2}^{n} k_i)}{\sqrt{2p_1^o 2p_2^o 2p_1'^o \prod_{i=2}^{n} 2k_i^o}} G(\vec{p}_1, \vec{p}_1', \vec{k}_2,\dots, \vec{k}_n). \qquad (5)$$

Equation (2) is valid for the amplitudes of any multiparticle processes. For two-particle processes and diffraction dissociation processes one may introduce an impact parameter representation. For diffraction dissociation processes $h_1+h_2\to h_1+X_n$, when particle system X_n is formed from the decay process of an excited state h_2^* with mass $W^2=(p_1+p_2-p_1')^2$, hadron h_1 is a fast particle with the fraction of the initial hadron longitudinal mometum $x=2(p_1')_{\parallel}/\sqrt{s}\sim 1$. In that case impact parameter appears to be a conjugated quantity to the variable $p_\perp=(p_1')_\perp$ and has the same physical meaning as for two-particle scattering processes.

To describe an elastic scattering process, we use variables s and t; and variables s, $p_\perp$, W^2, ξ_n - for the diffraction dissociation processes. Here ξ_n stands for variables related to the momenta $\vec{k}_2,\dots, \vec{k}_n$.

In the impact parameter representation eqs. (1) and (2) are reduced to the algebraic equations. Performing the Fourier-Bessel transform we get

$$f(s, b) = g(s, b) + i\rho(s)k(s, b)f(s, b) \qquad (1')$$

$$f(s,b,W^2, \xi_n)=g(s,b,W^2, \xi_n)+i\rho(s)k(s,b)f(s,b,W^2, \xi_n). \qquad (2')$$

Solving eqs.(1') and (2') we find

$$f(s,b,W^2, \xi_n) = \frac{g(s,b,W^2, \xi_n)}{g(s,b)}f(s, b), \qquad (6)$$

We introduce now the functions $u(s,b)$ and $u(s,b,W^2, \xi_n)$ through relations

$$u(s,b) = g(s,b)\left[\frac{1-iu(s,b)}{1-ik(s,b)}\right], \qquad (7)$$

$$u(s,b,W^2, \xi_n) = g(s,b,W^2, \xi_n)\left[\frac{1-iu(s,b)}{1-ik(s,b)}\right]. \qquad (8)$$

Then for the amplitudes $f(s,b)$ and $f(s,b,W^2,\xi_n)$ one comes to the following forms:

$$f(s,b) = \frac{u(s,b)}{1-iu(s,b)}\,, \tag{9}$$

$$f(s,b,W^2,\xi_n) = \frac{u(s,b,W^2,\xi_n)}{1-iu(s,b)}\,. \tag{10}$$

As it follows from eq. (9), the function $u(s,b)$ is a generalized reaction matrix, whereas $u(s,b,W^2,\xi_n)$ may be interpreted as a generalized reaction matrix for the diffraction dissociation processes.

The function $u(s,b,W^2,\xi_n)$ describes diffraction excitation of h_2^*-state and its subsequent decay to a system of particles X_n. It is natural therefore to choose for $u(s,b,W^2,\xi_n)$ an expression in a factorized form

$$u(s,b,W^2,\xi_n) = \tilde{u}(s,b,W^2)\,\phi(s,b,W^2,\xi_n), \tag{11}$$

where the function $\tilde{u}(s,b,W^2)$ is to be related to the dynamics of the quasi-elastic scattering process, and the function $\phi(s,b,W^2,\xi_n)$ -to the decay $h_2^* \to X_n$.

To fix on the forms of the functions $u(s,b)$ and $u(s,b,W^2,\xi_n)$, which determine dynamics of elastic scattering and diffraction dissociation processes, respectively, we assume the following basic features for quark structure of hadrons. It is generally accepted to treat hadrons consisting of quarks, the number and flavours of valence quarks are determined by the composition over quantum numbers. However, examination of hadronic reaction dynamics allows to conclude, that hadron structure seems more complicated, than it could be thought from the simple rules of hadron composition over quantum numbers. For instance, it seems that proton structure is not exhausted by three valence quark state. It seems natural to consider that the wave function of a hadron has components which contain extra quark-antiquark pairs in addition to its basic configuration of valence quarks. Then, for example, in wave function of proton there will exist components $|uud, n(q\bar{q})\rangle$, $q = u,d,s,c,b,t$, $n \geqslant 1$ together with it basic state $|uud\rangle$. Similar states were first explored in ref.[4] for description of open charm production processes. The inference on probability of various Fock states of proton wave function is to be obtained from the analysis of the corresponding reactions where these components contribute. Corresponding estimations can also be obtained with the help of bag model[4], which, gives, for instance, the value of 0.02 for probability of the component $|uud, c\bar{c}\rangle$. The quarks from additional $q\bar{q}$-pairs are treated analogously to conventional valence constituents in hadron-hadron interaction process.

We assume that valence quarks and $q\bar{q}$-pairs are located at the center of a hadron. In hadron-hadron interaction valence quarks conventional and from $q\bar{q}$-pairs) are scattered in a quasi-independend way in some effective potential field V_{eff}. This field is generated at collision of two hadrons h_1 anf h_2. The scattering amplitude results from valence quarks scattering by the potential V_{eff}.

In the model hadron state with an extra $q\bar{q}$-pair is responsible for the production of particles containing the quark of q-flavour. For example, the $|uud, c\bar{c}\rangle$ state of proton is responsible for difraction production of Λ_c and D, D and $\bar{D}$-mesons as well as for production of the states with hidden charm.

We consider now a specific method for construction of the functions $u(s,b)$ and $u(s,b,W^2,\xi_n)$ with account of above stated assumption. In constructing expression for the function $u(s,b)$, which determines the elastic scattering dynamics, we should only take into consideration hadron states with conventional valence quarks that are independently scattered by the potential V_{eff}. The contribution of the states with extra $q\bar{q}$-pairs is negligible in that case. Therefore it is natural to choose an expression in the factorized form for the function $u(s,b)$ [5]:

$$u(s,b) = \prod_{i=1}^{n_1} f_i(s_i,b) \prod_{j=1}^{n_2} f_j(s_j,b), \qquad (12)$$

where f_i are the amplitudes of quark scattering on the potential V_{eff}, n_1 and n_2 are the numbers of conventional valence quarks in hadrons h_1 and h_2, respectively.

Assumption on factorization of separate quark amplitudes at hadron-hadron amplitude composition is a distinguished feature of factorized quark model, which was first proposed in ref. [6], where quark amplitudes dependent on variables s and θ were used.

The use of the same value of impact parameter for all quark scattering amplitudes corresponds to the picture when valence quarks are localized in the center of hadrons, i.e. they are at a small distance from each other. Therefore, assuming that quark interactions at small distances are rather weak, one can consider the quarks to be independent and use eq. (12). Factorization of quark interactions in the impact parameter representation can be thus argued by absence of quark interactions at small distances.

It is natural to consider that quarks in a moving hadron have the same velocity, and therefore their momenta $p_i \sim m_i$, where m_i is the mass of i-quark. Then $s_i = s\dfrac{m_i^2}{(\sum_i m_i n_i)^2}$, where n_i is the number of valence quarks of flavour i.

Consider now the method of reconstructing the expression for the function $u(s,b,W^2, \xi_n)$ that determines the dynamics of diffraction dissociation. The diffraction dissociation process is treated in the model as a result of excitation and subsequent decay of the states of hadrons with additional $q\bar{q}$-pairs. Extraction from the total amplitude of the component related to the projection, for instance, on the five-quark state of a proton $|uud, q\bar{q}>$, implies the account of the probability factor $\alpha_q(s, W^2)$. By analogy with the case of elastic scattering for the function $\tilde{u}(s,b,W^2)$ we take the following factorized form

$$\tilde{u}(s,b,W^2) = \alpha_q(s,W^2) \prod_{i=1}^{n_1} f_i(s_i,b) \{ \prod_{j=1}^{n_2} f_j(\tilde{s}_j,b) f_q(s_q,b) f_{\bar{q}}(s_{\bar{q}},b) \}, \quad (13)$$

where $f_{q,\bar{q}}$ are the scattering amplitudes in the field V_{eff} for quarks belonging to $q\bar{q}$-pair. It is natural to admit here that the quark scattering amplitude f_q coinsides with the antiquark scattering amplitude $f_{\bar{q}}$.

Analytical properties of the scattering amplitude over $\cos\theta$ allow one to derive definite conclusions on the explicit form of the generalized reaction matrix/5/ and taking into account eqs. (12) and (13), on the form of the function $f_i(s_i,b)$. Moreover, the natural assumption about the short range character of the potential V_{eff} results in the exponential decrease of the quark amplitude with the impact parameter/7/. The above mentioned arguments imply the following form for the quark amplitude:

$$f_i(s_i, b) = g_i(s_i)\exp(-m_i b). \quad (14)$$

We assume the relevant interaction radius is determined by the quark mass m_i. With account for the polynomial boundedness of the U-matrix and the increase of the total cross-sections for the functions $g_i(s_i)$ we take the expression $g_i(s_i) = g_i(s_i/s_o)^{\lambda}$, where $\lambda > 0$, $s_o = 1$ GeV^2.

Since the total cross-section of the inelastic interactions $\sigma_{inel}(s)$ increases logarithmically in the framework of this approach/8/ and $\sigma_{diff}(s) < \sigma_{inel}(s)$, where σ_{diff} is a diffraction dissociation cross-section, then as it will be confirmed in what follows one has to appoint an energy-independent form for the scattering amplitude for the quarks from $q\bar{q}$-pair.

Having constructed an expression for the functions $u(s,b)$ and $u(s,b,W^2, \xi_n)$ and solving equations (1) and (2), we get for the amplitudes of hadron elastic scattering and diffraction dissociation $F(s,t)$ and $F(s,p_\perp^2, W^2, \xi_n)$ the following representations:

$$F(s,t) = \frac{s}{\pi^2} \int_0^\infty b db J_0(b\sqrt{-t}) \frac{u(s,b)}{1-iu(s,b)}, \quad (15)$$

$$F(s,p_\perp, W^2, \xi_n) = \frac{s}{\pi^2} \int_0^\infty b db J_0(bp_\perp) \frac{u(s,b,W^2, \xi_n)}{1-iu(s,b)}, \qquad (16)$$

where the functions $u(s,b)$ and $u(s,b,W^2, \xi_n)$ are defined by eqs. (11)-(14).

In summary, the "modeling" of the approach discussed consists in exploitation of some quark features of hadron structure and hadron interactions under construction of generalized reaction matrix (U-matrix) which is a kernel of the dynamical equations for the amplitude:

1. Hadrons consist of quarks, and moreover in their wave functions with a non-zero probability the components containing extra $q\bar{q}$-pairs $q=u,d,s,c,b,t$ in addition to the conventional valence quarks are presented.

2. The process of hadron-hadron scattering is imaged as a result of quasi-independent scattering of the valence quarks and quarks from the $q\bar{q}$-pairs by some effective potential V_{eff}, generated in overlapping of hadron structures.

3. The generalized reaction matrix is given in the impact parameter representation by the product:

$$u(s,b) \sim \prod_i f_i(s_i,b),$$

where $f_i(s_i, b)$ is the scattering amplitude of i-flavour quark by the potential V_{eff}.

Note, the calculation procedure of the model is defined unambiguously as soon as the form of the quark amplitude $f_i(s_i,b)$ is fixed

4. Proceeding from the analytical properties of the scattering amplitude over $\cos\theta$ we choose the dependence of the functions $f_i(s_i,b)$ on the impact parameter b in the form;

$$f_i(s_i, b) \sim \exp(-m_i b).$$

The radius of the i-quark interaction in the potential field V_{eff} is defined by an inverse value of its mass m_i.

In Section 2 we shall consider the application of this model in description of elastic scattering process. Section 3 is devoted to the analysis of diffraction dissociation and in Section 4 we discuss a possible mechanism of the polarization appearence of inclusively produced hyperons.

2. HADRON ELASTIC SCATTERING

The dynamics of elastic scattering of hadrons is defined by the generalized reaction matrix $u(s,b)$. A respective expression follows from eqs. (12) and (14) and has the form:

$$u(s,b) = iC\left(\frac{s}{M^2}\right)^{\lambda N} \exp(-Mb), \qquad (17)$$

where

$$N \equiv N_{h_1h_2} = \sum_i n_i = n_1 + n_2, \quad M \equiv M_{h_1h_2} = \sum_i m_i n_i,$$

$$C \equiv C_{h_1h_2} = \prod_{i=1}^{n_1} g_i m_i^{2\lambda} \prod_{j=1}^{n_2} g_j m_j^{2\lambda}, \qquad (18)$$

n_i is a number of valence quarks with the mass m_i. For simplicity the imaginary unit is singled out. The quantities N, M and C depend on the quark composition of hadrons h_1 and h_2 and valence quark masses.

From representation (15) and eq. (17) for the elastic scattering amplitude F(s,t) we obtain the following expression for the interaction radius of hadrons h_1 and h_2:

$$R_{h_1h_2}(s) = \frac{N}{M}\left[\lambda \ln\frac{s}{M^2} + \frac{\ln C}{N}\right]. \qquad (19)$$

The total cross section can be then presented in the form:

$$\sigma_{tot}^{h_1h_2}(s) = 4\pi\left(\frac{N}{M}\right)^2 \left[\lambda^2 \ln^2 \frac{s}{M^2} + 2\lambda \frac{\ln C}{N} \ln\frac{s}{M^2} + \frac{\ln^2 C}{N^2}\right]. \qquad (20)$$

Thus, the rate of the total cross-sections growth and their ratios are defined by the quark content of the interacting hadrons, the dependence on the masses of the constituent quarks enters through the quantity M.

Note, that the ratio $\ln C/N$ weakly depends on the quantities m_i and n_i.

The total cross-section of inelastic interactions has the form:

$$\sigma_{inel}^{h_1h_2}(s) = 16\pi\frac{N}{M^2}\left[\lambda\ln\frac{s}{M^2} + \frac{\ln C}{N}\right]. \qquad (21)$$

Relations (19)-(21) provide us with a number of consequences that can be verified experimentally. Let u- and d-quark masses be equal $m_q = m_{u,d}$. Then we get for the ratio of the total cross-sections

$$\frac{\sigma_{tot}(Kp)}{\sigma_{tot}(\pi p)} = \left(\frac{5}{m_s/m_q + 4}\right)^2.$$

Using the experimental data for the cross-sections $\sigma_{tot}(Kp)$ and $\sigma_{tot}(\pi p)$ for the ratios m_s/m_q we get the value 1.5. Fixing this value of the ratio m_s/m_q we obtain for the total cross-sections of hyperon-proton and proton-proton interactions the following ratios:

$$\frac{\sigma_{tot}(\Sigma p)}{\sigma_{tot}(pp)} = \left(\frac{6}{m_s/m_q + 5}\right)^2 = 0.85 \qquad (22)$$

and

$$\frac{\sigma_{tot}(\Xi p)}{\sigma_{tot}(pp)} = \left(\frac{6}{2m_s/m_q + 4}\right)^2 = 0.73.$$

These values agree rather well with respective experimental data[9]. Thus, the inequality $m_s > m_q$ leads to the following chains:

$$\sigma_{tot}(Kp) < \sigma_{tot}(\pi p),$$

$$\sigma_{tot}(\Xi p) < \sigma_{tot}(\Sigma p) < \sigma_{tot}(pp),$$

and allows one to obtain a good agreement with the experimental data on the cross-section ratios. Assuming the ratio m_Q/m_q be known it is easy to calculate values for the total cross-sections in the case of particle with valence heavy quarks. E.g. taking $m_c/m_q = 10$[10], we find for the cross-section ratio:

$$\frac{\sigma_{tot}(\Lambda_c p)}{\sigma_{tot}(pp)} = \left(\frac{6}{m_c/m_q + 5}\right)^2 = 0.16.$$

Thus, $\sigma_{tot}(\Lambda_c p)$ at $p_L \simeq 100 \div 200$ GeV/c will be approximately equal to 6 mb. Analogous estimation of the cross-section $\sigma_{tot}(\psi p)$ gives the value 1 mb. Asymptotically, the total cross-section grows as $\ell n^2 s$. The factor before $\ell n^2 s$ term is connected with a number of valence quarks and their masses. For the total cross-section of πN and NN interactions we have

$$\sigma_{tot}^{(\infty)}(s) = \frac{4\pi\lambda^2}{m_q^2} \ell n^2 s.$$

It is natural that taking into account the analyticity of the amplitude over $\cos\theta$ at the choice of the functions $f_q(s,b)$ and the identification of the quark interaction radius with its inverse mass allows us to confront our expressions with the corresponding rigorous bounds in QFT and thereby to obtain for m_q the estimation involving the mass of π-meson[11]:

$$m_q \geqslant 2\, m_\pi \lambda .$$

The value of the parameter λ is connected with the value of the power law index of the differential cross section at large angles[5]. The choice of $\lambda = 1/2$ leads to a good agreement with the πN- and NN-scattering data[12]. The quark mass boundary takes then a simple form $m_q \geqslant m_\pi$.

Note, that the asymptotic equality of the cross-sections $\sigma_{tot}(\pi N)$ and $\sigma_{tot}(NN)$ is related to the assumption on the equality of u- and d-quarks masses. The account for the difference of the masses $m_u < m_d$ leads to the breaking of the asymptotic equality of the total πN- and NN-interaction cross-sections, as well as to, e.g., the inequality $\sigma_{tot}(pp) > \sigma_{tot}(pn)$, which is in agreement with the experimental data.

Consider now the scattering at the angles $\theta \neq 0$. Expressions (15) and (16) lead to the amplitudes $F(s,t)$ and $F(s, p_\perp, W^2, \xi_n)$, which satisfy the unitarity condition. The behaviour of the amplitudes $F(s,t)$ and $F(s, p_\perp, W^2, \xi_n)$ is governed by the singularities in the impact parameter complex plane[13,14], i.e. by the poles, whose position is defined by the solution of the equation $1-i\,u(s,\beta) = 0$ ($\beta = b^2$), and the branching point at $\beta = 0$. The solution of this equation has the form $b_k(s) = R_{h_1h_2}(s)+ick$, $k = \pm 1, \pm 3, \dots$, where $R_{h_1h_2}$ is hadron interaction range and c is constant.

For the elastic scattering amplitude at $t \neq 0$ the calculation of the contributions of the poles in the impact parameter complex plane[13], which is the governing one for t fixed and $s \to \infty$, leads to to the following expansion of the amplitude in the series over the function $\tau(\sqrt{-t})$ decreasing with the growth of $|t|$:

$$F_{h_1h_2}(s,t) = s \sum_{k=1}^{\infty} \tau^k(\sqrt{-t})\, \Phi_k(R_{h_1h_2}(s), \sqrt{-t}), \qquad (23)$$

where

$$\tau(\sqrt{-t}) = \exp(-\frac{2\pi}{M}\sqrt{-t}), \text{ and } \Phi_k(R_{h_1h_2}(s), \sqrt{-t})$$

is a slowly varying function of its arguments.

From representation (23) for the scattering amplitude at sufficiently large transferred momenta ($|t| \gg 1$, $|t|/s \ll 1$) we obtain the following expression:

$$F_{h_1h_2}(s,t) = s\exp(-\frac{2\pi}{M}\sqrt{-t})\, \Phi_1(R_{h_1h_2}(s), \sqrt{-t}), \qquad (24)$$

wherefrom for the slope parameter $B = \frac{d}{d\sqrt{-t}} \ln \frac{d\sigma}{dt}$ find

$$B = 2\pi / \sum_i m_i n_i . \qquad (25)$$

Dependence (24) describes well the behaviour of the angular distributions in the elastic pp-scattering in the region beyond the second diffractive maximum[15]. Fixation for the u- and d-quark masses of the value $m_q = 150 \div 200$ MeV[10] allows to obtain a good agreement with the experimental data on the differential cross-section of the elastic pp-scattering. The ratio of the slopes of the angular distributions for various processes can easily be calculated by formula (25).

Consider now the scattering on large angles ($s \to \infty$, t/s fixed) The scattering amplitude in this kinematic region is determined by the behaviour of generalized reaction matrix at small impact patameter. As it was shown in[13] the analytical properties of the scattering amplitude over the variable $\cos\theta$ imply the singularity of the function $u(s, \beta)$ at $\beta = 0$. This results in the power law for the decrease of the large angle scattering amplitude. The presence of this singularity may be connected in model with localization of valence quarks at hadron center. It is natural that in the region of large angles the scattering amplitude depends essentially on the number of valence quarks. Using relation (17) and the results of paper[13] we obtain for the large angle scattering amplitude the following expression:

$$F_{h_1h_2}(s,\ t) = -\frac{is}{\pi^2}\frac{M}{C_{h_1h_2}|t|^{3/2}}\left(\frac{M^2}{s}\right)^{\lambda N}|f_{h_1h_2}(s,0)|^2. \qquad (26)$$

For the differential cross-section in the region of large angle scattering we have:

$$\left.\frac{d\sigma_{h_1h_2}}{dt}\right|_{\substack{s,t\to\infty\\ t/s\text{ - fix.}}} = \frac{32\pi}{C^2_{h_1h_2}M^4}\left(\frac{M^2}{s}\right)^{2\lambda N+3}|f_{h_1h_2}(s,0)|^4(1-\cos\theta)^{-3}. \qquad (27)$$

In formulae (26)-(27) $f_{h_1h_2}(s,0)$ is the value of the elastic scattering amplitude at $b = 0$. Note, that at $s \to \infty$ $f_{h_1h_2}(s,0) \to 1$. The obtained expressions imply that the exponent of decrease of the large angle scattering cross-sections is determined by the total number of the valence quarks and the value of the parameter λ . A proper fix on of the parameter λ ($\lambda = 0.5$-0.6) allows one to obtain the values for the exponent that agree with the values obtained by measuring the cross-sections of different processes of hadron elastic scattering at large angles[5].

In conclusion to this Section we mark, that the introduction of the non-zero phases of the quark scattering amplitudes allows one to describe the smooth behaviour of the angular distributions in the elastic pp-scattering in the momentum transferred region beyond the second diffraction maximum, as well as to explain the energy dependence of the spin-spin correlation parameters in the elastic pp-scattering at the angles close to 90^o[16].

3. DIFFRACTION DISSOCIATION PROCESSES

These processes are treated in the model as a result of excitation and the subsequent decay of the states containing extra $q\bar{q}$-pairs. The diffractive dissociation amplitude is defined by expression (16). The function $\tilde{u}(s,b,W^2)$, describing quasi-elastic scattering, may be represented with account for eqs. (12)-(14) in the form:

$$\tilde{u}(s,b,W^2) = a_q(s, W^2) f_q^2(b) u(s,b), \qquad (28)$$

the kinematic factors are included into the factor $a_q(s,W^2)$. Assuming reasonably that the function $\phi(s,b,W^2,\xi_n)$ rather weakly depends on the impact parameter at small b, expression (16) for the amplitude $F(s, p_\perp, W^2, \xi_n)$ may be presented in the following form:

$$F(s, p_\perp, W^2, \xi_n) = \frac{s a_q(s,W^2)}{\pi^2} \phi_o(s, W^2, \xi_n) \times$$

$$\times \int_o^\infty bdb\, f_q^2(b) \frac{u(s,b)}{1-iu(s,b)} J_o(bp_\perp), \qquad (29)$$

where $\phi_o(s, W^2, \xi_n) = \phi(s, b=0, W^2, \xi_n)$.

To calculate the amplitude in the kinematical region $p_\perp \neq 0$, $2p_\perp/\sqrt{s} \ll 1$ we use the method[13] based on the analysis of the singularities in the impact parameter complex plane. As a result we obtain the expansion of the function $F(s, p_\perp, W^2, \xi_n)$ over the parameter $r(p_\perp)$ decreasing with the growth of $p_\perp$:

$$F(s,p_\perp,W^2,\xi_n) \simeq a_q(s,W^2) g_q^2 s\, \phi_o(s, W^2, \xi_n) \exp\{-2m_q R_{h_1h_2}(s)\} \times$$

$$\times \sum_{k=1}^\infty r^k(p_\perp)\, \Phi_k(R_{h_1h_2}(s), p_\perp). \qquad (30)$$

Here and henceforth we denote the mass of the quarks from a $q\bar{q}$-pair by m_q (or m_Q), For sufficiently large values of $p_\perp$ in expansion (30) one can retain leading term only:

$$\frac{d\sigma_q}{dW^2 dp_\perp^2} \simeq 4\pi^5 \rho_q(s,W^2) e^{-\frac{4R_{h_1h_2}(s)}{R_q}} r^2(p_\perp) \,|\,\Phi_1(R_{h_1h_2}(s), p_\perp)|^2, \qquad (31)$$

with the notation $\rho_q(s, W^2) = g_q^4 a_q^2(s, W^2) \sum_{n\geq 3} \int d\Gamma_n\, |\phi_o(s,W^2,\xi_n)|^2$.

Here $d\Gamma_n$ is a phase space element written in terms of the variables ξ_n and $R_q = 1/m_q$ the radius of the interaction leading to diffractive dissociation. It follows from relation (31) that the cross-section $d\sigma_q/dW^2dp_\perp^2$ is suppressed as compared to the elastic scattering cross section by the factor $\exp[-4R_{h_1h_2}(s)/R_q]$, which has a clear geometrical meaning.

Proceed to the region of small transverse momenta. Specifically, we discuss here the diffractive production of the states containing heavy quarks Q. The interest in such production processes was inspired by discovery of a large value for cross-sections of the open charm particle production at ISR energies/17/, and indication that Λ_c are mainly produced diffractively, was obtained. In the region $p_\perp \gtrsim m_Q$ ($Q = c, b, t$) for the cross-section of the production of particles with heavy quarks we can directly use eq. (31).

From representation (29) in the kinematic region $p_\perp < m_Q$ we obtain

$$\frac{d\sigma_Q}{dW^2dp_\perp^2} \simeq \frac{\pi\rho_Q(s,W^2)}{4m_Q^2}|f_{el}(s,0)|^2 \frac{J_1^2(p_\perp/2m_Q)}{(p_\perp/2m_Q)^2}. \tag{32}$$

The factor $|f_{el}(s,0)|^2$ is related to the hadron matter density at b=0.

Relation (32) implies that the dependence of the cross-section on $p_\perp$ is rather weak at $p_\perp \lessdot m_Q$. The width of the distribution is determined by the quark mass. Thus, in the case of production of particles with heavier quarks the hadron h_1 distribution over the transverse momentum should be wider. Parallel to this the particles from the set X_n must have a similar distribution on the transverse component of their momenta, since the system of the particles X_n is produced diffractively. An explicit form of the function $\rho_Q(s,W^2)$ is unknown, and to define its dependence on W^2 we make use of empirical formula $d\sigma/dW^2 \sim 1/W^2$ obtained from analysis of the experimental data on diffractive dissociation/18/. Thus, taking $\rho_Q(s, W^2) = \rho_Q(s)/W^2$ we obtain for the total cross-section of the heavy particle production.

$$\sigma_Q(s) \cong \frac{\pi\rho_Q(s)}{16m_Q^2}|f_{el}(s,0)|^2 \int_{W_o^2}^{W_1^2} \frac{dW^2}{W^2} = \frac{\pi\rho_Q(s)}{16\ m_Q^2}|f_{el}(s,0)|^2 \ln\frac{s(1-x_1)}{W_o^2}, \tag{33}$$

where x_1 is the lower fractional momentum cut on the recoiling hadron h_1 ($x_1 \simeq 0.8$-0.9). The quantity W_o^2 is the threshold value for the associated production of a pair of hadrons containing Q-flavour quarks.

The asymptotic behaviour of $\sigma_Q(s)$ has the form:

$$\sigma_Q^{(\infty)}(s) \simeq \frac{\pi\rho_Q}{16m_Q^2} \ln s, \tag{34}$$

hence with account for the unitary condition and (21) we must thereby take $\rho_Q(s) \to \rho_Q = \text{const}$ at $s \to \infty$. This is guaranteed by the energy independence of the function $g_q(s)$ that has been previously taken into account in Section 1. At finite energies the growth of $\sigma_Q(s)$ may be more rapid due to the rise of the factor $\rho_Q(s)$ with the energy. This factor accounts for the threshold effects also.

Consequently, the growth of the cross-sections of the diffractive production of hadrons containing heavy quarks, is an asymptotically logarithmic one. The coefficient before the logarithm is defined, by the quantity m_Q^{-2}. Such a dependence of the cross-section on the quark mass is a consequence of the fact, that the quark interaction range is defined by the inverse value of its mass and results from the dynamical content of this approach. In contrast to [4] it is not a trivial consequence of the probability dependence of hadron states with heavy $Q\bar{Q}$-pairs on the mass of the Q-quark. The dependence of the cross-section on the quark mass $\sigma_Q \sim m_Q^{-2}$ has been previously obtained on the basis of the perturbation theory, which however, restricts in some way the kinematic region of applicability of these results. On the other hand, a number of models lead to the exponential decrease of the cross-section with m_Q. This explains the interest in the investigation of such a dependence in different approaches and to the analysis of the dynamical assumptions that are the basis of the dependence of the cross-section of the new heavy hadron production on the mass of the corresponding valence quark.

Using (34), we obtain for the ratio of the cross-sections of the production of particles with Q_1- and Q_2 - flavour quarks:

$$\frac{\sigma_{Q_1}(s)}{\sigma_{Q_2}(s)} \propto \left(\frac{m_{Q_2}}{m_{Q_1}}\right)^2 . \tag{35}$$

In the framework of the given model and at $s \to \infty$ we can obtain the following expression for the ratio of the cross-sections

$$\frac{\sigma_Q^{(\infty)}(s)}{\sigma_{tot}^{(\infty)}(s)} = \frac{\rho_Q}{64\,\lambda^2 N^2}\left(\frac{M}{m_Q}\right)^2 \frac{1}{\ln s} . \tag{36}$$

Relations (34)-(36) determine a relative order of the cross-section of the diffractive production of particles with heavy quarks, via the parameter m_Q^2. Then formula (35) implies that the cross-section Λ_b makes up about 10%, and the production cross-section Λ_t – $0.5 \div 1\%$ of the cross-section of the open charm particle production. It can easily be seen from eq. (36) that the asymptotical behaviour of the quantity $\sigma_Q^{(\infty)}(s)/\sigma_{tot}^{(\infty)}(s)$ is determined by the ratio of the light quark mass and the heavy quark mass.

The application of (33) for $\sigma_Q(s)$ and the experimental data in the production of charmed particles in pp-collisions at ISR energies allows one to calculate the energy that is necessary for the relation $\sigma_c(s)/\sigma_{tot}(s)$ to be maximum. This energy is E_{lab} = 70 GeV.

In fig. 1 the comparison of eq. (33) with the data available on the cross-section of the production of particles with open charm is shown. It is assumed that $|f_{el}(s,0)|^2=1$, and the quantity ρ_Q is energy independent. The calculations show that at the energies 50-70 GeV the cross-section of the open charm particle production will be 10÷30 μb. Fig. 2 shows the production cross-sections for Λ_b and Λ_t, obtained by formula (33) for the values of the masses of b- and t-quarks 5 and 20 GeV, respectively. In this, the quantities ρ_b and ρ_t were put to be equal to $\rho_b = \rho_t = \rho_c = 2.3$.

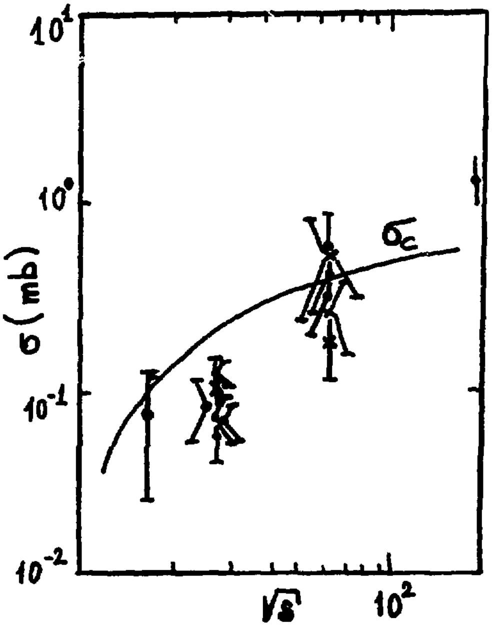

Fig. 1.

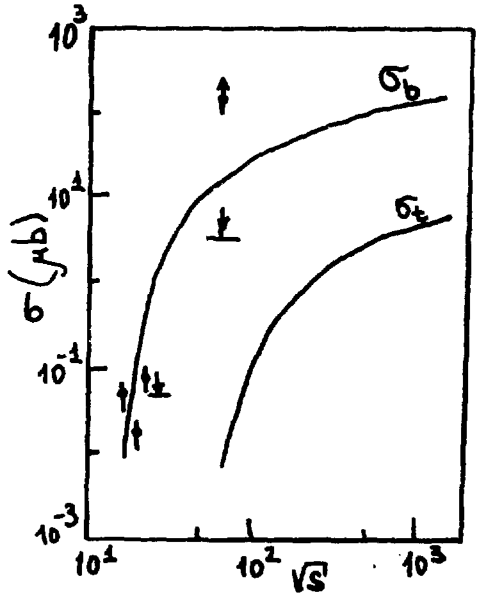

Fig. 2.

In paper [18] a possibility to separate kinematically the production processes of new heavy hadrons, in particular, to extract the processes of production of t-quark states from more intensive production states with c- and b-quarks, was discussed. It follows from (32), that a wider distribution of the hadron h_1 over the momentum $p_\perp$ in the reaction $h_1+h_2 \to h_1+X$ is related to the diffractive production of the states of X that contain heavier quarks. The width of the corresponding distribution in $p_\perp$ turns out to be of order m_Q. This conclusion may be useful for the analysis of the experimental data on the $p_\perp$-dependence of the cross-sections of the production of particles with open flavours. Note, that for $p_\perp \gtrsim m_Q$ the distributions are universal in $p_\perp$. We can consider similarly the processes of the double diffractive dissociation $h_1+h_2 \to X_1+X_2$. The cross-section $\sigma_{Q_1Q_2}(s)$ of the production of particles with the heavy quarks Q_1 and Q_2 in these processes is $\sigma_{Q_1Q_2}(s) \sim (\alpha_{Q_1}\alpha_{Q_2})^2/(m_{Q_1}+m_{Q_2})^2$. We also note that in the model parallel to the diffractive production of the open flavour particles, such as Λ_c and D, the hidden flavour particles, e.q. J/ψ, may be produced diffractively. The experimental measurements show however, that the production cross-section J/ψ is smaller than Λ_c by a factor of some orders of magnitude and, that in contrast to Λ_c, J/ψ is produced mainly in the central region. Consequently, the processes of the diffractive production of J/ψ must be suppressed, as compared to those of Λ_c.

Such a suppression in the model is caused by several reasons. In the first place note that the $c\bar{c}$-pair, has a vacuum quantum numbers and, hence, is in the $^3P_0$ state. That is why the χ-state alone can be generated diffractively, while the state J/ψ is a result of the radiation decay $\chi \to \gamma J/\psi$. This mechanism of J/ψ generation gives the suppression factor of $8\cdot10^{-3}$. Note also, that the probability of generating a particle with the $c\bar{c}$-pair from the five-quark state $|uud, c\bar{c}\rangle$ is one-fourth of the probability of generating particles containing only one quark from the $c\bar{c}$-pair. With account for this we obtain for the ratio of the diffractive production cross-sections: $\sigma_{J/\psi}/\sigma_c < 2\cdot10^{-3}$. Besides, there exist other reasons for the suppression of the diffractive J/ψ production cross-section. Thus, e.g., the χ-state will be generated if the value of the invariant mass of the $c\bar{c}$-pair is below the threshold of the production of the $D\bar{D}$-meson pair. But if the invariant mass $M_{c\bar{c}}$ is greater than the threshold of the $D\bar{D}$-pair production, then the open charm particles are mainly produced. The account for this circumstance, as well as for the colour degrees of freedom leads[4] to the factor of $\sim10^{-3}$. As a result, the diffractive J/ψ production cross-section is by a factor of about six orders of magnitude smaller than that of Λ_c and D. If we take the diffractive production cross-section of the D-meson at the FNAL energies[20] to be $\sigma(\pi N \to DX) \simeq 20\,\mu b$, then the value of the diffractive J/ψ production cross-section in the reaction will be at these energies $\sim$40 pb.

The above given dynamic reasons leading to the suppression of the diffractive J/ψ production, are equally valid for the production of particles with quarks of other flavours. Thus, e.g., the

ϕ-meson diffractive production should be strongly suppressed in comparison with the K-meson diffractive production which agrees with the experimental data on the pp-collisions.

In a number of papers/20/ arguments are given against the necessity to introduce the state with extra $q\bar{q}$-pairs. These arguments seem to be not quite convincing since the calculations on the basis of perturbation theory in the framework of QCD do not allow one to obtain the correct value of the production cross-section, e.g., Λ_c, and lead to the central distribution in x_F, which is in contradiction to the experimental data. This discrepancy with the experimental data is, probably, caused by neglecting non-perturbative effects. In connection with this, attempts to improve the situation by treating the Λ_c production as a two-step process/20/ seems to be artificial because of the small probability for these processes. Moreover, such an approach does not agree with a simple kinematic conclusion that a greater fraction of the momentum Λ_c is carried by the c-quark.

4. POLARIZATION IN INCLUSIVE HYPERON PRODUCTION

Measurements of the Λ-hyperon polarization in the reaction $pp \to \Lambda X$ in the case of the unpolarized target and beam brought us to an interesting result by discovering large polarization of Λ/21/ which is negative and increases in its modulus with the growth of $p_\perp$, reaching 20÷30% at $p_\perp$ = 1÷1.5 GeV/c. Any noticeable energy dependence up to the ISR energies was not found. The results of these, as well as of a number of experiments, made us appreciate the importance of a spin in hadron processes at high energies. In order to describe qualitatively large polarization in the inclusive production of Λ, and also of other hyperons (Σ, Ξ) several models were proposed, based on application of hadron quark structure. The point to be remarked is that some models require $P_\Lambda = 0$, which makes this experiment crucial for them. We calculate in this section the polarization asymmetry of diffractively produced Λ-hyperons. Since the initial protons are unpolarized, we assume that the constituent quarks should also be unpolarized. Note, that the spin and polarization of the Λ-hyperons are determined by the corresponding quantities of the s-quark.

The Λ-particle polarization in the model under consideration is caused by the mechanism of polarization of the $s\bar{s}$-quark pair which is in the state $^3P_0$ in the component $|uud, s\bar{s}\rangle$ of the proton wave function. This mechanism is determined by the structure of the effective potential V_{eff}, which indeed may have a rather complex dependence on the spin degrees of freedom. A possible form of its dependence on the spin and the orbital mometum was discussed in/16/.

Suppose that the structure of the potential is such that it extracts the state with a definite value of the spin projection $S_n=1$ of the $s\bar{s}$-pair on the direction of the normal to the scattering plane, whereas the states with $S_n=0,-1$ are being suppressed. In the scattering process the s and $\bar{s}$ quarks may either preserve the direction of their spins or change it for the opposite one. At this the $s\bar{s}$-pair retains its $^3P_0$ state.

The asymmetry arising here is the origin of the polarization of Λ. Note, that a possible reason of extracting the state with the fixed value of S_n is the non-zero orbital momentum of the $s\bar{s}$-pair. The absence of polarization in the inclusively produced protons in the reaction $pp \to pX$, when the u- and d-quarks are in the S-state, probably is in favour of this explanation.

Denote the non-spin flip amplitude of the $s(\bar{s})$-quark as $f_{so}(b)$, and the spin flip amplitude as $f_{sf}(b)$. Then the amplitude $F_\uparrow$ of the Λ production with the direction of spin in the positive one is determined by the amplitude $f_{so}(b)$, and the amplitude $F_\downarrow$ of the Λ production with the direction of spin in the negative one by the amplitude $f_{sf}(b)$. In the diffractive Λ-hyperon production $p_\perp^\Lambda \sim p_\perp$ and the Λ polarization in the process $p \xrightarrow{p} \Lambda$ are defined by the expression

$$P_\Lambda(p_\perp) = \frac{|F_\uparrow|^2 - |F_\downarrow|^2}{|F_\uparrow|^2 + |F_\downarrow|^2}. \tag{37}$$

With account for the aforementioned we find for the functions $u_{\uparrow(\downarrow)}(s,b,W^2,\xi_n)$

$$u_{\uparrow(\downarrow)}(s,b,W^2,\xi_n) = \alpha_s(s,W^2) f^2_{so(f)}(b) u(s,b)\, \phi(s,b,W^2,\xi_n), \tag{38}$$

and for the amplitudes $F_\uparrow$ and $F_\downarrow$, respectively,

$$F_{\uparrow(\downarrow)}(s,p_\perp,W^2,\xi_n) = \frac{s\alpha_s(s,W^2)}{\pi^2} \phi_o(s,W^2,\xi_n) \times$$

$$\times \int_0^\infty bdb f^2_{so(f)}(b) \frac{u(s,b)}{1-iu(s,b)} J_o(bp_\perp). \tag{39}$$

It is natural to use for the non-spin-flip amplitude the expression corresponding to the spinless case:

$$f_{so}(b) = g_o e^{-m_s b}. \tag{40}$$

The spin-flip amplitude is more central [16]. Taking this into account, we choose for it the expression

$$f_{sf}(b) = g_f e^{-\varkappa m_s b} \qquad (\varkappa > 1). \tag{41}$$

From the condition P_Λ $(p_\perp = 0) = 0$ for the parameters g_f, g_o and $\varkappa$ we obtain the relation $g_f = \varkappa g_o$.

For the region of small $p_\perp$ we obtain from (39):

$$F_\uparrow(s,p_\perp,W^2,\xi_n) \simeq \frac{s\alpha_s(s,W^2)g_o^2}{m_s^2}\phi_o(s,W^2,\xi_n)f_{el}(s,0)\frac{J_1(p_\perp/2m_s)}{(p_\perp/2m_s)}, \quad (42)$$

$$F_\downarrow(s,p_\perp,W^2,\xi_n) \simeq \frac{s\alpha_s(s,W^2)g_f^2}{\varkappa^2 m_s^2}\phi_o(s,W^2,\xi_n)f_{el}(s,0)\frac{J_1(p_\perp/2\varkappa m_s)}{(p_\perp/2\varkappa m_s)}.$$

Using for the Bessel function the representation $J_1(x) \simeq \frac{x}{2}e^{-x^2/8}$

which is valid at $0 \leqslant x < 1$, we get the following expression for the polarization

$$P_\Lambda(p_\perp) \simeq -\mathrm{th}[\frac{1}{8}(\frac{p_\perp}{2m_s})^2(1 - \frac{1}{\varkappa^2})]. \quad (43)$$

From here for the region of small $p_\perp$ we find:

$$P_\Lambda(p_\perp) \simeq -\frac{1}{8}\left(\frac{p_\perp}{2m_s}\right)^2(1 - \frac{1}{\varkappa^2}).$$

Note first, that the polarization parameter is a function of the ratio $p_\perp/2m_s$. The factors $\alpha_s(s,W^2)$ and $\phi_o(s,W^2,\xi_n)$ on which the amlitudes essentially depend are absent in the expression for $P_\Lambda(p_\perp)$.

It follows from (43) that in the region of rather small $p_\perp$, $P_\Lambda(p_\perp) < 0$, since $\varkappa > 1$. Note, that the Λ-hyperons polarization is energy-independent. The choice of the parameter $\varkappa = 3$ allows to obtain a satisfactory description of the experimental data on the parameter $P_\Lambda(p_\perp)$ at small $p_\perp$ /24/.

This discussion is simply be generalized for the case of the particles Λ_Q, with the heavy quark Q=c,b,t instead of the s-quark. Thus, for the polarization of Λ_Q at $p_\perp \leqslant m_Q$ we have

$$P_{\Lambda_Q}(p_\perp) \simeq -\mathrm{th}[\frac{1}{8}(\frac{p_\perp}{2m_Q})^2(1 - \frac{1}{\varkappa^2})]. \quad (44)$$

The polarization of, for example, Λ_c at $p_\perp^{\Lambda_c} \simeq 1 \div 1.5$ GeV/c must be, thereby, on the level of 1%. It is clear that the polarization

$P_{\Lambda_Q}(p_\perp)$ will be considerable at $p_\perp^{\Lambda_Q} = \frac{m_Q}{m_s} p_\perp^{\Lambda}$. Note also, that the polarization parameter should be negative at $p_\perp \leqslant m_Q$.

In the region of large $p_\perp^{\Lambda_Q}$ it is possible to obtain the expression for the polarization of Λ_Q using the method of /13/ for calculation of the contributions of singularities in the impact parameter plane. At sufficiently large $p_\perp$ polarization may oscillate and be non-trivially energy dependent.

Note that the expressions describing polarization imply a possibility for the polarization parameter to reach 1. This is connected with our initial assumption that the structure of the potential V_{eff} is such that the states with $S_n = 0,-1$ are completely suppressed. If these states are not completely suppressed, then the maximum value of the polarization is $|P_\Lambda|_{max} < 1$ and it may reach the value $1-\omega(S_n=0)-2\omega(S_n=-1)$, where $\omega(S_n=0)$ and $\omega(S_n=-1)$ being the probability of the states with $S_n=0, -1$.

In conclusion of this Section we give qualitative results on the behaviour of polarization in the production of $\Sigma^{\overset{+}{o}}$ and Ξ-hyperons in pp-collisions. Since the probability of the state in which the s-quark spin is antiparallel to the spin $\Sigma^{\overset{+}{o}}$, is twice as much as the probability of the state with parallel spins, then the amplitude $F_\uparrow(\Sigma)$ will be contributed mainly by spin flip amplitude f_{sf}, and the amplitude $F_\downarrow(\Sigma)$ - by the non-spin-flip amplitude of the s-quark f_{so}. The polarization sign $\Sigma^{\overset{+}{o}}$ must be, thereby, opposite to that of Λ. In the case of the Ξ-production we assume that the proton has two extra $s\bar{s}$-pairs. Surely, the probability of such a state is small.

However, the polarization of Ξ may be essential, since the expression for $P_\Xi(p_\perp)$ does not involve the corresponding probability factor $a_s^2(s, W^2)$. The spins of the two s-quarks in Ξ are parallel, and as a result the sign of the polarization of Ξ will coincide with the sign of the polarization of Λ in the reaction $pp \to \Lambda X$. The $\bar{\Lambda}$, s polarization in the $\bar{p}p \to \bar{\Lambda} X$ process must be equal to the Λ's polarization in the $pp \to \Lambda X$ process.

Note, that in the model considered the value of the polarization of different particles does not depend on the SU(6) group Clebsh-Gordon coefficients, the latter ones being similarly included into both amplitudes $F_\uparrow$ and $F_\downarrow$.

5. GEOMETRIC CHARACTERISTICS OF MODEL. CONCLUSION

In refs. /5,13/ as well in Section 2 and 3 of the present paper the calculation method based on the analysis of the singularities of the scattering amplitude in the impact parameter representation was used. Consider a possible interpretation of these singularities in the framework of the model discussed in the previous sections. The position of the k-th pole in the complex b-plane is defined by the expression $b_k(s) = R_{h_1 h_2}(s) + 2i\pi r(k+1/2)$, $k = 0,1,2...$, $r = 1/\sum_i m_i m_i$. The real part of the function $b_k(s)$ coincides with the interaction radius of the hadrons h_1 and h_2, which defines the asymptotic behaviour of the total cross-sections and the slope parameter of the diffraction cone. The imaginary part of the function $b_k(s)$ is defined by the quantity $r = 1/\sum_i m_i n_i$, which may be considered

as the interaction radius of the N-quark state with the potential V_{eff}. Such a multiquark state arises in the process of interaction of the initial hadrons. Note, that the quantities $R_{h_1 h_2}(s)$ and r are related by the equation

$$R_{h_1 h_2}(s) = \lambda r N \ln(r^2 s) + r \ln C.$$

The energy dependence of the first term accounts for the effects related to existence of the hadron structures(of the virtual clouds). The quantity r determines the behaviour of the elastic scattering amplitude in the region of the momentum transfer $|t| \gg 1$, $|t|/s \ll 1$. In contrast to $R_{h_1 h_2}(s)$ the radius r is energy-independent, that

leads to the energy-independent slope parameter of the angular distributions in the Orear's region.

The model discussed in this paper describes the basic experimental facts occuring in the elastic scattering and diffractive dissociation process at high energies. The notions on the quark structure of hadrons that are input of the model, make it possible to explain the value of the ratio of the total cross-sections for different interacting hadrons as well as to relate the values of other observables with the valence quark numbers and their masses.

The model is based on the solution of the single-time dynamic equations in quantum field theory that follows from the explicitly unitary representation of S-matrix. The composite hadron structure is taken into account at the choice of the kernels of these equations. Besides, certain assumptions have been made. These assumptions seem to be natural if one supposes that quarks weakly interact at small distances, and takes into account analytical properties of the scattering amplitude.

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