Chiral quark model and hadron scattering

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Hadron scattering at large and small distances is simultaneously described in the framework of the unitarization method based on a generalized reaction matrix. The assumed dynamics of hadron interactions is inspired by the ideas of the chiral quark models. The proposed model for hadron interactions is based on geometrical notions of the interactions of constituent quarks similar to those used by Chou and Yang for the description of hadronic interactions. This model allows us to reproduce the $\ln^2 s$ growth of the total cross sections as well as regularities in the behavior of differential cross sections, in particular, its power-law falloff at fixed angles. It also allows us to get the correct values for the total cross-section ratios of hyperon-proton and proton-proton interactions and to estimate the total cross-section values for interactions of the hadrons containing heavy quarks.

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I. INTRODUCTION

The knowledge of hadron structure and hadron interaction dynamics is the ultimate goal of strong interaction theory. Nowadays quantum chromodynamics (QCD) is generally accepted as such a theory. Perturbative expansion in running the coupling constant used for calculations of the observables [1] is based on the asymptotic freedom of QCD. These calculations rely on the parton model [2] where the hadron is considered to be composed of noninteracting massless pointlike constituents. Perturbative QCD (PQCD) enables one to describe spin-averaged observables in hadron interactions at short distances. However, PQCD meets certain problems when the data for spin observables are included in the analysis [3].

The availability of the regular calculational method for the hard process inspired the prompt application of the parton model to calculations of cross sections for the soft processes also. The danger of this extension is notorious. It is related both to the increase of the effective coupling constant and the phenomena of spontaneous breaking of the chiral symmetry.

Because of the latter, a hadron structure already at the distances of 0.10 fm diverges from the parton model picture [4]. The chiral-symmetry-breaking results, in particular, in the generation of quark masses and in the appearance of quark condensates. Quark masses are comparable with the hadron mass scale. Therefore the hadron is often represented as a loosely bound system of constituent quarks. These observations on the hadron structure lead to the understanding of several regularities observed in hadron interactions at large distances. Such a picture provides also reasonable values for the static characteristics of the hadrons, for instance, their magnetic moments [5].

The success of the chiral quark models naturally leads to attempts to enlarge the region of their applicability and, in particular, to use these models for the description of hadron interactions at small distances. It should be stressed here that significant spin effects observed in hard processes might have their interpretations in terms of nonperturbative hadron dynamics at small distances. Notice also that the chiral quark models appear to provide a transparent explanation of the results on spin structure function measurements in deep-inelastic scattering of polarized electrons and muons on a polarized proton target [6].

In this paper we attempt to incorporate some of the chiral quark model ideas in the context of the geometrical approach to the simultaneous description of hadron scattering at large and small distances. We depart from the Chou-Yang model which considers the nucleon as a blob of hadronic matter [7]. We use the ideas of this model while considering nucleon structure motivated by the chiral quark model.

II. STRUCTURE OF HADRONS

QCD in the nonperturbative regime should provide two important phenomena: confinement and spontaneous breaking of chiral symmetry. The scales relevant to these phenomena are characterized by the parameters $\Lambda_{\rm QCD}$ and Λ_χ, respectively [4]. The values of these parameters are

$$\Lambda_{\rm QCD} = 100 - 300 \quad \text{MeV}$$

and

$$\Lambda_\chi \simeq 4\pi f_\pi \simeq 1 \quad \text{GeV},$$

where f_π is the pion decay constant. Chiral $SU(3)_L \times SU(3)_R$ symmetry is spontaneously broken in the range between these two scales. The mechanism of chiral symmetry breaking, as already mentioned, generates the quark masses [4].

Most of the hadron interactions occur at the distances where chiral symmetry is spontaneously broken. But

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since those distances are less than the radius of confinement, constituent degrees of freedom appear to be relevant for the description of hadron interactions.

There are a number of models which incorporate chiral symmetry breaking and imply definite constraints for hadron structure. It is usually considered that those models realize the nonperturbative properties of QCD. They are based on the effective Lagrangian approach. These effective Lagrangians are supposed to be derived from nonperturbative QCD, e.g., in the limit of $N_c \to \infty$. A recent survey of chiral models was presented in Ref. [8].

We will try to incorporate, at least qualitatively, different aspects of hadron dynamics, and, following Ref. [9], represent the effective Lagrangian as a sum of the three terms:

$$\mathcal{L} = \mathcal{L}_\chi + \mathcal{L}_I + \mathcal{L}_C.$$

Here $\mathcal{L}_\chi$ is the term responsible for the spontaneous chiral symmetry breaking. We will consider as possible candidates for the description of that part of interaction the σ model [13] and Nambu–Jona-Lasinio model [14]. The term $\mathcal{L}_\chi$ is turned on first; it provides masses to quarks and leads to the appearance of the quark condensate. Constituent quarks in these models are not pointlike objects. In addition to effective masses their strong interaction dynamics provides them finite size also.

In general, chiral models represent the baryon as an inner core carrying baryonic charge and an outer cloud surrounding this core [8]. The presence of a hadron core is confirmed by the experimental data at low [10] and high [11] [12] energies as well. The observed $p_\perp$ distribution of high-mass muon pairs may be explained under the assumption that a proton core consists of valence quarks and its radius is estimated as 0.20 ± 0.03 fm.

Following these observations it is natural to represent a hadron consisting of the inner region where valence quarks are located and the outer region filled with condensates. This picture incorporates the features of the models with dynamical chiral symmetry breakdown and takes into account the presence of valence quarks and quark condensates inside a hadron.

Valence quarks in this approach are extended objects. They are described by their size and quark matter distribution. To account for the constituent quark interaction and confinement the terms $\mathcal{L}_I$ and $\mathcal{L}_C$ are introduced. The terms $\mathcal{L}_I$ and $\mathcal{L}_C$ do not affect the internal structure of constituent quarks.

Let us turn to the specific form of the chiral-symmetry-breaking term $\mathcal{L}_\chi$. Various possibilities may be considered. For example, the recent one was developed in [13] and is based on the σ model which takes into account the quark interaction with the scalar field $\xi(x)$. This field is a nonvanishing function in the outer region of the nucleon. This model considers an $SU(3)_L \times SU(3)_R$ symmetric scheme with local gauged invariance. In the model the cloud of interacting quarks appears in the outer region of the soliton which is inherent to the nonlinear σ model. The quark interaction is mediated by the scalar field

$$-g\xi(x)[\,\bar\psi_L(x)\psi_R(x) + \bar\psi_R(x)\psi_L(x)\,]. \qquad (1)$$

If this interaction is strong enough, massless quarks may be coupled into the states with zero momentum and spin, and the quark system can suffer a transition to the new ground state containing quark condensates. The new ground state thus is the superposition of correlated quark pair states.

For the first time such a mechanism was introduced in Ref. [14] where an analogy with the theory of superconductivity was used. The original version of the Nambu–Jona-Lasinio (NJL) model considers the nucleon fields as fundamental ones. The NJL model was reformulated for the quark fields with the advent of QCD.

We will refer to the version of this model which takes into account the $U(1)_A$-symmetry-breaking term [15]:

$$\begin{aligned}\mathcal{L}_\chi = \bar\psi(i\gamma\cdot\partial - \hat m^0)\psi + \sum_{a=0}^{8}\frac{1}{2}G[\,(\bar\psi\lambda_a\psi)^2\\ +(\bar\psi i\lambda_a\gamma_5\psi)^2\,]\\ +K[\det\bar\psi(1-\gamma_5)\psi + \det\bar\psi(1+\gamma_5)\psi],\end{aligned} \qquad (2)$$

where the quark field ψ has three colors (N_c=3) and three flavors (N_f=3) and the matrix $\hat m^0 = \mathrm{diag}(m_u^0,\, m_d^0,\, m_s^0)$ is composed from the current quark masses. Equation (2) may be considered as a minimal effective Lagrangian which reflects some of the basic properties of nonperturbative QCD. The last term in (2) obeys the chiral $SU(3)_L \times SU(3)_R$ invariance, but it breaks the unwanted $U(1)_A$ symmetry. The four-fermion Lagrangian of the NJL model reveals this symmetry in the $N_f \geq 3$ case. The first two terms describe the well-known NJL Lagrangian. These terms ensure the dynamical breaking of $SU(3)_L \times SU(3)_R$ chiral symmetry when the coupling constant G is large enough. It has been shown that chiral symmetry is broken dynamically and the quark acquires a mass when the coupling constant G is beyond its critical value.

The Lagrangian (2) in addition to the four-fermion interaction of the original NJL model includes a six-fermion $U(1)_A$-breaking term. The constituent quark masses, when the Lagrangian has the form (2), have been calculated in Refs. [15] as the solution of the coupled gap equations

$$\begin{aligned} m_u &= m_u^0 - 2G\langle\bar u u\rangle_0 - 2K\langle\bar d d\rangle_0\langle\bar s s\rangle_0,\\ m_d &= m_d^0 - 2G\langle\bar d d\rangle_0 - 2K\langle\bar u u\rangle_0\langle\bar s s\rangle_0,\\ m_s &= m_s^0 - 2G\langle\bar s s\rangle_0 - 2K\langle\bar u u\rangle_0\langle\bar d d\rangle_0. \end{aligned} \qquad (3)$$

In this approach massive quarks appear as quasiparticles, i.e., as a current quark and the surrounding cloud of quarks which consists of a mixture of left and right quarks of different flavors. Therefore, in addition to its mass, the quark acquires the complex internal structure and a finite size. Quark radii are determined by the radii of the quark-antiquark clouds surrounding it. Of course, the notion of quark radius depends on the probing field by which it is measured.

We assume that the strong interaction radius of quark r_q is determined by its mass:

$$r_q = \xi/m_q. \qquad (4)$$

Equation (4) was used earlier in the phenomenological analysis of the experimental data in Ref. [16]. Similar results for the quark radius were discussed recently in the framework of the NJL model [17]. Now we can proceed to discussions of hadron interactions.

III. HADRON INTERACTIONS

The above picture for hadron structure implies that the overlapping and interaction of outer clouds in hadron collisions occur at the first stage. Under this, the condensate is being excited and as a result the quasiparticles, i.e., massive quarks, appear in the overlapping region. The part of the hadron energy carried by the outer clouds of condensates being released in the overlapping region goes to the generation of massive quarks. The number of such quarks fluctuates. The average number of these quarks in the framework of these geometrical approaches can be represented as:

$$N(s,b) \propto N(s) D_c^A \otimes D_c^B, \tag{5}$$

where $\otimes$ denotes a convolution integral

$$\int D_c^A(\mathbf{x}) D_c^B(\mathbf{b}-\mathbf{x}) d^2\mathbf{x}.$$

The function D_c^H describes the condensate distribution inside the hadron H and b is the impact parameter of colliding hadrons A and B. In (5) the function $N(s)$ depends on the fraction of energy carried by the valence quarks and is determined by the dynamics of the phase transition of the condensate into constituent quarks. To estimate the function $N(s)$ we can use the maximal possible value

$$N(s) = \frac{[1-k]\sqrt{s}}{m_q},$$

where m_q is the constituent quark mass and k is the average fraction of hadron energy carried by valence quarks. In what follows we will see that such a dependence on s corresponds to the expected asymptotic regimes in the total and differential cross-section behavior.

Thus, the $N(s,b)$ quarks appear in addition to N ($N = N_A + N_B$) valence quarks of the initial hadrons. Note that valence quarks as well as the quarks produced by the excitation of chiral condensates are massive or dressed quarks. The hadronization of produced quarks provides the secondary hadrons in multiple production processes. In elastic scattering they suffer transformation into the condensate clouds of the final hadrons. In terms of the above picture, the generation of the effective field and the formation of the final hadrons should be described by the term $\mathcal{L}_C$ in the Lagrangian $\mathcal{L}$.

In the model, the valence quarks located in the central part of the hadron are supposed to scatter in a quasi-independent way by the external field generated by the intermediate quarks and by the self-consistent field of valence quarks themselves. This part of the interaction should be attributed to the term $\mathcal{L}_I$.

In accordance with the quasi-independence of valence quarks we represent the basic dynamical quantity of this approach in the form of the product [16,18]:

$$U(s,b) = \prod_{q=1}^{N} \langle f_q(s,b) \rangle \tag{6}$$

in the impact parameter representation. The factors $\langle f_q(s,b)\rangle$ correspond to the averaged individual quark scattering amplitude. The assumption of a factorized representation [20] for a U matrix in the form (6) implies that all valence quarks are scattered in the effective field simultaneously; i.e., there are no spectator valence quarks. Such a mechanism resembles the Landshoff mechanism of quark-quark independent scattering [21]. However, in our case we refer not to the pair interaction of valence quarks from the two colliding hadrons but in some sense to the Hartree-Fock approximation for the scattering of valence quarks in an effective field. The functions $\langle f_q(s,b)\rangle$ describe the scattering of valence quarks in an effective field.

The function $U(s,t)$ is given by the Fourier-Bessel transform

$$U(s,t) = \frac{s}{\pi^2} \int_0^\infty \prod_{q=1}^{N} \langle f_q(s,b) \rangle J_0(b\sqrt{-t}) b db.$$

The scattering amplitude F is related to the function U by the equation

$$F = U + iUDF, \tag{7}$$

which we present here in operator form. This relation allows one to obey unitarity provided the inequality $\mathrm{Im}\, U(s,b) \geq 0$ is satisfied. The function $U(s,b)$ is a generalized reaction matrix [19]. The amplitude $F(s,b)$ itself does not reveal the independence of the quark scattering (it cannot be represented in a factorizable form) and due to this fact we make reference to quasi independence.

Keeping in mind the above assumptions on quark scattering in such a field we represent the function $\langle f_q(s,b)\rangle$ in the form

$$\langle f_q(s,b) \rangle = [N(s,b) + N - 1]\, V_q(b). \tag{8}$$

In this representation $V_q(b)$ corresponds to a single quark-quark amplitude. Equation (8) implies that each valence quark is being scattered by all other $N-1$ valence quarks belonging to the same hadron as well as to the other hadron and by the $N(s,b)$ quarks produced by the excitation of the chiral condensates. Therefore the single scattering amplitude is multiplied by the factor $[N(s,b)+N-1]$.

For the single quark-quark amplitude $V_q(b)$ we apply the ideas of the Chou-Yang model since quarks in our approach are extended objects. Therefore $V_q(b)$ may be written as the convolution integral

$$V_q(\mathbf{b}) = \int D_q(\mathbf{x}) f_{qq}(\mathbf{b}+\mathbf{x}-\mathbf{y}) D_q(\mathbf{y}) d^2\mathbf{x} d^2\mathbf{y}, \tag{9}$$

where f_{qq} stands for the basic amplitude of a nonperturbative "parton-parton" interaction, and $D_q(\mathbf{x})$ and $D_q(\mathbf{y})$ are the two-dimensional quark matter distributions inside the interacting quarks.

As was noted, massive quarks are quasiparticles. Our basic "parton-parton" interaction is of a contact type. Hence we approximate the amplitude $f_{qq}(\mathbf{b}+\mathbf{x}-\mathbf{y})$ by the corresponding δ function in impact parameter space $\delta^2(\mathbf{b}+\mathbf{x}-\mathbf{y})$. Therefore, $V_q(b)$ is the two-dimensional Fourier transform of the product of two quark form factors. The quark form factor can be taken in the form $(\mathbf{q}^2+m_q^2/\xi^2)^{-2}$ [22]. The corresponding quark matter distribution is of the form $D_q \propto \exp(-m_q r/\xi)$. Then from (9) we obtain, for the function $V_q(b)$,

$$V_q(b) \propto \exp(-m_q b/\xi). \tag{10}$$

The b dependence of the function $V_q(b)$ may be more complicated owing to more complex expression for f_{qq}. The account for this is essential for a quantitative description of the experimental data, however, for qualitative considerations Eq. (10) is quite sufficient.

It should be noted that due to the peripherality of the condensate distribution inside a hadron the convolution $D_c^A \otimes D_c^B$ has a less steep b dependence than the function $V_q(b)$. Keeping in mind this one may neglect, the b dependence of the condensate convolution in comparison with the b dependence of $V_q(b)$. Thus we can approximate $N(s,b)$ in (8) by $N(s,0)$.

IV. DIFFERENTIAL CROSS SECTIONS

Given with the explicit form of the U matrix we can obtain the scattering amplitude as a solution of (7). We consider three kinematical regions: namely, the region of small momentum transfers,

$$|t| < m_p^2, \quad |t|/s \ll 1,$$

the region of intermediate t,

$$|t|/s \ll 1, \quad |t| \geq m_p^2$$

and the region of large t or fixed angle scattering, when

$$|t|,\ s \gg m_p^2, \quad |t|/s \text{ fixed}.$$

We will discuss qualitative results.

In impact parameter representation the scattering amplitude may be written in the form

$$F(s,b) = U(s,b)[1-iU(s,b)]^{-1}. \tag{11}$$

This is the solution of Eq. (7) at $s \gg 4m^2$. Note that the more familiar way to provide the direct channel unitarity consists in the representation of the scattering amplitude in the eikonal form

$$F(s,b) = \frac{i}{2}\left(1-e^{i\chi(s,b)}\right),$$

where $\chi(s,b)$ is the eikonal connected with the function $U(s,b)$ by the relation

$$e^{i\chi(s,b)} = \frac{1+iU(s,b)}{1-iU(s,b)}.$$

The generalized reaction matrix $U(s,b)$ as follows from the above considerations may be presented in the form

$$U(s,b) = ig(N-1)^N\left[1+\alpha\sqrt{s}/m_q\right]^N \exp(-Mb/\xi), \tag{12}$$

where $M = \sum_{q=1}^N m_q$. At very high energies we could neglect the energy-independent term in (12) and rewrite the expression for $U(s,b)$ in the form

$$U(s,b) = i\tilde{g}\left(s/m_q^2\right)^{N/2}\exp(-Mb/\xi). \tag{13}$$

The calculation of the scattering amplitude is based on the impact parameter representation

$$F(s,t) = \frac{s}{2\pi^2}\int_0^\infty d\beta F(s,\beta) J_0(\sqrt{-\beta t}), \quad \beta = b^2,$$

and on an analysis of the singularities of $F(s,\beta)$ in the complex β plane. Notice that the possibility of writing down an exact impact parameter or Fourier-Bessel representation for the elastic scattering amplitude valid for all energies and scattering angles was proved in Ref. [23].

The amplitude $F(s,\beta)$ has poles in the complex β plane,

$$\beta_n(s) = \frac{\xi^2}{M^2}\left\{\ln\left[\tilde{g}\left(s/m_q^2\right)^{N/2}\right] + i\pi n\right\},$$

$$n = \pm 1, \pm 3, \ldots, \tag{14}$$

and the branching point at $\beta = 0$. The amplitude $F(s,t)$ may be written then as a sum:

$$F(s,t) = F_p(s,t) + F_c(s,t), \tag{15}$$

where F_p is the pole contribution and $F_c(s,t)$ is the cut contribution. The pole contribution has the form

$$F_p(s,t) = -\frac{2is\xi^2}{\pi^2 M^2}\sum_n \sqrt{\beta_n(s)} K_0[\sqrt{t\beta_n(s)}] \tag{16}$$

and the cut contribution is determined as

$$F_c(s,t) = \frac{-isM}{\pi^2\xi\tilde{g}(s/m_q^2)^{N/2}}\,\frac{1}{(M^2/\xi^2-t)^{3/2}}. \tag{17}$$

From Eqs. (16) and (17) it is easily seen that the cut contribution dominates at

$$|t| > |t_0| = \frac{M^2}{\pi^2\xi^2}\ln^2[\tilde{g}(s/m_q^2)^{N/2}]. \tag{18}$$

At small t all terms of series (16) are essential and at $t \sim 0$ the differential cross section has the t dependence

$$\frac{d\sigma}{dt} \propto \exp[B(s)t], \quad B(s) \propto \ln^2 s \tag{19}$$

at $s \to \infty$.

In the region of medium t it is sufficient to take into account only a few or even one of the terms of series (16). The differential cross section in this region has the familiar Orear-type behavior

$$\frac{d\sigma}{dt} \propto \exp\left(-\frac{2\pi\xi}{M}\sqrt{-t}\right). \tag{20}$$

It is well known that the dependences (19) and (20) are in a good agreement with experiment. Equation (20) allows one to estimate the value of the inverse quark radius: $m_q/\xi = 150$–200 MeV. Thus, to reproduce the standard constituent quark masses the value of the parameter ξ is close to 2.

At large momentum transfers ($s \to \infty$, $|t|/s$ fixed) the contribution from the branching point ($\beta = 0$) is a dominating one. The angular distribution in this region has the power dependence

$$\frac{d\sigma}{dt} \propto \left(\frac{1}{s}\right)^{N+3} f(\theta). \tag{21}$$

Expression (21) is in rather good agreement with the experimental data though in general the values of the power index are different from those obtained via dimensional counting rules and in perturbative QCD. Thus, for pp-elastic scattering we have index 9 instead of 10 in PQCD; however, for πp scattering the power index coincides with that predicted by PQCD.

V. TOTAL AND INELASTIC CROSS SECTIONS

In this section we discuss behavior of the total and inelastic cross sections. In addition to the energy dependence of these observables we will emphasize its dependence on the geometrical characteristics of nonperturbative quark interactions.

The total cross section has the following energy and quark mass dependences:

$$\sigma_{\rm tot}(s) = \frac{\pi\xi^2}{\langle m_q\rangle^2}\Phi(s,N), \tag{22}$$

where $\langle m_q\rangle = \frac{1}{N}\sum_{q=1}^{N} m_q$ is the mean value of the constituent quark masses in the colliding hadrons and the function Φ does not depend on the quark masses. It can be represented in the two extreme cases as

$$\Phi(s,N) = \begin{cases} 8g^N(N-1)^N/N^2, & s \ll s_0, \\ \ln^2 s, & s \gg s_0. \end{cases} \tag{23}$$

The value of s_0 is determined by the equation $|U(s_0, b = 0)| = 1$. At high energies ($s \gg s_0$) the total cross-section ratios for different hadron pairs only depend on the mean value of quark mass in the colliding hadrons:

$$\frac{\sigma_{\rm tot}(h_1h_2)}{\sigma_{\rm tot}(h_3h_4)} = \frac{\langle m_q\rangle^2_{h_3h_4}}{\langle m_q\rangle^2_{h_1h_2}}. \tag{24}$$

Equation (22) provides the cross-section ratios for hyperon-proton interactions in agreement with the experimental data [24] if one assumes the standard value for the ratio of strange to light constituent quark mass, $m_s/m_q = 1.5$ ($m_q = m_u = m_d$):

$$\frac{\sigma_{\rm tot}(\Sigma p)}{\sigma_{\rm tot}(pp)} = 0.85, \quad \frac{\sigma_{\rm tot}(\Xi p)}{\sigma_{\rm tot}(pp)} = 0.73. \tag{25}$$

The ratio of the total cross sections of the Kp and πp interactions is also in agreement with experimental data:

$$\frac{\sigma_{\rm tot}(Kp)}{\sigma_{\rm tot}(\pi p)} = 0.83. \tag{26}$$

Equation (22) allows us to get estimations of the cross sections for the processes with heavy quarks. Thus, for the total cross section $\sigma_{\rm tot}(\psi p)$ we have

$$\frac{\sigma_{\rm tot}(\psi p)}{\sigma_{\rm tot}(\pi p)} = 0.15 \tag{27}$$

in the case of $m_c/m_q = 5$. It is interesting to note that (24) provides

$$\frac{\sigma_{\rm tot}(\pi p)}{\sigma_{\rm tot}(pp)} \to 1$$

at $s \to \infty$. So we could expect the deviation at higher energies of the above ratio from the value 2/3, predicted by the additive quark model [25] and also obtained in the model where the Pomeron couples to the single quark like a $C = +1$ photon [26,27].

However, for hadrons with strange and heavy quarks the above asymptotical value differs from unity; e.g.,

$$\frac{\sigma_{\rm tot}(Kp)}{\sigma_{\rm tot}(pp)} \to 0.83.$$

The explicit form for the inelastic cross section is

$$\sigma_{\rm inel}(s) = \frac{8\pi\xi^2}{N^2\langle m_q\rangle^2}\ln\left[1 + g(N-1)^N\left(1+\frac{\alpha\sqrt{s}}{m_q}\right)^N\right]. \tag{28}$$

And at asymptotically high energies the inelastic cross-section growth is

$$\sigma_{\rm inel}(s) = \frac{4\pi\xi^2}{N\langle m_q\rangle^2}\ln s. \tag{29}$$

It is clearly seen that the ratio of the inelastic cross sections for different hadrons depends on the number of quarks in the initial hadrons N and $\langle m_q\rangle$ at $s \to \infty$:

$$\frac{\sigma_{\rm inel}(h_1h_2)}{\sigma_{\rm inel}(h_3h_4)} = \frac{\langle m_q\rangle^2_{h_3h_4}N_{h_3h_4}}{\langle m_q\rangle^2_{h_1h_2}N_{h_1h_2}}. \tag{30}$$

At $s \gg s_0$ the dependence of the hadron interaction radius $R(s)$ and the ratio $\sigma_{\rm el}/\sigma_{\rm tot}$ on the parameters $\langle m_q\rangle$ is given by the equation:

$$R(s) = \frac{\xi}{2\langle m_q\rangle}\ln s, \tag{31}$$

$$\frac{\sigma_{\rm el}(s)}{\sigma_{\rm tot}(s)} = 1 - \frac{4}{N \ln s}. \tag{32}$$

It is useful to note here that such a behavior of the ratio $\sigma_{\rm el}/\sigma_{\rm tot}$ and $\sigma_{\rm inel}(s)$ is a result of self-damping of inelastic channels [28] at small impact distances. The quantity $\mathrm{Im}U(s,b)$ can be represented in the form [29]

$$\mathrm{Im}U(s,b) = \sum_{n\geq 3} \bar{U}_n(s,b), \tag{33}$$

where

$$\bar{U}_n(s,b) = \int d\Gamma_n |U_n(s,b,\{\zeta_n\})|^2$$

and $d\Gamma_n$ is an element of n-particle phase space volume. The functions $U_n(s,b,\{\zeta_n\})$ are determined by the dynamics of the $2 \to n$ process. According to (33) $\mathrm{Im}U(s,b)$ arises as a shadow of inelastic interactions. However, an increase of the function $\mathrm{Im}U(s,b)$ due to an increase of the contribution of inelastic channels leads to a decrease of the inelastic overlap function

$$\eta(s,b) = \frac{\mathrm{Im}U(s,b)}{|1 - iU(s,b)|^2}$$

at small impact parameters and feedback elastic channel. We can observe here the transition from shadow to antishadow scattering which could occur beyond the energy of 2 TeV [30] as was suggested recently by the new data [31] from the Collider Detector at Fermilab (CDF) Collaboration on the total, elastic, and diffraction cross sections.

VI. CONCLUDING REMARKS

Proceeding from nonperturbative dynamics the above model allows one to reproduce the main regularities in elastic scattering at all values of the momentum transfers. It should be emphasized that the observed power-law behavior of the differential cross sections at large angles results from nonperturbative dynamics. This implies that the power-law behavior of the cross sections at fixed angles alone is not sufficient for a conclusion on the validity of perturbative QCD.

One should note the feasibility of the nonpertubative approach to large angle hadron scattering. In Ref. [32] the key role was attributed to confinement effects. Recently another approach was developed where different pictures for soft and hard Pomerons were proposed. It was suggested [33] to use for soft Pomeron the picture with extended constituent quarks and respectively for hard Pomeron the PQCD-inspired picture.

In this paper the main role is attributed to chiral symmetry breaking. Under considerations of hadron structure and hadron interaction dynamics we rely on the conclusions derived from the models with spontaneously broken chiral symmetry. Unitarity in the direct channel was explicitly taken into account and it was shown how hadron interactions at both large and small distances may be included in the framework of a unified approach.

The nonperturbative origin of the hadron dynamics allows one to include quark helicity flip interactions and generalization of the above model for the spin case allows one to get a description of a number of significant spin effects observed in elastic scattering [34].

Of course, future experiments at higher energies and $p_\perp$ values will clarify the dynamics of elastic scattering. Among these experiments polarization measurements would play a most significant role. The crucial fact here is that one-spin transverse asymmetries in PQCD have a higher-twist origin. The present experiments do not show trends in vanishing of one-spin asymmetries [35]. However, in elastic scattering the independent quark scattering mechanism [21] could provide the helicity flip of the hadron despite the helicity conservation at the quark level [36]. Therefore to reveal the nature of hard scattering dynamics the polarization experiments in elastic scattering at high energies and $p_\perp$ should be performed together with the study of nuclear analyzing power. The nuclear target here will be used as a filter to eliminate components of a proton wave function with a large transverse separation of constituents [36].

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