Spin phenomena in high-energy hadron scattering

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A unified approach to the investigation of high-energy small-angle scattering of particles with arbitrary spin is developed. In the case of a weak energy dependence of the spin effects a mechanism of "spin dynamics" of the strong interactions appears as a basically new possibility at superhigh energies. The manifestation of the "spin" mechanism is investigated in the framework of a definite dynamical model that takes into account the structure of hadrons at large distances. The predictions of this and a number of other models for proton–proton and proton–antiproton elastic scattering up to $\sqrt{s} = 40$ TeV are considered.

INTRODUCTION

The achievements of the theory of elementary particles in recent years are to a large degree associated with the development of quantum chromodynamics (QCD), which is based on the idea of colored quarks[1] and the principle of local gauge invariance.[2] In QCD, the effective charge, introduced on the basis of renormalization-group considerations,[3] tends to zero at short distances. This phenomenon, called asymptotic freedom, makes it possible to use the well-developed methods of perturbation theory. At distances of the order of the hadron diameter, around 1 F, these methods are invalid, so that at the present time it is not possible to develop a microscopic theory of hadron physics in this interaction region.

The basic features in the behavior of physical quantities in the asymptotic region are investigated through the general properties of the S matrix.[4] For the study of specific processes, including interactions at large distances, many models,[5] which use different ideas about hadron structure, the contribution played by the interaction of constituents, etc., have been proposed. The study of spin phenomena in high-energy hadron reactions has become one of the important ways of testing the viability of model approaches and the basic propositions of QCD.

The development of experimental techniques has led to the appearance of rich experimental material, and this poses for the theoreticians very difficult questions. We mention the existence of a large polarization in the inclusive production of the majority of hyperons (while $\overline{\Lambda}$ and p have polarization near zero), the energy independence in the $pp \to \Lambda^0 X$ reaction of the polarization of the Λ^0 particles in a remarkably wide region (from 12 to 2000 GeV), and the large π^0 asymmetry.[6]

The nucleon–nucleon reactions have the greatest interest, since it is here that we have rich experimental material and the highest energies have been achieved.[7]

The polarization measurements in elastic proton–proton scattering, which in the region of small momentum transfers initially confirmed the idea of a power-law decrease of the spin-flip amplitude,[8] gave an unexpected result in the region $-t \sim 1$–3 GeV^2 at $p_L = 150$ and 300 GeV—a polarization of about 20% that does not disappear with increasing energy.[9] The ZGS Argonne data and the AGS Brookhaven data at $p_L = 16$ and 28 GeV show that the transverse polarization in elastic pp scattering becomes quite appreciable on the transition to large-angle scattering. A sharp increase in the spin correlation parameter A_{NN} (up to 50%) has been found near 90°.[10] In the next few years additional extremely interesting data will be provided by using a beam of polarized neutrons (with p_L up to 16 GeV) and polarized beams of protons and antiprotons in various accelerators. For example, at Serpukhov experiments will continue on the measurement of the asymmetry in meson production. In 1987 a start should be made at CERN on the $S\bar{p}pS$ of measurements of $\Delta\sigma_L$, of the asymmetry of $\pi^\pm$ and Λ^0 production in the region $0.5 \leqslant x_F \leqslant 0.9$ and $0.2 \leqslant p_\perp < 1.5$ GeV, and of the asymmetry of the π^0 mesons in beams of protons of opposite polarization in the region $x_F = 0$ and $p_\perp$ up to 4 GeV.

At the present time, there is no single theoretical model capable of combining all the varied experimental material and revealing the physical processes underlying it (see, for example, Ref. 11). It is known that the QCD vector interaction for massless quarks leads to small spin effects at high energies and large momentum transfers. This means that either sufficiently high energies have not yet been reached experimentally or that it is necessary to determine more accurately or modify the theoretical picture of hadron interactions.

Various ideas are invoked in theoretical models to explain the large spin effects at high energies.[12] For example, it has been noted[13] that the presence of quark constituents of hadrons may lead to the appearance of spin density currents within polarized particles. The droplet model for particles with spin has been generalized by means of this concept. As a result, the transparency of different parts of one of the colliding particles with respect to the parts of the other depends on the relative velocity of their motion. This leads to spin effects in hadron–hadron scattering that tend to a certain constant limit with increasing energy.

The notion of spin currents was developed in Ref. 14, in which two possibilities were considered: currents associated with rigid rotation of the particle as a whole and currents due to soft rotation of the matter within a hadron with angular velocity that decreases with increasing impact parameter. The natural absence of an energy dependence of the angular velocity of the hadron matter leads to spin effects that do not decrease with increasing energy. Subsequently, this eikonal model was combined with a developed variant of a "massive" quark model.[15] In it, the introduction of meson exchanges between the interacting quarks made it possible to obtain a fairly rich spin structure of the quark–quark inter-

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action. On the basis of the interference between the "soft" eikonal scattering amplitude and the "rigid" scattering amplitude, appreciable polarization effects were obtained for scattering in the region of large momentum transfers, and it was predicted that the asymmetry coefficient $A(90°)$ should tend to the value 0.97 at superhigh energies.

The rapid-growth model[16] leads to a weak energy dependence of the polarization in elastic hadron reactions. The model is based on the assumption of complete saturation of the partial waves at superhigh energies. The helicity-flip amplitude increases as $\ln s$, and this leads to large spin effects at high energies.

When an ISR experiment revealed that the differential cross section of elastic pp scattering had no alternating maxima and minima, this was explained in the u-matrix method[17] by a significant contribution of amplitudes with double helicity flip in the region $-t \geqslant 2\ \mathrm{GeV}^2$. In the quark–parton u-matrix model[18] it is assumed that the interaction of the hadrons produces a certain field $V_{\rm eff}$, in which the valence quarks are scattered in a quasi-independent manner. If certain assumptions are made about the structure and the interaction of the quarks within the resulting superposition of hadrons and it is assumed that this formation has a nonvanishing angular momentum, this model predicts appreciable and nonvanishing (in the limit $s \to \infty$) polarization in elastic scattering in the region of fixed angles up to 90°.

The effects of an anomalous color moment and their confinement in a hadron of finite volume can also lead to a spin-flip amplitude that does not decrease with increasing s.[19]

It was shown in Ref. 20 that the true parameter determining the polarization and asymmetry in QCD may be the mass of the polarized hadron. Physically, this is equivalent to renormalization of the mass of a quark as it moves in the external gluon field of the hadron.

Calculations of the quark magnetic moments on the basis of measured magnetic moments of hyperons[21] indicated extremally large nonstatic contributions (orbital or configuration mixing, exchanges of current quarks or meson states). In Ref. 22, the same data were used to calculate the proton and neutron form factors. It was shown that there is an appreciable contribution of d waves in the nucleon, and this created the possibility of obtaining a large asymmetry coefficient A_{NN} in the region of large angles.

Most of these considerations suggest that the spin effects at high energies and limited momentum transfers are intimately related to the dynamics of the strong interactions at large distances.

In this region, the coupling constant of the strong interaction increases, and QCD perturbation theory becomes invalid. There has recently appeared a hope of investigating the large-distance interaction by means of sum rules in QCD,[23,24] which have made it possible to obtain important results in the study of the dynamical properties of hadrons.[25-27] However, elastic particle scattering has not yet been investigated in the framework of this method.

In the theory of strong interactions, an important part is played by the Logunov–Tavkhelidze quasipotential method,[28] which combines the rigor of the fundamental principles of quantum theory with the possibility of using heuristic considerations about the nature of high-energy hadron interaction. The generalization of the quasipotential method to the case of particles with spin[29-35] opened up the possibility of effective use of the method to study real physical systems. In conjunction with the hypothesis of smoothness of the quasipotential,[36-38] which is intimately related to hadron interaction dynamics at large distances, the use of the quasipotential method made it possible to understand the basic features of high-energy small-angle particle scattering.[39-42] It was shown[41,42] that smoothness of the interaction results in an eikonal nature of the small-angle hadron scattering. The scattering amplitude has the representation

$$T(s,t)=\frac{is}{16\pi^3}\int d^2\rho e^{i\Delta\rho}\left(1-e^{i\chi(\rho,s)}\right). \tag{0.1}$$

Information about the large-distance interaction dynamics is contained in (0.1) in the eikonal phase $\chi(\rho,s)$. The eikonal nature of the scattering means that in small-angle scattering the particles move along nearly straight paths, and the particle wave function differs little from the plane-wave representation. The eikonal representation has been justified in quantum field theory in studies based on the approximation of straight paths.[43] In the framework of QCD the validity of the eikonal representation has been proved up to the eighth order of perturbation theory.[44]

The quasipotential theory of strong interactions makes it possible to develop a unified approach to the investigation of small- and large-angle scattering[45,46] by separating the contributions of the small and large distances in a smooth local quasipotential given by a representation of the form

$$V(s,t=-\Delta^2)=g(s)\int_0^\infty dx\rho_t(s,x)\,e^{xt}. \tag{0.2}$$

For hard processes the important region is $0<x<1\ \mathrm{GeV}^{-2}$, and for soft processes it is the region $x \geqslant 1\ \mathrm{GeV}^{-2}$. By means of the method, one can construct a representation for the scattering amplitude[45] that makes it possible to use effectively the separation between the contributions of the small and large distances made in the interaction quasipotential. This representation is a basis for studying high-energy scattering in different regions of the kinematic variables. It should be noted that the derivation of the representation was based on a quasipotential equation written in general form and valid for the scattering of particles of arbitrary spin; this makes it possible to perform the calculations for specific processes in the final stage by choosing an appropriate matrix structure of the equation and of the interaction quasipotential. The representation for the scattering amplitude has provided the basis for the development of an effective method of summation at large distances in hard hadron scattering. This is extremely important, since the great strength of the strong interaction at distances of the order of the hadron diameter requires complete summation of the contributions of the "soft" region to the scattering amplitude. Allowance for the contributions of the large distances in hard scattering leads to corrections that decrease as powers with increasing energy.[46] These corrections result in the appearance of an appreciable (up to 50%, for $\cos\theta = -0.4$ to -0.5) polarization in $\pi^\pm p$ scattering and a polarization of order 15%–30% in pp and np scattering at large angles at energies $E_L \sim 10$ GeV.

The results obtained in Refs. 45 and 46 made it possible to develop in the quasipotential theory of strong interactions unified methods for investigating the small-angle scattering of particles with spin; these are presented in the present re-view. On the basis of the expression for the scattering amplitude of particles with spin,[45] which is valid in the complete region of kinematic variables, we obtain a unified eikonal representation for the amplitude of small-angle scattering of particles of arbitrary spin.[47] We investigate the criteria for its applicability, these making it possible to prove the validity of the eikonal representation for the scattering amplitude in the case of spin effects that decrease in a power-law manner with increasing energy.[48]

However, the physical phenomena noted above and the model approaches require study of high-energy hadron scattering in the presence of a weak energy dependence of the spin effects. The summation of the large-distance contributions made in this case in the quasipotential method leads to a possibility of modifying the eikonal representation, in which terms that increase as $\sqrt{s}$ appear in the eikonal phase, which determines the amplitude without spin flip.[49] This effect can change the strong-interaction dynamics at superhigh energies.

To study high-energy small-angle scattering, we have developed a dynamical model of the interaction of hadrons[50,51] that takes into account their large-distance structure. The model yields naturally small spin-flip amplitudes that do not vanish with increasing energy,[48,49] and which we call anomalous. The model has been used to describe meson–nucleon and nucleon–nucleon scattering[52,53] in the diffraction region. We consider in detail the spin effects determined by the anomalous terms of the scattering amplitude that appear in the model. At ISR energies, these terms make it possible to reproduce fairly accurately the picture of the polarization in nucleon–nucleon scattering without giving a significant contribution to the differential cross sections.[49,54] At superhigh energies, allowance for the anomalous contributions has made it possible to predict quantitatively the behavior of the $p\bar{p}$ scattering cross sections at the energies of the CERN $p\bar{p}$ collider. The self-consistent picture obtained for the polarization and the scattering cross sections is an indication of the first manifestations of the mechanism of "spin" dynamics of the strong interactions at superhigh energies.

Comparison of the results of the model with the predictions of a number of other models has revealed the most interesting critical regions of the kinematic variables for the undertaking of future experiments.

1. UNIFIED EIKONAL REPRESENTATION FOR THE SCATTERING AMPLITUDE OF PARTICLES WITH SPIN

We consider the elastic scattering of strongly interacting particles at high energies and fixed momentum transfers:

$$s \to \infty, \, |t/s| \ll 1.$$

Assuming the existence of a smooth local quasipotential, we obtain in the framework of the quasipotential approach a unified eikonal representation for the scattering of particles of arbitrary spin. We choose the quasipotential equation in this case in the general form[45]

$$\hat{T}(s, \mathbf{p}, \mathbf{k}) = \hat{V}(s, \mathbf{p}, \mathbf{k}) + \int \frac{d^3 q \hat{V}(s, \mathbf{p}, \mathbf{q})\, \hat{A}(s, q)\, \hat{T}(s, \mathbf{q}, \mathbf{k})}{(E^2(\mathbf{q}) - E^2 - i0)}, \tag{1}$$

where $E(\mathbf{q}) = \sqrt{m_1^2 + \mathbf{q}^2} + \sqrt{m_2^2 + \mathbf{q}^2}$; $E = \sqrt{s} = E(\mathbf{p}) = E(\mathbf{k})$ is the total c.m.s. energy, m_1 and m_2 are the masses of the colliding particles, and $\hat{A}(s,q)$ is a matrix whose form and rank depend on the spin of the interacting particles. To solve Eq. (1), one can choose a quasipotential that satisfies the representation (0.2) and is an analytic function of t in the half-plane Re $t \leqslant 0$.

We seek a solution of Eq. (1) by successive iteration. For the $(n+1)$ iteration term of the scattering amplitude, we have

$$\hat{T}_{n+1}(s, \mathbf{p}, \mathbf{k}) = \int dx_1 \ldots dx_{n+1} \times \exp\left\{ t \Big/ \sum_{j=1}^{n+1} \frac{1}{x_j} \right\} \hat{J}_n(x_1, \ldots, x_{n+1}),$$

$$\hat{J}_n(x_1, \ldots, x_{n+1}) = [g(s)]^{n+1} \int \prod_{j=1}^{n} \frac{d\Delta}{E^2(\Delta_j + \lambda_j) - E^2 - i0} \times \exp[-C_{ij}\Delta_i\Delta_j]$$
$$\times \hat{\rho}\left(s, x_1, \frac{\mathbf{p}+\lambda_1+\Delta_1}{2}\right) \hat{A}(s, \Delta_1+\lambda_1)\, \hat{\rho}\left(s, x_2, \frac{\lambda_1+\lambda_2+\Delta_1+\Delta_2}{2}\right)$$
$$\ldots \times \hat{A}(s, \Delta_n + \lambda_n) \times \hat{\rho}\left(s, x_{n+1}, \frac{\lambda_n + \Delta_n + \mathbf{k}}{2}\right), \tag{2}$$

where

$$C_{ij}\Delta_i\Delta_j = x_1\Delta_1^2 + \sum_{h=2}^{n} x_h(\Delta_h - \Delta_{h-1})^2 + x_{n+1}\Delta_n^2;$$

$$\lambda_i = \mathbf{l} + \left[1 - 2\left(\sum_{j=1}^{i} \frac{1}{x_j}\right) \Big/ \left(\sum_{j=1}^{n+1} \frac{1}{x_j}\right)\right] \mathbf{r};\ \mathbf{l} = \frac{\mathbf{p}+\mathbf{k}}{2};\ \mathbf{r} = \frac{\mathbf{p}-\mathbf{k}}{2}.$$

The representation (2) was used to study large-distance effects in hard hadron scattering.[46] It makes it possible to simplify appreciably the investigation of the small-angle scattering of particles with spin. In this region, the small-distance effects are unimportant, and the strong-interaction dynamics is determined by the large distances. One can show[46] that the density function corresponding to the soft part of the quasipotential is concentrated in a region of finite x. This permits the conclusion that all the shifts λ_i in (2) are of the order of a large momentum $\lambda_i \sim \mathbf{l}$, and the dynamical factors in (2) can be replaced by their expansions near the point $\lambda_i \sim \mathbf{l}$. As a result, we obtain for the leading term of the scattering amplitude[57]

$$\hat{T}_{n+1}(s, t) = \frac{1}{(2\pi)^3}\left(\frac{i}{8p}\right)^n \int d^2\rho e^{i\Delta\rho} \int_{-\infty}^{\infty} dz_1 \ldots dz_{n+1} \theta(z_1 - z_2) \times \theta(z_2 - z_3) \times \ldots \times \theta(z_n - z_{n+1})\, \hat{V}(\rho, z_1, \mathbf{l}) \times \hat{A}(s, \mathbf{l})\, \hat{V}(\rho, z_2, \mathbf{l}) \ldots \hat{A}(s, \mathbf{l})\, \hat{V}(\rho, z_{n+1}, \mathbf{l}), \tag{3}$$

where $\hat{V}(\rho, z_i, \mathbf{l})$ is the Fourier transform of the quasipotential:

$$\hat{V}(\rho, z, \mathbf{l}) = g(s) \int_0^\infty dx \left(\frac{\pi}{x}\right)^{3/2} \hat{\rho}(s, x, \mathbf{l}) \exp\left(-\frac{1}{4x}(\rho^2 + z^2)\right).$$

The expression (3) generalizes the eikonal representation for the scattering amplitude of particles of arbitrary spin. To calculate the scattering amplitude in each particular case, it is necessary to specify explicitly the quasipotential $\hat{V}(\rho, z)$ and the matrix $\hat{A}(s,\mathbf{l})$. As an example, we consider the scattering of spinless particles. In this case, $\hat{A}(s,\mathbf{l})$ has the form

$$\hat{A}(s,\ \mathbf{l})=\frac{4}{\sqrt{m^2+\mathbf{l}^2}}\approx\frac{4}{|\mathbf{p}|},$$

and the quasipotential is a scalar function of its arguments. From (3) we then obtain

$$T_{n+1}(s,\ t)=\frac{is}{16\pi^3}\int d^2\rho e^{i\Delta\rho}\frac{1}{(n+1)!}\Big[\frac{2i}{s}\int\limits_{-\infty}^{\infty}dzV(\rho,\ z)\Big]^{(n+1)},\tag{4}$$

and this leads to the standard eikonal representation (0.1), the eikonal phase being related to the quasipotential by

$$i\chi(\rho,\ s)=\frac{2i}{s}\int\limits_{-\infty}^{\infty}dzV(\rho,\ z,\ s).\tag{5}$$

The smoothness of the quasipotential ensures smallness of the corrections $\sim 1/p$. We now obtain eikonal representations for meson–nucleon scattering amplitudes. In the high-energy limit, we obtain for the matrix $\hat{A}(s,\mathbf{l})$, which is determined in Ref. 34,

$$\hat{A}(s,\ \mathbf{l})\underset{s\to\infty}{\sim}2\hat{n}(\mathbf{l}),\ \hat{n}(\mathbf{l})=\gamma_0-\gamma\mathbf{l}/|\mathbf{l}|.\tag{6}$$

As the quasipotential, we use the amplitude for scattering of a meson by a nucleon in invariant form:

$$\hat{T}=a+\frac{1}{\sqrt{s}}b\hat{Q},\ \hat{Q}=(\hat{q}_1+\hat{q}_2)/2.\tag{7}$$

Here, q_1 and q_2 are the meson momenta in the initial and final states. We have introduced $1/\sqrt{s}$ to make the amplitudes a and b have the same dimensions. In the region of high-energy small-angle scattering ($s\to\infty$, t fixed)

$$\hat{T}=a+b\hat{n}(-\mathbf{l}).\tag{8}$$

To go over to the r space, we must make a Fourier transformation with respect to the momentum transfer. We then obtain a matrix form for the quasipotential of the meson–nucleon interaction:

$$\hat{V}(s,\ r,\ \mathbf{l})=a(s,\ r)+b(s,\ r)\hat{n}(-\mathbf{l}),\tag{9}$$

where the quasipotentials $a(s,r)$ and $b(s,r)$ are the Fourier transforms of the amplitudes $a(s,t)$ and $b(s,t)$.

Calculating the matrix elements in (8), we obtain for the helicity amplitudes of meson–nucleon scattering[57]

$$\left.\begin{aligned}T_{++}(s,\ t)&=\frac{ip}{4\pi^3}\int d^2\rho e^{i\Delta\rho}[1-e^{\chi_0(\rho)}],\\ T_{+-}(s,\ t)&=-\frac{i\Delta}{16\pi^3}\int d^2\rho e^{i\Delta\rho}\chi_1(\rho)e^{\chi_0(\rho)}.\end{aligned}\right\}\tag{10}$$

Here

$$\chi_0(\rho)=\frac{i}{p}\int\limits_{-\infty}^{\infty}dzb(\rho,\ z);$$

$$\chi_1(\rho)=\frac{i}{p}\int\limits_{-\infty}^{\infty}dza(\rho,\ z)$$

$$\times\frac{1}{2}\Big[\exp\Big\{-\frac{i}{p}\int\limits_{z}^{\infty}dz'b(\rho,\ z')\Big\}+\exp\Big\{-\frac{i}{p}\int\limits_{-\infty}^{z}dz'b(\rho,\ z')\Big\}\Big].$$

The corresponding expressions for the interaction quasipotentials and the helicity amplitudes of NN scattering are given in Ref. 48.

We now show that a unified eikonal representation is valid, subject to certain restrictions on the energy dependence of the quasipotential. For this, we use an equation for the wave function equivalent to Eq. (1):

$$(E^2(-i\nabla)-E^2)\hat{\psi}_{\mathbf{p}}(\mathbf{r})=\hat{A}(-i\nabla)\hat{V}(E,\ r)\hat{\psi}_{\mathbf{p}}(\mathbf{r}).\tag{11}$$

The quantities that occur here have already been defined above.

The solution of Eq. (11) in the form of a weakly distorted plane wave,[41]

$$\psi_{\mathbf{p}}(\mathbf{r})=e^{ipz}\hat{F}_{\mathbf{p}}(\mathbf{r}),\tag{12}$$

where $\hat{F}_p(\mathbf{r})$ is a slowly varying function, corresponds to an eikonal nature of high-energy small-angle hadron scattering. By means of operator expansions, we obtain the equation

$$\Big(4p^2-E^2-8ip\frac{\partial}{\partial z}\Big)\hat{F}_{\mathbf{p}}(\mathbf{r})=\hat{A}(\mathbf{p})\hat{V}^+(E,\ r)\hat{F}_{\mathbf{p}}(\mathbf{r}).\tag{13}$$

One can show that a solution of the form (12) is possible only when we have the restriction

$$\hat{A}(\mathbf{p})\hat{V}^+(E,\ r)\leqslant E.\tag{14}$$

Then Eq. (13) becomes

$$-8ip\frac{\partial\hat{F}_{\mathbf{p}}(\mathbf{r})}{\partial z}=A(\mathbf{p})V^+(E,\ r)F_{\mathbf{p}}(r).$$

The solution of this equation can be written in the form

$$\hat{F}_{\mathbf{p}}(r)=1+\sum_{n=1}^{\infty}\Big(\frac{i}{8p}\Big)^n\int\limits_{-\infty}^{\infty}dz_1\ \dots\ dz_n\theta(z-z_1)\ \dots$$

$$\dots\ \theta(z_{n-1}-z_n)\hat{A}(\mathbf{p})\hat{V}^+(\rho,\ z_1)\ \dots\ \hat{A}(\mathbf{p})\hat{V}^+(\rho,\ z_n),\tag{15}$$

and this leads to the unified eikonal representation (3) for the scattering amplitude.

Using the asymptotic behavior of the matrix $\hat{A}(\mathbf{p})$,[45] we immediately obtain from (14) corresponding restrictions on the growth with the energy of quasipotentials that are consistent with eikonal small-angle hadron scattering:

$$\left.\begin{aligned}&\hat{V}_{0,0}(E,\ r)\leqslant E^2;\ \hat{V}_{0,\frac{1}{2}}(E,\ r)\leqslant E;\\ &\hat{V}_{\frac{1}{2},\frac{1}{2}}(E,\ r)\leqslant\text{const}.\end{aligned}\right\}\tag{16}$$

We note that the experimental data indicate a logarithmic growth of the total cross sections in the region $\sqrt{s}\lesssim 100$ GeV. It is easy to show that such behavior corresponds to the maximally permitted growth of the quasipotentials that determine the amplitudes without spin flip (see Ref. 48), for which the interaction range must increase logarithmically. In this case, we obtain restrictions on the ratios of the amplitudes with and without spin flip:

$$\left.\begin{aligned}&\frac{|M^{MN}_{+-}(s,\ t)|}{|M^{MN}_{++}(s,\ t)|}\leqslant\frac{1}{\sqrt{s}};\\ &\frac{|M^{NN}_{++,\ +-}(s,\ t)|}{|M^{NN}_{++,\ ++}(s,\ t)|}\leqslant\frac{1}{\sqrt{s}};\\ &\frac{|M^{NN}_{++,\ --}(s,\ t)|}{|M^{NN}_{++,\ ++}(s,\ t)|}\leqslant\frac{1}{s}.\end{aligned}\right\}\tag{17}$$

Accordingly, we have restrictions on the energy dependence of the physically observable parameters of the polarization and differential cross sections:

$$\mathcal{P}(s,\,t)\leqslant\frac{1}{\sqrt{s}};\quad \frac{d\sigma}{dt}=\frac{d\sigma}{dt}\Big|_{\text{asympt}}+\frac{d\sigma}{dt}\Big|_{\text{spin}},\quad \frac{d\sigma}{dt}\Big|_{\text{spin}}\leqslant\frac{1}{s}, \tag{18}$$

and these are consistent with the experiments.[8]

Thus, the requirement that the eikonal nature of high-energy hadron scattering be valid leads to upper bounds on the energy dependence of the quasipotentials and the helicity amplitudes and, as a result, the spin effects in the region of small scattering angles must die out rapidly with increasing energy.

2. PARTICLE SCATTERING IN THE CASE OF QUASIPOTENTIALS THAT INCREASE RAPIDLY WITH THE ENERGY

For the simple example of scattering of spinless particles, we now consider the consequences of violation of the restriction (16). We show that in this case the eikonal representation can be modified.

In high-energy small-angle scattering, a quasipotential equation is equivalent to a Schrödinger equation with an energy-dependent potential. In this case, the restriction on the growth of the potential with the energy has the form

$$V_{\text{eff}}(E,\,r)\leqslant E. \tag{19}$$

We investigate the possibility of rapid growth of the effective interaction potential with an anomalous term that violates the restriction (19):

$$V_{\text{eff}}(E,\,r)=p^2a(r)+pb(r). \tag{20}$$

Solutions of the Schrödinger equation

$$(\nabla^2+p^2)\,\phi_p(r)=-V_{\text{eff}}(E,\,r)\,\phi_p(r) \tag{21}$$

can now no longer be found in the form of a weakly distorted plane wave (12), since the right- and left-hand sides of the equation increase in the same way with the energy. Therefore, we seek the solution in the form

$$\phi_p(r)=\exp\left(ip\varkappa(r)\right)F(r) \tag{22}$$

with the boundary condition $\phi_p(r)\to_{z\to-\infty}\exp(ipz)$.

We require that in the high-energy limit the transverse momentum components in the scattering process be small compared with the longitudinal momentum of the incident particle. This condition means that we restrict ourselves here to small scattering angles. As a result, we obtain the system of equations

$$\left.\begin{aligned}&\left(\frac{d\varkappa}{dz}\right)^2=1+a(r);\\ &2\frac{d\varkappa}{dz}\frac{dF}{dz}+\left(\frac{d^2\varkappa}{dz^2}+\frac{1}{i}b(r)\right)F(r)=0.\end{aligned}\right\} \tag{23}$$

From this we readily find for the wave function the solution

$$\phi_p(r)=e^{ipz}\exp\left[ip\int_{-\infty}^{z}(\sqrt{1+a(r)}-1)\,dz-\frac{1}{2i}\int_{-\infty}^{z}\frac{b(r)\,dz}{\sqrt{1+a(r)}}-\frac{1}{4}\ln(1+a(r))\right] \tag{24}$$

and the representation corresponding to it for the scattering amplitude[63]:

$$T(\mathbf{p},\,\mathbf{k})=\frac{p}{(2\pi)^3}\int d^2\rho e^{i\Delta\rho}\int_{-\infty}^{\infty}dz\,\Big\{(pa(r)+b(r))\exp\Big[ip\int_{-\infty}^{z}(\sqrt{1+a(r)}-1)\,dz-\frac{1}{2i}\int_{-\infty}^{z}\frac{b(r)\,dz}{\sqrt{1+a(r)}}-\frac{1}{4}\ln(1+a(r))\Big]\Big\}, \tag{25}$$

which does not have an eikonal form. Thus, violation of the restriction (14) can lead to violation of the eikonal representation for the small-angle scattering amplitude.

We now consider how the particle scattering takes place in the presence of the anomalous term of the potential. In this case, the system is described by a wave function that differs appreciably from the free function.

Applying to this wave function the momentum operator, we find its longitudinal and transverse components:

$$\left.\begin{aligned}&p_z\sim p\sqrt{1+a(r)},\\ &p_\perp\sim p\int_{-\infty}^{z}\frac{da(r)}{d\rho}\Big/\left(2\sqrt{1+a(r)}\right)dz.\end{aligned}\right\} \tag{26}$$

Thus, within the interaction region there is a change in the longitudinal momentum of the particle, and we obtain a nonvanishing transverse component, which distorts the front of the plane wave. Within the interaction range the effect of the potential is so great that there is a certain deflection of the particle from its straight path. This necessarily leads to an increase in the effective range of the interaction and, therefore, of the total cross sections. We shall show this below.

We note here that the solutions which we have found for the wave function (24) and for the scattering amplitude are, in general, valid in the case when the transverse momentum component is small compared with the longitudinal component. Using (26), we obtain

$$\frac{p_\perp}{p_z}\sim\frac{\int_{-\infty}^{z}\frac{da(r)}{d\rho}\Big/\left(2\sqrt{1+a(r)}\right)dz}{\sqrt{1+a(r)}}\ll 1,$$

a result that holds when the anomalous term of the quasipotential is small:

$$a(r)\ll 1. \tag{27}$$

In this case, the deflection of the particle from the straight path will be small, and the eikonal nature of the scattering process must be preserved approximately. Indeed, when the condition (27) is satisfied, the expression (25) goes over into the eikonal representation for the scattering amplitude:

$$T(\mathbf{p},\mathbf{k})=\frac{ip}{4\pi^3}\int\times d^2\rho e^{i\Delta\rho}\left[1-\exp\left[-\frac{1}{2ip}\int_{-\infty}^{\infty}dz\,(p^2a(r)+pb(r))\right]\right]. \tag{28}$$

In the presence of an anomalous term in the quasipotential, the eikonal phase in (28) contains the term $p^2a(r)$, which increases rapidly with the energy and begins to be dominant at superhigh energies. If in this region the anomalous term of the eikonal phase has for large impact parameters the form

$$-\frac{p}{2i}\int_{-\infty}^{\infty} dza(r) \sim -p\lambda e^{-(\rho^2/R^2)^{\gamma}},$$

then for the total cross sections we obtain the estimate

$$\sigma_{\text{tot}} \sim R^2 [\ln | p\lambda |]^{1/\gamma}. \tag{29}$$

Therefore, anomalous behavior of the eikonal phase leads to a growth of the total cross sections. At the same time, the Froissart bound permits only square-root branch points ($\gamma=\frac{1}{2}$) of the eikonal phase, and this corresponds to quasipotentials that can be represented in the form of a superposition of Yukawa potentials. The absence of rapid growth of the total cross sections in the region of ISR energies indicates once more a fairly small value of the anomalous term.

As a result, it may be concluded that we have obtained restrictions on the energy dependence of the quasipotentials such that when they are satisfied the eikonal representation for the small-angle scattering amplitude is valid. A consequence of the violation of these restrictions is a modification of the eikonal representation. Our analysis indicates smallness of the possible anomalous terms of the quasipotential.

We now consider the scattering of particles with spin. The quasipotential that describes meson scattering by nucleons has the matrix form (9). Using (9), we can readily show that spin effects with a weak energy dependence arise if the quasipotential $a(s, r)$ has the following growth with the energy:

$$a(s, r) \sim (s/2)\,\alpha(r).$$

The quasipotential equation for particles with spins 0 and 1/2 can be reduced to a two-component equation of the form[55]

$$(\nabla^2+p^2)\hat{\phi}_p(r) = -\hat{V}_{\text{eff}}(s, r)\hat{\phi}_p(r),$$

where $\hat{V}_{\text{eff}}(s, r)$ has a complicated form, but its principal asymptotic term with respect to s is the scalar potential[56]

$$\hat{V}^{\text{as}}_{\text{eff}}(s, r) = -p^2\alpha^2(r).$$

As a result, we arrive at an equation of Schrödinger type with an effective potential that increases anomalously fast with increasing energy. For it, all the results obtained above hold. This permits the conclusion that in the presence of a weak energy dependence of the spin effects the eikonal representation can be modified. A detailed investigation of meson–nucleon scattering in this case was made in Ref. 63. Here, we consider in more detail NN scattering. The quasipotential equation for the wave function of two particles with spins 1/2 and masses m in the center-of-mass system can be written in the form[32]

$$[E-\hat{I}\otimes\hat{H}(-i\nabla)-\hat{H}(i\nabla)\otimes\hat{I}+\gamma_0\otimes\gamma_0\hat{V}(E, r)]\hat{\psi}_p(r)=0, \tag{30}$$

where

$$\hat{H}(i\nabla) = m\gamma_0 + i\alpha\nabla;$$
$$E=\sqrt{s}=2\sqrt{m^2+p^2}\simeq 2p;$$

$\mathbf{p}$ is the momentum of the particle in the initial and final states, and $\hat{\mathbf{I}}$ is the 4×4 unit matrix. We take the quasipotential $\hat{V}(E, r)$ on the basis of the relativistically invariant form for the scattering amplitude in the high-energy limit and take into account anomalous terms leading to a weak energy dependence of the spin effects:

$$\hat{V}(E, r) = \frac{\sqrt{s}}{2}\alpha(E, r) + a(E, r) + [\hat{I}\otimes\hat{n}(-\mathbf{l}) + \hat{n}(\mathbf{l})\otimes\hat{I}] \times\left[\frac{\sqrt{s}}{2}\beta(E, r)+b(E, r)\right] + \hat{n}(\mathbf{l})\otimes\hat{n}(-\mathbf{l})\,d(E, r). \tag{31}$$

To solve Eq. (30), we shall seek the 16-component wave function $\hat{\psi}_{\mathbf{p}}(\mathbf{r})$ in the form

$$\hat{\psi}_{\mathbf{p}}(\mathbf{r}) = \begin{bmatrix} f_1(\mathbf{r}) \\ f_2(\mathbf{r}) \\ f_3(\mathbf{r}) \\ f_4(\mathbf{r}) \end{bmatrix} \chi(\mathbf{p})\otimes\chi(-\mathbf{p}), \tag{32}$$

where the matrices f_1, f_2, f_3, and f_4 have four rows each; $\chi(\mathbf{p})$ and $\chi(-\mathbf{p})$ are free two-component spinors corresponding to particles 1 and 2.

For the functions f_i we can write down a system of equations, from which we find the corresponding equation for each of the components. Thus, for f_i, chosen in the form (22), we have[49]

$$\left[\tilde{A}_1p+\tilde{A}_0-(\tilde{B}_1p+\tilde{B}_0)\frac{1}{A_1p+A_0}(B_1p+B_0)\right]F(r)=0, \tag{33}$$

where

$$A_1 = 2\left(1+\frac{\alpha}{2}-\beta-\varepsilon_+\varphi_+\right);$$
$$A_0 = 2(m-b)-(a+d)+2i(\partial_z\varphi_+)+4i\varphi_+\partial_z+4(d-b)\varphi_+ + 2(d-2a)\varphi_+^2-([\sigma^{(1)}, \nabla]_z\varphi_+)-([\sigma^{(2)}, \nabla]_z\varphi_+);$$
$$B_1 = 2\varepsilon_+\varphi_-\sigma_z^{(1)}\sigma_z^{(2)};$$
$$B_0 = -i\varphi_-(\sigma_z^{(1)}\sigma^{(2)}\nabla+\sigma_z^{(2)}\sigma^{(1)}\nabla)-i(\sigma_z^{(1)}\sigma^{(2)}\nabla+\sigma_z^{(2)}\sigma^{(1)}\nabla)\varphi_+ -[d+2(d-b)\varphi_-+2(d+b)\varphi_++2(a+2b)\varphi_+\varphi_-]\sigma_z^{(1)}\sigma_z^{(2)} -\frac{1}{2}\beta(\varphi_+-\varphi_-)(\sigma_\perp^{(1)}\sigma_z^{(2)}-\sigma_\perp^{(2)}\sigma_z^{(1)})\Delta;$$
$$\varphi_\pm = \varepsilon_\pm/(2-\alpha),\ \varepsilon_\pm=(\partial_z\varkappa)\pm\beta.$$

Here

$$\sigma^{(1)}=\begin{pmatrix}1&0\\0&1\end{pmatrix}\otimes\sigma,\quad \sigma^{(2)}=\sigma\otimes\begin{pmatrix}1&0\\0&1\end{pmatrix},$$

where σ are the Pauli matrices. The operators $\tilde{A}_1, \tilde{A}_0, \tilde{B}_1, \tilde{B}_0$ are obtained from A_1, A_0, B_1, B_0, respectively, by the substitution

$$\beta\to-\beta,\quad b\to-b,\quad m\to-m,\quad \sigma^{(1)}\leftrightarrow\sigma^{(2)}.$$

Expanding the left-hand side of Eq. (33) in powers of $1/p$, equating to zero the coefficients of the two leading powers, and taking into account the smallness of the anomalous terms of the quasipotential, $\alpha\ll 1$, $\beta\ll 1$, we obtain for the functions $\varkappa$ and F the equations

$$\frac{d\varkappa(r)}{dz} = 1-\frac{\alpha^2(r)}{8}-2\beta^2(r); \tag{34}$$

$$\frac{dF(r)}{dz} = \frac{1}{2i}\left\{4d(r)-[\sigma^1+\sigma^{(2)}, \mathbf{n}_\rho]_z\frac{d\beta(r)}{d\rho}\right\}F(r). \tag{35}$$

From (34) and (35) we have[49]

$$\varkappa(r) = z - 2\int_{-\infty}^{z}\left(\beta^2(\rho, z') + \frac{\alpha^2(\rho, z')}{16}\right) dz';$$

$$F(\mathbf{r}) = \exp\left[\frac{2}{i}\int_{-\infty}^{z} d(\rho, z')\, dz'\right]$$

$$\times\left\{1 + \frac{1}{4}\left[J^{(1)}(\beta; \mathbf{r}) + J^{(2)}(\beta; \mathbf{r})\right]\right\}, \qquad (36)$$

where

$$J^{(1,2)}(\beta; \mathbf{r}) = \frac{2}{i}[\sigma^{(1,2)}, \mathbf{n}_\rho]_z \int_{-\infty}^{z} \frac{d\beta(\rho, z')}{d\rho} dz'; \quad \mathbf{n}_\rho = \frac{\rho}{|\rho|}.$$

As a result, the solution of Eq. (30) can be expressed in the form

$$\hat{\psi}_p(\mathbf{r}) \simeq \frac{1}{2}\begin{bmatrix} 1 \\ \sigma_z^{(1)}\left(1 + \frac{\alpha}{2} + 2\beta - \frac{1}{p} J^{(1)}(d; \mathbf{r})\right) \\ -\sigma_z^{(2)}\left(1 + \frac{\alpha}{2} + 2\beta - \frac{1}{p} J^{(2)}(d; \mathbf{r})\right) \\ -\sigma_z^{(1)}\sigma_z^{(2)}\left(1 + 4\beta - \frac{1}{p}(J^{(1)}(d; \mathbf{r}) + J^{(2)}(d; \mathbf{r}))\right) \end{bmatrix}$$
$$\times e^{ip\varkappa(r)} F(\mathbf{r})\, \chi(\mathbf{p}) \otimes \chi(-\mathbf{p}). \qquad (37)$$

Determining the scattering amplitude by the relation

$$T(\mathbf{p}', \mathbf{p}) = \frac{1}{4\pi}\int d\mathbf{r} \bar{\psi}_{\mathbf{p}'}^{(0)}(\mathbf{r})\, \hat{V}(r)\, \psi_{\mathbf{p}}(\mathbf{r}),$$

we obtain in the limit $p \to \infty$ by means of the free solution of Eq. (30),

$$\bar{\psi}_{\mathbf{r}}^{(0)} \simeq \frac{\chi^+(\mathbf{k}) \otimes \chi^+(-\mathbf{k})}{2}\Big[1, -\sigma_z^{(1)} + \frac{1}{p}\sigma_\perp^{(1)}\Delta_\perp, \ \sigma_z^{(2)} - \frac{1}{p}\sigma_\perp^{(2)}\Delta_\perp,$$
$$-\sigma_z^{(1)}\sigma_z^{(2)} + \frac{1}{p}(\sigma_z^{(1)}\sigma_\perp^{(2)} + \sigma_z^{(2)}\sigma_\perp^{(1)})\,\Delta\Big]\, e^{-i\mathbf{kr}},$$

the scattering amplitude in the form

$$T(\mathbf{k}, \mathbf{p}) = \chi^+(\mathbf{k}) \otimes \chi^+(-\mathbf{k})[T_{++,++}(\mathbf{k}, \mathbf{p}) + i\sigma_y^{(1)} T_{++,+-}(\mathbf{k}, \mathbf{p})$$
$$+ i\sigma_y^{(2)} T_{++,-+}(\mathbf{k}, \mathbf{p})]\,\chi(\mathbf{p}) \otimes \chi(-\mathbf{p}),$$

where

$$\left.\begin{aligned} T_{++,++}(\mathbf{k}, \mathbf{p}) &= \frac{i}{2\pi^2}\int \rho\, d\rho J_0(\rho\Delta)\,[1 - e^{\chi_0(s,\rho)}]; \\ T_{++,+-}(\mathbf{k}, \mathbf{p}) &= T_{++,-+}(\mathbf{k}, \mathbf{p}) \\ &= -\frac{1}{2\pi^2}\int \rho\, d\rho J_1(\rho\Delta)\,\chi_1(s, \rho)\, e^{\chi_0(s,\rho)}; \\ \chi_0(s, \rho) &= -\frac{2}{i}\int_{-\infty}^{\infty}\left\{d(\rho, z) - \frac{\sqrt{s}}{2}\left[\beta^2(\rho, z) + \frac{\alpha^2(\rho, z)}{16}\right]\right\} dz; \\ \chi_1(s, \rho) &= \frac{1}{2i}\int_{-\infty}^{\infty}\frac{d\beta(\rho, z)}{d\rho} dz. \end{aligned}\right\} \qquad (38)$$

Note that here we have found only the leading terms of the scattering amplitude without flip and with a single flip of the spin. If the energy dependence of the interaction quasipotential has the form (31), the ratio of the amplitudes with and without flip has the form

$$\frac{|T_{++,+-}(s, t)|}{|T_{++,++}(s, t)|} \sim \text{const}; \quad \frac{|T_{++,--}(s, t)|}{|T_{++,++}(s, t)|} \sim \frac{1}{\sqrt{s}}.$$

Thus, we have obtained modified eikonal representations for the nucleon–nucleon scattering amplitudes (38) valid for small anomalous terms of the quasipotential (31).

These expressions have the same form as the standard eikonal representation. The difference is that the eikonal phase χ_0 has a term that increases as $\sqrt{s}$ and which contains the small squares of the anomalous terms of the quasipotential. These terms have a spin nature. Indeed, it is easy to show that the quasipotentials β and α determine the leading terms of the amplitudes with spin flip of one or two particles. The amplitude without spin flip can be obtained by multiple scattering by the quasipotential $d(s, r)$, and also by twofold scattering by the quasipotentials β and α. In the normal case, such terms are suppressed in a power-law manner, but in the anomalous case they become decisive in the limit $s \to \infty$.

We note that eikonal phases containing squares of the quasipotentials were obtained earlier in Refs. 55 and 56. In them, increasing terms also occur in the case of spin effects that depend weakly on the energy.

Anomalous terms of this kind in the eikonal phase are unimportant at low energies. But at superhigh energies they lead to a rapid growth of the total cross sections and to a number of other effects that we shall discuss in detail below.

3. INTERACTION DYNAMICS AND HADRON STRUCTURE AT LARGE DISTANCES

By means of the expressions found for the physical quantities, it is possible to make a detailed investigation of the small-angle scattering of particles with spin and thus obtain information about strong-interaction dynamics at high energies. For this, using model representations, it is necessary to determine the soft quasipotential, which is the Born term of the scattering amplitude. To solve the posed problem, it is necessary to understand the structure of the t-channel contributions, which determine the dynamics of high-energy hadron interactions at large distances.

We represent the Born term of the scattering amplitude as a sum of the t- and u-channel contributions:

$$T(s, t, u) = T(s, t) + T(s, u), \qquad (39)$$

where $T(s, t)$ and $T(s, u)$ have peaks for forward and backward scattering, respectively. The first term of the sum (39) has the form (0.2), and for the second we have the representation

$$T(s, u) = g(s)\int_0^\infty dy \rho_u(s, y)\, e^{yu}. \qquad (40)$$

In the region $|t|/s \ll 1$, it is natural to consider the contribution to the amplitude (39) of only the t-channel term.

By an inverse Laplace transformation it is possible to calculate the density function $\rho_t(s, x)$ in (0.2):

$$g(s)\,\rho_t(s, x) = \frac{1}{2\pi i}\int_{-i\infty}^{i\infty} dt e^{-xt}\, T(s, t). \qquad (41)$$

Closing the contour of integration in the region of positive t, we obtain

$$g(s)\,\rho_t(s, x) = \frac{1}{\pi}\int_{t_0}^{\infty} dt e^{-xt}\, \mathrm{Im}_t\, T(s, t)$$
$$+ (\text{contribution of pole terms}). \qquad (42)$$

The representation (42) for the density function can be readily obtained from dispersion relations for the scattering amplitude at fixed s.

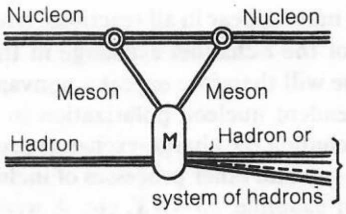

FIG. 1. Contribution of two-meson exchange to hadron–hadron processes at large distances.

The representation (42) has a form closely similar to the sum rules of Ref. 24. It can be regarded as a sum rule for determining the parameters of the soft part of the density function by means of the calculated imaginary part of the amplitude $T(s, t)$ on the t-channel cut.

Note that in the study of high-energy small-angle hadron scattering one can consider in (42) only the contributions of the leading asymptotic terms of the scattering amplitude. In this case the pole terms in (42) can be ignored. As was shown above, in the region of fixed t, which is determined by the particle interaction at distances of the order of the hadron diameter, values of x satisfying $x \gtrsim 1\ \mathrm{GeV}^{-2}$ are important in the density function. The presence of the exponential factor e^{-xt} in (42) makes possible a restriction in this case to the contributions of the nearest singularities in the t channel. Therefore, for the calculation of the imaginary part of the scattering amplitude we take into account in the unitarity condition only the elastic contributions corresponding to the 2π cut. Taking into account in the scattering amplitudes in the unitarity condition the pole term and the asymptotic term, we find that only the triangle singularities in the t channel shown in Fig. 1 contribute significantly to $\mathrm{Im}_t T(s, t)$. They are associated with allowance for the pion cloud around the hadrons and, thus, are determined by the particle structure at large distances.

For the πN scattering amplitude, the form of which is important for calculating the contribution of the diagrams in Fig. 1 to the integral (42), we use the very simple analytic form

$$T_{\pi N}(s, t) \sim is \exp[-b\sqrt{m^2 - t}]$$

with parameters determined from the experimental data on high-energy πN scattering at momentum transfers $|t| \leqslant 1\ \mathrm{GeV}^2$. The contribution of the triangle singularity to the density function $\rho_t^{NN}(s, x)$ found by means of the sum rule (42) decreases rapidly (exponentially) at small and large x and has a broad maximum at $x \sim 3\ \mathrm{GeV}^{-2}$ (Fig. 2). This function can be approximated by

$$\rho_t^{NN}(s, x) \sim is \exp\left[-\frac{c^2}{4x} - d^2 x\right]$$

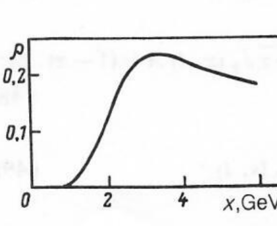

FIG. 2. Contribution of the triangle singularity to the density function $\rho_t^{NN}(s,x)$.

with the parameter values

$$c \simeq 5\ \mathrm{GeV}^{-1}, \quad d \simeq 0.7\ \mathrm{GeV}.$$

This contribution completely determines the behavior of the eikonal phase of pp scattering at large impact parameters $\rho \gtrsim 5\ \mathrm{GeV}^{-1}$.

Naturally, the 2π exchange effects taken into account in the sum rule cannot completely reproduce the soft part of the interaction in the complete range of variation of the impact parameter, since for $\rho \leqslant 1/2\mu_\pi \sim 4\ \mathrm{GeV}^{-1}$ it is also necessary to calculate diagrams containing, for example, inelastic contributions with the exchange of four, six, etc., mesons, and also the exchange of heavier states. It is difficult to take into account all these effects. However, they have a central nature, and their contribution can be separated in the form of a certain function that possesses free parameters.

Thus, by means of the sum rule (42) we have shown that the triangle diagrams with pion exchange in the t channel make the main contribution to the density function in the region of large distances. The same results are obtained in the dynamical model of hadron interaction of Ref. 50, in which the effects of the meson clouds around the particles are taken into account and the contribution of the diagrams in Fig. 1 to the hadron scattering amplitude are calculated.

The model is based on fairly simple dynamical ideas, close to those of Ref. 58, in which the hadron is regarded as an object consisting of a hard central part surrounded by a cloud of quark–antiquark pairs. The model made it possible to explain on a unified basis all the basic properties of high-energy hadron scattering in a wide range of the momentum transfer at ISR energies.[59,60] The ideas on which the model is based lead naturally to a slow change in the slope of the diffraction peak at small momentum transfers, to the presence of a unique diffraction dip, and to a small slope beyond the second diffraction maximum, which is determined by the effective radius of the central interaction region, this being ~0.5 F. We emphasize that the calculations made for $\pi^- p$ scattering predicted that the diffraction dip at $p_L = 200$ GeV arises in the region $|t| \sim 3.5$–$4.2\ \mathrm{GeV}^2$[2,61] in agreement with the experiment of Ref. 62 (Fig. 3). In the spinless case, the model is equivalent to the sum-rule contributions considered above and admits a simple generalization to the case of particles with spin.

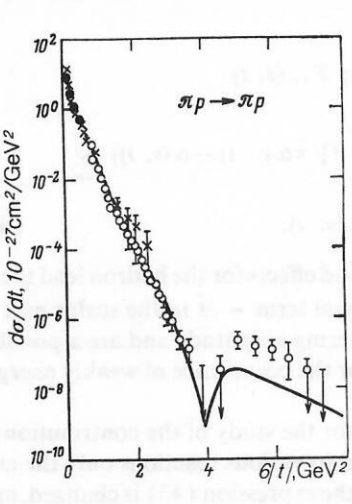

FIG. 3. Differential cross sections of $\pi^- p$ scattering at energy $E_L = 200$ GeV.

The satisfactory description given by this dynamical model for a large group of hadron scattering processes suggests that its picture of hadron structure at large distances corresponds to physical reality. For this reason, its generalization to the interaction of particles with spin is extremely topical.

We consider a hadron process with participation of a nucleon. The contribution of the meson-cloud effects to the scattering amplitude can be represented by a diagram (see Fig. 1) in which M is the amplitude of the meson–hadron subprocess and depends on the variable s'.

The corresponding integral for the scattering amplitude can be represented in the form

$$T_{sp}(s,\ t)=\frac{g^2}{i(2\pi)^4}\int d^4qM(s',\ t)\,\varphi[(k-q)^2,\ q^2]\,\varphi[(p-q)^2,\ q^2]$$
$$\times\frac{\bar{u}(p)\,\gamma_5\,(\hat{q}+m)\,\gamma_5 u(k)}{[q^2-m^2+i\varepsilon]\,[(k-q)^2-\mu^2+i\varepsilon]\,[(p-q)^2-\mu^2+i\varepsilon]},$$
$$s'=[k+k'-q]^2. \tag{43}$$

Here, m and μ are the masses of the nucleon and the meson. The matter distribution within the nucleon is taken into account by means of the functions φ. For simplicity, we have averaged over the nucleon spins in the amplitude M in (43). In this case, it is a scalar quantity. Going over to light-cone variables in (43), $q_+=q_0+q_z=xk_+$, $q_-=q_0-q_z$, $\mathbf{q}_\perp$, and integrating over q_-, we obtain for the matrix structure in the numerator of (43) the expression

$$\hat{q}-m=\frac{1}{2}xk_+\gamma_-+\frac{1}{2}\frac{\mathbf{q}_\perp^2+m^2}{xk_+}\gamma_+-\mathbf{q}_\perp\boldsymbol{\gamma}_\perp-m. \tag{44}$$

Calculations[53] show that the πN scattering amplitude obtained by means of (43) can be transformed to the form

$$T_{\pi p}(s,\ t)=\bar{u}(\mathbf{p})\,[\sqrt{s}\,\alpha(s,\ t)+a(s,\ t)+b(s,\ t)\,\hat{n}(-\mathbf{l})]\,u(\mathbf{k});$$
$$\hat{n}(-\mathbf{l})=\gamma_0+\boldsymbol{\gamma}\mathbf{l}/|\mathbf{l}|;\quad \mathbf{l}=(\mathbf{p}+\mathbf{k})/2;\quad \alpha\sim a\sim b. \tag{45}$$

Note that the first and last terms of the expression (44) contribute to $\alpha(s,t)$. The contribution contains the mass of the intermediate nucleon state, and also the masses of the scattering nucleons, which are also dimensional parameters determining the anomalous (containing an additional $\sqrt{s}$) terms of the scattering amplitude.

For the matrix elements of MN scattering we have from (45)

$$T_{++}(s,\ t)\underset{s\to\infty}{\sim}b(s,\ t);\ T_{+-}(s,\ t)$$
$$\sim\sqrt{\frac{|t|}{s}}\,(\sqrt{s}\,\alpha(s,\ t)+a(s,\ t))\underset{s\to\infty}{\sim}$$
$$\sim\sqrt{|t|}\,\alpha(s,\ t). \tag{46}$$

Thus, the meson-cloud effects for the hadron lead to the appearance of an additional term $\sim\sqrt{s}$ in the scalar part of the meson–nucleon scattering amplitude and are a possible dynamical mechanism for the occurrence of weakly energy-dependent spin effects.

We emphasize that for the study of the contribution of the hadron's meson cloud to various reactions only the amplitude M (see Fig. 1) in the expression (43) is changed, and this does not change the structure of the expression (43). It may be concluded from this that anomalous terms containing an additional factor $\sqrt{s}$ must appear in all reactions, irrespective of the structure of the t-channel exchange in the amplitude M of Fig. 1. One will therefore expect a nonvanishing weakly energy-dependent nucleon polarization in a large class of processes, including the charge-exchange processes $\pi^-p\to\pi^0n$, $\pi^-p\to\eta n$ and other processes of inclusive nucleon scattering, for example, $pp\to\Lambda X$, etc. A weak energy dependence of the polarization has been confirmed in the last process, but in this direction further theoretical and experimental investigations are needed.

4. PERIPHERAL EFFECTS IN NUCLEON–NUCLEON SCATTERING AT $\sqrt{s}\leqslant 100$ GeV

We consider in more detail nucleon–nucleon scattering. The study of this process is very topical, since the experimental material is richest here and the highest energies have been achieved.

For nucleon–nucleon scattering, we must substitute for the amplitude in (43) the matrix element of the meson–nucleon scattering amplitude:

$$T_{\pi N}(s',\ t)=\bar{u}(-\mathbf{p})\left[\frac{\sqrt{s'}}{2}A(s',\ t)+\left(\frac{\hat{p}+\hat{k}}{2}-\hat{q}\right)B(s',\ t)\right]$$
$$\times u(-\mathbf{k}). \tag{47}$$

Since the anomalous terms of the meson–nucleon scattering amplitude are small,[49] they can be ignored in (47). As a result, it can be assumed that the invariant amplitudes A and B have a weak energy dependence.

Assuming that the matter distribution within the nucleon is the same as the charge distribution, we use for the function φ in (43) the standard dipole approximation

$$\varphi(l^2,\ q^2=m^2)=\beta^4/(l^2-\beta^2-i\varepsilon)^2,\quad \beta^2=0.71\ \mathrm{GeV}^2,$$

which can be made since the nucleon in the intermediate state is near the mass shell.

Calculating the corresponding integrals, we find that the matrix element of the peripheral part of the nucleon–nucleon scattering amplitude in the region of high energies and small momentum transfers can be represented in the form[65]

$$T_p(s,\ t)$$
$$=\bar{u}(-\mathbf{p})\otimes u(\mathbf{p})\,\{\tilde{a}_p(s,\ t)+\tilde{b}_p(s,\ t)\,[\hat{n}(\mathbf{l})\otimes\hat{I}+\hat{I}\otimes\hat{n}(-\mathbf{l})]$$
$$+d_p(s,\ t)\,\hat{n}(\mathbf{l})\otimes\hat{n}(-\mathbf{l})\}\,u(-\mathbf{k})\otimes u(\mathbf{k}). \tag{48}$$

The amplitudes in (48) have the form

$$\tilde{a}_p(s,\ t)=\alpha_p(s,\ t)\frac{\sqrt{s}}{2}+a_p(s,\ t); \tag{48a}$$

$$\alpha_p(s,\ t)=-\frac{4g^2\beta^8}{(2\pi)^3}M\int_0^1 dxx^5(1-x)^{3/2}J_0(x,\ t)\,A(s(1-x),\ t) \tag{48b}$$

$$a_p(s,\ t)=\frac{2g^2\beta^8}{(2\pi)^3}\Delta^2\int_0^1 dxx^5\sqrt{1-x}\,J_2(x,\ t)\,A(s(1-x),\ t); \tag{48c}$$

$$\tilde{b}_p(s,\ t)=\beta_p(s,\ t)\frac{\sqrt{s}}{2}+b_p(s,\ t); \tag{49a}$$

$$\beta_p(s,\ t)=-\frac{2g^2\beta^8}{(2\pi)^3}M\int_0^1 dxx^5(1-x)^2J_0(x,\ t)\,B(s(1-x),\ t); \tag{49b}$$

$$b_P(s,\ t)=\frac{2g^2\beta^8}{(2\pi)^3}\int_0^1 dx x^5\left\{\frac{\Delta^2}{2}J_2(x,\ t)(1-x)B(s(1-x),\ t)\right.$$
$$+\frac{1}{4}\left[M^2\frac{(1-x^2)}{x}J_0(x,\ t)+\Delta^2 J_2(x,\ t)\right.$$
$$\left.\left.+\frac{1}{x}J_1(x,\ t)\right]A(s(1-x),t)\sqrt{1-x}\right\};\qquad(49c)$$

$$d_P(s,\ t)=\frac{g^2\beta^8}{(2\pi)^3}\int_0^1 dx x^5\left[M^2\frac{(1-x^2)}{x}J_0(x,\ t)+\Delta^2 J_2(x,\ t)\right.$$
$$\left.+\frac{1}{x}J_1(x,\ t)\right](1-x)B(s(1-x),\ t).\qquad(49d)$$

The integrals $J_i(x,t)$ are determined in Ref. 53.

The quasipotentials $\tilde{a}_p$ and $\tilde{b}_p$ have anomalous terms that lead to a weak energy dependence of the spin effects. Calculating the integrals (48) and (49) by means of a Gaussian shape of the meson–nucleon scattering amplitude, we find that $\alpha\sim\beta$. The quasipotential β can be approximated by an expression convenient for concrete calculations:

$$|\beta(r)|\sim H\exp(-M\sqrt{B^2+r^2})\qquad(50)$$

with parameters $H=132$, $M=0.80$ GeV, $B=9.7\ \mathrm{GeV}^{-1}$. This approximation shows that the anomalous term of the quasipotential is small compared with the normal terms. As shown above, in this case contributions that grow as $\sqrt{s}$ and contain the square of the quasipotential β can appear in the eikonal phase. Note that for consistent determination of the interaction quasipotential it is, in general, necessary to calculate the second Born term of the scattering amplitude.[28] This is an extremely complicated problem, since its solution requires calculations in field theory of all the three-loop diagrams which occur. Therefore, we shall use below the amplitudes (38) obtained for interaction quasipotentials determined in the Born approximation. We shall assume that in the general case the form of the expressions (38) remains the same. The criterion for testing the predicted effects will be comparison of the theory with the experimental data. As a result, for the quasipotentials corresponding to (48) and (49) and with allowance for the fact that $\alpha^2/16\ll\beta^2$ we obtain the relations (38) for the NN scattering amplitudes with eikonal phases

$$\chi_0(s,\ \rho)=-\frac{2}{i}\left[\int_{-\infty}^{\infty}d_P(s,\ r)\,dz-\frac{\sqrt{s}}{2}\int_{-\infty}^{\infty}dz\beta^2(s,\ r)\right]$$
$$=\chi_P(s,\ \rho)+\frac{\sqrt{s}}{2}\chi_{\mathrm{an}}(s,\ \rho);$$
$$\chi_1(s,\ \rho)=\frac{1}{2i}\frac{d}{d\rho}\int_{-\infty}^{\infty}\beta(s,\ r)\,dz=\frac{d}{d\rho}\tilde{\chi}_1(s,\ \rho).\qquad(51)$$

It should be noted that at high energies only the amplitude $B(s,t)$ of meson–nucleon scattering is known with a sufficient degree of accuracy. It is this amplitude that determines the leading asymptotic terms of the nucleon–nucleon amplitudes (48) and (49) calculated in the model. The amplitude $A(s,t)$ makes it possible to determine the peripheral contributions to the spin-flip NN scattering amplitudes, which decrease with the energy. As shown above, they are important at relatively low energies $E_L\lesssim 100$ GeV.

Estimates show that the contribution of the anomalous terms to the scattering cross sections becomes important only at energies $\sqrt{s}\gtrsim 100$ GeV, except for the region of the

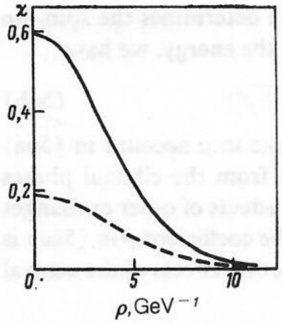

FIG. 4. Contribution of the region of large distances to the eikonal phases. The continuous curve represents $2\chi_0(s,\rho)$; the broken curve, $2\chi_1(s,\rho)$.

diffraction dip, where their influence can be noted somewhat earlier.

The peripheral parts of the eikonal phases $\chi_p(s,\rho)$ and $\chi_1(s,\rho)$ calculated in the model for $\sqrt{s}=53$ GeV are given in Fig. 4.

In actual calculations analyzing experimental data on pp scattering we approximated the calculated phases by

$$-\chi_P(s,\ \rho)=h_P(b_P^2+\rho^2)\left[\exp(-\mu_P'\sqrt{b_P^2+\rho^2})\right.$$
$$\left.+h_{\mathrm{as}}\exp(-\mu_P''\sqrt{b_P^2+\rho^2})\right];\qquad(52a)$$
$$\chi_{\mathrm{an}}(s,\ \rho)=4(1-i)H^2(B^2+\rho^2)^{1/2}K_1(M\sqrt{B^2+\rho^2}).\qquad(52b)$$

In the anomalous term, we have introduced the factor $(1-i)$ to take into account the $s\to u$ crossing-symmetric diagram in the scattering amplitude. This leads to asymptotic equality of the pp and $p\bar{p}$ cross sections as $s\to\infty$.

There will also be contributions to the total eikonal phase from the interaction of the central parts of the hadrons; we determine these phenomenologically in the form

$$-\chi_c(s,\ \rho)=h_c\exp(-\mu_c(s)\sqrt{b_c^2+\rho^2}).\qquad(53)$$

where $\mu_c(s)$ and b_c are the effective mass and effective interaction range of the central parts of the hadrons.

To determine the energy dependence of the effective masses and interaction ranges we use the hypothesis of geometrical scaling,[64] which must be approximately satisfied in the region of ISR energies, and local dispersion relations.[65,66] For $\chi_p(s,\rho)$ we then obtain

$$\chi_P(s,\ \rho)=\chi_P(\rho^2/b^2(\varkappa_1(s)));\quad \varkappa_1(s)=\left[1+\alpha_P\left(\ln s-i\frac{\pi}{2}\right)\right]^{1/2}.\qquad(54)$$

But for $\chi_c(s,\rho)$ we consider the energy dependence only in the effective mass of the interaction:

$$\mu_c(s)=\mu_c^0/\varkappa_2(s);\quad \varkappa_2(s)=\left[1+\alpha_c\left(\ln s-\frac{i\pi}{2}\right)\right]^{1/2}.\qquad(55)$$

In the model, the form of the peripheral part of the $1/\sqrt{s}$ term of the eikonal phase is close to (52). This made it possible to use for the $1/\sqrt{s}$ term the same form of the eikonal phase. For the total eikonal phase of the amplitude without spin flip we obtain as a result

$$\chi_0(s,\ \rho)=\chi_c(s,\ \rho)(1+\gamma\chi_c(s,\ \rho))+g_1\chi_P(s,\ \rho)$$
$$+g_2^2\chi_{\mathrm{an}}(s,\ \rho)\frac{\sqrt{s}}{2}$$
$$+(A+iB)/\sqrt{s}\,(\chi_c(s,\ \rho)+g_1\chi_P(s,\ \rho)),\qquad(56a)$$

and for the phase $\chi_1(s, \rho)$, which determines the spin-flip amplitude that varies slowly with the energy, we have

$$\chi_1(s, \rho) = \rho H g_2 M K_0 (M \sqrt{B^2+\rho^2}). \qquad (56b)$$

The coefficients g_1 and g_2 take into account in (56a) and (56b) the possible deviation from the eikonal phases calculated in the model due to the effects of other exchanges (see, for example, Ref. 67), and the coefficient γ in (56a) is determined by the inelastic interaction effects of the central parts of the hadrons.

For the terms of the spin-flip amplitude that decrease with the energy we use the simplest Gaussian approximation:

$$|T(s, t)|^2 = \frac{|t|}{s} h^2_{1/\sqrt{s}} e^{-2b_{1/\sqrt{s}}|t|}. \qquad (56c)$$

The parameters in (56a)–(56c) were determined by analyzing the experimental data on the differential cross sections of pp scattering at energies $4.5 \leqslant \sqrt{s} \leqslant 540$ GeV.[68–71,75] A $\chi^2/\bar{\chi}^{2} = 0.98$ quantitative description of the experimental data for $N = 1227$ experimental points (see also Table I and Fig. 5) was obtained for the parameter values $\mu_c^0 = 0.687 \pm 0.005$ GeV, $b_c = 2.185 \pm 0.015$ GeV^{-1}, $\alpha_p = 0.188 \pm 0.007$, $\alpha_c = 0.031 \pm 0.002$, $A = 1.82 \pm 0.03$ GeV, $B = 4.02 \pm 0.04$ GeV, $h_c = 4.67 \pm 0.05$, $\gamma = 0.219 \pm 0.002$, $g_1 = 1.0$, $g_2 = 0.85$.

It can be seen from Table II that in the region of ISR energies the assumptions made in many models of approximate geometrical scaling are indeed satisfied:

$$\frac{B(0)}{\sigma_{tot}} \simeq \text{const}; \quad \frac{d\sigma/dt_{\,2\text{nd max}}}{\sigma^2_{tot}} \simeq \text{const}.$$

However, at energy $\sqrt{s} = 540$ GeV we already note a significant difference, particularly for the second relation. Note that, as is found experimentally,[88] an appreciable deviation from KNO scaling is also observed at this energy.

Despite the fact that in the analysis of the experimental data by means of the expression (56) the contribution of the anomalous term was felt only in the region $0.8 \leqslant |t| \leqslant 3$ GeV^2, the constant of this term can be determined reliably with 5% accuracy, confirming the presence of such a term in the eikonal phase. Its contribution to the differential cross sections in the region of the diffraction dip at $\sqrt{s} = 60$ GeV was $\sim 1 \cdot 10^{-32}$ cm^2/GeV^2, i.e., about 30%. This result makes it possible to consider the predictions of the model at much higher energies, where the contribution of the anomalous term becomes larger, indeed dominant in the region of the diffraction dip.

An important property of the model described above is the possibility of using it for $p\bar{p}$ scattering at sufficiently low energies. Indeed, the amplitude of $p\bar{p}$ scattering obtained

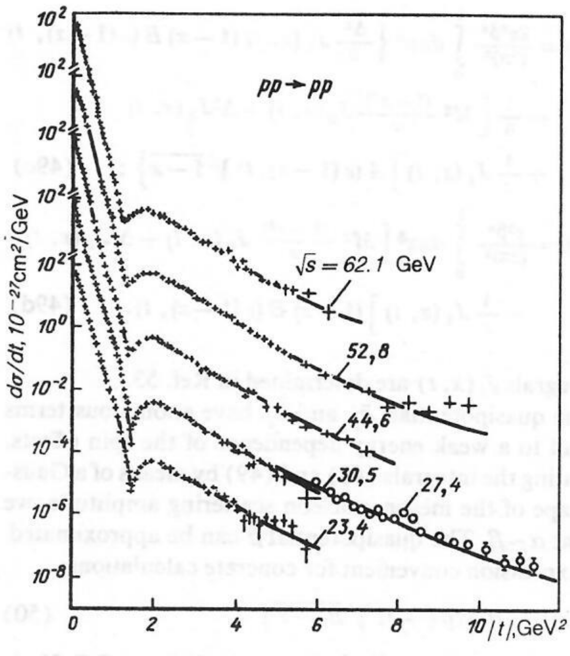

FIG. 5. Differential cross sections of *pp* scattering in the region of small and large momentum transfers (the theoretical curves at all energies pass effectively through the experimental points, including the regions of the diffraction dip and maximum).

from (56) by means of $s \to u$ crossing agrees well with the recent experimental data on $p\bar{p}$ scattering at $p_L = 50$ GeV (Ref. 69) and $p_L = 1850$ GeV (Ref. 72) (Fig. 6). The presence of a clearly expressed diffraction structure in $p\bar{p}$ scattering at $p_L = 50$ GeV is explained by the fact that in the region of the diffraction dip the decisive contribution to the differential cross section is made by the real part of the amplitude without spin flip, which in this case is approximately two times less[60] than for pp scattering. This is a consequence of the relationship of the total cross sections, which determines the coefficients in the $1/\sqrt{s}$ term of χ_0 as $B \simeq 2.2A$.

The proximity of the predictions of the model to the experimental data on $p\bar{p}$ scattering at sufficiently low energies is an indication of the small contribution of the spin-flip amplitudes to the differential cross sections of hadron scattering in a wide range of momentum transfers.

We note that for ISR energies the differential cross sections of pp and $p\bar{p}$ scattering calculated in the model are close to each other, in good agreement with recent experimental data obtained at CERN[73] in the region of the diffraction peak.

In the region of small momentum transfers the model leads to a smooth decrease in the slope of the diffraction peak

TABLE I. Description of experimental data on pp scattering at different energies.

Interval of t, GeV^2	0.03–0.5	0.07–1.03	0.035–2.85	0.02–3.5	0.05–14.2	0.05–10.0	0.03–0.5	0 (σ_{tot})
Interval of $\sqrt{s}$, GeV	4.5–9.3	9.78	13.6–18.17	19.4	23.4–30.5	44.5–62.1	540	4.5–62.1
$\chi^2/\bar{\chi}^2$	1.16	0.79	0.9	1.02	0.95	1.05	0.92	0.87
Number of experimental points	95	58	123	170	316	290	84	91

TABLE II. Verification of the hypothesis of geometrical scaling.

$\sqrt{s}$, GeV	9.78	13.4	19.4	30.5	62.1	540
$\frac{B(0)}{\sigma_{tot}}$	0.29	0.3	0.31	0.32	0.32	0.28
$\frac{\left.\frac{d\sigma}{dt}\right\vert_{2nd\ max}}{\sigma_{tot}^2}\ 10^8$	5.9	4.0	3.15	3.0	3.3	48.3

with increasing $|t|$. Such behavior is universal for different hadron reactions at high energies, and is confirmed experimentally at $p_L = 70$ GeV (Serpukhov) (Ref. 74), $p_L = 200$ GeV (Fermilab), and $\sqrt{s} = 540$ GeV (CERN).[75]

Thus, the dynamical model developed here, which takes into account the effects of the meson cloud of the hadron and the spins of the interacting particles, makes it possible to reproduce completely all the known properties of high-energy hadron scattering in a wide range of momentum transfers. The model has as a consequence anomalous terms of the scattering amplitude that are small in value and lead to spin effects that vary slowly with increasing energy.

We emphasize that a self-consistent picture of the scattering can be obtained only for a large interaction range of the anomalous quasipotential $\beta(s, r)$. This confirms the conclusion drawn in the model of a peripheral nature of the part of the interaction quasipotential that increases rapidly with the energy.

5. EFFECTS OF THE NEW MECHANISM OF "SPIN" DYNAMICS OF THE STRONG INTERACTIONS AT SUPERHIGH ENERGIES

The results obtained above indicate the possible existence of a new mechanism of "spin" dynamics of the strong interactions at high energies.

Indeed, in the presence of an anomalously increasing $\sqrt{s}\beta(s, t)$ contribution to the NN scattering amplitude (31) we arrive at a weak energy dependence of the spin effects in the Born term of the helicity amplitudes:

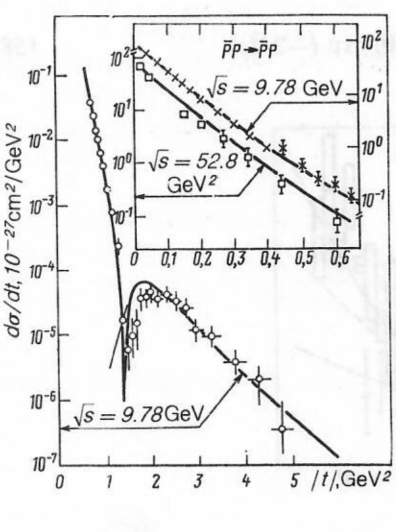

FIG. 6. Predictions of the model for the differential cross sections of $\bar{p}p$ scattering for small and large momentum transfers.

$$T^B_{++,++}(s, t) \sim d(s, t);$$
$$T^B_{++,+-}(s, t) \sim \sqrt{\frac{|t|}{s}}\sqrt{s}\beta(s, t);$$
$$\left.\frac{|T^B_{++,+-}(s, t)|}{|T^B_{++,++}(s, t)|}\right|_{t\ \mathrm{fixed}} \sim f(\ln s). \quad (57)$$

Such behavior of the spin-flip amplitudes is characteristic of the models of Refs. 17, 58, and 77–80. It also arises in the dynamical model of hadron interaction with allowance for the particle spins developed above.

Analysis of the expressions (51)–(53) permits the conclusion that the strong-interaction dynamics at high and asymptotically high energies may have different mechanisms.

Thus, at energies $\sqrt{s} \leqslant 100$ GeV, at which the terms of the form (52b) do not make a significant contribution to the eikonal phase, the helicity amplitudes of NN scattering have the standard structure. In this case, the amplitude without spin flip is determined by rescatterings on the quasipotential $d(s, r)$, and the amplitude with spin flip is determined by single scattering by the quasipotential $\beta(s, r)$ with subsequent rescatterings by $d(s, r)$.

With increasing energy, the terms with twofold spin flip in χ_0 begin to be dominant, and the helicity amplitudes are completely determined by the contributions of the quasipotential $\beta(s, r)$. As noted above, the contributions to the eikonal phase that increase with the energy have a spin nature, since they can arise only in the case of the interaction of particles with spin. Therefore, at asymptotically high energies the mechanism of the strong-interaction dynamics may have a spin nature.[54] The magnitude of the physical effects produced by the spin mechanism at high energies can be determined by means of the anomalous term $\beta(s, r)$ of the quasipotential found in the previous section.

We consider first the predictions of the model for polarization phenomena at $p_L \gtrsim 100$ GeV. As noted above, the spin effects that decrease with the energy become unimportant in this region, and the behavior of the spin-flip amplitude (38), (51) and, therefore, of the polarization is to a large degree determined by the anomalous term $\beta(s, r)$.

The predictions obtained in the model for the polarization of pp scattering based on the parameter values determined by analyzing the differential cross sections are given in Fig. 7(a) (for $p_L = 200$ and 300 GeV). Overall, they correctly reproduce the behavior of the experimental data.[8] We note that the polarization is positive at small $|t|$ and changes sign at $|t| \simeq 0.4$–0.5 GeV^2. A change in the sign of the polarization in this region of momentum transfers is observed experimentally already at $p_L = 45$ GeV and persists at high energies.[76]

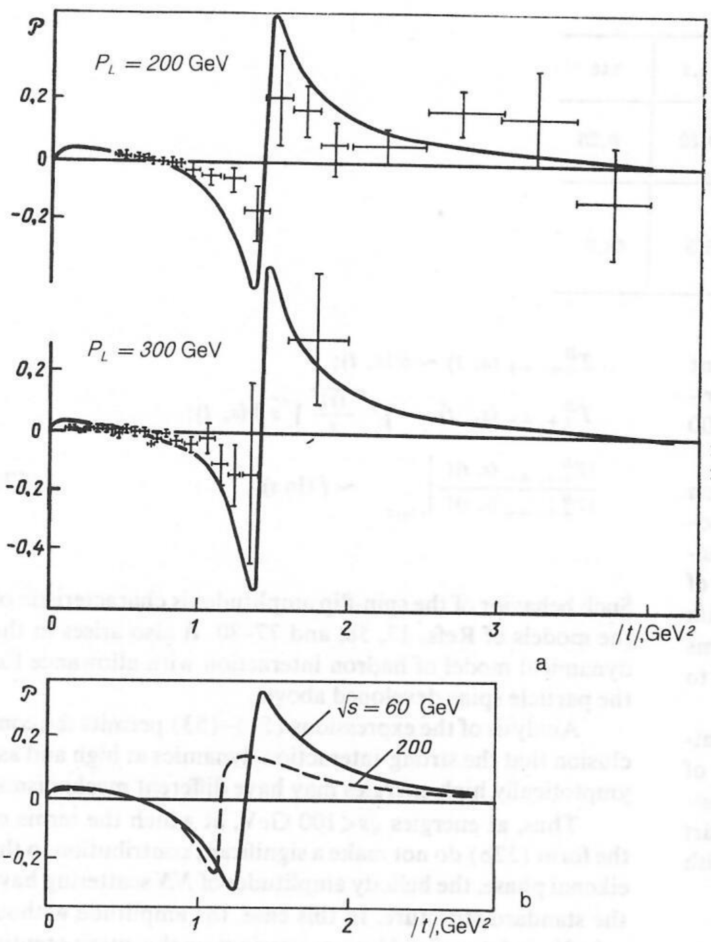

FIG. 7. Comparison of predictions of the dynamical model for the polarization with experiment at high energies (a) and predictions for the polarization at higher energies (b).

In the considered model, this effect is associated with the passage through zero of the leading asymptotic term of the spin-flip amplitude; this term is largely determined by the anomalously increasing part of the quasipotential.

In the region of low energies, the contribution of the $1/\sqrt{s}$ terms of the spin-flip amplitude leads to positive polarization at $|t| \leqslant 1$ GeV^2.

A similar picture of polarization phenomena at $p_L \gtrsim 100$ GeV is observed in a number of other models.[80-82]

The predictions of the model for the polarization at higher energies are given in Fig. 7(b). It should be noted that the weak energy dependence of the polarization at fixed $|t| = 0.18$ GeV^2 (Fig. 8) is associated with the anomalous behavior of the spin-flip amplitude.[52]

Thus, the dynamical model leads to a number of predictions about the behavior of the polarization at high energies, the verification of which requires experiments at energies $E_L \gtrsim 500$ GeV. The results of the model agree with the data in the experimentally studied region. This permits the conclusion that the quasipotential $\beta(s, r)$ in the model has been correctly calculated. For it determines the spin contribution to the eikonal phase χ_0 (52) of the amplitude without spin flip, and this can significantly change the behavior of the scattering cross sections at superhigh energies.

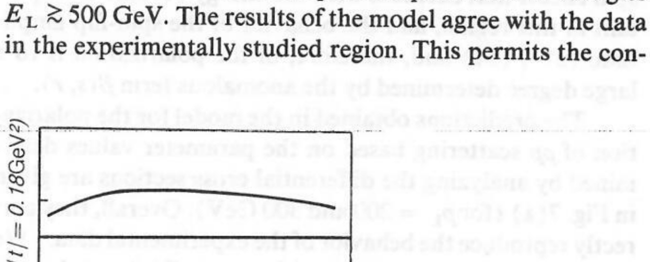

FIG. 8. Energy dependence of the polarization obtained in the model for $|t| = 0.18$ GeV^2.

We determine the effect of the mechanism of the spin dynamics at the energies of the CERN $S\bar{p}pS$. We consider first the growth of the total cross sections. A quasipotential $\beta(s, r)$ of the form (50) leads to an exponential asymptotic behavior of the anomalous term of the eikonal phase at large impact parameters:

$$\chi_{an}(s, \rho) \sim \chi_{an}(0) \exp(-2M\rho). \quad (58)$$

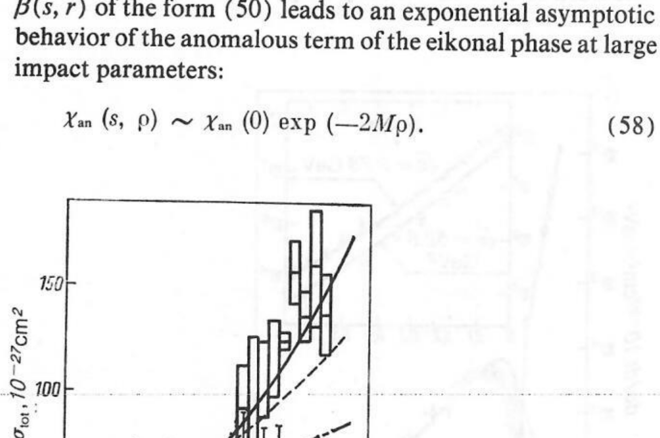

FIG. 9. Total cross sections of pp and $\bar{p}p$ scattering calculated in the model with the spin mechanism (continuous curve) and without it (chain curve) and Amaldi's calculations (broken curve).

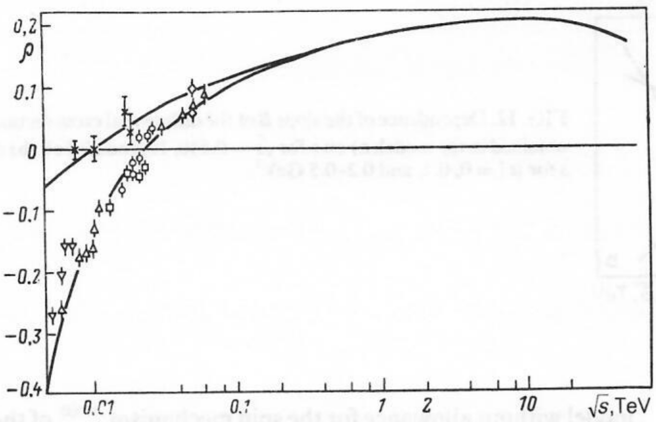

FIG. 10. The ratio $\rho = \mathrm{Re}\, T(s, 0)/\mathrm{Im}\, T(s, 0)$ for pp and $\bar{p}p$ scattering obtained in the model [the upper curve is $\rho(\bar{p}p)$].

In this case, we have for the total cross sections the asymptotic behavior

$$\sigma_{\mathrm{tot}}(s) \underset{s\to\infty}{\simeq} \frac{\pi}{M^2} \ln^2\left(\frac{\sqrt{s}}{2}\,|\chi_{\mathrm{an}}(0)|\right). \qquad (59)$$

Thus, a consequence of our model is a rapid growth of the cross sections at superhigh energies, which, however, does not contradict the Froissart bound and has a spin nature. The results obtained for the total cross sections up to energies $\sqrt{s} = 40$ TeV are given in Fig. 9. At energies $\sqrt{s} \leqslant 100$ GeV, the main contribution to the growth of the total cross sections is made by the standard mechanism associated with the growth of the effective interaction range. At higher energies, the spin mechanism becomes important, and its contribution to σ_{tot} at $\sqrt{s} = 540$ GeV is about $6\cdot 10^{-27}$ cm^2.[54]

The ratio

$$\rho = \frac{\mathrm{Re}\, T(s, 0)}{\mathrm{Im}\, T(s, 0)}$$

obtained in the model for pp and $\bar{p}p$ scattering is in good agreement with the experiments at $\sqrt{s} \gtrsim 7$ GeV (Fig. 10). The presence of the anomalous terms hardly changes the energy dependence of this quantity.

We note that the model predicts the following values for the slope of the diffraction peak at different momentum transfers and $\sqrt{s} = 540$ GeV:

$B(0) = 16.9\,\mathrm{GeV}^{-2}$; $B(0.1) = 15.2\,\mathrm{GeV}^{-2}$;

$B(0.3) = 13.6\,\mathrm{GeV}^{-2}$.

For comparison, we give the data of the UA-4 experiment[84]:

$B(0.1) = (15.2 \pm 0.2)\ \mathrm{GeV}^{-2}$;

$B(0.3) = (13.4 \pm 0.13)\ \mathrm{GeV}^{-2}$.

At the same time, the model describes well the experimental material at small momentum transfers in a fairly wide energy region from $\sqrt{s} = 4.5$ GeV to $\sqrt{s} = 540$ GeV (see Fig. 11 and Table II).

We note the nonstandard behavior of the slope of the diffraction peak when s and t vary (Fig. 12). It can be seen that for t near zero B increases with the energy approximately as $\ln s$ at $\sqrt{s} < 10$ TeV. At higher energies, $B(0)$ begins to increase as $\ln^2 s$. In the region of large momentum transfers $t \gtrsim 0.1$ GeV^2 the slope of the diffraction peak has an interesting energy dependence for $\sqrt{s} > 1$ TeV. Thus, for $-t \sim 0.1$ GeV^2 the value of B increases as $\ln^2 s$, while for $-t \sim 0.2$–0.5 GeV^2 the value of B begins to decrease. Such behavior of the slope is due to the manifestation in the region of small momentum transfers of the anomalous spin term of the eikonal phase, which has a peripheral nature.

As already noted, the anomalous term is manifested in the region of the diffraction dip at much lower energies. Indeed, a consequence of the rapid growth with the energy of the peripheral term of the eikonal phase (51) is a rapid change in the differential cross sections in the region of the diffraction dip and the maximum when $\sqrt{s} \geqslant 100$ GeV. Allowance for the spin contribution in the amplitude without spin flip increases by an order of magnitude the differential cross sections in the region $1 \leqslant |t| \leqslant 1.5$ GeV^2 at $\sqrt{s} = 540$ GeV. As a result, the model predicts that the diffraction structure at the energies of the $S\bar{p}pS$ collider at CERN completely disappears[54,83] and a "shoulder" appears in the differential cross sections (Fig. 13). For comparison, we give in the figure the data of the UA-4 experiment[84] made with the $S\bar{p}pS$ collider; they are in complete agreement with the predictions of the model. The predictions of the dynamical

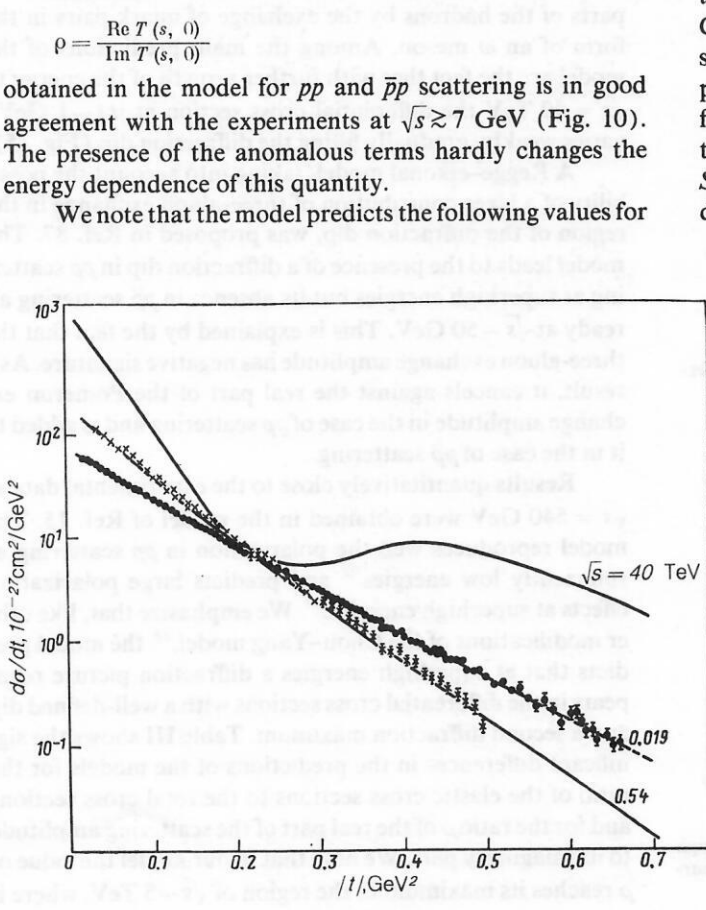

FIG. 11. Differential cross sections in the region of small momentum transfers for pp and $\bar{p}p$ scattering (continuous curve) obtained in the model.

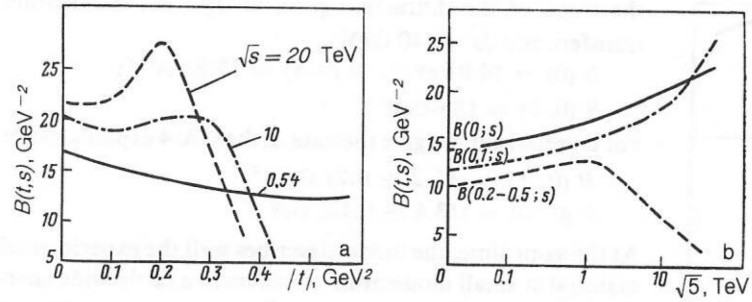

FIG. 12. Dependence of the slope B of the differential cross sections obtained in the model: a) on t for $\sqrt{s} = 0.540$, 10, and 20 TeV; b) on s for $|t| = 0$, 0.1, and 0.2–0.5 GeV^2.

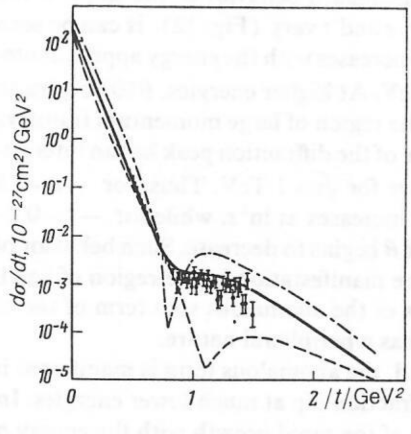

FIG. 13. Predictions of models for elastic $\bar{p}p$ scattering at $\sqrt{s} = 540$ GeV. The continuous curve is for our model with allowance for the spin contribution,[83] the broken curve is for our model without allowance for the spin contribution (geometrical scaling) and the chain curve is for the Chou–Yang model (factorized eikonal).[85]

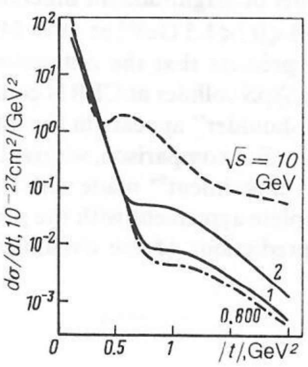

FIG. 14. Manifestations of the mechanism of spin dynamics in the differential cross sections at superhigh energies.

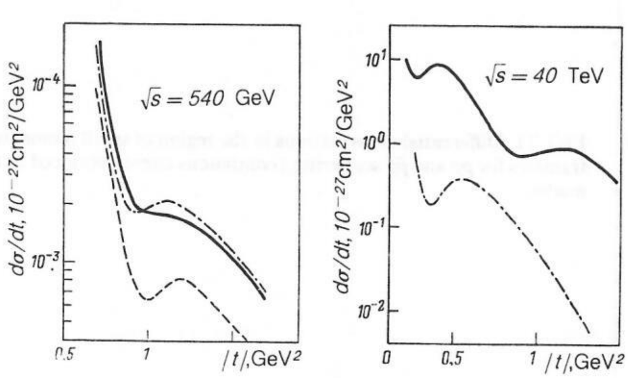

FIG. 15. Predictions of different models for $\bar{p}p$ scattering at superhigh energies. The continuous curve is for our model, the chain curve for Bourrely and Soffer, and the broken curve for Islam.

model without allowance for the spin mechanism,[57,60] of the standard model of factorized eikonal type,[85] and of a number of other models (see, for example, Ref. 83) differ from the experiment at $|t| \sim 1.0–1.2$ GeV^2 by one or two orders of magnitude. The influence of the mechanism of the spin dynamics at different energies can be traced in Fig. 14.

To what extent is this influence uniquely determined, and in what case can one regard the existence of the new spin mechanism of hadron interaction as proved? For this, we compare our predictions with the predictions of some of the other most viable models.

At the present time various models, as a rule with modification of some ideas, have yielded theoretical values for the differential cross sections of $p\bar{p}$ scattering at $\sqrt{s} = 540$ GeV. For example, in the model of Ref. 86 a picture qualitatively close to the experiment is obtained with a weakly expressed diffraction dip and maximum. The model is based on the idea that two mechanisms operate in the region of small scattering angles: a "soft" diffraction mechanism and a "hard" mechanism, corresponding to the interaction of the central parts of the hadrons by the exchange of quark pairs in the form of an ω meson. Among the main predictions of the model are the fact that with further growth of the energy to $\sqrt{s} = 40$ TeV the differential cross section at $|t| \sim 1$ GeV^2 varies weakly, gradually filling the diffraction dip (Fig. 15).

A Regge–eikonal model, taking into account the possibility of a large contribution of three-gluon exchange in the region of the diffraction dip, was proposed in Ref. 87. The model leads to the presence of a diffraction dip in pp scattering at superhigh energies but its absence in $p\bar{p}$ scattering already at $\sqrt{s} \sim 50$ GeV. This is explained by the fact that the three-gluon exchange amplitude has negative signature. As a result, it cancels against the real part of the Pomeron exchange amplitude in the case of pp scattering and is added to it in the case of $p\bar{p}$ scattering.

Results quantitatively close to the experimental data at $\sqrt{s} = 540$ GeV were obtained in the model of Ref. 15. The model reproduces well the polarization in pp scattering at sufficiently low energies[70] and predicts large polarization effects at superhigh energies.[71] We emphasize that, like other modifications of the Chou–Yang model,[88] the model predicts that at superhigh energies a diffraction picture reappears in the differential cross sections with a well-defined dip and a second diffraction maximum. Table III shows the significant differences in the predictions of the models for the ratio of the elastic cross sections to the total cross sections and for the ratio ρ of the real part of the scattering amplitude to its imaginary part. We note that in our model the value of ρ reaches its maximum in the region of $\sqrt{s} \sim 5$ TeV, where it

TABLE III. Comparison of predictions for $\bar{p}p$ scattering in our model and in Ref. 15.

$\sqrt{s}$, TeV		0.540	1	2	10	20	40
σ_{tot}, 10^{-27} cm^2	I	62	67.9	76.1	98.4	109.4	121.2
	II	60.67	69	81	123	146	175
σ_{el}, 10^{-27} cm^2	I		14.9	17.9	26.8	31.4	36.4
	II	12.43	15.4	20.7	42.6	54.1	66.1
$\frac{\sigma_{el}}{\sigma_{tot}}$	I	0.21	0.22	0.235	0.27	0.287	0.3
	II	0.205	0.22	0.256	0.35	0.37	0.38
ρ	I	0.13	0.128	0.128	0.122	0.118	0.114
	II	0.151	0.175	0.197	0.195	0.195	0.157
$\frac{d\sigma}{dt}$, 10^{-24} cm^2/GeV^2	I	1.92	3.47	11.2	38.6	49.9	53.7
$\|t\|=1,0$ GeV^2	II	1.7	7.6	38.8	350.0	440.0	750.0
$B(s, t)$, GeV^{-2}	I	15.6			18.6	19.6	20.7
$0\leqslant\|t\|\leqslant 0.15$GeV^2	II	15.56	16.14	16.54	19.1	21.8	26

Note. I is for Ref. 15; II is for our predictions.

differs appreciably from the predictions of the model of Ref. 15.

Thus, systematic allowance for the spin structure of interacting hadrons on the basis of the unified eikonal representation leads to the possibility of a new mechanism of spin interaction dynamics, and at superhigh energies this becomes dominant in the region of small scattering angles. On the basis of this mechanism and in the framework of a model that considers the internal nucleon structure, we have explained from a unified point of view all the physical phenomena of hadron elastic scattering in a wide range of momentum transfers for variation of the scattering energy in the interval $4.5\leqslant\sqrt{s}\leqslant 540$ GeV. This permits the conclusion that in the $S\bar{p}pS$ collider at CERN the first manifestations of the mechanism of the strong-interaction spin dynamics may have been detected.

Such a mechanism leads to a number of clearly expressed effects for the physical quantities at superhigh energies, and these could be tested in the accelerators of the next generation (UNK, $S\bar{p}pS$, Fermilab).

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