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Предасимптотические эффекты в нуклон-нуклонном рассеянии на большие углы

Предасимптотическое поведение упругих нуклон-нуклонных амплитуд изучено в пределе высокоэнергетического рассеяния на большие углы. Полученные формулы использованы для совместного описания и интерпретации экспериментальных данных по упругому pp- и pn -рассеянию.

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Preasymptotic Effects in Nucleon-Nucleon Large-Angle Scattering

The preasymptotic behaviour of elastic nucleon-nucleon amplitudes is studied in the limit of high-energy large-angle quasipotential scattering. Formulae obtained are used for a common description of experimental data on pp and pn elastic scattering.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

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1. INTRODUCTION

The Logunov-Tavkhelidze quasipotential approach proved to be one of the most efficient methods in studying the asymptotic and preasymptotic behaviour of differential cross sections of high-energy large-angle hadron scattering for the following three reasons: First, the quasipotential equations are essentially two-particle ones and thus are most suitable for the description of hadron-hadron elastic scattering. Second, the quasipotential equations themselves enable us to account self-consistently for preasymptotic effects in transition energy regions where the dominant interaction mechanism is changed. Third, when considered in the framework of quantum-field-theory models, the quasipotential exhibits naturally the dependence on the relative coordinate of colliding particles and on their energy, as well. This indicates that both the form of the quasipotential and magnitude of its parameters can change with growing energy.

The structure of phenomenological quasipotentials which account correctly for the experimental data on high-energy elastic hadron scattering was discussed in refs. [1,2]. Their result suggests that all the requirements are met by analytic quasipotentials given by the integral representation:

$$\hat{g}(s, \vec{\Delta}) = \int_0^\infty dx\, \hat{\rho}(s, x)\, e^{-x\vec{\Delta}^2}; \quad t, u = -\vec{\Delta}^2$$

(s, t ,u are the usual Mandelstam variables of two-particle reaction).

The small-angle scattering is dominated by the effects connected with the global structure of hadron as a soft object of finite size that is rendered correctly by the Gaussian quasipotentials. The typical density function here is the δ-function:

$$\hat{\rho}(s, x) = \hat{b}\,\delta(x - a), \qquad (1.1)$$

where parameters a and $\hat{b}$ may be slowly varying functions of s.

As to the large-angle scattering, the behaviour of differential cross sections is controlled by details of the inner structure of hadron at short distances that is observed from the behaviour of $\hat{\rho}(s, x)$ near $x = 0$. Thus, the automodel asymptotics of differential cross sections of large-angle exclusive scattering

$$\frac{d\sigma}{dt} \sim \frac{1}{s^{2M}}\, f(t/s); \quad s \to \infty; \quad t/s = \text{const}, \qquad (1.2)$$

originally understood within the assumption about the presence of point-like constituents within hadrons [3,4], can be obtained in the framework of quasipotential approach provided the weak limit for the function $\hat{\rho}(s, x)$

$$\lim_{s\to\infty} s^{M-1}\hat{\rho}(s, x = \eta/s) = \hat{\Psi}(\eta); \quad 0 < \eta < \infty; \quad M > 0 \qquad (1.3)$$

does exist.

In that way we see, that in high-energy limit the dynamics of small and large-angle scattering is governed by different mechanisms, and accordingly two asymptotic quasipotentials can be constructed. But as far as finite-energy large-angle scattering is considered, certain interference of two mechanisms takes place. That is, the "soft" component of interaction connected with large distances and given by the quasipotential with density function (1.1) generates corrections to the asymptotic amplitude, determined by short-distance interaction that is predominant in this region of momentum transfers. These corrections decrease with growing energy and lead, in particular, to the deviation from strict automodelity (1.2):

$$\frac{d\sigma}{dt} \sim \frac{1}{s^{2M}}\left(f(t/s) + 1/s\, f_1(t/s)\right); \quad s \to \infty; \quad t/s = \text{const}. \qquad (1.4)$$

They also break the γ_5-invariance of the amplitude even for γ_5-invariant interaction that manifests itself in nonzero polarization:

$$P \sim 1/s\; \mathcal{P}(t/s). \qquad (1.5)$$

In this paper we shall apply the method previously developed in refs. [5,6] for the quantitative investigation of preasymptotic effects (1.4), (1.5) in nucleon-nucleon large-angle scattering. In so doing, the assumption about the charge independence of strong interactions will enable us to describe both reactions of pp and pn elastic scattering with one set of parameters. Appropriate helicity amplitudes with correction terms of two leading orders in $1/p$, where p is the c.m.s. momentum of colliding particles, are presented in Sec. 2. In Sec. 3 the formulae obtained are used for a detailed analysis of available experimental data.

2. THE DESCRIPTION OF LARGE-ANGLE NUCLEON-NUCLEON SCATTERING AT MODERATE ENERGIES

The quasipotential equations for a system of two particles with spin 1/2 have been derived in a number of papers (see, e.g., refs. [7,8]), and we use the equation from ref. [8]. In the momentum space it is of the form:

$$\hat{T}(s, \vec{p}, \vec{k}) = \hat{g}(s, \vec{p}, \vec{k}) + \int d\vec{q}\, \hat{g}(s, \vec{p}, \vec{q})\, \frac{\hat{A}(s, \vec{q})}{E^2(\vec{q}) - E^2 - i0}\, \hat{T}(s, \vec{q}, \vec{k}). \qquad (2.1)$$

Here $\hat{T}(s, \vec{p}, \vec{k})$ is the off-mass-shell matrix scattering amplitude, $\vec{p}$ and $\vec{k}$ are the c.m.s. momenta of particles before and after the collision, $E(\vec{q}) = 2\sqrt{\vec{q}^2 + m^2}$; $E = \sqrt{s} = E(\vec{p}) = E(\vec{k})$ and

$$\hat{A}(s,\vec{q}) = [\frac{E^2 - \frac{1}{2}E^2(\vec{q})}{E} + \hat{H}^{(1)}(\vec{q}) + \hat{H}^{(2)}(-\vec{q}) + \frac{2}{E}\hat{H}^{(1)}(\vec{q})\hat{H}^{(2)}(-\vec{q})],$$

$\hat{H}^{(1,2)}(\vec{q})$ are the energy operators of the first and second particles, resp.:

$$\hat{H}^{(1,2)}(\vec{q}) = m\gamma_0^{(1,2)} + \gamma_0^{(1,2)}\vec{\gamma}^{(1,2)}\vec{q}.$$

If the charge independence of strong interactions is assumed in the framework of isospin symmetry, the differential cross sections both of pp and pn scattering can be expressed via one matrix amplitude $\hat{T}(s,\vec{p},\vec{k})$ as follows:

$$\frac{d\sigma^{NN}}{dt} \sim \sum_{spin} M^{NN}(\vec{p},\vec{k})M^{+NN}(\vec{p},\vec{k})\Big|_{\substack{s=4(\vec{p}^2+m^2)=4(\vec{k}^2+m^2)\\ t=-(\vec{p}-\vec{k})^2}},$$

where

$$M^{pp}(\vec{p},\vec{k}) = \langle\bar{\Psi}_1^{\sigma_1}(\vec{p})\bar{\Psi}_2^{\sigma_2}(-\vec{p})|\hat{T}(s,\vec{p},\vec{k})|\Psi_1^{\sigma_1'}(\vec{k})\Psi_2^{\sigma_2'}(-\vec{k})\rangle - \qquad (2.2)$$

$$- \langle\bar{\Psi}_1^{\sigma_1}(\vec{p})\bar{\Psi}_2^{\sigma_2}(-\vec{p})|\hat{T}(s,\vec{p},-\vec{k})|\Psi_1^{\sigma_2'}(-\vec{k})\Psi_2^{\sigma_1'}(\vec{k})\rangle$$

and

$$M^{pn}(\vec{p},\vec{k}) = \langle\Psi_1^{\sigma_1}(\vec{p})\Psi_2^{\sigma_2}(-\vec{p})|\hat{T}(s,\vec{p},\vec{k})|\Psi_1^{\sigma_1'}(\vec{k})\Psi_2^{\sigma_2'}(-\vec{k})\rangle. \qquad (2.3)$$

To solve the equation (2.1), we are to choose the explicit form of the quasipotential $\hat{g}(s,\vec{p},\vec{k})$; it is natural to represent it as a sum of "soft" and "hard" components corresponding to the asymptotic quasipotentials for the interaction at large and short distances:

$$\hat{g}(s,\vec{p},\vec{k}) = \hat{g}_s(s,\vec{p}-\vec{k}) + \hat{g}_h(s,\vec{p},\vec{k}).$$

In high-energy scattering with small momentum transfers the spin-flip amplitudes are small as compared to the spin-non-flip ones. This requirement is met by the following matrix structure of the "soft" quasipotential [9]:

$$\hat{g}_s(s,\vec{p}-\vec{k}) = \gamma_0^{(1)}\gamma_0^{(2)}g_s(s,\vec{p}-\vec{k}).$$

As to the "hard" quasipotential, we shall choose its matrix structure under the requirement of γ_5-invariance of interaction at large energies and momentum transfers [10]:

$$\hat{g}_h(s,\vec{p},\vec{k}) = \gamma_\mu^{(1)}\gamma^{\mu(2)}g_{1h}(s,\vec{p},\vec{k}) +$$
$$+ \gamma_\mu^{(1)}\gamma_5^{(1)}\gamma^{\mu(2)}\gamma_5^{(2)}g_{2h}(s,\vec{p},\vec{k}).$$

The data available on pp small-angle scattering are well reproduced by a simple purely imaginary Gaussian potential with the density function:

$$\rho_s(s,x) = 2ig\delta(x-a)$$

and the parameters being [1]:

$$i\chi(0) = -0.5;\quad a = 2.5\,(\mathrm{GeV/c})^{-2},$$

where $i\chi(0) = -2\pi^2 g/a$ is the eikonal phase at the zero impact parameter.

The "hard" component of the quasipotential is to account also for the exchange forces in the nucleon-nucleon system, so we take:

$$g_{1h}(s,\vec{p},\vec{k}) = \int_0^\infty dx\,\rho_{1h}(s,x)e^{-x(\vec{p}+\vec{k})^2};$$

$$g_{2h}(s,\vec{p},\vec{k}) = \int_0^\infty dx\,\rho_{2h}(s,x)e^{-x(\vec{p}-\vec{k})^2}$$

and the density functions $\rho_{1h}(s,x)$ and $\rho_{2h}(s,x)$ are approximated as follows:

$$\rho_{1h}(s,x) = \frac{Ce^{-2i\chi(0)}}{\Gamma(\nu+1)s^{M-\nu-1}}x^\nu e^{-dx};$$

$$\rho_{2h}(s,x) = \frac{De^{-2i\chi(0)}}{\Gamma(\gamma+1)s^{M-\gamma-1}}x^\gamma e^{-dx}.$$

In ref. [6] we have obtained the helicity amplitudes of elastic pp scattering with the corrections of two leading orders in $1/p$ for arbitrary local vector-vector and axial-axial "hard" quasipotentials. Then, generalizing that result to the case of the exchange vector-vector potential, we get for the helicity amplitudes of elastic pp and pn scattering, the relation between them being fixed in (2.2), (2.3):

$$T^{pp}_{++,++}(s,t) = (1+z)\{C(s,u)F_1(\nu,-z) + D(s,t)F_1(\gamma,z) -$$

$$-0.44(p_0/p)^2(C(s,t) + D(s,u)\};$$

$$T^{pp}_{++,--}(s,t) = -(1-z)\{C(s,t)F_1(\nu,z) + D(s,u)F_1(\gamma,-z) -$$

$$-0.44(p_0/p)^2(C(s,u) + D(s,t))\};$$

$$T^{pp}_{+-,+-}(s,t) = 2\{C(s,t)[F_1(\nu,z) + (0.1(\nu+1) - 0.22(1+z))(p_0/p)^2] +$$

$$+C(s,u)[F_1(\nu,-z) + (0.1(\nu+1) - 0.22(1-z))(p_0/p)^2] -$$

$$-D(s,t)[F_1(\gamma,z) + (0.1(\gamma+1) + 0.22(1-z))(p_0/p)^2] -$$

$$-D(s,u)[F_1(\gamma,-z) + (0.1(\gamma+1) + 0.22(1+z))(p_0/p)^2]\};$$

$$T^{pp}_{++,+-}(s,t) = -\sqrt{1-z^2}[(C(s,u) - C(s,t))F_2(\nu) +$$

$$+(D(s,t) - D(s,u))F_2(\gamma)]; \qquad (2.4)$$

$$T^{pn}_{++,++}(s,t) = (1+z)[C(s,u)F_1(\nu,-z) + D(s,t)F_1(\gamma,z)];$$

$$T^{pn}_{+-,+-}(s,t) = 2\{C(s,u)[F_1(\nu,-z) + (0.1(\nu+1) - 0.22(1-z))(p_0/p)^2] -$$

$$-D(s,t)[F_1(\gamma,z) + (0.1(\gamma+1) + 0.22(1-z))(p_0/p)^2]\};$$

$$T^{pn}_{++,+-}(s,t) = -\sqrt{1-z^2}\{C(s,u)F_2(\nu) + D(s,t)F_2(\gamma)\};$$

where $p_0 = 1$ (GeV/c); $z = \cos\theta$ is the cosine of the c.m.s. scattering angle and

$$C(s,t) = \frac{C}{s^M}\left(\frac{s}{|t|+d}\right)^{\nu+1}; \quad C(s,u) = \frac{C}{s^M}\left(\frac{s}{|u|+d}\right)^{\nu+1};$$

$$D(s,t) = \frac{D}{s^M}\left(\frac{s}{|t|+d}\right)^{\gamma+1}; \quad D(s,u) = \frac{D}{s^M}\left(\frac{s}{|u|+d}\right)^{\gamma+1};$$

$$F_1(\nu,z) = 1 - i\cdot 0.365\nu(p_0/p) + (0.0363\nu^2 + 0.0182\nu -$$

$$-0.9617 - 0.2\frac{(\nu+1)^2}{1-z+\frac{d}{2p^2}})(p_0/p)^2;$$

$$F_2(\nu) = 0.469(p_0/p) - i\cdot 0.167(\nu+1/2)(p_0/p)^2.$$

The exponent M determines the rate of the power fall-off of differential cross sections with growing energy; the quark counting rules [3] predict the value $M = 5$. We remark, that the formalism developed can be applied for arbitrary asymptotic quasipotentials. The Gaussian and power potentials exploited above possess the advantage of being characterized by only few free parameters, that is, three real (C, D, d) and two integer (ν,γ) ones. Nevertheless, they will enable us to fit the data on two reactions in a wide range of energies and scattering angles.

3. COMPARISON WITH EXPERIMENT

Let us proceed now to the analysis of experimental data on the differential cross sections of pp and pn large-angle scattering in the energy region of $p_L \gtrsim 7$ (GeV/c) [11].

The differential cross sections are expressed in terms of helicity amplitudes as follows:

$$\frac{d\sigma^{NN}}{dt} = |T^{NN}_{++,++}|^2 + |T^{NN}_{++,--}|^2 + |T^{NN}_{+-,+-}|^2 + 4|T^{NN}_{++,+-}|^2 .$$

The results of the fit are presented in the table and in Figs. 1,2 and show good agreement of theoretical curves with experiment. It should be noticed that the consideration of corrections enables us to achieve better description of the data as compared to the fits neglecting corrections [12], the number of fitted parameters being the same.

Table

	C	ν	D	γ	d	χ^2	$\bar{\chi}^2$	$\chi^2/\bar{\chi}^2$
pp & pn	-8033 ±63	0	610 ±16	3	2.379 ±0.034	309	159	1.94
pp	-8166 ±66	0	634 ±16	3	2.415 ±0.033	137	79	1.75

We also remark that when describing pp scattering only, the formulae (2.4) require practically unchanged values of the parameters (see the table). This is a strong evidence in favour of the chosen relation between the amplitide of two reactions which is fixed by (2.2), (2.3).

The same conclusion is supported by comparing the ratio

$$R(p_L) = \frac{d\sigma^{pn}}{dt}(90°) / \frac{d\sigma^{pp}}{dt}(90°)$$

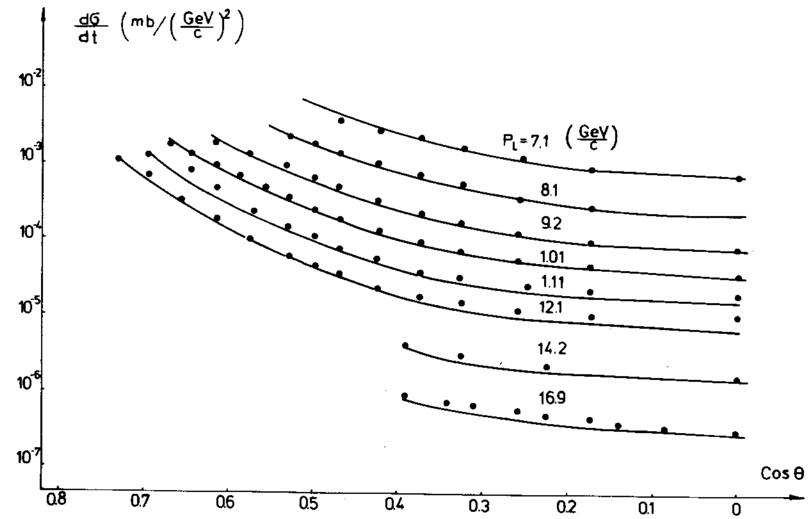

Fig. 1. $\frac{d\sigma}{dt}$ for pp scattering.

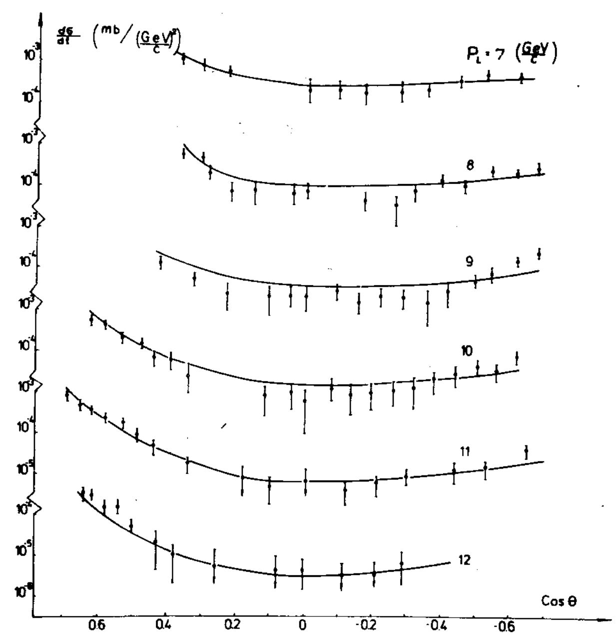

Fig. 2. $\frac{d\sigma}{dt}$ for pn scattering.

predicted by our model with that derived directly from the data. The appropriate curves are depicted in Fig. 3, where the dashed lines are predictions of the statistical model [13], by Fishbane and Quigg [14] and by Wu and Yang [15].

As was mentioned above, the consideration of corrections results in the deviations from the exact automodelity (1.2) that can be quantitatively described by the effective power N_{eff} (s, z):

$$\frac{d\sigma^{NN}}{dt}\Big|_{z = const} \sim s^{-N_{eff}^{NN}(s,\, z)} ,$$

where $N_{eff}(s,z) \to 10$ when $s \to \infty$. The smooth approximations of differential cross sections obtained allow us to predict the effective powers $N_{eff}(s,z)$. Their values as functions of laboratory momentum for different scattering angles are shown in Figs. 4,5. The increase of $N_{eff}(p_L)$ at small p_L

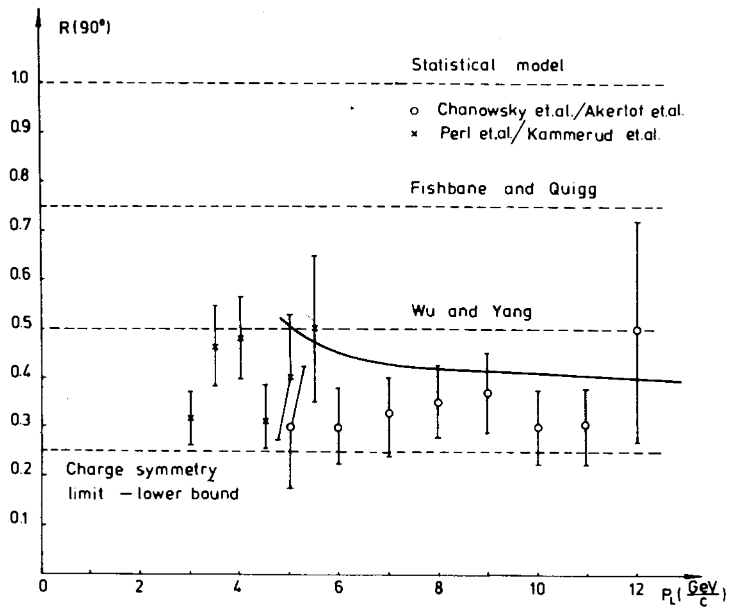

Fig. 3. Energy dependence of the ratio of the pn to pp scattering cross section at $\theta = 90°$.

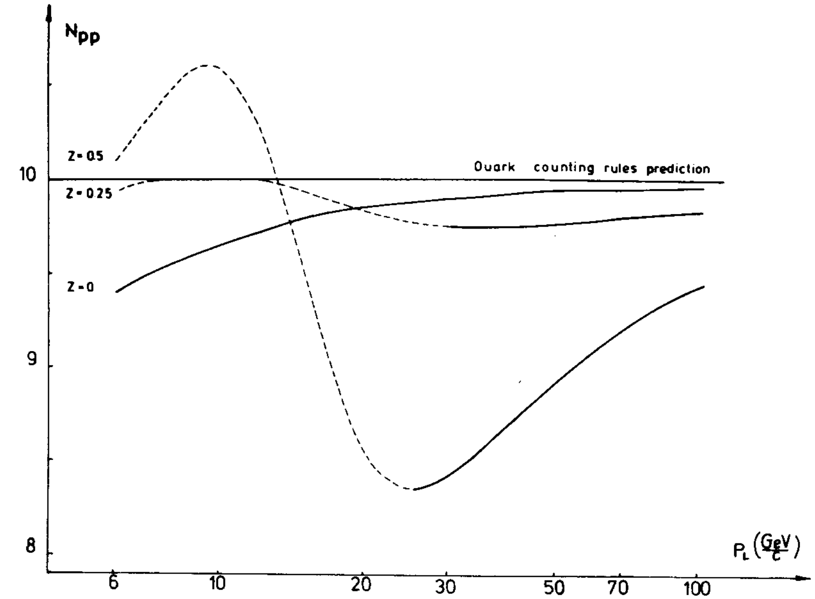

Fig. 4. $N_{eff}(p_L)$ for pp scattering.

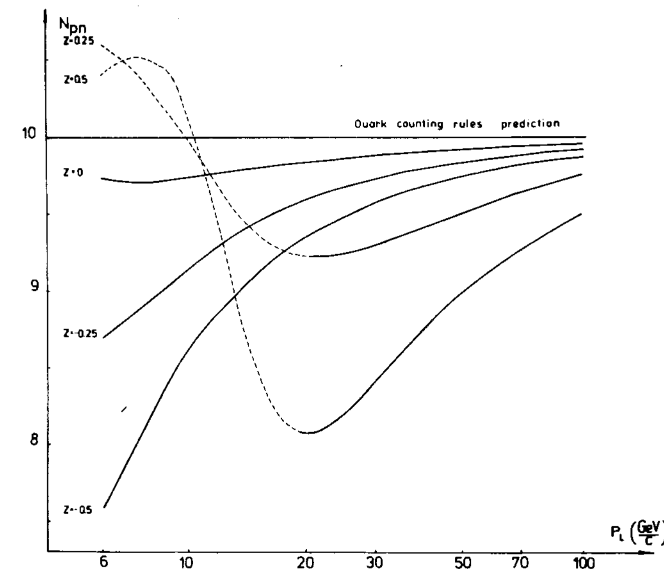

Fig. 5. $N_{eff}(p_L)$ for pn scattering.

(these parts of curves are dashed lines) is caused by the growth of corrections which are compatible in this energy region with the leading asymptotic term. The angular dependences of N_{eff}^{pn} for different energies are plotted in Fig. 6. They indicate that the approximation developed is self-consistent throughout the whole angular range only for sufficiently high energies ($p_L \geq 20$ (GeV/c)). Nevertheless, the formulae obtained reproduce correctly the experimental data even for $p_L \sim 7$ (GeV/c).

As to the numerical values of effective powers, the deviations from the quark counting rules are larger for pn scattering, whereas for pp scattering the corrections to different helicity amplitudes considerably compensate each other. For the interval 6 (GeV/c) $< p_L <$ 12 (GeV/c)

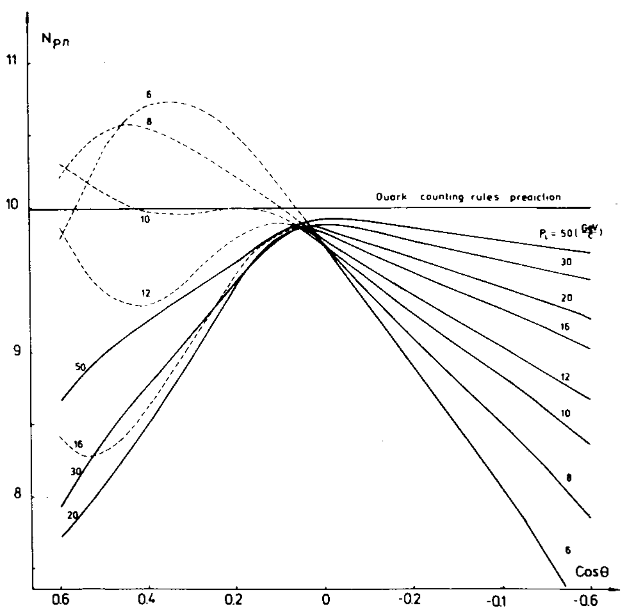

Fig. 6. $N_{eff}(z)$ for **pn** scattering.

our model predicts the following average values of N_{eff}^{pn}:

$$N_{eff}^{pn}(90^\circ) = 9.74; \; N_{eff}^{pn}(120^\circ) = 8.32$$

that is in good agreement with the values derived directly from the data [11]. It is worth mentioning that the deviation of $N_{eff}^{pp}(90^\circ)$ from $N_{as} = 10$ is practically negligible for all energies $p_L \geq 6$ (GeV/c). This prediction is supported by recent experimental results [16].

The power behaviour is observed also for the inclusive production of particles with high $p_\perp$, where the dependences of effective powers on energy are analogous to that discussed above. As it was shown in ref. [17], these energy-dependent effective powers appear only if the deviations from the Bjorken scaling in deep inelastic lepton-hadron processes are taken into account.

The nonzero mass of interacting particles breaks the γ_5-invariance of the amplitude regardless of the γ_5-invariance of interaction, and this results in nonzero polarization decreasing as s^{-1} with growing energy. The polarizations predicted by the model discussed for the processes of pp and pn scattering are plotted in Figs. 7,8. In sign and order of magnitude, they are compatible with the experimentally measured polarizations.

Thus, the above analysis shows that in the energy regions where the experimental data on exclusive nucleon-nucleon large-angle scattering are available, taking into account corrections resulting from the interaction at large distances considerably improves the description of the data. Analogous results can be obtained if the parameters of the "hard" quasipotential are assumed to be logarithmic functions of energy [12], that brings, of course, an extra arbitrariness. But our method enables us to avoid this arbitrariness in a constructive way, as the correction parameters are determined from the data corresponding to the region of momentum transfers where the hadron dynamics is of a different nature.

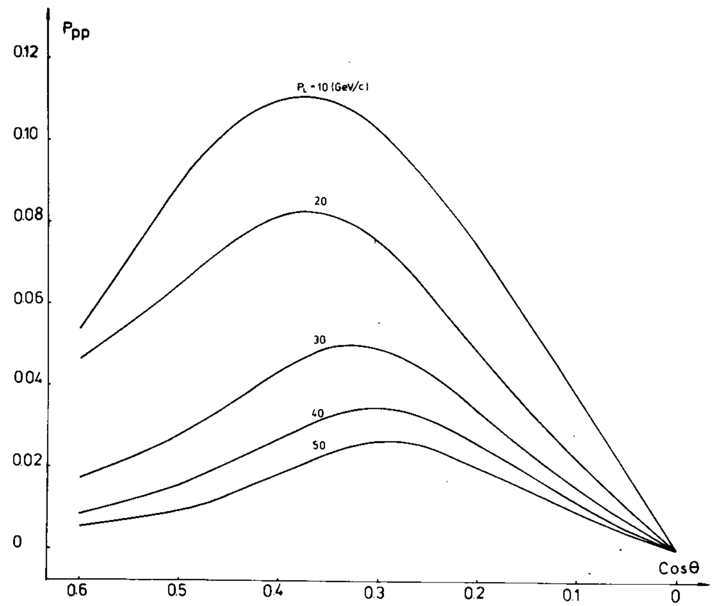

Fig. 7. Predictions for pp polarization.

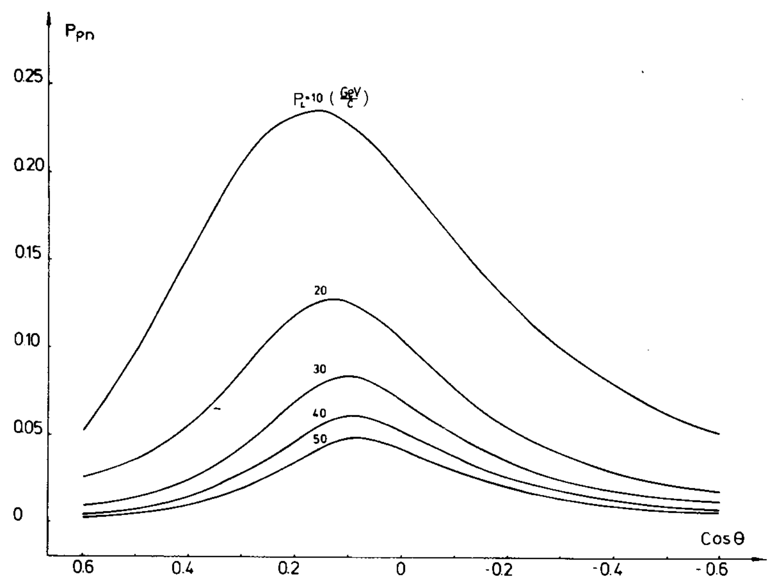

Fig. 8. Predictions for pn polarization.

Summarizing the above discussion, we stress that the considered preasymptotic effects in nucleon-nucleon large-angle scattering are characteristic of the domain of energies and momentum transfers where the hadron dynamics is dominated by the interaction of point-like constituents of hadrons at short distances, but the effects of the global hadron structure, i.e., of large distances are still pronounced. The magnitude of these effects makes it certainly necessary to take them into account in describing and interpreting the data.

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