Nucleon-nucleon scattering in the 0–6 GeV range and the relativistic optical model based on deep attractive forbidden state potentials

V. G. Neudatchin and N. P. Yudin
Institute of Nuclear Physics, Moscow State University, Moscow, U.S.S.R.

Yu. L. Dorodnykh
Institute for Nuclear Research, U.S.S.R. Academy of Sciences, Moscow, U.S.S.R.

I. T. Obukhovsky
Institute for Nuclear Physics, Moscow State University, Moscow, U.S.S.R.

(Received 9 April 1990)

The optical relativistic model based on deep attractive forbidden state quasipotentials describes well the angular and energy dependence of differential *np*- and *pp*-scattering cross sections and polarizations, including the transition from the *U*-shaped form of the angular distributions of *np* scattering to forward scattering at the energy $E_{\text{lab}} \geq 1$ GeV, when the scattering *P*-phase shifts, equal to π at low energies, pass through $\pi/2$ (the triplet and singlet *S*-phase shifts start from the values 2π and π, respectively). The higher scattering phase shifts ($L \geq 2$) are small everywhere. The potentials are determined from the scattering phase shifts in the low-energy region $E_{\text{lab}} < 1$ GeV. They have been chosen in the simple Gaussian form and for different partial waves they differ in depth and width (from $V_0 = 0.73$ GeV, $a = 0.85$ fm to $V_0 = 2.40$ GeV, $a = 0.45$ fm). A few negative phase shifts $^3D_1$, $^3F_3$, $^3G_3$, which reflect the peripheral repulsion due to the spin-orbital and tensor interactions, are calculated by means of the one-boson-exchange potential periphery matched to the central attraction. The imaginary part W of the optical potential $V + iW$ is determined by the value $\sigma_{\text{tot}}/\sigma_{\text{el}}$ and grows rapidly with increase in the energy E_{lab}, so that the phase-shift values lose sensitivity to the real part V of optical potential if E_{lab} exceeds 5 GeV. Finally, the quantum chromodynamics (QCD) effects are discussed which may underlie the potentials considered.

I. INTRODUCTION

There are many indications of the manifestation of quark degrees of freedom in atomic nuclei, starting from the two-nucleon system, but all of them, unfortunately, are very indirect, which brings about the feeling of dissatisfaction. For example, in the description of deuteron magnetic form factors, the quark many-body theory[1,2] and the traditional hadron theory, with such elements as the Bonn potential and meson-exchange currents,[3] may compete with each other if different form factors, soft or hard, are used at the meson-nucleon vertices. Next, the data on cumulative effects may be described on the basis of the hypotheses about both cold[4] and hot[5] fluctuations. Especially demonstrative is the problem of the nucleon-nucleon interaction in the short range $r \leq 1$ fm. The absence of a common standpoint here is reflected well by the fact that, now, six or seven different approaches are being developed, each of which, apart from a certain empirical success, has appreciable physical disadvantages. We touch upon two of them. Specifically, the concept of "realistic" meson potentials with short-range repulsion,[6,7] which is associated in ideology with the vector-meson exchange,[8] fails to describe the *S*- and *P*-phase shifts, apart from the weak point that here is a very great number of variable parameters (as a cost for empirical success), and the authors of the above-mentioned papers invoke here either the purely empirical parametrization of potentials[7] or the phenomenological form factors at the meson-nucleon vertex,[6] which now belong to the quark description. There are also such aspects of the *NN* interaction as cannot altogether be reflected by the meson exchange; namely, the formation of a string in the virtual exchange of color between nucleons, etc.

The quark approaches, which rest most actively, i.e., microscopically, on the quark-nucleon structure and deal with the six-quark (6*q*) problem, represent the method of expansion of the wave function in the nucleon overlap region in terms of the quark shell-model 6*q* configurations[9–11] and the resonating-group-method (RGM) components.[2,11,13] In the former method, consideration is given to quark configurations in the spherical 6*q* Massachusetts Institute of Technology (MIT) bag model which expresses the nonperturbative effects in a simplified manner in terms of the external pressure of the expelled vacuum. The levels in this bag model are split in energy by the perturbative quark-gluon interaction. In this way, a useful result has been obtained,[2,9–11] namely, in the description of the deuteron and *NN* interaction, the important role of the nontrivial quark configuration $s^4p^2[42]_x[42]_{CS}$ was revealed, which occurs along with the simple configuration $s^6[6]_x[2^3]_{CS}$ and becomes most favorable energetically if the strong color magnetic *qq* interaction is prevailing. The excited configuration mentioned above, as has long been noted,[11] is a microscopic basis for deep attractive forbidden state NN potentials[14–17] (FSP's) which we use in the present paper.

The qualitative conclusions outlined above were confirmed by the RGM calculations.[11,10] In this scheme, use was made of the frozen internal structure of the nucleon s^3 of three constituent quarks and of the simple phenomenological confinement (such as a "funnel") in their interaction. It has also been elucidated that an important role is played by the inclusion of internal excitations of negative parity in each of the nucleons of the NN system, and the corresponding RGM calculation, on a very broad basis for the first time, has given the NN attraction in the intermediate range $r \simeq 1$ fm.[18] This argument certainly is similar in spirit to the above-mentioned important role of excited quark configurations, which is also associated with the formation of a repulsive core. Namely, the destructive interference of the excited configuration s^4p^2 and the simple one s^6 (they have been mentioned above) leads to the fact that, as the RGM calculations show,[11,10] the wave function for the relative motion of nucleons $\Phi(r_{NN})$ in the S wave dies out in the short range, and the interaction between nucleons, as follows from most of the RGM calculations, acquires a soft repulsive core[11] which is so habitual in the meson "realistic" potentials. The above-mentioned destructive interference of the quark configurations s^6 and s^4p^2 is quite compatible with the electron-scattering form factors for the deuteron[1,2] if, as has been noted above, one assumes soft meson form factors. Yet, since there is noticeable arbitrariness here, we need a wider range of data. Such data, it would appear, are involved in the momentum distribution of nucleons in the deuteron, which is extracted from several types of experiments, both exclusive and inclusive. The corresponding theoretical procedure[19,20] takes into account the final-state interaction but ignores the important feature of composite particles, which is well known in the theory of nucleon clustering in atomic nuclei:[21] the significant role may be played by the amplitudes of deexcitation of virtual orbitally excited nucleons, which is characteristic of the quark configuration s^4p^2. Therefore, traditional conclusions[19,20] that the momentum distribution of nucleons in the deuteron corresponds to the repulsive core potentials should be regarded as some preliminary orientation in the intricate problem. With all this in mind, we sum up the main features of the quark microscopic picture of the NN interaction (for a wider discussion, see the review in Ref. 22). Namely, although we still virtually try, in a very naive phenomenological manner, to fit the fundamental quantum chromodynamics (QCD) results to the confinement region, very difficult for this theory, we yet understand in this case that, in reality, the question concerns a rather unstable balance of the quark configurations s^6 and s^4p^2, with which the existence or nonexistence of the short-range NN repulsion is associated. Therefore, the repulsion cited is not at all a stable fundamental property of the NN system and may, for example, be absent even in the ground state, or, if the repulsion is still present in the ground state, then it may be replaced by the nucleon-nucleon attraction with increase in the particle energy (with the accompanying appearance of a wave-function node characteristic of the FSP concept) since, in this case, obviously the configuration balance is shifted towards the excited quark configuration s^np^m. It can look like an "invisible" jump up (equal to π) of the scattering phase shift when crossing the corresponding positive-energy bound state.

In such a situation, on the one hand, it is necessary to enhance the reliability of microscopic calculations of the NN system, taking into account more and more adequately the nonperturbative interaction between constituent quarks.

In the nonperturbative QCD vacuum, the initial chiral invariance of the theory is spontaneously broken, being realized in the Goldstone mode; that is, light quarks (u,d) acquire constituent masses $m_q \simeq \frac{1}{3}m_N$ (Ref. 23) and new carriers of the interaction between quarks (π, σ, etc.) appear. The resultant change in the quark-quark potential may not at all be reduced to only some renormalization of constants of the perturbative quark-gluon interaction, as it is sometimes assumed empirically.[24] Moreover, the assumption itself of the pairwise interaction between constituent quarks is problematic. Furthermore, all this is complicated by the effects of a strong external meson field.

The problem of developing a microscopic quark theory of the NN interaction is complicated by the fact that the quark-nucleon radius is small, being, as is known, only 0.5–0.6 fm, and laborious search for convincing, experimentally observed, quark effects, which might constitute the field of action of such a theory, represents the second aspect of the possible plan of action under discussion.

In the present work, extending and developing the line of investigations,[14–17] we show that the nucleon-nucleon scattering data (differential cross sections and polarizations), extended in energy to include the values $E_{\rm lab}=6$ GeV (for one to be able to judge the short-range NN interaction), seem just to reflect the desired quark effect and convincingly demonstrate the dominant role of the excited quark configurations s^np^m ($n=4$, $m=2$ for S waves; $n=3=m$ for P waves). This is reflected by the efficiency of forbidden state potentials and by the corresponding very peculiar behavior of the scattering phase shifts, which is exhibited precisely in a wide energy interval of 0–6 GeV. Experiments are discussed which are necessary for the comprehensive verification of the existence of a node of the deuteron wave function at the point of the conventional repulsive core; that is, for an independent testing of the balance of quark configurations.

II. OPTICAL MODEL

For the analysis of the problems of interest to us, we use an optical model similar to that which is long known in low-energy nuclear physics in the description of nucleon-nucleus and cluster-cluster scattering. Such a model has not been previously employed here.

We rely on one of the variants of the relativistic quasipotential equation,[25,8] which takes the following form for the partial amplitudes:

$$T^J_{LS}(p,p')=V^J_{LS}(p,p') + \int V^J_{LS}(p,k)g(k,\sqrt{s})T^J_{LS}(k,p')dk \ , \quad (1)$$

$$V^J_{LS}(p,p')=(p\varepsilon_p)^{1/2}(p'\varepsilon_{p'})^{1/2} \times \int Y^*_{LM}(\hat{\mathbf{n}})V^J_{LS}(\mathbf{q})Y_{LM}(\hat{\mathbf{n}}')d\hat{\mathbf{n}}d\hat{\mathbf{n}}' \ ,$$

$$V^J_{LS}(\mathbf{q})=(2\pi)^{-3/2}\int \exp(i\mathbf{q}\cdot\mathbf{r})V^J_{LS}(\mathbf{r})d\mathbf{r} \ ,$$

$$g(k,\sqrt{s})=\frac{2k}{\frac{1}{4}s-\epsilon_k^2} \ ,$$

$$\epsilon_k^2=m^2+k^2 \ ,$$

$$\mathbf{q}=\mathbf{p}-\mathbf{p}' \ ,$$

$$\mathbf{p}=\mathbf{n}p \ ,$$

where the potentials V^J_{LS} are complex. These equations were solved using the Noyes-Kowalsky regularization[26] by the matrix inversion method.[26,27] In the numerical integration, use was made of the Gauss quadrature formulas. The relation between the partial amplitude and the scattering phase shift is given by the expression[28]

$$T^J_{LS}(p,p)=\frac{1}{2\pi i}(1-S^J_{LS})\equiv\frac{1}{2\pi i}(1-\eta^J_{LS}e^{2i\delta^J_{LS}}) \ , \quad (2)$$

where $4p^2=s-4m^2$, and, if the potential is real, then

$$T^J_{LS}(p,p)=-\frac{1}{\pi}\sin\delta^J_{LS}e^{i\delta^J_{LS}} \ .$$

The differential scattering cross section is given by the expression[8]

$$d\sigma/d\Omega=\frac{1}{4}\pi^2\lambda\!\!\!^-{}^2\sum_{lS}B_l(S)P_l(\cos\theta) \ ,$$

$$B_l(S)=\sum_{L_1L_2J_1J_2}[\bar{Z}(L_1J_1L_2J_2{:}Sl)]^2\mathrm{Re}(T^{J_1}_{L_1S}T^{J_2*}_{L_2S}) \ , \quad (3)$$

$$\bar{Z}(L_1J_1L_2J_2{:}Sl)=\hat{L}_1\hat{J}_1\hat{L}_2\hat{J}_2W(L_1J_1L_2J_2{:}Sl) \ , \qquad \hat{L}=\sqrt{2L+1} \ . \quad (4)$$

For the unpolarized initial nucleons, the degree of polarization of the nucleon emitted at an angle θ is given by the expression[28]

$$P(\theta)=\frac{Sp\{\sigma_y TT^+\}}{Sp\{TT^+\}} \ , \quad (5)$$

where σ_y is the Pauli matrix (the y axis is chosen to be perpendicular to the reaction plane).

In the denominator of Eq. (5) there occurs the differential cross section [see Eq. (3)] and the numerator appears as

$$S_p\{\sigma_y TT^+\}=\frac{\sqrt{3}}{8}\pi^2\lambda\!\!\!^-{}^2\sum_{LSL'J'l}(-1)^{L'+J'+1}\begin{Bmatrix}L & J & 1\\ J' & L' & 1\end{Bmatrix}\begin{Bmatrix}L & 1 & J\\ L' & 1 & J'\\ l & 1 & l\end{Bmatrix}\mathrm{Im}\{T^{J'}_{L'1}T^{J*}_{L1}\}\bar{P}^1_l(\cos\theta) \ , \quad (6)$$

where

$$\begin{Bmatrix}L & 1 & J\\ L' & 1 & J'\\ l & 1 & l\end{Bmatrix} \ ,$$

$$\begin{Bmatrix}L & J & 1\\ J' & L' & 1\end{Bmatrix}$$

are the $9j$ and $6j$ symbols, and the polynomials $\bar{P}^1_l(\cos\theta)$ are defined in the monograph.[28]

Finally, the total elastic cross section and the total cross section for all the processes are determined by the expressions[29]

$$\sigma_{\mathrm{el}}=\tfrac{1}{4}\pi^2\lambda\!\!\!^-{}^2\sum_{LSJ}(2J+1)|1-\eta^J_{LS}e^{2i\delta^J_{LS}}|^2 \ , \quad (7)$$

$$\sigma_{\mathrm{tot}}=\tfrac{1}{2}\pi^2\lambda\!\!\!^-{}^2\sum_{LSJ}(2J+1)(1-\eta^J_{LS}\cos 2\delta^J_{LS}) \ . \quad (8)$$

III. SCATTERING PHASE SHIFTS, DIFFERENTIAL CROSS SECTIONS, AND POLARIZATIONS

In the present paper we intend to investigate the efficiency of the FSP concept for describing the differential cross sections and the polarizations in the scattering of unpolarized particles. The quark-cluster model[9–12] imposes rather remarkable constraints onto the NN-interaction potential parameters. For instance, the S-wave-function node at $r\simeq 0.5$ fm in the NN channel, as it follows from the quark configuration s^4p^2 and from the baryon excitation spectrum parameter $\hbar\omega\simeq 500$ MeV,[9,10] requires the potential depth $V_0\simeq 1.2$ GeV and the width $a\simeq 0.5$ fm. This NN potential being refined empirically has one forbidden $0S$ state and no forbidden states in the highest even partial waves $D,G,I,\ldots,$ in perfect agreement with the requirements of the quark configuration s^4p^2.

Turning to odd waves, we automatically obtain that the P-wave-function node, as it follows from the quark configuration s^3p^3 and from the same $\hbar\omega$ value, should lie at $r\simeq 0.9$ fm (the P-wave "hard-core" position). There should be one forbidden $1P$ state and no forbidden states in $F,H,\ldots,$ waves. As it will be shown in what follows, the phase-shift running versus energy confirmed by the experimental data gives a rather sound substantiation of all these interconnections.

The pronounced J splitting of the ^3P_J phase shifts in the energy range of $E_{\mathrm{lab}}=0.3-3$ GeV, which will also be demonstrated, shows a joint effect of short-range spin-orbital and tensor forces (for some microscopical quark-gluon estimations, see Ref. 30). The large spin-orbital contribution to this splitting, presented by the "spin-orbital" combination of the ^3P_J phase shifts,[31] can be directly connected, in the usual phenomenological treatment, with the great gradient value of our central potential, $dV/dr \simeq 2$ GeV/fm at $r=a$.

Proceeding here in the simplest instructive way, we introduce the J dependence ($\mathbf{J}=\mathbf{L}+\mathbf{S}$) of the deep potentials characterizing the partial waves with given L, S, and J values ($L \leq 4$). This phenomenological strong J dependence accumulates the above effect of spin-orbital and tensor interactions in a simplified form (with no mixing angles), which is just reflected by the polarization.[32]

It proves to be sufficient to parametrize most of the partial-wave potentials by the simplest Gaussian form

$$V(r) = -V_0 \exp(-r^2/a^2) . \quad (9)$$

These potentials are listed in Table I and are determined from the familiar scattering phase shifts at energies of $E_{\text{lab}} < 1$ GeV.[33] On this basis, the scattering phase shifts,

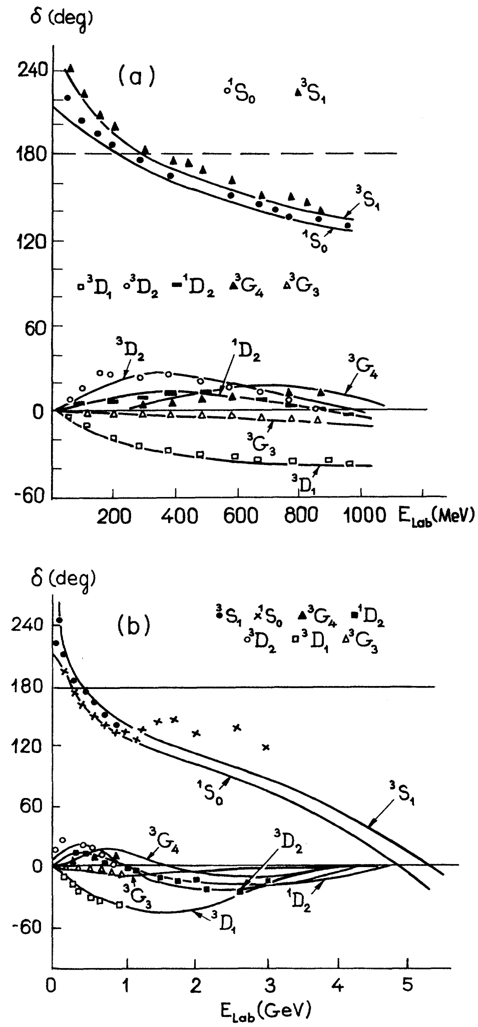

FIG. 1. (a) and (b) Phase shifts for even waves of NN scattering in the optical model [the solid line, our calculation; the dashed line, experimental data (Refs. 33 and 37)].

TABLE I. Parameters of the NN potential for different partial waves.

	$^1S_0$	$^3S_1$	$^1P_1$	$^3P_0$	$^3P_1$	$^3P_2$	$^1D_2$[a]	$^3D_1$[a]	$^3D_2$
V_0 (GeV)	1.52	1.61	1.08	2.15	0.75	2.39	1.26	1.40	1.56
a (fm)	0.65	0.65	0.88	0.42	0.71	0.45	0.78	0.42	0.56

	$^3D_3$	$^1F_3$	$^3F_2$	$^3F_3$[a]	$^3F_4$	$^1G_4$	$^3G_3$[a]	$^3G_4$	$^3G_5$
V_0 (GeV)	1.32	1.31	0.98	1.80	1.30	0.71	1.10	0.63	0.67
a (fm)	0.52	0.61	0.73	0.43	0.57	0.61	0.54	0.51	0.46

[a]See text for further details.

in turn, are extended to include, with no free parameters, the intermediate-energy region of 1–6 GeV (see Figs. 1 and 2). The addition to potential (9) in the region of low energies $E_{\text{lab}} < 0.4$ GeV of π-meson exchange tail is of no essential importance (of course, we have no intention to compete with the quality of "realistic" potential fit[6,7] of low-energy data).

However, the singlet D-wave requires, to describe the low-energy phase-shift running, a little bit more of a flexible deep attractive potential,

$^1D_2$:

$$V(r) = -V_0 \exp(-r^2/a^2) - V_1 \exp(-r^2/a_1^2) , \tag{10}$$

where values of the parameters V_0, a are listed in Table I and $V_1 = 0.71$ GeV, $a_1 = 0.44$ fm. Further, three partial waves $^3D_1$, $^3F_3$, and $^3G_3$ have small negative values of low-energy phase shifts due to the peripheral repulsion produced by the tensor and spin-orbital forces within the one-boson-exchange potential (OBEP) approach. We represent this repulsion by the simplest OBEP-like terms[34] matched to the universal short-range attraction discussed above,

$^3D_1$, $^3F_3$, $^3G_3$:

$$V(r) = -V_0 \exp(-r^2/a^2) + \theta(r - 1.5a) V_{\text{OBEP}}(r) , \tag{11}$$

where the V_0, a values (see Table I) and $V_{\text{OBEP}}(r)$ characteristics (see Ref. 34) do depend on the channel quantum numbers L, S, J. The π-meson exchange potentials are no longer of importance and lose their meaning at energies of $E_{\text{lab}} > 0.4$ GeV, etc.

For all the partial waves, naturally, we go over to complex optical potentials, adding to each real deep attractive potential the imaginary part which is proportional to the real one, in which case, the proportionality factor rapidly increases with increasing energy (Fig. 3). This is necessary to ensure the experimental relations $(\sigma_{\text{el}}/\sigma_{\text{tot}}) = 0.9$, 0.62, 0.44, 0.40 at the energies $E_{\text{lab}} = 0.6$, 1.0, 2.2, 3.2 GeV and so on.[35] The main features of the E dependence of the scattering phase shifts, as presented in Figs. 1 and 2, are as follows. The FSP description of the properties of the deuteron and scattering phase shifts at low energies of $E_{\text{lab}} < 500$ MeV is characterized by its compactness:[19,20] the number of parameters is four or five times smaller than that for the realistic meson potentials.[6,7] However, the main peculiar feature of the FSP concept, which distinguishes this concept from the others, is a specific behavior of the scattering phase shifts as a function of energy. Namely, all the lowest phase shifts $(L = 0, 1)$ are positive. In accordance with the generalized Levinson theorem, the S- and P-phase shifts, with increase in energy from zero, start from the value $\delta_L = \pi$ (and the triplet S-phase shift is 2π at zero) while the higher phase shifts $(L \geq 2)$ start from zero. The Born region $\delta_L < 1$, where all the phase shifts "converge" with increasing energy, lies at the energies $E_{\text{lab}} \sim 5$ GeV. In order to "recognize" this physical pattern, it is necessary to involve, in consideration, a wide energy range from zero to a value of just about 5–6 GeV.[13] In fact, the quark-gluon interaction of nucleons is more fundamental than the meson exchange and covers the region of smaller distances and higher energies. The principal importance

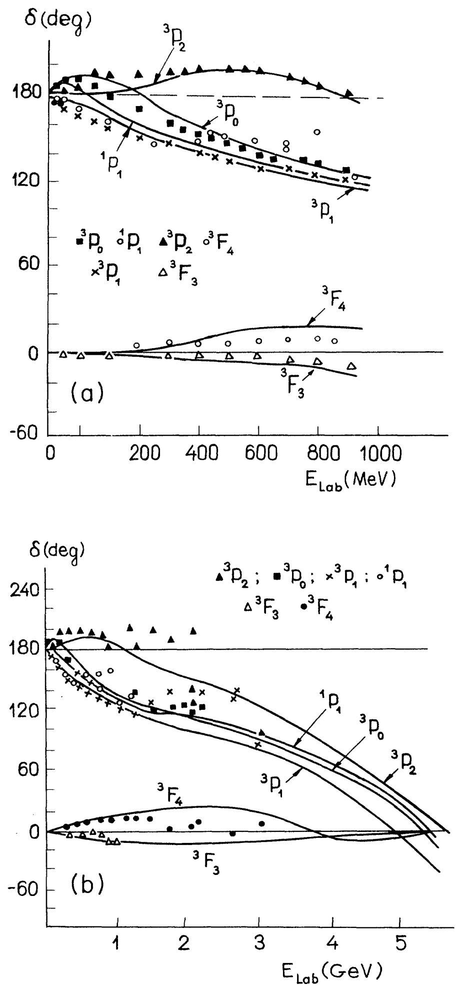

FIG. 2. (a) and (b) Phase shifts for odd waves of NN scattering in the optical model [the solid line, our calculation; the dashed line, experimental data (Refs. 33 and 37)].

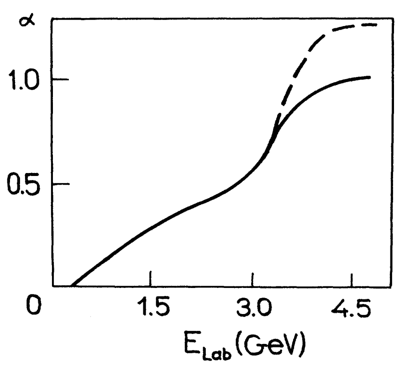

FIG. 3. The E dependence of the coefficient α, $W_{LS}^J(r) = \alpha V_{LS}^J(r)$ in each partial wave LSJ, α is dependent on L parity, too (the solid line, L even; the dashed line, L odd).

of precisely considering a wide energy range was previously demonstrated by the example of "recognition" of the $\alpha\alpha$ interaction.[13]

As a matter of fact, the experimental data on the NN scattering throughout the region $E_{\text{lab}}=0-5$ GeV have, so far, not been interpreted theoretically; that is, have not been adopted by theorists, and the present study, as far as we know, is the first attempt to propose the unified physical picture for this wide region. Let us also make some more particular remarks.

First, at the energies $E_{\text{lab}} \geq 1$ GeV, the odd phase shifts, while approaching the values of $\pi/2$ at $L=1$, become just as active in scattering as the even S- and D-phase shifts. This immediately implies the disappearance of the U-shaped form of the angular distributions in the np scattering in the transition from the energies $E_{\text{lab}} \leq 700$ MeV, when the odd waves have no effect, in practice, on the scattering, to the energies $E_{\text{lab}} \geq 1$ GeV, when the destructive interference of even and odd waves in the backward hemisphere is very significant. Figure 4 shows that the FSP concept explains well this phenomenon that has long been experimentally established,[35,8] whereby, in theory, in the explanation of both the shape of the curves and the absolute cross section, there is not a single variable parameter: all the parameters are determined by the description of scattering at $E_{\text{lab}} \leq 1$ GeV and by the values of the inelasticity $\sigma_{\text{el}}/\sigma_{\text{tot}}$ in the region under consideration, 600 MeV $< E_{\text{lab}} < 6$ GeV. The evolution of the differential pp scattering cross section[35] (Fig. 5) is also well accounted for. All these circumstances justify our reconstruction of scattering phase shifts as presented in Figs. 1 and 2.

Second, the small positive phase shifts ($^1D_2$, $^1F_3$, and others) reverse sign because of absorption at energies of about 1 GeV and the small negative phase shifts $^3D_1$, $^3F_3$, $^3G_3$, on the contrary, do not reverse their sign owing to the absorption at the same energies. Here, the only result of the short-range attraction is that the negative phase shifts mentioned above die out very fast as energy increases above 1 GeV. Now we point out the significant E dependence of the moduli $|S_L| \equiv \eta_L$, which is presented for our optical model in Fig. 6. This dependence reflects well the influence of absorption, which increases with energy. We see that, at the values of $E_{\text{lab}} > 4-5$ GeV,

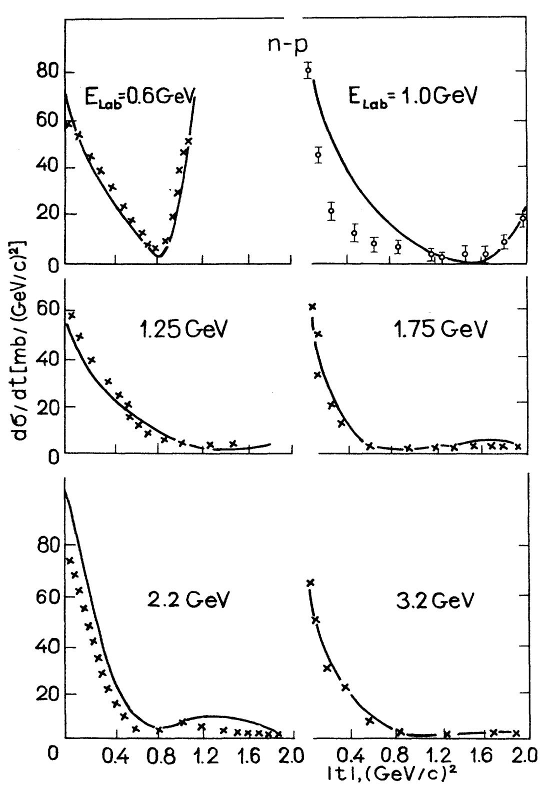

FIG. 4. Description of the differential np scattering cross sections at different energies. The solid line, our calculation; the experimental data are from Ref. 35.

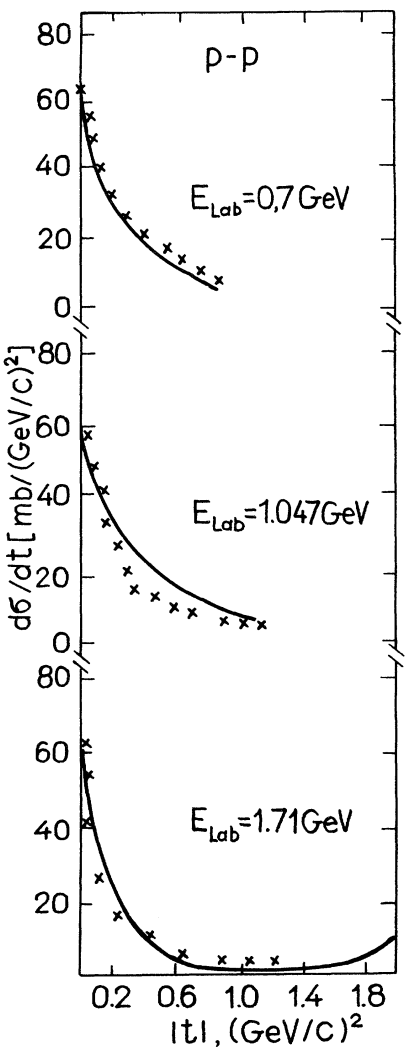

FIG. 5. Description of the differential pp scattering cross sections at different energies. The solid line, our calculation; the experimental data are from Ref. 35.

values of $|S_L|$ are small for the low partial waves $L=0$, 1, 2; that is, the most informative lower scattering phase shifts become unstable here from the viewpoint of empirical reconstruction, and the "transparency window" is closed, although the higher scattering phase shifts remain sensitive to the peripheral part of the real potential, also, at high energies. On the whole, fortunately, the "transparency window" turns out to be sufficiently wide, and the energy interval $E_{\rm lab}=0-6$ GeV reflects quite well the pronounced individual features of the E dependence of scattering phase shifts, which are characteristic of the FSP concept. Figure 7 gives an evidence that our description of differential cross sections gives the reasonable results even at the high-energy border of the "transparency window." Now we turn to the analysis of the polarization data, which, combined with the differential sections, just makes it possible to draw some noteworthy conclusions.

If the amplitudes T^J_{LS} are independent of J (at given L and S), then the polarization $P(\theta)$ becomes zero. Our potentials, markedly corrected in comparison with the preliminary work[17] (see Table I), are strongly dependent on J, and the polarization data in the wide energy region considered are a verification of these potentials, which are independent of the differential cross sections. The results of calculations of the polarization in the energy region $E_{\rm lab}=0.5-5$ GeV are displayed in Figs. 8–10 against the experimental data.[36] It can be seen from these figures that there is a stable agreement between the calculated and measured polarizations $P(\theta)$, and only the 5-GeV data show some influence of neglected partial waves $L\geq 5$. The even-L–odd-L splitting of the absorption coefficient α, shown in Fig. 3, is introduced to take into account the empirical data[37] on the energy dependence of the reflection coefficients $\eta_L\equiv|S_L|$ (see Fig. 6).

Thus, our description has an advantage over the attempt[38] to consider the polarization in the region of $E_{\rm lab}\simeq 1.5$ GeV from the conventional standpoint on the basis of exchange of vector mesons. The J splitting of phase shifts and the corresponding J splitting of potentials, which we analyze, can be interpreted in terms of the

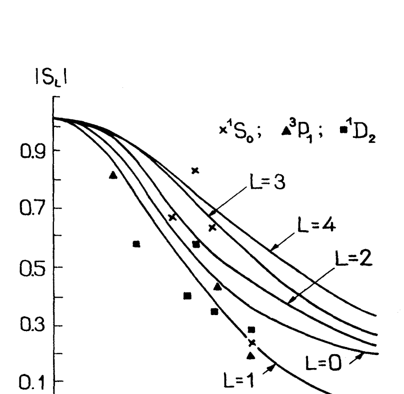

FIG. 6. The E dependence of moduli of the S matrix $\eta_L=|S_L|$ in the optical model. The experimental data are given after Ref. 37.

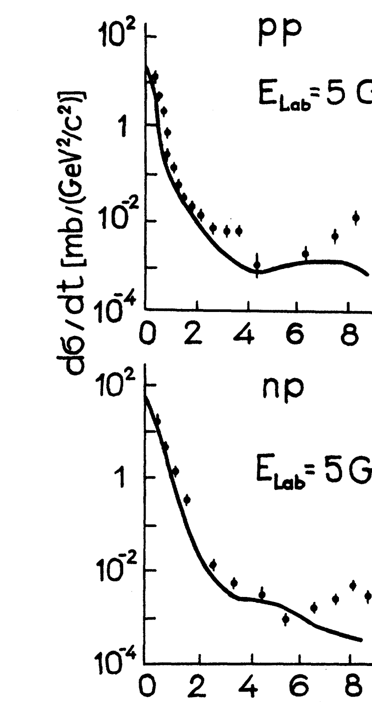

FIG. 7. Dependence of the pp and np scattering cross sections on the momentum transfer value at the energy $E_{\rm lab}=5$ GeV. The experimental data are from Refs. 35 and 36.

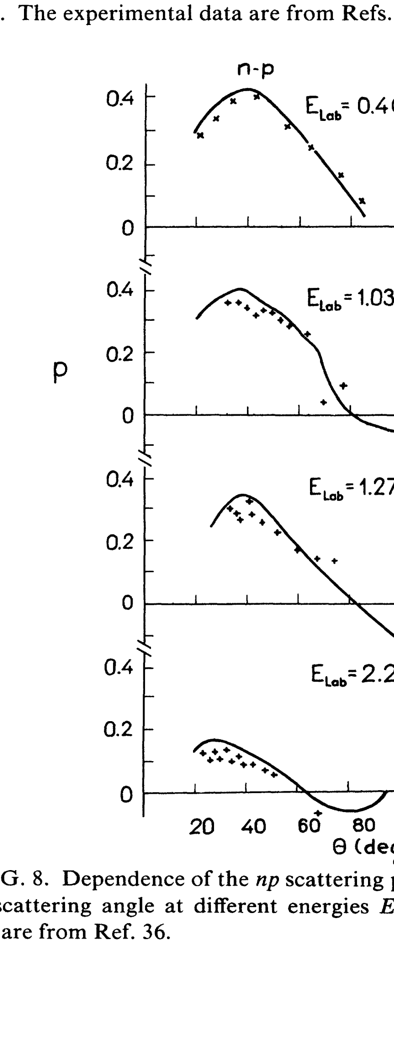

FIG. 8. Dependence of the np scattering polarization $P(\theta)$ on the scattering angle at different energies E. The experimental data are from Ref. 36.

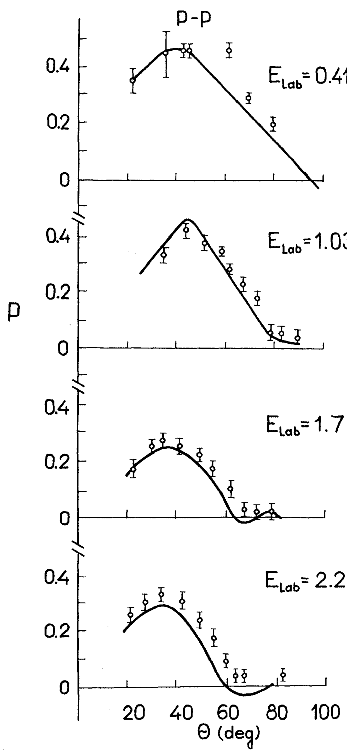

FIG. 9. Dependence of the pp scattering polarization $P(\theta)$ on the scattering angle at different energies E. The experimental data are from Ref. 36.

spin-orbit and tensor interactions [in the pure spin-orbit interaction, the phase shift $\delta(^3P_1)$ would lie between the phase shifts $\delta(^3P_2)$ and $\delta(^3P_0)$ but, as we see, the real situation is far from this]. The total number of adjustable parameters (which are fixed in our theory entirely by the low-energy data $E_{\rm lab} \leq 800$ MeV) can probably be diminished even more if we try to unify the description of singlet waves by means of more flexible potentials than

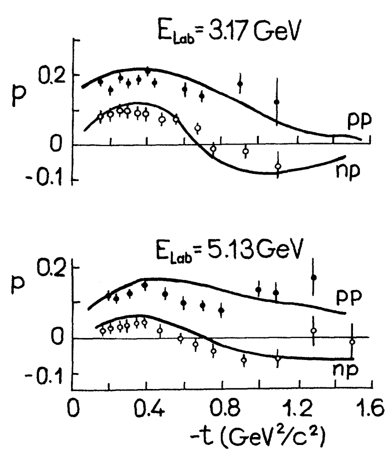

FIG. 10. Dependence of the pp and np scattering polarization $P(t)$ on the momentum-transfer value at the energies $E_{\rm lab}=3$ and 5 GeV. The experimental data are from Ref. 36.

Gaussian ones.

We now compare the results of our calculation of the scattering phase shifts (Figs. 1 and 2) with the tentative phase-shift analysis by Hoshizaki.[37] For this purpose, obviously, it is necessary to move its triplet and singlet S-phase shifts by a value of 2π and π upward, respectively, and all the P-phase shifts by π. As a result, we see (Figs. 1 and 2) the undoubtful general correspondence, against the background of which narrow dibaryon resonances can occur, which Hoshizaki attempts to separate out (this, possibly, reflects the influence of the opening inelastic channels and also the decay of some specific multiquark states[39]). In addition, one can see the general instability of the empirical phase-shift analysis, which comes from the fact that the total number of significant scattering phase shifts is large and the phase shifts themselves are complex quantities as a result of absorption.

IV. DISCUSSION AND PROSPECTS

Thus, a new general picture of the NN interaction becomes visible in outline, how it is exhibited in the scattering phase shifts. The above-mentioned picture is such that we observe separation of the scattering phase shifts into two groups: the "large" S- and P-phase shifts (they are equal to 2π or π at zero energy) and the "small" ones which include all the others. This result of Ref. 17 and of the present study is just in agreement with the quark configuration $s^4p^2[42]_x[42]_{CS}$ for the deuteron ground state. It is well known in the theory of cluster phenomena ($\alpha\alpha$ scattering, etc.)[40,41] that such a configuration corresponds to the wave-function node for the S wave instead of the repulsive core in the NN channel, which is the cause of displacement of the S-phase shift by the value π upward (the generalized Levinson theorem[40]). The configuration cited has long been discussed in the literature as the important component of the wave function for the NN system,[9–13] and the more interesting are the results we have obtained now in favor of precisely this configuration that, at the present level of knowledge, there are no arguments as to why the component $s^4p^2[42]_x$ may suppress all the others (it would be the only one, just as in the t-h system,[40] should there be no color). Next, if the quark configurations below s^4p^2 are not realized dynamically, then, for the P wave of NN scattering, we shall have the quark configuration $s^3p^3[33]_x[33]_{CS}$ if we proceed from the principle of maximum CS symmetry,[9–11,14] i.e., in the P wave there will be a node at the radius $r \simeq 0.9$ fm and the P-phase shifts will start at low energies just from the value π.[16] A change of the absolute scale of the S- and P-phase shifts is well substantiated in the good and natural matching of our phase-shift curves to the Orear diffraction cone regime[41] at the high-energy border of $E_{\rm lab} > 6$ GeV, where all the phase shifts take on small negative values due to the very predominant absorption. This matching rules out the alternative interpretation (possible if there is no strong absorption) when the asymptotic values of S- and P-phase shifts are equal to $-\pi$ owing to the specific shape of the hard core.[42]

The main features of the differential cross section and polarizations are formed primarily by the profound effect of a mutually consistent falloff with increase in the energy of six large phase shifts $L=0,1$, and the peculiarity of the picture lies in that, although the S-phase shifts move downward fairly rapidly, the P-phase shifts, due to the centrifugal barrier influence, are "sticking" near the value $\delta_L=\pi$ up to the energies $E_{\text{lab}}=300$–500 MeV. Hence, there is no longer question concerning the observation that, in some P channels, the systems exhibit attraction while in the others, repulsion. We see that, in all the P channels, there is short-range attraction; that is, the Majorana exchange forces do not manifest themselves in the short range while, in the long-range region of pion exchange, the interaction in odd waves may be absent owing to a possible structure of pion exchange $\frac{1}{2}(1+P^x)$,[8] which, however, should be substantiated anew since, previously, this fact was deduced mainly from the U-shaped symmetry of differential np scattering cross section mentioned above.

It turns out that the potentials for singlet waves $L\leq 3$ have remarkable features in common. We hope these potentials can be unified by means of some flexible construction like the extended expression (9) separately for even waves ($[f]_x=[42]$) and odd waves ($[f]_x=[33]$). For triplet waves, such unification and compactness of the description will be associated, obviously, with the incorporation in the explicit form of the tensor and spin-orbit interactions with account taken of the qualitative possibilities of the quark approach to this problem.[36]

The above results give us an important evidence for the bright quark effects in the nucleon-nucleon scattering based on the composite nucleon structure. We see no alternative interpretation of the broad energy range data discussed above. However, to reinforce our arguments in favor of the FSP concept, it is necessary to undertake independent investigations into the existence of the excited quark configurations and the wave function node in the deuteron at $r\simeq 0.5$ fm.

First, the excited quark configuration s^4p^2 in the deuteron gives rise to virtual excited nucleons N^* and N^{**} (with one or two oscillator quanta $\hbar\omega$ of the internal excitation) which can be observed as spectators in the exclusive quasielastic knockout experiments $^2\text{H}(e,e'p)N^*$, $^2\text{H}(e,e'p)N^{**}$, etc.,[43] with the expected probability of $w\sim 10^{-3}$ as compared to the known $^2\text{H}(e,e'p)n$ cross section. It is reasonable here to have the final-state relative p-N^* (or p-N^{**} or p-n) momentum value p_{rel} large enough, $p_{\text{rel}}\geq 1$ GeV/c, to avoid the meson-exchange corrections, etc.[44]

Second, here one should keep in mind the calculation of the momentum distribution of nucleons in the deuteron taking into account those "cluster deexcitation" peculiarities[21] which were outlined in the Introduction.

Third, one should theoretically and experimentally investigate the deuteron photodisintegration $\gamma+d\rightarrow n+p$ in the energy region $E_\gamma\sim 500$–1000 MeV, i.e., where the P-phase shifts pass through $\pi/2$ when decreased. Use of polarized photons enables one to separate out the $E1$ absorption.[45] The quark configurations appear in the theory of the process in question in a quite different manner than in the electron form factors for the deuteron (both elastic and spin-isospin-flip inelastic form factors). So, here we shall have independent valuable information. Let us now discuss other important, albeit more indirect, possibilities for verifying the description of NN scattering with the aid of the FSP concept.

In the foregoing we have pointed out that the peripheral repulsion, associated with the spin-orbit and tensor interactions, outbalances the effect of central attraction in the case of the phase shifts $^3D_1$, $^3F_3$, $^3G_3$, etc., which are negative at all energies. The general consistency of the short-range spin-orbital and tensor interactions (including mixing angles which were not involved in our present investigation) can be verified by analysis of polarization correlation experiments.

Next, the mutual transparency of nucleons should lead to marked changes in the properties of nuclear matter. In particular, due to large virtual momenta of nucleons in the region of their overlap, the microscopic picture of compressibility of the nucleus should change.

It is also interesting to elucidate to what extent the concept of deep attractive potentials is applicable to other hadron pairs. As regards the $N\bar{N}$ interaction, where no *a priori* prerequisite arose about the existence of a repulsive core associated with vector mesons, it has long been admitted that here the optical potential with the attractive real part is applicable which, according to modern estimates,[46] is qualitatively similar to what we have in the NN system, but the strongly bound $N\bar{N}$ states are not necessarily dynamically forbidden, although if they do exist they seem to be greatly smeared in energy.

However, certainly, the most consistent approach, as has been noted in the Introduction, lies in a further quark shell-model investigation of the nonperturbative interaction between constituent quarks in the confinement region of partial nucleon-nucleon overlap.

The results outlined above indicate that significantly more powerful carriers of the NN interaction are required that those which are usually considered in the quark models of nuclear forces. In fact, if we have high-power NN interaction potentials 1–2 GeV deep, this is comparable to the energy of a gluon condensate in the volume of the two-nucleon system.[47] In this respect, of interest is Ref. 48 where the interaction between quarks via the exchange of glueballs or heavy colored gluonic bunches is discussed. The masses of glueballs and bunches have the desired energy scale $\simeq 1$ GeV and the interaction they produce will take place just in the region of significant nucleon overlap, which corresponds to low partial waves $E_{\text{lab}}=1$–5 GeV. Our results suggest the symmetry $(\lambda\lambda)(\sigma\sigma)$ for this color-exchange interaction.[9–11]

The authors express gratitude to Prof. V. V. Anisovitch, Prof. V. I. Efimov, Prof. A. M. Green, Prof. F. Cross, Prof. A. D. Jackson, Prof. V. I. Kukulin, Prof. F. Myhrer, Prof. E. M. Myman, Prof. D. O. Riska, Prof. M. G. Ryskin, Prof. V. I. Savrin, and Prof. J. Soffer for stimulating comments and advice.

[1]L. Ya. Glozman, N. A. Burkova, and E. I. Kuchina, Z. Phys. A **332**, 339 (1989).
[2]H. Ito and A. Faessler, Nucl. Phys. **A470**, 626 (1987).
[3]W. Leidemann and H. Arenhoevel, Z. Phys. A **326**, 333 (1987).
[4]A. V. Efremov, V. T. Kim, and G. I. Lykasov, Yad. Fiz. **44**, 241 (1986); V. V. Burov, V. K. Lukyanov, and A. I. Titov, Fiz. Elem. Chastits At. Yadra [Sov. J. Part. Nucl.] **15**, 1249 (1984).
[5]S. Sh. Ameyev, E. S. Golubyatnikova, B. N. Kalinkin, G. A. Shakhanov, and V. L. Shmonin, Joint Institute for Nuclear Research Report JINR-P2-87-466, 1987.
[6]R. Machleidt, K. Holinde, and Ch. Estor, Phys. Rep. **149**, 1 (1987).
[7]M. Lacombe, B. Loiseau, J. M. Richard, R. Vihn Mau, J. Cote, P. Pires, and R. de Tourreil, Phys. Rev. C **21**, 861 (1980).
[8]G. E. Brown and A. D. Jackson, *Nucleon-Nucleon Interaction* (North-Holland, Amsterdam, 1976).
[9]I. T. Obukhovsky and A. M. Kusainov, Phys. Lett. B **238**, 142 (1990).
[10]A. M. Kusainov, V. G. Neudatchin, and I. T. Obukhovsky, Phys. Rev. C (to be published).
[11]I. T. Obukhovsky, V. G. Neudatchin, Yu. F. Smirnov, and Yu. M. Tchuvil'sky, Phys. Lett. **88B**, 231 (1979).
[12]M. Harvey, Nucl. Phys. **A352**, 326 (1981).
[13]V. G. Neudatchin, I. T. Obukhovsky, V. I. Kukulin, and N. F. Golovanova, Phys. Rev. C **11**, 128 (1975).
[14]V. G. Neudatchin, I. T. Obukhovsky, and Yu. F. Smirnov, Fiz. Elem. Chastits At. Yadra [Sov. J. Part. Nucl.] **15**, 1165 (1984).
[15]V. I. Kukulin, V. N. Pomerantsev, V. M. Krasnopol'sky, and P. B. Sazonov, Phys. Lett. **125B**, 20 (1984); Yad. Fiz. **43**, 559 (1986).
[16]V. M. Krasnopol'sky, V. I. Kukulin, and V. N. Pomerantsev, Izv. Akad. Nauk SSSR, Ser. Fiz. **51**, 898 (1987).
[17]Yu. L. Dorodnykh, V. G. Neudatchin, N. P. Yudin, and I. T. Obukhovsky, Yad. Fiz. **48**, 1796 (1988); **50**, 170 (1989).
[18]K. Maltman and N. Isgur, Phys. Rev. D **29**, 952 (1984).
[19]C. Chiofi degli Atti, E. Pace, and G. Salme, Phys. Rev. C **36**, 1208 (1987); C. Chiofi degli Atti, Nucl. Phys. **A463**, 127 (1987).
[20]H. Arenhoevel, Nucl. Phys. **A384**, 287 (1982).
[21]N. F. Golovanova, I. M. Il'in, V. G. Neudatchin, Yu. F. Smirnov, and Yu. M. Tchuvil'sky, Nucl. Phys. **A262**, 444 (1976); V. G. Neudatchin, Yu. F. Smirnov, and N. F. Golovanova, in *Advances in Nuclear Physics,* edited by J. Negele and E. Vogt (Plenum, New York, 1979), Vol. 11, pp. 1–132.
[22]F. Myhrer and J. Wroldsen, Rev. Mod. Phys. **60**, 629 (1988).
[23]E. V. Shuryak, Phys. Rep. **115**, 151 (1984); D. I. Dyakonov and V. Yu. Petrov, Zh. Eksp. Teor. Fiz. **89**, 361 (1985); **89**, 751 (1985).
[24]N. Isgur and G. Karl, Phys. Rev. D **11**, 2653 (1979).
[25]A. A. Logunov and A. N. Tavkhelidze, Nuovo Cimento **29**, 380 (1963); R. Blankenbeckler and R. Sugar, Phys. Rev. **142**, 1051 (1966).
[26]H. P. Noyes, Phys. Rev. Lett. **15**, 538 (1965); O. O. Kowalski, *ibid.* **15**, 798 (1965).
[27]M. I. Haftel and F. Tabakin, Nucl. Phys. **A158**, 1 (1970).
[28]A. M. Baldin, V. I. Goldansky, V. M. Maximenko, and I. L. Rosental, *Kinematices of Nuclear Reactions* (Atomizdat, Moscow, 1968).
[29]F. Nikitiu, *Phase-Shift Analysis in the Physics of Nuclear Interactions* (Mir, Moscow, 1983).
[30]Y. Suzuki and K. T. Hecht, Nucl. Phys. **A420**, 525 (1984); Y. He, F. Wang, and C. W. Wong, *ibid.* **A451**, 666 (1986).
[31]P. S. Signell, in *Two-Body Forces in Nuclei,* edited by S. M. Austin and G. M. Crawley (Plenum, New York, 1972), p. 9.
[32]G. R. Goldstein and M. J. Moravcsik, Can. J. Phys. **65**, 445 (1987).
[33]R. A. Arndt, L. D. Roper, R. A. Bryan, R. B. Clark, B. J. Wer West, and P. S. Signell, Phys. Rev. D **28**, 97 (1983).
[34]J. M. Eisenberg and W. Greiner, *Nuclear Theory* (North-Holland, Amsterdam, 1972), Vol. 3.
[35]D. V. Bugg, D. S. Salter, G. M. Stafford, R. E. George, K. F. Riley, and R. J. Tapper, Phys. Rev. **146**, 980 (1966); M. L. Perl, J. Cox, M. J. Longo, and M. N. Kreisler, Phys. Rev. D **1**, 1857 (1970).
[36]P. R. Robrish, O. Chamberlain, R. D. Field, R. Z. Fuzesy, W. Gorn, C. C. Morehouse, T. Powell, S. Rock, S. Shannon, G. Shapiro, H. Weisberg, and M. J. Longo, Phys. Lett. **31B**, 617 (1970); M. Borghini *et al.*, *ibid.* **B31**, 405 (1970); R. Diebold, D. S. Ayres, S. L. Kramer, A. J. Pawlicki, and A. B. Wicklund, Phys. Rev. Lett. **35**, 632 (1975).
[37]Y. Higuchi and N. Hoshizaki, in *High Energy Photonuclear Reactions,* Proceedings of the International Symposium at the Institute of Nuclear Study (University of Tokyo Press, Tokyo, 1983), p. 124.
[38]I. Hulthage and F. Myhrer, Phys. Rev. C **30**, 298 (1984).
[39]M. P. Locher, M. E. Sainio, and A. Svarc, Adv. Nucl. Phys. **17**, 47 (1986).
[40]V. I. Kukulin, V. G. Neudatchin, I. T. Obukhovsky, and Yu. F. Smirnov, *Clusters as Subsystems in Light Nuclei* (Vieweg, Braunschweig, 1983), Vol. 3; V. I. Kukulin, V. G. Neudatchin, and Yu. F. Smirnov, Nucl. Phys. **A245**, 429 (1975).
[41]J. Orear, Phys. Lett. **13**, 190 (1964), Phys. Rev. D **18**, 2484 (1978); A. A. Logunov and O. A. Khrustalev, Fiz. Elem. Chastits At. Yadra [Sov. J. Part. Nucl.] **1**, 70 (1970); D. Ito, Prog. Theor. Phys. Suppl. 626 (1965).
[42]F. Michel and G. Reidemeister, Z. Phys. A **329**, 385 (1988).
[43]L. Ya. Glozman, V. G. Neudatchin, I. T. Obukhovsky, and A. A. Sakharuk, Phys. Lett. B **252**, 23 (1990).
[44]M. A. Ignatenko and G. I. Lykasov, Yad. Fiz. **46**, 1080 (1987).
[45]V. P. Barannik *et al.*, Nucl. Phys. **A451**, 751 (1986).
[46]T. A. Shibata, Nucl. Phys. **A478**, 623 (1988); T. Hippchen, K. Holinde, and W. Plessas, in Proceedings of the International Conference on Few-Body and Quark-Hadronic Systems, Dubna, 1987, edited by V. K. Lukyanov [Joint Institute of Nuclear Research Report D4-87-692, 1987], p. 215.
[47]M. A. Shifman, A. I. Vainstein, and V. I. Zakharov, Nucl. Phys. **B147**, 385; **B448**, 519 (1979).
[48]V. V. Anisovitch, S. M. Gerasyta, and I. V. Keltula, Yad. Fiz. **46**, 1508 (1987); V. V. Anisovitch and S. M. Gerasyta, *ibid.* **44**, 174 (1986).