Elastic electron–deuteron scattering

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Theoretical and experimental studies of elastic electron–deuteron scattering published after 1974 are reviewed. The possible part played by meson and quark degrees of freedom in the deuteron is discussed in the light of recent achievements. A new consistently relativistic approach to the investigation of the electromagnetic structure of the deuteron, developed with the participation of the authors, is described. The results are compared with those of other authors. Polarization effects in elastic *ed* scattering are considered.

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INTRODUCTION

For a long time, elastic electron–deuteron scattering has been the subject of intense theoretical and experimental study. Some reasons for this go back to the fifties and sixties (the problem of describing the nucleon–nucleon interaction off the mass shell, the problem of relativistic description of composite hadronic systems, the problem of extracting electromagnetic form factors of the neutron from *ed* scattering data), while others relate to investigations of recent years (the validity of nonrelativistic methods in the theoretical investigation of relativistic composite systems, the quark structure of the deuteron, meson and baryon degrees of freedom in the deuteron, various problems in the physics of weak interactions). Therefore, there has recently been a definite increase in interest in the deuteron.

The general situation in electron–deuteron scattering is briefly as follows. In the experimental studies of recent years the results of earlier experiments have been reproduced, made more accurate, or extended to the region of larger momentum transfers. No experiments of new type (polarization) have been made. The general direction of the theoretical studies can be characterized, on the one hand, by the rapid development of the nonrelativistic approach, i.e., the construction of ever more complicated and cumbersome nucleon–nucleon potentials, and, on the other, by the increasing "penetration" (at least by corresponding attempts) of relativistic field-theoretical methods into the physics of few-nucleon systems in general and, in particular, into the apparently weakly bound and typically nonrelativistic system that is the deuteron. There have been some studies of the contributions of the relativistic effects, the meson and baryon degrees of freedom, the manifestation of the quark structure of few-nucleon systems, and other questions. But at the present time one can hardly say that the results match the efforts that have been made. For example, there are still no clear and unambiguous answers to old and "very simple" questions such as that of the precise contribution of the D state to the deuteron and the correct calculation of the static characteristics (magnetic and quadrupole moments of the deuteron). Nor is there complete clarity in the framework of the standard approaches with regard to the problem of the extraction of the charge form factor of the neutron and its slope at the origin from *ed* scattering data. Nevertheless, analyzing the results obtained since 1974 (the earlier results are discussed in the review of Ref. 1), we can say that detailed comparison of the modern theoretical approaches with the experiments requires much more ingenious and accurate experiments than those so far made.

Limitations of space have forced us to omit all questions relating to electrodisintegration of the deuteron, and also the problem of extracting the neutron charge form factor from *ed* scattering. Besides this, in elastic *ed* scattering we do not consider the problem of weak interactions, and restrict ourselves to electromagnetic and strong interactions.

1. PRESENT EXPERIMENTAL STATUS OF ELASTIC ELECTRON–DEUTERON SCATTERING

In the single-photon approximation, the cross section for elastic scattering of unpolarized electrons by unpolarized deuterons has the form[1)]

$$\frac{d\sigma}{d\Omega_e}=\left(\frac{d\sigma}{d\Omega_e}\right)_{\text{Mott}}[A(Q^2)+B(Q^2)\,\mathrm{tg}^2(\theta_e/2)], \tag{1}$$

where $(d\sigma/d\Omega_e)_{\text{Mott}}$ is the Mott scattering cross section, and Q^2 is the square of the 4-momentum transfer. The longitudinal and transverse parts of the cross section (1) can be expressed in terms of the charge $G_C(Q^2)$, magnetic $G_M(Q^2)$, and quadrupole $G_Q(Q^2)$ form factors, respectively, by means of the relations

$$\left.\begin{aligned} A(Q^2)&=G_C^2(Q^2)+\frac{8}{9}\eta^2G_Q^2(Q^2)+\frac{2}{3}\eta G_M^2(Q^2);\\ B(Q^2)&=\frac{4}{3}\eta(1+\eta)G_M^2(Q^2),\ \ \eta=Q^2/4M_d^2.\end{aligned}\right\} \tag{2}$$

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1)*Translator's Note.* The Russian notation for the trigonometric, inverse trigonometric, hyperbolic trigonometric functions, etc., is retained here and throughout the article in the displayed equations.

Since 1973, the results of six experiments[2-7] on elastic ed scattering have been published. The experiments of Refs. 2, 3, and 7 greatly extended the region with respect to Q^2 of measured values of $A(Q^2)$ and $B(Q^2)$. In Ref. 2, $A(Q^2)$ was measured up to $Q^2 = 154.1\ \mathrm{F}^{-2}$, which exceeded by almost 4.5 times the value of the square of the momentum transfer achieved earlier in elastic ed scattering, which was $Q^2 = 35.4\ \mathrm{F}^{-2}$. Naturally, the experimental errors grow strongly with increasing Q^2, since the value of $A(Q^2)$ itself then decreases strongly [in the interval measured in Ref. 2, from the point $Q^2 = 20.5\ \mathrm{F}^{-2}$ to the point $Q^2 = 154.1\ \mathrm{F}^{-2}$, the cross section (1) decreases by 2×10^5 times]. At the maximal value of the square of the momentum transfer achieved in Ref. 2, $Q^2 = 154.1\ \mathrm{F}^{-2}$, an estimate of the upper limit of the scattering cross section was obtained: $(d\sigma/d\Omega_e) \leqslant 5\times10^{-40}\ \mathrm{cm}^2/\mathrm{sr}$ (i.e., not more than one event per week at the maximal current of the 19-GeV electrons).

Indirect results of the determination of $A(Q^2)$ at large Q^2 up to $Q^2 = 8\ (\mathrm{GeV}/c)^2$ are given in Ref. 7. The reported experiment was a measurement of the cross section of the reaction of inclusive ed scattering at a value of order $W^2 \simeq M_d^2$ for the square of the invariant hadronic mass in the final state. The function $A(Q^2)$ was determined from these data by using the fact that for $W - M_d \lesssim 200$ MeV and $Q^2 \gtrsim 2\ (\mathrm{GeV}/c)^2$ the relation $W_2(Q^2, W^2)/A(Q^2) = \mathrm{const}(W^2)$ holds, where νW_2 is the structure function of deep inelastic scattering. Among the experiments at small Q^2, we mention the results obtained in Ref. 4 of precision measurements of the ratio of the cross sections of elastic ed and ep scattering in the interval $0.044 \leqslant Q^2 \leqslant 3.98\ \mathrm{F}^{-2}$.

Thus, the longitudinal part $A(Q^2)$ of the cross section of elastic ed scattering has now been measured in the very wide interval $0.044 \leqslant Q^2 \leqslant 205.4\ \mathrm{F}^{-2}$ of momentum transfers. The modern data for $A(Q^2)$ are given in Fig. 1 (see also Ref. 1, Fig. 2).

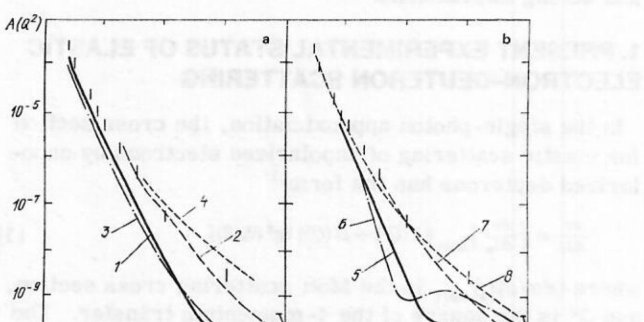

FIG. 1. Experimental data for the function $A(Q^2)$ compared with predictions of theoretical models: a) 1) results of calculations with Reid potentials with soft core, 2) the same with allowance for meson exchange currents, 3) results of calculations with Reid potentials with soft core and Gross-type relativization, 4) six-quark model[153]; b) 5) our results in the nonrelativistic case, 6) the same with complete allowance for the relativistic corrections to the deuteron charge form factor (Ref. 27), 7) phenomenological relativization (Ref. 45), 8) Kobushkin's six-quark model.[140]

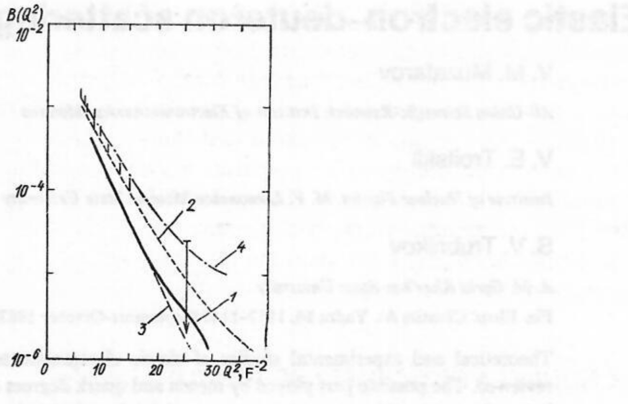

FIG. 2. Experimental data for the function $B(Q^2)$. The notation is the same as in Fig. 1a.

The elastic ed scattering experiments made in recent years have also made it possible to obtain new information about the function $B(Q^2)$. The experiments of Ref. 3 (the results of which are shown in Fig. 2) extended the region of measured values from $Q^2 = 14\ \mathrm{F}^{-2}$ to $Q^2 = 25.7\ \mathrm{F}^{-2}$. The value obtained for $B(Q^2)$ (admittedly, with a large error) at $Q^2 = 25.7\ \mathrm{F}^{-2}$ is $(0.59\pm1.20)\times10^{-5}$. Bearing in mind that the existing data for $G_M(Q^2)$ at small values of Q^2 are characterized by only moderate accuracy, we should mention the experiment of Ref. 6, in which the deuteron magnetic form factor was measured with high accuracy at $Q^2 = 0.031\ \mathrm{F}^{-2}$ (electron with initial energy $E_e = 56.4$ MeV scattered elastically through angle $\theta_e = 180°$). In order to determine more accurately the contribution of the meson exchange currents to the deuteron form factors, it would be desirable to make equally accurate measurements of $G_M(Q^2)$ for other values of Q^2. So far, $B(Q^2)$ has been measured with some degree of accuracy in the interval $0.28 \leqslant Q^2 \leqslant 25.7\ \mathrm{F}^{-2}$. The experimental data for $B(Q^2)$ are shown in Fig. 2 (see also Ref. 1, Fig. 3).

The most interesting polarization experiments, namely, elastic scattering of electrons by an aligned deuteron target or experiments in which the polarization tensor of the recoil deuterons is measured, have not so far been made at any value of the momentum transfer. The information available in the literature on the possibility and plans for such experiments is fragmentary and indirect. In Ref. 8, for example, the prospects for polarization measurements in elastic ed scattering at large Q^2 are rated very pessimistically. At low and medium Q^2 such experiments are more realistic. In this connection, a polarized deuteron target is mentioned in Ref. 9, while in Ref. 10 it is pointed out that an experiment to measure the polarization tensor of the recoil deuterons at $Q^2 = 6\ \mathrm{F}^{-2}$ is currently in the stage of practical examination. Further, Ref. 11 announces the experimental possibility of measuring the recoil-deuteron polarization tensor in elastic ed and πd scattering on the basis of the $^3\mathrm{He}(d,p)^4\mathrm{He}$ reaction in the interval of deuteron kinetic energies $T_d = 27$–47 MeV (it may be possible to extend this interval). On this basis, there is in Ref. 12 a preliminary discussion of the possibility of measuring the recoil-deuteron polarization tensor in elastic ed scattering in the region

$15 \leq Q^2 \leq 30$ F^{-2}. An immediate result of such an experiment would be the determination of the presence of a first zero in the deuteron charge form factor (see Secs. 3, 7, and 8) and (if it exists) precise localization of the position of this zero. Note that the existence of this zero is predicted by all potential models, whereas calculations of the deuteron charge form factor with allowance for an admixture (with probability β^2_{6q}) of a pure six-quark ($6q$) state indicate the absence of such a zero in the entire region of Q^2 for $\beta^2_{6q} \gtrsim 1\%$.

2. NONRELATISTIC IMPULSE APPROXIMATION

In the impulse approximation and the nonrelativistic limit the standard parametrization of the electromagnetic form factors of the deuteron in terms of the electromagnetic form factors of the nucleons and the deuteron wave function has the form

$$\left.\begin{array}{l} G_C = 2G^N_{ES}C_E,\ G_Q = 2G^N_{ES}C_Q, \\ G_M = (2G^N_{MS}C_S + G^N_{ES}C_L)\,M_d/M, \end{array}\right\} \quad (3)$$

where the structure functions C_E, C_Q, C_L, and C_S of elastic ed scattering are determined from the S- and D-wave deuteron functions by the well-known quadrature formulas

$$\left.\begin{array}{l} C_E(Q^2) = \int\limits_0^\infty [u^2(r) + w^2(r)]\, j_0(Qr/2)\, dr; \\ C_Q(Q^2) = \frac{3}{\sqrt{2}\,\eta} \int\limits_0^\infty w(r) \left[u(r) - \frac{w(r)}{\sqrt{8}} \right] j_2(Qr/2)\, dr; \\ C_L(Q^2) = \frac{3}{2} \int\limits_0^\infty w^2(r)\, [j_0(Qr/2) + j_2(Qr/2)]\, dr; \\ C_S(Q^2) = \int\limits_0^\infty \left\{ \left[u^2(r) - \frac{w^2(r)}{2} \right] j_0(Qr/2) \right. \\ \left. + \frac{1}{\sqrt{2}}\, w(r) \left[u(r) + \frac{w(r)}{\sqrt{2}} \right] j_2(Qr/2) \right\} dr. \end{array}\right\} \quad (4)$$

In (3) and (4), G^N_{ES} and G^N_{MS} are the charge and magnetic isoscalar nucleon form factors, $j_n(z)$ are spherical Bessel functions, M is the nucleon mass, and M_d is the deuteron mass.

In the static limit, we obtain from the relations (3)

$$\left.\begin{array}{l} G_C(0) = 1;\ G_Q(0) = M_d^2 Q_d;\ G_M(0) \\ = \frac{M_d}{M} \left[\mu_p + \mu_n - \frac{3}{2} P_d \left(\mu_p + \mu_n - \frac{1}{2} \right) \right], \end{array}\right\} \quad (5)$$

where $P_d = \int_0^\infty w^2(r)dr$ is the probability that the deuteron is in the D state, and μ_N is the nucleon magnetic moment.

A detailed discussion of the characteristics and comparative advantages of the various nucleon–nucleon potentials can be found, for example, in Refs. 13–17, and also in the original literature of recent years.[18–22] Without making any claim to completeness of the references (in view of the extent of the corresponding literature), we mention only some of the papers[16,23,24] in which the deuteron form factors have been calculated on the basis of the relations (3) and (4) for different nucleon–nucleon potentials.

The calculations made permit the following conclusions to be drawn. First, in the region of small Q^2 ($\lesssim 10$ F^{-2}) the results agree reasonably with the experimental data. At larger Q^2, the theoretical predictions for $A(Q^2)$ lie systematically below the experimental data, the discrepancy between theory and experiment increasing all the time with increasing Q^2; it is almost an order of magnitude at $Q^2 \simeq 154.1$ F^{-2} (see, for example, Refs. 2, 23, and 24). Second, at $Q^2 \gtrsim 25$ F^{-2} the change in the function $A(Q^2)$ due to admissible variation of the electromagnetic nucleon form factors, which are not sufficiently well measured in this region (especially the neutron charge form factor), has the same order as the corresponding changes due to the model choice of the potential. Therefore, for example, for a sufficiently exotic approximation of the neutron form factor in the region $Q^2 \gtrsim 10$ F^{-2} the agreement between theory and experiment can be greatly improved.[8] However, such a choice appears to contradict the results of determination of the neutron charge form factor from experiments on pion electroproduction on the proton.[25]

Thus, at large Q^2 the results of calculations of the form factors in the framework of the nonrelativistic impulse approximation do not agree with the experimental data for any potential model (the relatively good agreement between theory and experiment when the Feshbach–Lomon potential is used[2] must be regarded as fortuitous, since this potential is unrealistic[26] in the deuteron channel). Analogous calculations by two of the present authors[27] on the basis of a model-free dispersion approach (see Sec. 5) confirm this conclusion. The elimination of the existing discrepancy between theory and experiment—including the discrepancy for the static characteristics—must be attributed to the need to take into account fully the relativistic corrections, the contribution of the meson exchange currents, and the $6q$ configurations in the deuteron (see Secs. 4, 6, and 7, respectively).

We note that as regards the choice of the most realistic potential for the two-body nucleon–nucleon interaction the most informative investigations are those of nucleon–deuteron scattering and calculations of the characteristics of three-particle systems and nuclear matter. This is so because these characteristics are very sensitive to the behavior of the t matrix of nucleon–nucleon scattering off the mass shell, and the nearest singularity of the off-shell t matrix is the deuteron pole. Such investigations have been made intensively in recent years (see, for example, Refs. 28–33) with a view to determining the D-wave admixture P_d in the deuteron and determining more accurately the value ρ_d of the asymptotic D/S ratio of the deuteron wave functions. In particular, when allowance is made for the latest experimental results for ρ_d extracted from data on the measurement of the tensor component $T_{22}(\theta)$ of the process of elastic dp scattering, the following value[37] is obtained by averaging over the results of Refs. 34–37: $\rho_d = 0.0264 \pm 0.0003$.

3. RELATIVISTIC IMPULSE APPROXIMATION

The results so far obtained offer hope of obtaining some sort of accurate description of the relativistic structure of composite hadronic systems only for the special case of weakly bound composite systems.[38–41]

From this point of view, the ideal object for testing the various theoretical approaches is, as before, the deuteron.

The relativistic methods of describing the electromagnetic form factors of the deuteron proposed by different authors on the basis of physical assumptions and results fit into a single general scheme, but they can be nominally divided into the following groups in accordance with the formulation of the initial propositions:

a) calculations in the framework of the relativistic impulse approximation, with approximation of the exact solution of the Bethe–Salpeter equation for the *npd* vertex function[42];

b) calculations in the relativistic impulse approximation with exact solution of the approximate—with respect to the Bethe–Salpeter equation–quasipotential equation to take into account the *npd* vertex function[43];

c) allowance for relativistic recoil effects by means of semiphenomenological constructions that admit interpretation in the framework of the quasipotential approach[44];

d) calculations in the relativistic impulse approximation based on a phenomenological relativization (consistent with ideas about the composite quark structure of hadrons) of the nonrelativistic expressions for the *npd* vertex function[45];

e) disperion calculations in the relativistic impulse approximation.[27,46,47]

Following the general scheme, and considering in the relativistic impulse approximation the deuteron form factors $G^{\mu}(q^2)$ defined with respect to the matrix element of the electromagnetic current of the deuteron by[48]

$$G^{\mu}(q^2)=(2p^0 2p'^0)^{-1/2}\langle p',\ \xi'\,|j^{\mu}|\,p,\ \xi\rangle$$
$$=-e\Big\{G_C(Q^2)(\xi'^*\xi)(p+p')^{\mu}+G_Q(Q^2)[\xi^{\mu}(\xi'q)-\xi'(\xi q)]-G_M(Q^2)\frac{(\xi q)(\xi'^*q)}{2M^2}(p+p')^{\mu}\Big\}, \tag{6}$$

we can write down (Fig. 3) an integral representation of the form[49]

$$G^{\mu}(q^2)=\int\frac{d^4k}{(2\pi)^4}\mathrm{Tr}\{\overline{\Lambda}^{\beta}(p-k,\ k+q)[\tilde{k}-\tilde{q}+M]\times F^{\mu}(k,\ k+q)[\tilde{k}+M]\Lambda^{\alpha}(k,\ p-k)[\tilde{k}-\tilde{p}+M]\}\times\frac{1}{(k+q)^2-M^2}\frac{1}{k^2-M^2}\frac{1}{(k-p)^2-M^2}\xi_{\alpha}\xi_{\beta}'^*, \tag{7}$$

where ξ and ξ' are the deuteron polarization vectors in the initial and final states, $\tilde{k}=k^{\gamma}k_{\gamma}$, $\Lambda^{\alpha}(p_1,p_2)$ is the Bethe–Salpeter *npd* vertex function, and $F^{\mu}(k_1,k_2)$ is the off-shell isoscalar nucleon form factor.

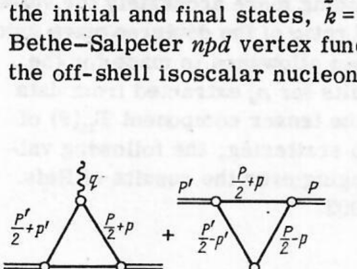

FIG. 3. Sum of diagrams determining the matrix element of the electromagnetic current of the deuteron in the impulse approximation.

For convenience and without loss of generality, we can go over to the frame of reference[50]

$$\left.\begin{array}{l}p=\left(d+\frac{M_d^2}{4d},\mathbf{O}_{\perp},d-\frac{M_d^2}{4d}\right),\quad q=\left(\frac{q_{\perp}^2}{4d},\ \mathbf{q}_{\perp},\ -\frac{q_{\perp}^2}{4d}\right);\\ k=\left(xd+\frac{k^2+k_{\perp}^2}{4xd},\ \mathbf{k}_{\perp},\ xd-\frac{k^2+k_{\perp}^2}{4xd}\right).\end{array}\right\} \tag{8}$$

In the new independent variables of integration $k_{\perp}$, k^2, and x, the propagators in the representation (7) have the form

$$[(p-k)^2-M^2]^{-1}=\left[-k^2\frac{(1-x)}{x}-\frac{k_{\perp}^2}{x}+(1-x)M_d^2-M^2\right]^{-1}; \tag{9a}$$

$$[(k+q)^2-M^2]^{-1}=[k^2-(1-x)q_{\perp}^2-2k_{\perp}q_{\perp}-M^2]^{-1}. \tag{9b}$$

Using the Feynman prescription for avoiding the pole singularities of the propagators, we can readily see that for $x\notin(0,1)$ the pole singularities of all the propagators with respect to k^2 (and also the admissible branch-point singularities of the vertex functions at the thresholds of the inelastic channels) lie in the lower half-plane. Therefore, only the integration over the region $0<x<1$ makes a nonvanishing contribution to the expression (7). In this region, the pole singularity of the propagator (9a) goes over to the upper half-plane, so that, calculating the integral (7) by means of the residues, we obtain

$$G^{\mu}(q^2)=-\frac{i}{(2\pi)^3}\int_{-\infty}^{\infty}\int_0^1\frac{d^2k_{\perp}\,dx}{2(1-x)}\times\Big\{\frac{1}{(k+q)^2-M^2}\frac{1}{k^2-M^2}\mathrm{Tr}\,[\overline{\Gamma}^{\beta}(k+q)(\tilde{k}-\tilde{q}+M)F^{\mu}(k,\ k+q)\times(\tilde{k}+M)\Gamma^{\alpha}(k)(\tilde{k}-\tilde{p}+M)]\Big\}_{(p-k)^2=M^2}\xi_{\alpha}\xi_{\beta}'^*+C^{\mu}_{\mathrm{inel}}, \tag{10}$$

where, following Ref. 46, we have introduced the notation $\Gamma^{\alpha}(p_1)$ for the half-off-shell *npd* vertex function, which is determined from $\Lambda^{\alpha}(p_1,p_2)$ by the substitution $p_2^2=M^2$; C_{inel} denotes the contribution of the admissible cuts of the vertex function at the thresholds of the inelastic channels.

At the subtraction point, the propagator factors in (10) have the form

$$(k^2-M^2)^{-1}=(1-x)[x(1-x)M_d^2-k_{\perp}^2-M^2]^{-1}; \tag{11a}$$
$$((k+q)^2-M^2)^{-1}=\left[xM_d^2-\frac{k_{\perp}^2+M^2}{1-x}-(1-x)q_{\perp}^2-2k_{\perp}q_{\perp}\right]^{-1}. \tag{11b}$$

Hence, by virtue of the assumed smallness of ε/M_d (where $\varepsilon=2M-M_d$ is the deuteron binding energy) it is easy to see that in the nonrelativistic limit, i.e., for $k_{\perp}\ll M_d$, $x\approx\frac{1}{2}$, the propagator terms $(k^2-M^2)^{-1}$ and $((k+q)^2-M^2)^{-1}$ have, respectively, the orders $O(M_d/\varepsilon)$ and $O(\max(M_d/\varepsilon, M^2/q_{\perp}^2))$, which ensures a dominant contribution of the nonrelativistic impulse approximation. The relative smallness of C^{μ}_{inel} is due to the fact that for the vertex function the nearest singularities of branch-point type in the inelastic channels are at $k^2\approx(M+m_{\pi}^2)$, $(k+q)^2\approx(M+m_{\pi})^2$. The general representation in the deuteron rest frame for $\Gamma^{\alpha}(p)$ in terms of the complete set of scalar functions of the invariant variables has the form

$$\Gamma^{\alpha}(p)=F(p^2)\gamma^{\alpha}-\frac{G(p^2)}{M_d}p^{\alpha}-\frac{M_d-\tilde{p}}{M_d}\left[H(p^2)\gamma^{\alpha}-\frac{I(p^2)}{M_d}p^{\alpha}\right]. \tag{12}$$

The concrete way in which the scalar functions occurring in this expansion are found distinguishes the different approaches to the calculation of the electromagnetic form factors of the deuteron in the relativistic impulse approximation with an on-shell spectator nucleon.

In the nonrelativistic approximation $\Gamma^{\alpha}(p)$ is completely determined by the functions $F(p^2)$ and $G(p^2)$, which are then related to the nonrelativistic deuteron wave function in the impulse approximation by the substitution

$$\frac{F(k^2)}{k^2-M^2}\to u(k)-\frac{w(k)}{\sqrt{2}},\quad \frac{G(k^2)}{k^2-M^2}\to\frac{3M_d^2}{k}\,\frac{w(k)}{\sqrt{2}}. \tag{13}$$

This connection can serve as the source of fairly reasonable approximations in the calculation of the relativistic deuteron form factors. Thus, in Ref. 45 Fernandez-Pacheco *et al.* proposed for the deuteron wave function phenomenological approximations (of rational form in the variables $k_{\perp}^2$ and x) consistent in the nonrelativistic limit with Hulthén's deuteron wave function and reflecting in the relativistic region the composite quark structure of the deuteron.

The results [with neglect of the D-wave contribution to the function $F(k^2)$ and with three free parameters in the approximation for $u(k)$ and $w(k)$] are shown in Figs. 1 and 2. It can be seen that there is satisfactory agreement between the theory and experiment [especially for the function $B(Q^2)$]. The systematic excess of the theoretical predictions with respect to the experimental data for $A(Q^2)$ in the region $10\le Q^2\le 60$ F^{-2}, which leaves no room for possible corrections for meson exchange currents, may be due to the fact that: 1) in the calculations, the D-wave contribution to the function $F(k^2)$ was ignored; 2) no investigation was made of the sensitivity of the results to the choice of approximations of the nucleon form factors more realistic than dipole fitting; 3) insufficiently realistic approximations were chosen for the deuteron wave function.

An alternative and more consistent way of finding the vertex function $\Gamma^{\alpha}(k)$ is based on the use of the Bethe–Salpeter equation (an excellent exposition of questions relating to the Bethe–Salpeter equation can be found in Ref. 51; see also Refs. 15 and 52) for the two-particle T matrix of the NN interaction (Fig. 4). In the impulse approximation, with approximation of the contribution of the irreducible diagrams V by the total contribution of the one-boson exchange diagrams, the Bethe–Salpeter equation simplifies somewhat, but its quantitative solution is still technically difficult to realize.[15,52]

Furthest developed is the technique of solving the quasipotentials obtained by replacing the two-particle propagator G by approximate propagators of, for example, the form[15]

$$g(P,k)=\frac{1}{(2\pi)^3}\int\limits_{4M^2}^{\infty}\frac{ds'}{s'-w-i0}f(s',w)$$
$$\times\,\delta^{(+)}\left[\left(\frac{P}{2}+k\right)^2-M^2\right]\delta^{(+)}\left[\left(P'-\frac{P}{2}-k\right)^2-M^2\right], \tag{14}$$

where $P'^{\mu}=(\sqrt{s'/w})P^{\mu}$, and from the requirement of fulfillment of the condition of two-particle unitarity in the elastic channel a unique restriction is imposed on

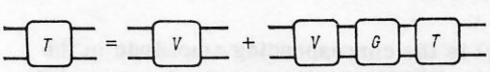

FIG. 4. Bethe–Salpeter equation for the T matrix of two-particle scattering.

the, in general arbitrary, function $f(s',w)$: $f(w,w)=1$. Choosing $f(s',w)=(\sqrt{w}+\sqrt{s'})/2\sqrt{w}$ or $f(s',w)=(\sqrt{w}+\sqrt{s'})/2\sqrt{s'}$, respectively, we arrive at the equations of Gross[53-55] and Kadyshevskii.[56,60]

Various papers (see Refs. 53–55, 57, and 58) have been devoted to the calculation of the vertex function $\Gamma^{\alpha}(k)$ on the basis of Gross's quasipotential equation with approximation of V by the sum of diagrams of π-, σ-, ρ-, and ω-meson exchange. The subject of investigation in the quoted papers is the four-component deuteron wave function $\Psi_d^{\mu}(k,p)$, which is determined from the vertex function $\Gamma^{\alpha}(k)$ in such a way that in the nonrelativistic limit the two nonvanishing components are equal to the nonrelativistic S and D wave functions, and the two other components (classified on the basis of the transformation properties as $^1P_1$ and $^3P_1$ wave functions) make a nonvanishing contribution to $\Gamma^{\alpha}(k)$ only in the relativistic region.[43,59] Using a simple analytic approximation of the relativistic deuteron wave function constructed in this manner,[58] corresponding calculations of the relativistic deuteron form factors were made in Ref. 8. For the most interesting[52] case of pseudovector coupling at the πNN vertex the results obtained in Ref. 8 are shown in Figs. 1 and 2. Analysis of them yields the following conclusions. First, the complete allowance for the relativistic corrections made in this way systematically worsens the agreement between theory and experiment. Second, the change in the results due to the arbitrariness in the choice of the model effective potential and the approximation of the nucleon form factors "masks" the sensitivity of the results to allowance for the relativistic corrections in the complete region of Q^2. Third, as is noted in Ref. 8, even approximate allowance for the fact that the spectator nucleon is off the mass shell shows that the corresponding corrections to the deuteron form factors are in order of magnitude equal to the crossed interference contributions of the small components of the relativistic deuteron wave function.

A more general description of the electromagnetic form factors of the deuteron with the npd vertex function found by solving (by means of Padé approximants) the Bethe–Salpeter equation is developed in Ref. 42. The kernel of the equation is approximated by the contribution of the diagrams of π-, η-, ε-, δ-, ρ-, and ω-meson exchange.[52] The results of the relativistic calculation are closest to the corresponding results obtained in the nonrelativistic description with the Reid potential. This may be an indication that the relativistic effects are not taken into account sufficiently correctly in the framework of Gross's formalism.

4. DISPERSION APPROACH TO THE CALCULATION OF THE ELECTROMAGNETIC STRUCTURE OF THE DEUTERON

Attempts at direct application of dispersion methods[61-64] to describe the structure of the deuteron do not give quantitative agreement with the experiments for $Q^2\gtrsim 6$ F^{-2} and hinder theoretical analysis because of the "mixing" of the internal structure of the interacting particles and the nucleon current in the total deuteron current.

The new dispersion approach presented in this section to the investigation of the electromagnetic structure of the deuteron takes into account explicitly the composite nature of the system and has the following advantages: a) it describes the structure of the composite system directly in terms of experimentally "observable" quantities—the scattering phase shifts and the single-particle form factors of the particles constituting the system; b) it admits a unified treatment in both the relativistic and nonrelativistic cases; c) it does not contain parameters fitted to the experimental data; d) it makes it possible to obtain a correct formulation of the potential-free inverse scattering problem, i.e., the problem of constructing the two-particle wave function directly from the scattering data; e) it gives a satisfactory description of the experimental data on elastic *ed* scattering in the complete range of momentum transfers; f) it makes it possible to describe in a self-consistent way the processes of elastic and quasielastic electron–deuteron scattering; g) it is consistent with the analytic properties of the deuteron form factor that follow from the Mandelstam representation. The basis of the method was developed by Shirokov in an investigation of the general question of connection between fields and particles.[90] In Ref. 1, the basic assumptions of the method are formulated and the results of some calculations are given. The approach was developed further in Refs. 27 and 65–83.

In the derivation of the integral representation that determines the electromagnetic form factors of the composite two-particle system in terms of the phase shifts and the single-particle form factors of the particles forming the system one uses the mathematical formalism of the solution of a Riemann boundary-value problem,[88,89] which is well known in quantum field theory (see, for example, Refs. 84–87). Omitting the details of the derivation (see Ref. 1), we give the final representation, obtained, in particular, for the deuteron charge form factor:

$$G_C(Q^2)=\Gamma^2\int_{4M^2}^{\infty}\int_{s_1}^{s_2}\frac{\Delta(s)\,f(s,\ t,\ s')\,\Delta(s')\,ds\,ds'}{(s-M_d^2)\,(s'-M_d^2)}+R(Q^2);$$

$$s_{1,2}(s,\ t)=2M^2+\frac{1}{2M^2}(2M^2-t)(s-2M^2)\pm$$

$$\pm\frac{1}{2M^2}[(-t)(4M^2-t)\,s\,(s-4M^2)]^{1/2}. \qquad (15)$$

Here, $f(s,t,s')$ is the multipole S-wave form factor of the triplet channel of *np* scattering corresponding to the multipole parametrization[91] of the matrix element of the electromagnetic current on the noninteracting *np* system.[93] The function $R(Q^2)$ takes into account the contribution of the so-called (see, for example, Refs. 1 and 92) unphysical cuts; $\Delta(s)$ denotes the discontinuity of the Jost function $f(k)$ [$k=\sqrt{s^2-4M^2}$; see Eq. (17) below] across the physical cut. The normalization constant Γ is determined from the condition $G_C(0)=1$.

Taking into account the unphysical cuts by a renormalization of the constant Γ, we used[93,70] the representation (15) to calculate the relativistic and [after transition in (15) to the nonrelativistic limit] nonrelativistic

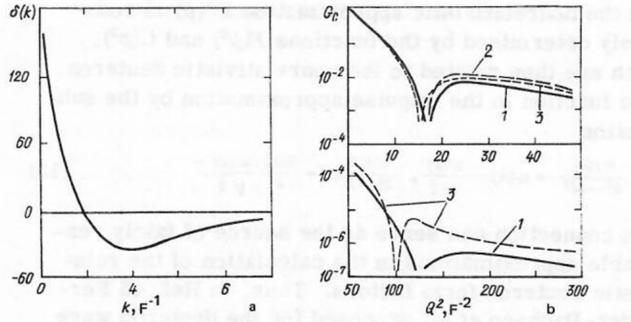

FIG. 5. Choice of the phase-shift functions for the $^3S_1$ wave (a) and the corresponding results for the nonrelativistic (1) and relativistic (2) deuteron form factors compared with the analogous results (3) for the Reid potential with soft core with allowance for the meson exchange currents (b).

deuteron form factors. The results are shown in Fig. 5. In the calculations, we set $G_{En}\equiv 0$ and used the dipole approximation for G_{Ep} and the scaling relations for G_{Mp} and G_{Mn}. The phase-shift function of the $^3S_1$ wave for $k\leqslant 2.2\ \mathrm{F}^{-1}$ was fixed on the basis of data of phase-shift analysis. In the region $4.2\leqslant k\leqslant 11.7\ \mathrm{F}^{-1}$ the phase-shift function is not well determined,[94] and we found it from amplitude Regge-pole fitting of the available experimental data on high-energy nucleon–nucleon scattering.[79] In the intermediate region, the phase-shift function was determined by continuity using the technique of Padé approximants. The possible change in the results from taking into account the D wave was estimated using McGee wave functions, and is not more than 15%.

The complete description of the electromagnetic form factors of the deuteron on the basis of a representation of the type (15) involves the solution of two auxiliary problems, namely, allowance for the contribution of the unphysical cuts and allowance for the tensor nature of the nucleon–nucleon potential. The contribution of the unphysical cuts to the deuteron form factors can be fixed exactly[69,77] on the basis of the following three general requirements: 1) consistency in the nonrelativistic limit of the dispersion description of electromagnetic processes involving the deuteron with the corresponding description in terms of wave functions; 2) the assumption that the amplitude of deuteron electrodisintegration satisfies the Mandelstam representation; 3) the fulfillment of the condition of elastic unitarity by the two-particle amplitude of nucleon–nucleon scattering.

As is shown in Refs. 69, 79, and 83, on the basis of these requirements, for the amplitude $M_{lsj}(k,q)$ of deuteron electrodisintegration[96,97] (in the multipole parametrization) the following integral representation can be obtained:

$$M_{lsj}(k,\ q)=M_{lsj}^B(k,\ q)+2k^{-1}\,\mathrm{Re}\ C_{lsj}(k,\ q)$$

$$-[\pi k f_{lsj}(k)]^{-1}\int_{\infty}^{\infty}\frac{\mathrm{Im}\,f_{lsj}(p)}{p-k-\mathrm{i}0}[pM_{lsj}^B(p,\ q)+2\,\mathrm{Re}\,C_{lsj}(p,\ q)]\,dp, \qquad (16)$$

where $M_{lsj}^B(k,q)$ is the corresponding amplitude in the Born approximation[98–102]; $C_{lsj}(k,q)$ is the function of the contribution of the unphysical cuts, and $f_{lsj}(k)$ is the Jost function of the uncoupled $^{2S+1}l_j$ channel determined from the S-matrix element by the quadrature formula[95, 103]

$$f_{lsj}(k)=\exp\left\{-\frac{1}{2\pi i}\int\limits_{-\infty}^{\infty}\frac{\ln S_{lsj}(p)}{p-k-i0}\,dp\right\} \qquad (17)$$

(here and in what follows, we omit all multipole indices or some of them in the cases when this does not cause confusion). Thus, to make quantitative calculations on the basis of the representation (16) it remains to determine the function $C_{lsj}(k,q)$. If we assume that the amplitude $M_{lsj}(k,q)$ satisfies the Mandelstam representation, then the answer to the posed question is known[69]: Under the additional requirement of equivalence of the dispersion and Schrödinger descriptions of the process of deuteron electrodisintegration, the function $C_{lsj}(k,q)$ is uniquely determined by a representation of the form

$$C_l(k,\ q)=-i\int\limits_{m_\pi}^{\infty}d\sigma\int\limits_{m_\pi/2}^{\sigma/2}d\rho\,\frac{\rho(\beta,\ \sigma)}{k-i\beta}\times[B_l(k-i\sigma,\ q)-B_l(i\beta-i\sigma,\ q)], \qquad (18)$$

where

$$B_l(k,\ q)=\int\limits_0^{\infty}w_l^{(-)}(kr)\,j_l(qr/2)\,u_0(r)\,dr.$$

Here, $w_l^{(-)}(z)$ is a Bessel–Hankel function, and $\rho(\beta,\sigma)$ is the spectral function of the unphysical cuts. The corresponding analytic expressions for the wave functions of np scattering have the form

$$v_{lsj}(k,\ r)=krj_l(kr)+a_{lsj}(k,\ r)-[\pi f_{lsj}(k)]^{-1}\int\limits_{-\infty}^{\infty}\frac{\mathrm{Im}\,f_{lsj}(p)}{p-k-i0}[prj_l(pr)+a_{lsj}(p,\ r)]\,dp, \qquad (19a)$$

where

$$a(k,\ r)=2\,\mathrm{Re}\int\limits_{m_\pi}^{\infty}d\sigma\int\limits_{m_\pi/2}^{\sigma/2}d\beta\,\frac{\rho(\beta,\ \sigma)}{k-i\beta}\times\{w_l^{(-)}(k-i\sigma)\,r-w_l^{(-)}[(i\beta-i\sigma)\,r]\}. \qquad (19b)$$

The presence of the bound state is taken into account by introducing a simple zero in the Jost function of the corresponding partial-wave channel. At the same time, the wave function of the bound state is found by taking the residue in the representation (19) at the point of the bound state.

Thus, in order to particularize fully the representations (18) and (19), it remains to specify a sufficiently reasonable approximation of the function $\rho(\beta,\sigma)$. To this end, it was proposed in Ref. 69 to use the normalization condition of the scattering wave functions, determining the free parameters in the chosen (for example, following Refs. 99 and 104) parametrization of the function $\rho(\beta,\sigma)$ by the requirement of minimality with respect to some norm of the "normalization-defect" function

$$\mathrm{Def}(k,\ p)=\int\limits_0^{\infty}v(k,\ r)^*\,v(p,\ r)\,dr-\frac{\pi}{2}\,\delta(k-p). \qquad (20)$$

A more general way of determining $\rho(\beta,\sigma)$ was developed in Ref. 77, in which, for a complete model-free determination of the function $\rho(\beta,\sigma)$, the normalization condition was regarded as a quadratically nonlinear integral equation for $\rho(\beta,\sigma)$. It was shown that on the introduction of a physically motivated regularization the corresponding problem can be properly posed in Tikhonov's sense,[105] and a convergent iterative procedure was constructed to find the solution of the regularized equation.

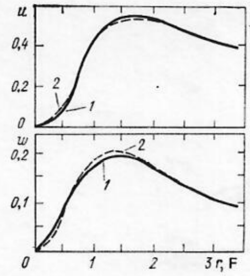

FIG. 6. Deuteron wave functions (1).[76] For comparison, the Reid wave functions with soft core (2) are shown.

A generalization with allowance for functional noncommutativity[106] of the above results to the case of channels coupled with respect to the quantum numbers can be found in Refs. 71, 76, 79, and 83. The results then obtained for the deuteron wave function (with realistic allowance for the D-wave admixture and complete allowance for the unphysical cuts) are shown (following Ref. 76) in Fig. 6. The deuteron form factors calculated on the basis of the constructed deuteron wave function, the longitudinal and transverse parts of the differential cross section of elastic ed scattering, and the tensor polarization are shown in Figs. 7–9.

Since the constructed deuteron wave function is completely determined by the np scattering data, it is of particular interest to investigate the sensitivity of the results to the choice of the scattering data in the high-energy region, in which there are no reliable data of phase-shift analysis. Corresponding calculations were made in Ref. 72 for the nonrelativistic charge form factor of the deuteron and led to the following conclusions: a) variation of the scattering data at high energies significantly changes the behavior of the deuteron wave function at small distances and the behavior of the charge form factor at large Q^2; b) the sensitivity of the results to the method used to take into account the unphysical cuts is much weaker than the dependence on the choice of the phase shifts; c) a stronger lowering to the region of negative values of the phase-shift func-

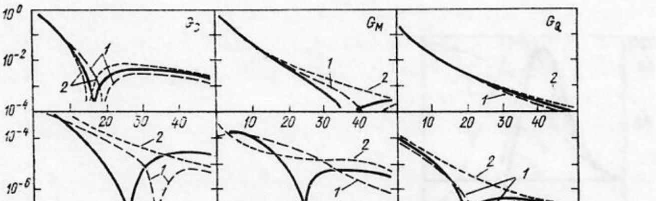

FIG. 7. Dependence for the nonrelativistic charge, magnetic, and quadrupole form factors of the deuteron. The continuous curves are our results[27]: 1) results of calculation for the Reid potential with soft core; 2) the same with allowance for the meson exchange currents.

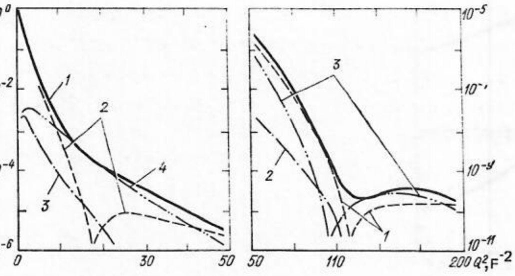

FIG. 8. Relative contribution to the function $A(Q^2)$ (curve 1) of the charge 2), magnetic 3), and quadrupole (4) form factors of this deuteron.

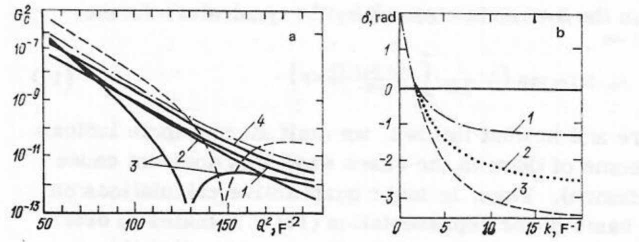

FIG. 10. Nonrelativistic deuteron charge form factor (a, curves 1–3) for different choices (b) of the phase-shift function of the $^3S_1$ wave in the region of high energies (curve 4 is the result of calculation with the McGee wave functions.

tion at high energies (corresponding to a stronger NN repulsion in the region of the core) leads to a shift in the positions of the "dips" of the charge form factor to lower Q^2 (Fig. 10).

Thus, in the framework of the developed dispersion approach one can hope to obtain a correct description of the relativistic electron form factors of the deuteron even with a fairly simple relativization consisting of relativistic allowance for the kinematic cut and allowance for the contribution of the unphysical cuts by a renormalization. The good agreement between the results that we obtained[27] in this way and the latest indirect experimental data on elastic ed scattering[7] encourages our hopes of subsequent construction of a consistent theory of the deuteron with complete allowance for the relativistic unphysical cuts.

5. THE PROBLEM OF THE MESON EXCHANGE CURRENTS

In the calculation of the electromagnetic characteristics of few-nucleon systems and, in particular, the deuteron, the need to include in the treatment corrections for the meson exchange currents is dictated by the fact that corresponding calculations with realistic NN potentials in the impulse approximation lead to violation of the gauge invariance of the theory. These considerations were the basis for the construction in the lowest orders of perturbation theory of a systematic nonrelativistic theory of meson exchange currents,[107] a detailed exposition of which can be found in the reviews of Refs. 108–111, the conference proceedings of Refs. 112 and 113, and the excellent collection of Ref. 114.

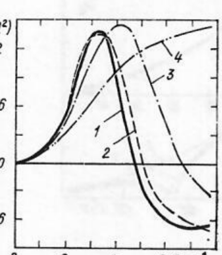

FIG. 9. Tensor polarization in elastic ed scattering: 1) our results[81]; 2) data for the Reid potential with soft core; 3) Hamada–Johnston data; 4) data for the Hulthén potential with soft core.

In the case of the deuteron, as noted, for example, in Ref. 23, the existing arbitrariness in the allowance for the meson exchange currents means that these corrections can always be introduced in such a way that the observed characteristics of electromagnetic processes with the participation of the deuteron are not changed (so-called minimal allowance for the meson exchange currents). From the point of view of the analysis of the experimental data on processes with the participation of few-nucleon systems the need to introduce corrections for the meson exchange currents in a nonminimal way is motivated by the following facts: 1) the ratios of the theoretically predicted values to the corresponding experimental data are 93.5% for the magnetic moments $\mu(^3H)$ and $\mu(^3He)$, 90.5% for the total cross section of thermal np capture,[115] and 95% for the Gamow–Teller β-decay matrix element for the ^3H nucleus[116]; 2) in the impulse approximation, the differential cross section of elastic ed scattering can be satisfactorily described only in the region $Q^2 \leqslant 35\ F^{-2}$. For $Q^2 \geqslant 35\ F^{-2}$, the discrepancy between theory and experiment increases with increasing Q^2, being almost an order of magnitude at $Q^2 \simeq 154.1\ F^{-2}$.[2,8] Allowance for the relativistic corrections in powers of Q^2/M^2 in the framework of the quasipotential approach does not improve the situation.[8] The relative errors in the description of the static characteristics of the magnetic and quadrupole moments change approximately in inverse proportion to each other to within 2% when different nucleon–nucleon potentials are used[15]; 3) the experimentally measured differential cross section of backward deuteron electrodisintegration near the threshold exceeds the corresponding theoretical prediction in the impulse approximation by 25% at $Q^2 = 1\ F^{-2}$ and by an order of magnitude for $Q^2 = 10\ F^{-2}$.[96]

Some of the diagrams due to the meson exchange currents in the lowest order of perturbation theory (Fig. 11) can be taken into account in the framework of

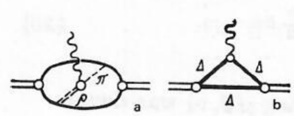

FIG. 11. The $\rho\pi\gamma$ diagram determining the meson-exchange-current contribution to the deuteron form factors in the lowest order of perturbation theory (a) and a typical diagram of the isobar contribution in the intermediate state to the matrix element of the electromagnetic current of the deuteron (b).

the quasipotential approach.[8,15] However, because of the nonuniqueness in the choice of the quasipotential equation, the relative contribution of the meson exchange currents to the characteristics of the above processes will depend on the choice of the particular modification of the quasipotential equation. One of the ways of overcoming this difficulty is to use data on weak-interaction processes in few-nucleon systems.[117]

Suppression of the relative contribution of the meson exchange currents of second order is proved by direct quantitative calculations.[118] The relative contribution to the deuteron form factors of meson exchange currents of even higher orders is suppressed quantitatively in the region $Q^2 \lesssim M^2$ by multiplicative factors $(q^2/M^2)^l$ with $l \geqslant 1$, and in the region $Q^2 \gtrsim M^2$, where the effects of the quark structure of nucleons are manifested, it is suppressed[119] in accordance with the Okubo–Zweig–Iizuka rule.[120]

Qualitatively, the suppression of the relative contribution of the first-order meson exchange currents with intermediate states of higher masses can be readily interpreted in the language of the dispersion methods of Ref. 79. Thus, it is natural to expect that the main contribution to the differential cross section of deuteron electrodisintegration will be made by the meson exchange current due to the $\pi\pi$ current (see Fig. 11). The analogous contributions from the $\omega\pi$, $\rho\eta$, and $\rho\rho$ currents characterize the resonance contributions of the many-pion currents. The Feynman diagrams corresponding to these currents have a higher order than the $\pi\pi$-current diagram. It is well known[121] that with increasing order of a diagram its singularities with respect to all the invariant variables become less pronounced. Therefore, in the phenomenological allowance for the meson exchange currents the contributions of the $\omega\pi$, $\rho\eta$, and $\rho\rho$ currents, and also the higher-mass currents enter in a "regular manner" and in a finite region with respect to Q^2 make a relatively small contribution to the deuteron form factors, which can be taken into account effectively by a small variation of the meson–nucleon form factors that determine the matrix element of the $\pi\pi$ current. For the same reasons, the contribution of the $\rho\pi$ current to the differential cross section of elastic *ed* scattering[122] and of the $\pi\pi$ current to the total cross section of thermal *np* capture[123] are dominant.

The calculations so far made to take into account the corrections for the meson exchange currents for elastic *ed* scattering[124] and also for deuteron electrodisintegration[96,125] and thermal *np* capture[123] do indeed lead to better agreement between theory and experiment. However, the problem of the correct allowance for the meson exchange currents is still, when considered as a whole, in a confused state. This is due to the following circumstances.

First, the currently available approaches have a strongly phenomenological nature and do not possess predictive power—the appearance of new experimental data in the region of large Q^2 constantly leads to the need to review the already existing methods of taking into account the contribution of the meson exchange currents.[2] The contributions of the first-order meson exchange currents of the lowest-mass currents are calculated on the basis of the hypothesis of a partially conserved axial current (PCAC) and the use of current algebra. Because of this, the inaccuracy in the determination of these currents for $Q^2 \neq 0$ cannot be estimated. It is even harder to estimate the accuracy with which allowance is made for the corrections for the first-order meson exchange currents of higher-mass currents, which make a nonvanishing contribution even to the charge form factor of the deuteron.[118] In addition, the meson–nucleon form factors that determine the contributions of the meson exchange currents are not sufficiently well known for $Q^2 \neq 0$ and are poorly determined even in the static limit $Q^2 = 0$.

Second, the correct description of the dynamics of the nucleon–nucleon interaction and complete allowance for the relativistic corrections lead to a significant improvement in the agreement between theory and experiment even without the introduction of meson-exchange-current corrections—for elastic *ed* scattering this was shown in Refs. 27, 45, and 70, and for deuteron electrodisintegration in Refs. 92 and 101. This fact can be regarded as an indication that in the construction of a self-consistent procedure for taking into account meson exchange currents there may for some processes be a mutual compensation of the contributions of the different meson-exchange-current diagrams. With regard, for example, to the estimate of the influence of the axial meson exchange currents on the doublet rate of capture in the reaction $\mu^- + d \to 2n + \nu_\mu$, this last conclusion received exact quantitative confirmation in Ref. 126.

Third, the calculations in the literature of concrete meson-exchange-current diagrams have been made in the nonrelativistic formalism,[8,131] and this is a serious defect of the existing approaches. It would be very desirable to make a relativistic calculation of the meson exchange currents, since in the region in which it is necessary to "unfreeze" the meson degrees of freedom (i.e., at $Q^2 \approx M^2$) the contribution of the relativistic effects may be appreciable.

The problem of taking into account an admixture of isobars in the deuteron is related to the problem of taking into account the meson exchange currents and is in as confused a state. Experimentally, the upper limit for the relative contribution of the isobar admixtures in the deuteron is not well determined.[125] The corresponding estimates vary in a wide range, depending on the choice of the phenomenological *NN* potential. Even for the most acceptable one-boson exchange potentials (OBEP), the estimates do not inspire much confidence. This is due to the fact that because of the duality relations[127] the contribution of the additive corrections to the OBEP determined, for example, by the potential of 2π-meson exchange with a contribution of the isobars in the intermediate state (see Fig. 11) is partly effectively imitated by the OBEP parameters.

6. DEUTERON AS A SIX-QUARK SYSTEM

A discussion of the possible part played by quark degrees of freedom in nuclei can be found, for example, in Refs. 128–167. We note, first, that, in our opinion, the present experimental status of elastic *ed* scattering is such that it does not yet necessarily call for inclusion of quark degrees of freedom for its description. Nevertheless, it is to be expected that at large Q^2 (how large is not yet clear[8,168]) some variant of a quark approach to the deuteron will become unavoidable.

In speaking of the behavior of the hadron form factors at large Q^2, we must above all mention the pioneering papers of Ref. 133, in which the rules of dimensional quark counting lead to the following asymptotic behavior for the spin-averaged form factor of any hadron h:

$$F_h(Q^2) \sim (Q^2)^{1-n},\ Q^2 \to \infty, \tag{21}$$

where n is the number of elementary constituents of hadron h (see also Refs. 174 and 175). It is generally assumed that in accordance with (21) the results of measurement of the function $A(Q^2)$ (Refs. 2 and 7) in the region $0.8 \leqslant Q^2 \leqslant 8\ (\mathrm{GeV}/c)^2$ reveal a transition to an asymptotic behavior of the type (21): $F_d(Q^2) \sim (Q^2)^{-2\pm0.5}$ (see the discussion in Refs. 2, 7, 135, 137, and 139).

An elegant analysis of the $6q$ state in the bag model is made in Ref. 139. This takes the simplest case of a static spherically symmetric cavity in which all six quarks are in the same energy level with $J^P = 1/2^+$. It was found that in the analysis of the possible dibaryonic states of the $6q$ bag it is very important to take into account the color degree of freedom, which was first introduced in the pioneering studies of Ref. 169. The upshot established in Ref. 139 was that the contributions of different baryonic compositions of the $6q$ bag with the deuteron quantum numbers are about 11% for the proton–neutron state, about 9% for the $\Delta\Delta$ pair, and about 80% for the baryon pair B_cB_c with hidden color. In other words, in the simplest model the contents of the deuteronlike $6q$ bag have little in common with the real deuteron. It is possible that further development of the theory of quark bags will "bring" the $6q$ state closer to the real deuteron. An analysis of the various spectroscopic properties of the deuteronlike $6q$ bag (for example, the "old" question of the admixture of isobar states in the deuteron) has been made in many other studies (Refs. 142, 143, 152, 156, 161, and 164).

At the present time, we must regard as more realistic the intermediate so-called (following the terminology of Ref. 152) hybrid model of the deuteron (see Ref. 139), which is developed more fully in Refs. 140, 149, and 153. In this model, the quarks in a nucleon are bound in a region of space $r \leqslant r_N \sim 1$ F, and with a distance of order $r \sim 1.5-2$ F between the centers of mass of the two three-quark clusters (nucleons) the deuteron is described in terms of nucleon degrees of freedom. At small r, the neutron and proton partly overlap (as one says, the np system "collapses" into a $6q$ system) and then the fundamental variables are the quark fields. Accordingly, the main assumption of the hybrid model is that the deuteron wave function can be represented as a superposition of np and $6q$ states:

$$|d\rangle = \alpha\,|np\rangle + \beta\,|6q\rangle, \tag{22}$$

the wave function of the np system falling off at distances less than the nuclear-core radius r_c, while the $6q$ state makes a vanishingly small contribution when $r > r_c$. The normalization condition for the wave function (22) has the form $\alpha^2 + \beta^2 + 2\alpha\beta I_0 = 1$, where I_0 is the overlap integral of the np and $6q$ states. According to the estimates made in Ref. 140, $I_0 < 0.1\%$. In accordance with (22) the deuteron form factors can be represented as

$$G_i(Q^2) = \alpha^2 G_i^{(np)} + \beta^2 G_i^{(6q)} + 2\alpha\beta G_i^{(\mathrm{int})}(Q^2). \tag{23}$$

Methods of calculating the first term in (23) were discussed above, in Secs. 3–6. The second term requires special treatment in the framework of $6q$ models. The contribution of the last term is very small, and because of this we ignore it.

Before we turn to a discussion of the results obtained in such an approach, we enunciate the two most delicate questions that arise in the analysis of the representation (22). First, is there not a problem of double allowance for the same dynamics in the deuteron in the constructive realization of such a program? Second, from what arguments can one reliably hope to extract the value of β^2? It seems to us that there are no unambiguous or reasonably convincing answers to either of these questions. The first question is discussed, for example, in Refs. 158 and 167 (enthusiasts for the hybrid model simply avoid this question). With regard to the second question, each of the authors of Refs. 140, 144, 149, and 153 develops his own prescription. The results do not agree at all well (β^2 may vary in the interval 0.42 $\leqslant \beta^2 \leqslant 7\%$, while according to the estimates of Ref. 167 the value of β^2 may actually be near zero: $\beta^2 \simeq 0.05\%$). Both these questions need further study.

In Ref. 140, results obtained with the Reid potential with a soft core and Gross-type relativization[8] were used for the deuteron form factors, while the form factor of the $6q$ state for the $(OS)^6$ configuration was calculated in the nonrelativistic three-dimensional oscillator model. The final result is

$$G_C^{6q}(Q^2) = F_{6q}(Q^2),\ G_M^{(6q)}(Q^2) = \mu_{6q} F_{6q}(Q^2),\ G_Q^{6q}(Q^2) \equiv 0; \tag{24}$$

$$F_{6q}(Q^2) = \exp\left(-\frac{5}{24}\frac{Q^2}{\omega}\right), \tag{25}$$

where ω is the parameter of the oscillator potential. In the model of a relativistic oscillator (with the additional assumption of ρ dominance for the deuteron form factors) the nonrelativistic $6q$ form factor (25) is replaced by the relativistic $F_{6q}^R(Q^2)$ in accordance with

$$F_{6q}^R(Q^2) = \left(1 + \frac{Q^2}{m_\rho^2}\right)^{-1}\left(1 + \frac{Q^2}{4M^2}\right)^{-4} F_{6q}\left(Q^2, 1 + \frac{Q^2}{4M^2}\right), \tag{26}$$

where the pentapole factor in (26) is introduced "by hand." The numerical value $\omega = 7.9\ \mathrm{F}^{-2}$ is chosen by fitting of the nucleon form factors. The weight of the $6q$ state was determined from the deviation of the experimental value of the deuteron magnetic moment from the value calculated in accordance with the standard formula:

$$\begin{aligned}\mu_d &= \mu_p + \mu_n - \frac{3}{2}P_d\left(\mu_p + \mu_n - \frac{1}{2}\right)\\ &+ \beta^2\left[\mu_{6q} - \mu_p - \mu_n + \frac{3}{2}P_d\left(\mu_p + \mu_n - \frac{1}{2}\right)\right]\\ &\equiv \mu_{np} + \beta^2(\mu_{6q} - \mu_{np}).\end{aligned} \tag{27}$$

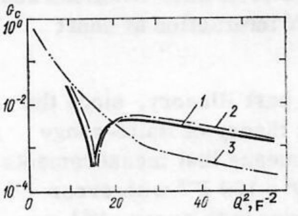

FIG. 12. Deuteron charge form factor: 1) for the Reid potential with soft core and with allowance for meson exchange currents; 2) our results with complete allowance for the relativistic corrections to the deuteron charge form factor; 3) Kobushkin's results with allowance for the admixture of the $6q$ state in the deuteron.[140]

Substituting $P_d = 6.5\%$, and $\mu_{6q} = 1.9\mu_{\text{nuc}}$, we obtain $\beta^2 = 1.4\%$ [note that in Ref. 140 there is evidently an error in the expression for μ_d—the term $\beta^2(3P_d/2)(\mu_p + \mu_n - \frac{1}{2})$ has been omitted, and as a result the value $\beta^2 = 2\%$ was obtained]. The results of calculation of the form factors in accordance with Eqs. (23)–(26) are given in Figs. 1 and 12. It can be seen that for $Q^2 \geqslant 1$ (GeV/c)^2 the contribution of the $6q$ state determines the behavior of $A(Q^2)$ and eliminates[7] the disagreement between theory and experiment that is characteristic of calculations in the framework of potential models. Note also that the calculation of the deuteron charge form factor with neglect of the contribution of the $6q$ state reveals (see Secs. 3–5) oscillations of the form factor, i.e., the existence of zeros of the form factor for $Q^2 > 0$. The first zero, which is stable with respect to the choice of the NN potential, lies in the region $Q^2 \approx 16$–20 F^{-2}. Allowance for a $6q$ admixture in the deuteron results in the disappearance of the zeros. Therefore, polarization experiments at $Q^2 > 20$ F^{-2} and determination of the sign of the deuteron charge form factor will give strong arguments either for or against hybrid quark models.

Finally, we make two remarks about the prescription for the calculation of β^2 adopted in Ref. 140. First, as was shown above, the relativistic effects and the meson exchange currents lead to an appreciable complication of the nonrelativistic formula $\mu_d = \mu_{np}$, and in principle one can "saturate" μ_d without quark degrees of freedom (see, for example, Ref. 170). Second, even if the validity of (27) is accepted without further reservation, one must bear in mind the tendency to systematic lowering of P_d. For example, in contrast to the previously adopted most probable value $P_d = (6$–$7)\%$ for one of the currently most widely used potentials—the potential of the Paris group[21]—$P_d = 5.77\%$, which reduces the weight of the $6q$ state immediately by a factor 1.5: $\beta^2 \approx 0.9\%$. Moreover, modern data on NN scattering do not rule out (and for the description of the deuteron photodisintegration experiments even prefer) the value $P_d = 4\%$,[171] which leaves no room at all for a significant contribution of the $6q$ state.

Results analogous to those of Ref. 140 were obtained in Refs. 149 and 153, in which the model of a relativistic harmonic oscillator was also taken as the basis. The main difference from Ref. 140 is in the different way of determining the parameters of the contribution of the $6q$ state to the deuteron form factors. The expressions obtained in Ref. 153 are

$$\left.\begin{aligned} G_C^{(6q)}(Q^2) &= [(1 + 7Q^2/4m^2)/(1 + Q^2/2m^2)]\, F_{6q}^R(Q^2), \\ G_M^{(6q)}(Q^2) &= 2F_{6q}^R(Q^2),\ G_Q^{(6q)}(Q^2) \equiv 0; \end{aligned}\right\} \quad (28)$$

$$F_{6q}^R(Q^2) = (1 + Q^2/2m^2)^{-5} \exp\left[-\frac{5}{4\alpha_{6q}} \frac{Q^2}{1 + Q^2/2m^2}\right]. \quad (29)$$

The value $\alpha_{6q} = 2\sqrt{2}\,\alpha_{3q} = 1.1$ (GeV/c)^2 is found, as in Ref. 140, from best fitting of the nucleon form factors. The two remaining free parameters θ and m are fixed by the condition of best fitting of $A_{6q}(Q^2)$ to the experimental data for $A(Q^2)$ when $Q^2 \geqslant 1$ (GeV/c)^2. The values found are $m = 1.3$ GeV and $\sin^2\theta = 0.05$ (i.e., the contribution of the $6q$ state is found to be 3.6 times greater than in Ref. 140). The results of the calculations of the functions $A(Q^2)$ and $B(Q^2)$ are close to those obtained in Ref. 140. For comparison, we mention that in the previous paper of the same authors, Ref. 149, in which the deuteron spin was not taken into account, the same parameters were found to be $\alpha_{6q} \approx 1.4$ (GeV/c)^2, $m \approx 1.22$ GeV, $\sin^2\theta = 0.07$. It is interesting to note that the same value for β^2 is also obtained from qualitative arguments[128,139] associated with the possibility of tunneling of the neutron and proton below the barrier of the nucleon–nucleon repulsive core and fusing into a $6q$ system.

In Refs. 144 and 145, the $6q$ state at large Q^2 is regarded as a constructive realization of the idea, put forward some time ago,[172] of the possible existence of fluctuations of nuclear matter in small nuclear volumes (Blokhintsev's so-called flucton). The scheme of the calculations and the results [the representation (22), (23), the use of the model of a relativistic oscillator, the determination of the weight of the $6q$ state in the deuteron, and the filling of the zeros of the deuteron charge form factor] are close to those discussed above. In Ref. 144, there is a more careful analysis of the interference term in (22). It is found that both at comparatively small [$Q^2 \leqslant 1$ (GeV/c)^2] and at larger momentum transfers the contribution of this term to $|G_C(Q^2)|$ can be appreciable. In Ref. 144, the same value for the contribution of the $6q$ state as in Refs. 135, 139, and 149 was obtained: $\beta^2 = 0.07$.

In Refs. 158, 159, and 167, the following results were obtained. Calculations were made of the exchange corrections to the deuteron charge form factor, i.e., the contribution of the $6q$ state to $G_C(Q^2)$ due to antisymmetrization with respect to quarks from different $3q$ clusters in the region of their overlap was determined. These exchange corrections are not related to any particular dynamical mechanism of the interquark interaction. The contribution of the quark exchange corrections is small. Further, in connection with the "double counting" problem mentioned above, the wave functions of the np and $6q$ states in the deuteron were orthogonalized. In contrast to earlier studies, in which only the $(OS)^6$ configuration was investigated, a larger basis of $6q$ states was considered in Refs. 158, 159, and 167. As a result of this spectroscopic investigation it was found that there is a compensation of the contributions of the different states in the complete $6q$ basis. Therefore, if there is no resonant enhancement of any one of the six-quark configurations, the weight of the $6q$ state in the deuteron may be decreased compared with the previous estimates by two orders of magnitude: $\beta^2 \approx 0.05\%$.

Thus, at the present time the uncertainty factor for the contribution of the $6q$ state (like its method of determination) is very large. Intensive investigations are currently being made in this direction. Actually, there is at present a certain contradiction between the desire to "see" quarks in the deuteron already at $Q^2 \approx 1$ (GeV/c)^2 and, for example, the circumstance that at such Q^2 there is already a reasonably good description of $A(Q^2)$ in the framework of nonrelativistic potential models.[173]

7. POLARIZATION IN ELASTIC ELECTRON-DEUTERON SCATTERING

The traditional subject of investigation—the polarization tensor of the recoil deuteron or of the deuteron target in elastic ed scattering with a beam of polarized electrons—has been calculated and discussed in many papers. Thus, in addition to the references given in Ref. 1, from the last seven years we must also add Refs. 8–10, 16, 22, 24, 33, 176–178, 180, and 181. In our view, the most compact and physically transparent among all these studies is still Ref. 182.

Measurement of the components of the polarization tensor is needed for the separate determination of the contributions of the charge and quadrupole form factors to the function $A(Q^2)$. This problem can be solved by measuring the tensor $T(Q^2)$, which is equal to

$$T(Q^2) = \frac{4\sqrt{2}}{3}\eta \frac{G_C(Q^2)G_Q(Q^2) + \frac{1}{3}\eta G_Q^2(Q^2)}{G_C^2(Q^2) + \frac{8}{9}\eta^2 G_Q^2(Q^2)}. \qquad (30)$$

A characteristic graph of $T(Q^2)$ (taken from Ref. 8) is shown in Fig. 9.

The usefulness of the expression (30) is determined by two circumstances, which are revealed in the nonrelativistic formalism. First, we have the relation $T(Q^2) \sim \eta G_Q(Q^2)$, so that, following Refs. 24 and 179, it can be hoped that a measurement of $T(Q^2)$ in the region $4 \leqslant Q^2 \leqslant 9$ F^{-2}, in which $\eta G_Q(Q^2)$ reaches a maximum according to realistic estimates, makes it possible to determine fairly accurately the D-wave admixture P_d in the deuteron. However, it is pointed out in Refs. 10 and 16 that such an analysis may be rendered difficult by inadequate knowledge of other quantities (such as the neutron charge form factor, the meson-exchange-current contribution and the admixture of isobars in the deuteron, and relativistic effects) and an anomalously high accuracy in the measurement of $T(Q^2)$ may be required. Second, it follows from (30) that $T(Q^2)$ depends in practice only on the ratio $x(Q^2) = G_Q(Q^2)/G_G(Q^2)$ of the quadrupole and charge form factors. In the nonrelativistic approximation, it is readily seen that $T(Q^2)$ does not depend on the nucleon form factors and is determined solely by the deuteron wave function. It is on the basis of this that proposals were made in Refs. 33, 178, and 179 to use measurement of $T(Q^2)$ in different regions of Q^2 (in Ref. 33, the region $36 \leqslant Q^2 \leqslant 100$ F^{-2} was investigated; in Ref. 78, the region near the point $Q^2 = 20$ F^{-2}) to choose realistic deuteron wave functions and obtain information about the NN interaction at short distances.

Such optimistic hopes are in part illusory, since the admissibility in nonrelativistic theory of finite-range unitary transformations[10,23,33] means that measurements of $T(Q^2)$ in even a wide range $Q^2 \leqslant 100$ F^{-2} with error ±10% (and in individual cases even with error ±1%) cannot distinguish a large class of unitarily equivalent deuteron wave functions. Besides this, allowance for the relativistic corrections "entangles" the contributions of the isoscalar charge and magnetic form factors of the nucleon in the deuteron form factors (see, for example, Refs. 8 and 68). An even greater arbitrariness in the determination of the deuteron wave function from data on elastic polarization ed scattering arises when meson and baryon degrees of freedom are included. All these questions require further investigation (we note here that the influence of relativistic corrections on $T(Q^2)$ was also considered in Refs. 33 and 183).

Having in mind separate experimental determination of $G_C(Q^2)$ and $G_Q(Q^2)$, it is also helpful to analyze elastic scattering of electrons by a polarized deuteron target. It is well known[33,182] that if for fixed θ_e the deuteron target is polarized at right angles to the scattering plane, then the cross section for scattering by such a target is

$$\left(\frac{d\sigma}{d\Omega_e}\right)_{T_2} = \left(\frac{d\sigma}{d\Omega_e}\right)_{\mathrm{Mott}} \frac{1}{3}\left[G_C^2 + \eta(1+\eta)G_M^2 \operatorname{tg}^2(\theta_e/2)\right]. \qquad (31)$$

Separating in (31) the contribution G_M (or choosing $\theta_e \leqslant 10°$), one can determine $G_C(Q^2)$, and in conjunction with (2) this makes it possible to find $G_Q(Q^2)$.

In Ref. 16, which continues the investigations of the same authors,[24,179] there is a detailed study of the dependence of the various elements [including ones different from (30)] of the polarization tensor T_{ik} of the recoil deuteron in elastic ed scattering on the choice of the NN interaction potential. Seven potential models of the NN interaction with different behavior at short distances and different relative contributions of the tensor forces are considered. However, the important question of the influence of relativistic effects in the impulse approximation on T_{ik} was not, in fact, considered. It would be desirable to fill this gap.

In Refs. 180 and 181, a new important step forward in the study of polarization phenomena in elastic ed scattering was made. It was here proposed to return from the polarization tensor to the polarization vector of the recoil deuteron at the new level made possible by the appearance of beams of polarized electrons. The new result is that in the elastic scattering of longitudinally polarized electrons by unpolarized deuterons the polarization vector p_x of the recoil deuterons is no longer zero. The polarization in the scattering plane in the direction perpendicular to the momentum of the recoil deuteron is the most interesting. In accordance with Ref. 180, this polarization is

$$p_x = -\frac{4}{3F_d^2}\sqrt{\eta(1+\eta)}\, G_M\left(G_C + \frac{1}{3}\eta G_Q\right)\operatorname{tg}(\theta_e/2), \qquad (32)$$

so that measurement of p_x also solves the problem of

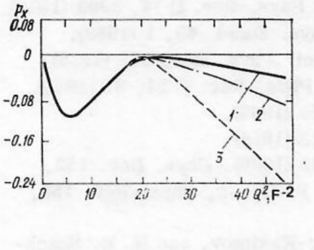

FIG. 13. Vector polarization $p_x(Q^2)$ in elastic *ed* scattering (for $\theta_e = 40°$): 1) for the Reid potential with soft core; 2) the same with Gross-type relativization; 3) data with the Lomon–Feshbach potential for $P_d = 4.5\%$.

separating G_C and G_Q. It is very important for future measurements that $p_x \sim \sqrt{\eta}$. We recall that in the currently attained kinematics $\eta \ll 1$, and all tensor polarizations are proportional to η. In Ref. 180, the polarizations are calculated in the relativistic impulse approximation for several deuteron wave functions (Fig. 13).

One more important result obtained in Refs. 180 and 181 relates to the old problem of measuring $G_{En}(Q^2)$. In the scattering of longitudinally polarized electrons by unpolarized neutrons, the vector polarization of the recoil neutron analogous to (32) is

$$\left. \begin{array}{l} I_0 p_x = -2\sqrt{\tau(1+\tau)}\, G_{Mn} G_{En}\, \mathrm{tg}(\theta_e/2); \\ I_0 = G_{En}^2 + \tau G_{Mn}^2 [1 + 2(1+\tau)\,\mathrm{tg}^2(\theta_e/2)], \end{array} \right\} \qquad (33)$$

where $\tau = Q^2/4M^2$. In other words, $p_x \sim G_{En}$ (and, in addition, $p_x \sim \sqrt{\tau}$). Therefore, measurements using a beam of polarized electrons give an interesting possibility of direct determination of $G_{En}(Q^2)$. The same result was obtained in Refs. 184–186 in a study of deuteron electrodisintegration.

8. CONCLUSIONS

Our analysis in the previous sections of this review of the experimental and theoretical studies demonstrate the breadth of modern investigations of elastic *ed* scattering. We end by listing the main experimental and theoretical problems whose solutions will lead, in our view, to a further development of the theory of the deuteron in terms of both the traditional hadronic (nucleon, meson, baryon) and the quark degrees of freedom, and also the acquisition of new information about the neutron form factors.

The main experimental tasks are:

1. Direct measurements of $A(Q^2)$ for $Q^2 > 4\,(\mathrm{GeV}/c)^2$, extraction of $A(Q^2)$ from the inclusive cross section of *ed* scattering at $Q^2 > 8\,(\mathrm{GeV}/c)^2$, and also more accurate determination in this way of the values of $A(Q^2)$ in the interval $6 \le Q^2 \le 8\,(\mathrm{GeV}/c)^2$ (if counting-rate restrictions do not permit direct measurements of the cross section of elastic *ed* scattering).

2. Measurements of the transverse part of the cross section $B(Q^2)$ for $Q^2 > 1\,(\mathrm{GeV}/c)^2$. The corresponding proposals are discussed in Refs. 8 and 187. In Ref. 8 there is a discussion of the possibility of measuring $B(Q^2)$ with detection in the final state of the scattered electrons at $\theta_e = 180°$ in the indicated region of large Q^2. In Ref. 187, it is proposed to measure the cross section of elastic *ed* scattering with detection in the final state of recoil deuterons emitted forward ($\theta_d = 0°$). Counting-rate restrictions make it possible to measure it up to $Q^2 \le 3\,(\mathrm{GeV}/c)^2$. The main advantages of such a method are the following: a) only one particle is detected in the final state; b) there is no need to construct Rosenbluth graphs, which lowers the accuracy of the results. We note here that the magnetic form factor of ^{3}He has already been measured[188] up to $Q^2 \le 32$ F^{-2}.

3. Measurement of the polarization tensor $T(Q^2)$ of the recoil deuterons at all values of Q^2.

4. Measurement of the polarization vector $p_x(Q^2)$ of the recoil deuterons in the case of a beam of polarized electrons.

In the theory of the deuteron, it is desirable to develop the following directions:

1. The study of the arbitrariness which arises when allowance is made for off-shell effects, and also in the determination of the position and contributions of the unphysical cuts in the relativistic case.

2. Relativization of the calculations of the meson exchange currents.

3. Analysis of the possible part played by quark degrees of freedom, and especially the search for ways of more reliable determination of the weight of the $6q$ state in the deuteron.

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