Relativistic random-phase-approximation response function for quasielastic electron scattering in local density approximation

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The longitudinal and transverse response functions for quasielastic electron scattering on ^{12}C, ^{40}Ca, and ^{48}Ca are calculated in the relativistic $\sigma\omega\rho$ model in random-phase approximation. We use space-dependent baryon densities and effective mass of the nucleons, as obtained from mean-field theory of closed-shell nuclei, in local-density approximation. The longitudinal response function agrees with experimental data quite well, while the transverse response function underestimates data. Our results are also compared with other calculations of the response functions in relativistic models.

INTRODUCTION

Relativistic field theories of nucleons interacting with scalar and vector meson fields have been very successful in describing nuclear matter and closed-shell nuclei.[1] The main features of this type of model are already present in the simplest version, the $\sigma\omega$ model, also called QHD-I in Ref. 1. Large attractive and repulsive potentials arise from the scalar σ and vector ω meson, respectively, and the effective mass of nucleons in nuclear matter is reduced to about one-half of its free value. Basically, we use the $\sigma\omega\rho$ model (called QHD-II in Ref. 1), which includes the vector-isovector ρ meson, and discuss the effect of the ρ meson by comparison with the simpler $\sigma\omega$ model.

Quasielastic electron scattering is a valuable tool for exploring the properties of nucleons in the nuclear environment. Cross section measurements in the domain of the quasielastic peak which do not specify different excitation modes (longitudinal and transverse, respectively) can be very well interpreted in terms of the Fermi motion of independent nucleons inside the nucleus. Only some years ago, however, measurements of the separated longitudinal and transverse response functions have become available, and since then they have been a challenge to any theory of nuclear structure. As yet it has not been possible to explain longitudinal and transverse response simultaneously. Nonrelativistic calculations[2] have indeed shown that a careful treatment of the mean field and of the p-h interactions in the excited states contribute essentially to the q dependence of the separated response functions. Quasielastic scattering consequently qualifies as a sensitive test of the q dependence of residual interactions, in a rather wide q range ($1 < q < 2.5$ fm^{-1}).

Quasielastic electron scattering on symmetric nuclear matter has been investigated in the relativistic random-phase approximation (RPA) in the $\sigma\omega$ model in Ref. 3 and the $\sigma\omega\rho$ model in Ref. 4. In this paper we generalize the work of Refs. 3 and 4 to asymmetric nuclear matter and use this formalism to calculate the response function of closed-shell nuclei in the local-density approximation. This is in contrast to Nishizaki, Kurasawa, and Suzuki,[5] who take the structure of the nucleus fully into account, but calculate its interaction with the photon in the Hartree approximation. There is no free parameter in our calculation, the local effective mass and baryon densities are from mean-field theory as described in Ref. 6, and we use the coupling constants and masses determined there from the saturation properties of nuclear matter and the rms charge radius of ^{40}Ca.

We compare our results with data from Saclay[7–9] for the nuclei ^{12}C, ^{40}Ca, and ^{48}Ca, and momentum transfer $|\mathbf{q}|$ between 300 and 550 MeV. The relation to other work on quasielastic electron scattering is discussed.

HARTREE APPROXIMATION

The cross section for electron scattering with momentum transfer $q=|\mathbf{q}|$ and energy transfer $\omega=q_0$ is given in terms of the Mott cross section σ_M by ($q_\lambda^2=\omega^2-q^2$)

$$\frac{d^2\sigma}{d\Omega\,dE}=\sigma_M\left[\left[\frac{q_\lambda^2}{q^2}\right]^2 S_L(q,\omega)+\left[-\frac{q_\lambda^2}{2q^2}+\tan^2\frac{\theta}{2}\right]S_T(q,\omega)\right]. \quad (1)$$

In nuclear matter the longitudinal (transverse) response function S_L (S_T) is proportional to the volume V and the imaginary part of the longitudinal (transverse) polarization propagator Π_L (Π_T):

$$S_{L,T}(q,\omega)=\frac{V}{\pi}\,\mathrm{Im}\Pi_{L,T}(j_{L,T},j_{L,T};k_{F_\mathrm{p}},k_{F_\mathrm{n}},M^*;q,\omega)\ . \quad (2)$$

The polarization propagator depends on the baryon current j, the Fermi wave numbers k_{F_p} and k_{F_n} of protons and neutrons, the effective mass M^* of the nucleons (the same for protons and neutrons), and on the kinematical variables q and ω.

The baryon current in terms of the Dirac spinor $\psi(\mathbf{x})$, the mass M of the free nucleon, and the anomalous moments $\lambda_\mathrm{p}=1.79$ and $\lambda_\mathrm{n}=-1.91$ reads

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$$j_\mu(q,\omega)=F_1(q_\lambda^2)\int d^3x\, e^{i\mathbf{q}\cdot\mathbf{x}}\bar{\psi}(\mathbf{x})\gamma_\mu\psi(\mathbf{x}) + F_2(q_\lambda^2)\int d^3x\, e^{i\mathbf{q}\cdot\mathbf{x}}\bar{\psi}(\mathbf{x})\frac{\lambda}{2M}i\sigma_{\mu\nu}q^\nu\psi(\mathbf{x}) . \tag{3}$$

For $\mu=0$ this is j_L, and we choose the coordinate system such that the transverse current j_T is given by j_2. F_1 and F_2 are the Dirac and Pauli form factors of the nucleons to be discussed later. By using this effective baryon current, the pion is partly taken into account.[1]

In the Hartree approximation (one particle-hole pair) proton and neutron do not interfere and their Hartree polarization propagators just add (we suppress the dependence on M^*, q, and ω),

$$\Pi_T^H(j_T,j_T;k_{F_\mathrm{p}},k_{F_\mathrm{n}})=\Pi_T^H(j_{T_\mathrm{p}},j_{T_\mathrm{p}};k_{F_\mathrm{p}})+\Pi_T^H(j_{T_\mathrm{n}},j_{T_\mathrm{n}};k_{F_\mathrm{n}}) , \tag{4}$$

and analogously for the longitudinal response. The Hartree response of proton and neutron matter can be evaluated exactly as in Ref. 3 for symmetric nuclear matter. We define

$$c_\mathrm{p}=\frac{\lambda_\mathrm{p}}{2M},\quad c_\mathrm{n}=\frac{\lambda_\mathrm{n}}{2M} \tag{5}$$

(to be multiplied later by the appropriate form factor) and, as in Ref. 3 [except for an additional factor $(2\pi)^{-3}$ multiplying ρ_s and I_n]

$$\rho_s(k_F)=\int_0^{k_F}\frac{d^3k}{(2\pi)^3}\frac{4}{(M^{*2}+\mathbf{k}^2)^{1/2}} , \tag{6}$$

$$I_n(q,k_F)=-2\int\frac{d^4p}{(2\pi)^3}\frac{\theta_p}{E_p}\delta(p_0-E_p)\left[\frac{(2p_0+\omega)^n}{(p+q)^2-M^{*2}+i\epsilon}+\frac{(2p_0-\omega)^n}{(p-q)^2-M^{*2}+i\epsilon}\right] -2\pi i\int\frac{d^4p}{(2\pi)^3}\frac{\theta_p}{E_p}\frac{\theta_{p+q}}{E_{p+q}}\delta(p_0-E_p)\delta(p_0+\omega-E_{p+q})(2p_0+\omega)^n$$
$$[\theta_p=\theta(k_F-|\mathbf{p}|),\ E_p=(\mathbf{p}^2+M^{*2})^{1/2}] , \tag{7}$$

$$\Pi_v(k_F)=\tfrac{1}{2}[\rho_s(k_F)-q^2I_0(k_F)+I_2(k_F)] , \tag{8}$$

$$\Pi_T(k_F)=-\left[\frac{\omega^2+q^2}{2q^2}\right]\rho_s(k_F)-\frac{1}{2}\left[\frac{q_\lambda^2}{q^2}I_2(k_F)+(4M^{*2}+q_\lambda^2)I_0(k_F)\right] . \tag{9}$$

For different Fermi momenta of protons and neutrons, the result for the Hartree polarization propagator reads (again we suppress the dependence of I_n on q):

$$\Pi_L^H(j_{L_\mathrm{p}},j_{L_\mathrm{p}};k_{F_\mathrm{p}})=\Pi_v(k_{F_\mathrm{p}})-2M^*c_\mathrm{p}q^2I_0(k_{F_\mathrm{p}})-c_\mathrm{p}^2[\omega^2\rho_s(k_{F_\mathrm{p}})+\tfrac{1}{2}q_\lambda^2I_2(k_{F_\mathrm{p}})+2M^{*2}q^2I_0(k_{F_\mathrm{p}})] , \tag{10a}$$

$$\Pi_L^H(j_{L_\mathrm{n}},j_{L_\mathrm{n}};k_{F_\mathrm{n}})=-c_\mathrm{n}^2[\omega^2\rho_s(k_{F_\mathrm{n}})+\tfrac{1}{2}q_\lambda^2I_2(k_{F_\mathrm{n}})+2M^{*2}q^2I_0(k_{F_\mathrm{n}})] , \tag{10b}$$

$$\Pi_T^H(j_{T_\mathrm{p}},j_{T_\mathrm{p}};k_{F_\mathrm{p}})=\tfrac{1}{2}\Pi_T(k_{F_\mathrm{p}})-2M^*c_\mathrm{p}q_\lambda^2I_0(k_{F_\mathrm{p}})-\tfrac{1}{2}c_\mathrm{p}^2q_\lambda^2[\rho_s(k_{F_\mathrm{p}})+\Pi_T(k_{F_\mathrm{p}})+(4M^{*2}+q_\lambda^2)I_0(k_{F_\mathrm{p}})] , \tag{10c}$$

$$\Pi_T^H(j_{T_\mathrm{n}},j_{T_\mathrm{n}};k_{F_\mathrm{n}})=-\tfrac{1}{2}c_\mathrm{n}^2q_\lambda^2[\rho_s(k_{F_\mathrm{n}})+\Pi_T(k_{F_\mathrm{n}})+(4M^{*2}+q_\lambda^2)I_0(k_{F_\mathrm{n}})] . \tag{10d}$$

In the special case $k_{F_\mathrm{p}}=k_{F_\mathrm{n}}$ we recover the result in Ref. 3.

RANDOM-PHASE APPROXIMATION

The relativistic RPA was developed for symmetric nuclear matter in Ref. 3 and is again easily extended to asymmetric nuclear matter. Writing

$$\Pi_L^{\mathrm{RPA}}=\Pi_L^H+\delta\Pi_L^{\sigma,\omega}+\delta\Pi_L^\rho , \tag{11a}$$

$$\Pi_T^{\mathrm{RPA}}=\Pi_T^H+\delta\Pi_T^{\sigma,\omega}+\delta\Pi_T^\rho , \tag{11b}$$

and with the definitions (m_s,m_v,m_ρ and g_s,g_v,g_ρ are the masses and coupling constants of σ,ω,ρ), again as in Ref. 3, but generalized to asymmetric nuclear matter,

$$\chi_s=\frac{g_s^2}{m_s^2-q_\lambda^2},\quad \chi_v=\frac{g_v^2}{m_v^2-q_\lambda^2} ,\qquad \chi_\rho=\frac{g_\rho^2}{m_\rho^2-q_\lambda^2},\quad \tilde{\chi}_{v,\rho}=\frac{q_\lambda^2}{q^2}\chi_{v,\rho} , \tag{12}$$

$$\Pi_v=\Pi_v(k_{F_p})+\Pi_v(k_{F_n}),\quad \Pi_T=\Pi_T(k_{F_p})+\Pi_T(k_{F_n})\ , \tag{13a}$$

$$\Pi_s=-\tfrac{1}{2}[\rho_s(k_{F_p})+\rho_s(k_{F_n})]+(2M^{*2}-\tfrac{1}{2}q_\lambda^2)[I_0(k_{F_p})+I_0(k_{F_n})]\ , \tag{13b}$$

$$\Pi_{sv}=M^*[I_1(k_{F_p})+I_1(k_{F_n})]\ , \tag{13c}$$

$$\Pi_{-1}=(M^*+\tfrac{1}{2}c_p q_\lambda^2)I_1(k_{F_p})+\tfrac{1}{2}c_n q_\lambda^2 I_1(k_{F_n})\ , \tag{13d}$$

$$\Pi_0=\Pi_v(k_{F_p})-M^*q^2[c_p I_0(k_{F_p})+c_n I_0(k_{F_n})]\ , \tag{13e}$$

$$\Pi_2=\tfrac{1}{2}\Pi_T(k_{F_p})-M^*q_\lambda^2[c_p I_0(k_{F_p})+c_n I_0(k_{F_n})]\ , \tag{13f}$$

we obtain, by solving Dyson's equation,

$$\delta\Pi_L^{\sigma,\omega}=\frac{\chi_s(1-\widetilde{\chi}_v\Pi_v)\Pi_{-1}^2+\widetilde{\chi}_v(1-\chi_s\Pi_s)\Pi_0^2+2\chi_s\widetilde{\chi}_v\Pi_{sv}\Pi_{-1}\Pi_0}{(1-\widetilde{\chi}_v\Pi_v)(1-\chi_s\Pi_s)-\chi_s\widetilde{\chi}_v\Pi_{sv}^2}\ , \tag{14a}$$

$$\delta\Pi_L^{\rho}=\frac{\widetilde{\chi}_\rho}{4-\widetilde{\chi}_\rho\Pi_v}\{\Pi_v(k_{F_p})-q^2M^*[c_p I_0(k_{F_p})-c_n I_0(k_{F_n})]\}^2\ , \tag{14b}$$

$$\delta\Pi_T^{\sigma,\omega}=\frac{\chi_v}{1-\frac{1}{2}\chi_v\Pi_T}(\Pi_2)^2\ , \tag{14c}$$

$$\delta\Pi_T^{\rho}=\frac{\chi_\rho}{4-\frac{1}{2}\chi_\rho\Pi_T}\{\Pi_T(k_{F_p})-q_\lambda^2M^*[c_p I_0(k_{F_p})-c_n I_0(k_{F_n})]\}\ . \tag{14d}$$

NUCLEON FORM FACTOR

At the four-momentum transfer relevant for quasielastic electron scattering, the suppression of the nuclear response due to the nucleon form factors is very important. Just as the authors of Ref. 5, we therefore use the parametrization[10]

$$\begin{aligned}
F_{1p}(q_\lambda^2)&=f_D(q_\lambda^2)\left[\frac{1-(q_\lambda^2/4M^2)(1+\lambda_p)}{1-q_\lambda^2/4M^2}\right]\ ,\\
F_{2p}(q_\lambda^2)&=f_D(q_\lambda^2)\frac{1}{1-q_\lambda^2/4M^2}\ ,\\
F_{1n}(q_\lambda^2)&=0\ ,\\
F_{2n}(q_\lambda^2)&=f_D(q_\lambda^2)=\left[1-\frac{q_\lambda^2}{0.71\ \mathrm{GeV}^2}\right]^{-2}\ ,
\end{aligned} \tag{15}$$

and also compare with the simpler, but less accurate parametrization

$$F_{1p}=F_{2p}=F_{2n}=f_D,\quad F_{1n}=0\ . \tag{16}$$

The structure functions are taken into account in our formulae by multiplying c_p by F_{2p}, c_n by F_{2n}, and all quantities proportional to the proton charge by F_{1p}.

LOCAL DENSITY APPROXIMATION

In this approximation we assume that the response of a nucleus is just the sum of the responses of its volume elements, each characterized by its local Fermi wave numbers $k_{F_p}(r)$, $k_{F_n}(r)$, and local effective mass $M^*(r)$ (we only consider closed-shell nuclei), and treated as nuclear matter with these parameters. From (2) we obtain the response per unit volume and then integrate over the nucleus

$$S_{L,T}(q,\omega)=4\int_0^\infty r^2\,\mathrm{Im}[\Pi_{L,T}^{\mathrm{RPA}}(k_{F_p}(r),k_{F_n}(r),M^*(r);q,\omega)]dr\ . \tag{17}$$

The local parameters are calculated by solving self-consistently the equations of the $\sigma\omega\rho$ model in mean-field theory for the nucleus under investigation.[6] As an example, in Fig. 1 we show the effective mass and Fermi wave numbers as a function of the radius for ^{40}Ca. At the center the neutron density is slightly larger than the proton density, because Coulomb interaction is included.

RESULTS

We discuss the longitudinal response function first. In Fig. 2 we have plotted S_L for ^{40}Ca at $q=410$ MeV with data taken from Ref. 7. Comparing our theoretical results in the random-phase approximation with the data, we see that the maximum occurs at too low an energy transfer and is too high by about 25%. As in nuclear matter, the RPA response is lower then the response in the Hartree approximation[3,4] and describes data better. We have made explicit the contribution of the ρ meson to this reduction by comparing with the longitudinal response in the $\sigma\omega$ model. As shown in Fig. 3, our Hartree result in the local-density approximation is quite similar to the result obtained in Hartree approximation with the structure of the nucleus taken fully into account,[5] but somewhat higher and shifted to lower energy transfer. Thus it may

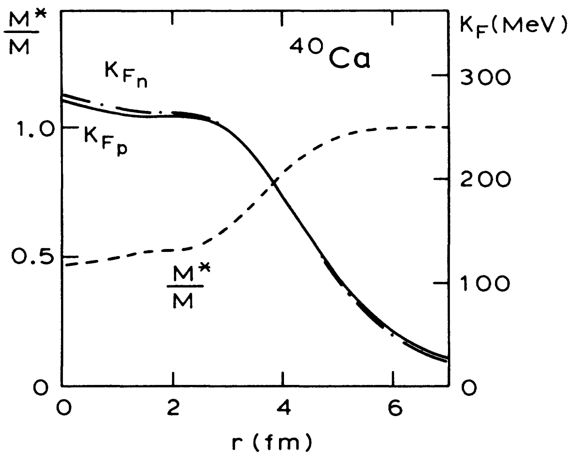

FIG. 1. Local effective mass M^* (dashed line) and Fermi wave numbers for protons (solid line) and neutrons (dotted-dashed line) for ^{40}Ca.

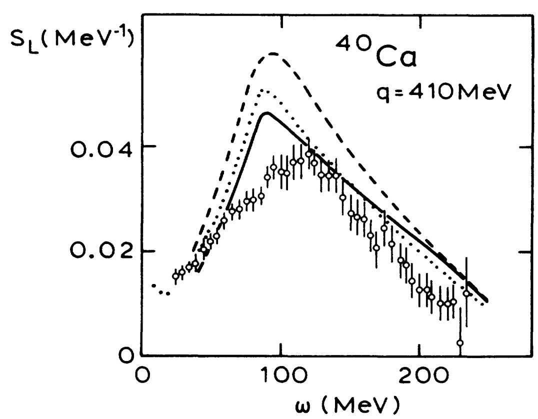

FIG. 2. Longitudinal response function for ^{40}Ca at $q = 410$ MeV. The RPA result in the $\sigma\omega\rho$ ($\sigma\omega$) model is shown by the solid (dotted) line, our Hartree result by the dashed line; data are from Ref. 7.

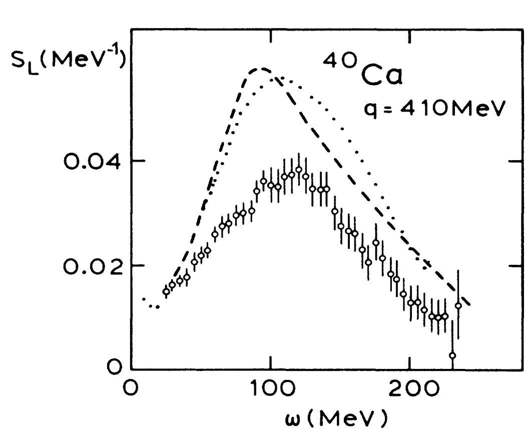

FIG. 3. Same as Fig. 2, the dashed line is our Hartree result, but now the dotted line is the Hartree result as given in Ref. 5.

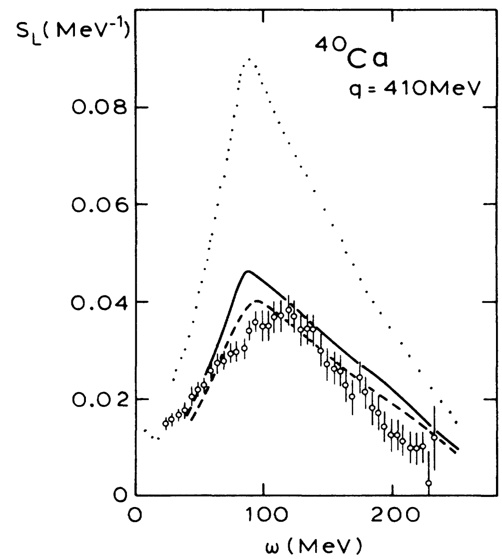

FIG. 4. Importance of the nucleon form factors for the longitudinal response function of ^{40}Ca at $q = 410$ MeV. The RPA result with the better parametrization (15) of the form factors is shown by the solid line, the parametrization (16) by the dashed line. The dotted line shows the RPA result with no form factor at all.

be expected that a RPA calculation along the lines of Ref. 5 improves the agreement with the data.

In Fig. 4 we make clear the importance of the use of the correct structure function of the nucleons. It reduces the nuclear response by about one-half, and even more at higher momentum transfer. If we use the approximate parametrization (16) instead of the better one, (15), S_L is lower by about 15%. The same behavior was also observed in Ref. 5. But even the better parametrization is not perfect. At momentum transfers relevant for quasielastic electron scattering, the electric form factor of the proton, for example, is known with an uncertainty of

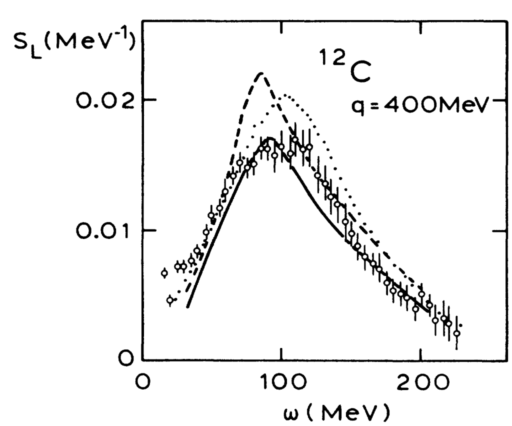

FIG. 5. Longitudinal response function for ^{12}C at $q = 400$ MeV. The solid (dashed) line shows our RPA (Hartree) result, the dotted line the Hartree result as given in Ref. 5.

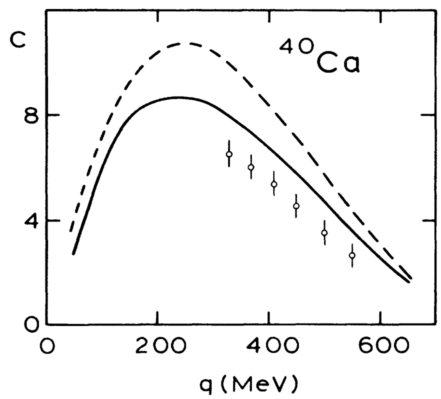

FIG. 6. Sum rule for ^{40}Ca. Data are from Ref. 7; the solid (dashed) line shows our RPA (Hartree) results.

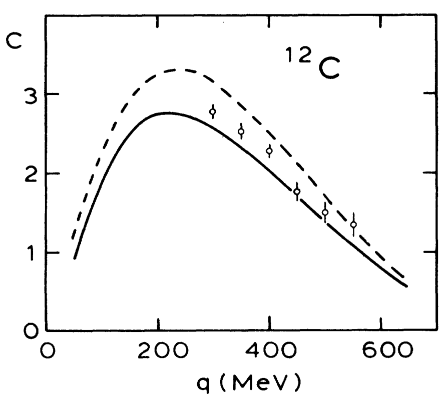

FIG. 7. Same as Fig. 6, but for ^{12}C with data from Ref. 8.

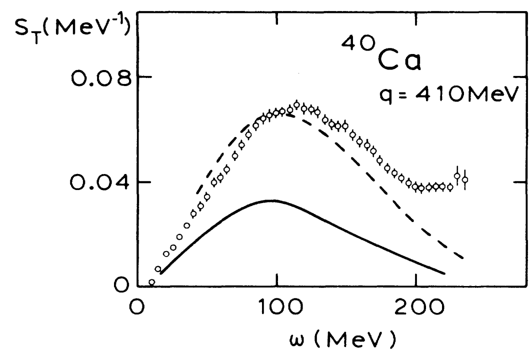

FIG. 9. Transverse response function for ^{40}Ca at $q=410$ MeV. The data are from Ref. 9; our RPA result is shown by the solid line, the Hartree result as given in Ref. 5 by the dashed line.

$\pm 3\%$ and is below the values given by (15) by about 5–10%.[10] The experimental information on the neutron structure function is much poorer. In view of the importance of the structure functions for our analysis (they enter quadratically), we should keep these uncertainties in mind.

The longitudinal response function for ^{12}C at $q=400$ MeV is shown in Fig. 5. The maximum of our RPA calculation in the local-density approximation occurs again at too low an energy transfer, but now its absolute value agrees with the data. Comparing with the results of Nishizaki, Kurasawa, and Suzuki,[5] we observe the same trends as for ^{40}Ca, but somewhat stronger.

In Figs. 6 and 7 we show the sum rule, which has recently been discussed in connection with the problem of the "missing charge,"[7] defined by

$$C(q)=\int_0^q S_L(q,\omega)d\omega \tag{18}$$

for ^{40}Ca and ^{12}C. The data are from Refs. 7 and 8. The

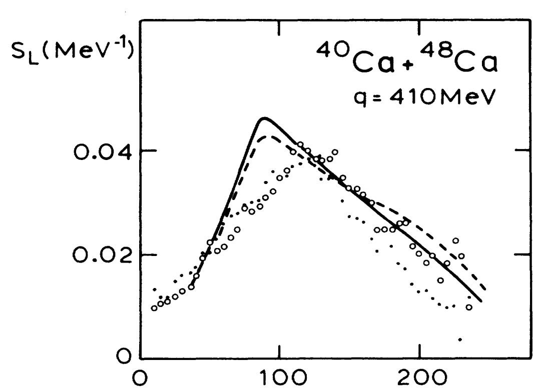

FIG. 8. Longitudinal response function of ^{40}Ca and ^{48}Ca at $q=410$ MeV. The solid (dashed) line shows the RPA result for ^{40}Ca (^{48}Ca); the solid circles (open circles) are data for ^{40}Ca (^{48}Ca) from Ref. 7. For clarity we have omitted the error bars.

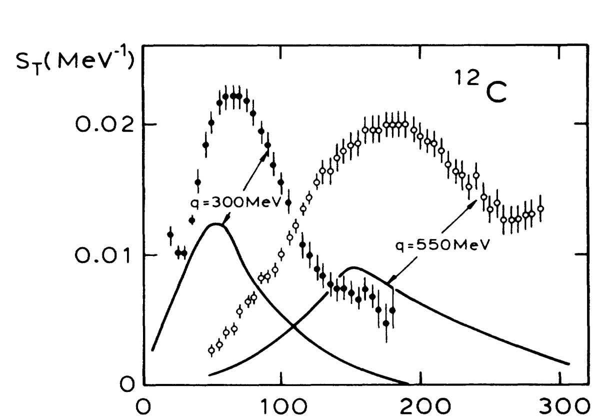

FIG. 10. Transverse response function for ^{12}C at $q=300$ and 550 MeV. The data, solid and open circles, respectively, are from Ref. 8; the solid (dashed) line is our RPA result for $q=330$ MeV (550 MeV).

sum rule is slightly overestimated for ^{40}Ca and slightly underestimated for ^{12}C for all values of the momentum transfer measured.

In Fig. 8 we compare our RPA results at $q=410$ MeV for ^{40}Ca and ^{48}Ca with data from Ref. 7. If we go from ^{40}Ca to ^{48}Ca in the data, the maximum tends to be higher, and it is shifted to larger energy transfer, while the maximum slightly decreases in our calculation and occurs at about the same energy for both isotopes. However, the status of the experiment seems to be not yet settled. Recent data from Bates[11] show significant deviations from the Saclay data.[7] For ^{40}Ca the response is higher by as much as 20% for $50<\omega<100$ MeV and smaller for larger energy transfer. For ^{48}Ca at the maximum of the response this discrepancy amounts to about 50%. While this does not improve the overall agreement with our calculation, it indicates the not yet settled status of the experiments, and should be a warning against preliminary conclusions.

We now turn to a discussion of the transverse response function. In Fig. 9 we show our results for ^{40}Ca at $q=410$ MeV in the RPA and the local-density approximation, together with data from Ref. 9 and the Hartree result, but with the structure of the nucleus taken fully into account, from Ref. 5. Our Hartree results are the same as our RPA results within a few percent. The calculation reported in Ref. 5 is in much better agreement with data than ours, which might indicate that the detailed structure of the nucleus is much more important for the transverse than for the longitudinal response. The observed similarity of Hartree and RPA results is in agreement with nuclear matter calculations,[3] but differs significantly from nonrelativistic calculations in the local-density approximation including particle-hole interactions, as given in Ref. 12. The transverse response is also very insensitive to details of the parametrization, (15) or (16), of the form factor of the free nucleon, although its effect is to reduce the transverse response, as in the longitudinal case, by about one-half.

For ^{12}C we find the same discrepancy of the transverse response with the data. This is shown in Fig. 10 for the example of momentum transfers $q=300$ and 550 MeV. The trend with energy is described correctly, but the absolute value is too low by about 50%. The same applies to the comparison of the transverse response of ^{40}Ca and ^{48}Ca.

DISCUSSION

Our observation that the longitudinal response function is in reasonable agreement with the data, while the transverse one underestimates the data, agrees with other calculations in the relativistic $\sigma\omega$ model. In Ref. 13 the same problem was encountered with a rough estimation of the effects of the large scalar and vector potentials on the quasielastic response. Our more refined analysis does not resolve this discrepancy. We have already discussed the results of Ref. 5, where the response was calculated in Hartree approximation with the realistic one-body density of the finite nucleus. Our Hartree result in the local-density approximation is in quite good agreement with theirs for the longitudinal response, and the relativistic RPA correlations reduce the response and improve the agreement with the data. For the transverse response our results are much lower than the results of Ref. 5; the detailed structure of the nucleus seems to be important in this case. In Ref. 14 medium-modified form factors of the nucleons in the nucleus are invoked to soften the longitudinal response. Their result fits the data very well, but at the cost of stronger disagreement for the transverse response. However, as we have seen, the RPA correlations significantly reduce the longitudinal response, and we have reason to believe that a realistic RPA calculation for a finite nucleus along the lines of Ref. 5 may yield even better agreement with data for the longitudinal response without such strong assumptions.

As to the transverse response, in this channel we expect important contributions from pion exchange and ρ meson tensor coupling. Quite generally, in this type of model the isovector channel is not well understood, as is exemplified by calculations of the magnetic moments of closed shell ± 1 nuclei.[15] While the isoscalar moments are in reasonable agreement with the data, the isovector moments are too large. The two problems are intimately connected and a better understanding of the role of the pion in relativistic models is needed to resolve them.

CONCLUSION

We have calculated the longitudinal and transverse response functions of closed-shell nuclei in the local-density and relativistic random-phase approximations, using the relativistic $\sigma\omega\rho$ model. We have no adjustable parameters. Our work shows that RPA correlations reduce the longitudinal response and constitutes, after Ref. 5, a further step toward a complete relativistic RPA calculation of finite nuclei. The local-density approximation is a good approximation for the longitudinal response, but seems to be insufficient for the transverse response. The missing strength in the transverse response is, in our opinion, due to the fact that in the calculation employed here the pion is only partly included by using the effective baryon current (3).

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