Gauge and Lorentz invariant pionic correlations in quasi-elastic electron scattering

M.B. Barbaro

Dipartimento di Fisica Teorica and INFN - Università di Torino, Via P. Giuria1, I- 10125
Turin, Italy
1. – Introduction

The problem of accounting for relativistic dynamics in nuclear physics is a daunting one and far from being solved (see, for example, Refs. [1, 2]). However in modern experimental studies of electron scattering from nuclei the typical values of energy and momentum transfer are comparable to, or even larger than, the scale set by the nucleon mass: accordingly one must expect relativistic effects to be important.

One reason for going to high energy and momentum transfers is the possibility of extracting information on the nucleon’s form factors, in particular on the strange and axial ones, which are measured via polarized electron scattering. Due to the large number of form factors involved in the process, the scattering on a single proton is not sufficient to disentangle the interesting quantities and heavy targets are needed to yield complementary information: hence the necessity of controlling the nuclear dynamics in such kinematical conditions. Besides this practical motivation there is a more interesting one from the nuclear physics’ point of view: the scattering of polarized electrons off complex targets can shed light on some nuclear correlations which are not accessible by unpolarized electrons, as will be shown in Section 3.

The traditional approaches to wave functions and operators, used in most calculations to describe the high-energy regime, involve leading-order expansions of the electroweak currents: such an approach is highly constrained to work only at relatively low energies and momenta, so that an exact treatment of relativistic effects is required.

There are two simple general principles which ought to be respected by any consistent treatment of the nuclear problem, independently of the details of the nucleon-nucleon interaction and of the theoretical framework adopted to deal with nuclear correlations (Hartree-Fock, RPA, Brueckner-Hartree-Fock, etc):

1. Lorentz covariance: the nuclear current must transform as a four-vector under a Lorentz boost;

2. Gauge invariance: the nuclear current must be conserved.

The simplest model in which the above fundamental requirements can be accomplished is the relativistic Fermi gas model (RFG), i.e., a system of nucleons moving freely inside the nucleus with relativistic kinematics. When focus is placed on the quasielastic region, where high-energy knockout of nucleons is kinematically favored, this model, while undoubtedly too simple to encompass the aspects of nuclear dynamics, is nevertheless a convenient place to start in such explorations.

In the quasielastic regime it is reasonable to expect pions to play a role that differs from the dynamics typically occurring near the Fermi surface, where one expects other mesons (σ and ω in particular) to dominate. For quasielastic scattering the residual interaction of relevance is principally that between a low-energy hole and a very high-energy particle, and for this the pion is expected to play an important role. Accordingly, as the next step after the basic RFG of non-interacting nucleons, the most important ingredient of the quasi-elastic nuclear responses is believed to be the one-pion exchange (OPE) potential.
Gauge and Lorentz invariant pionic correlations in quasi-elastic electron scattering

In this context a consistent first-order operator, embodying all Feynman diagrams built out of nucleons and pions with one exchanged pion and one photon attached to all the possible lines can be set up to represent the two-body current. Importantly, the latter is gauge invariant [3].

This fully-relativistic operator includes both the meson-exchange currents (MEC) and the so-called correlation currents. The latter are often not included in model calculations because they give rise to contributions already accounted for in the initial and final nuclear wave functions. However, the present approach is based on an uncorrelated relativistic Fermi gas whose states are Slater determinants built out of (Dirac) plane waves. Within a perturbative scheme one is free to consider the one-pion correlation contributions to the responses as arising either explicitly in the wave functions or from an appropriate current operator acting on unperturbed states: here the choice will be the latter. Clearly, should it be possible to sum up the whole perturbative expansion, then the results obtained starting with the true “correlated” wave function would be exactly recovered.

In what follows I shall illustrate the application of the pionic model to the unpolarized (Section 2) and polarized (Section 3) quasi-elastic electron scattering, introduce a “semi-relativistic” expansion which allows to “mimick” relativity at any value of energy and momentum transfers (Section 4) and briefly illustrate the scaling and “superscaling” behavior of the model (Section 5).

2. – Parity-conserving electron scattering

The formalism for the unpolarized, inclusive \((e,e')\) process is developed in great detail in a variety of papers (see, e.g., Refs. [4, 5, 6]). Hence here only the basic formulae are reported and particular emphasis is placed on the above mentioned Lorentz- and gauge-invariance issues.

In the extreme relativistic limit (ERL), in which the incident electron energy \(\varepsilon \gg m_e\), the cross section reads

\[
\frac{d\sigma}{d\Omega'_e d\omega} = 2\alpha^2 Q^4 \left(\frac{\varepsilon'}{\varepsilon}\right) \eta_{\mu\nu} W^{\mu\nu} = \sigma_M \left[ v_L R_L(q,\omega) + v_T R_T(q,\omega) \right],
\]

where \(\Omega'_e = (\theta_e, \Phi_e)\) is the scattered electron solid angle, \(\alpha\) the fine structure constant, \(Q_{\mu} = (\omega, q)\) the transferred four-momentum, \(\eta_{\mu\nu}\) and \(W^{\mu\nu}\) the leptonic and hadronic tensor, respectively, \(\sigma_M\) the Mott cross section and \(v_{L,T}\) are kinematical factors (defined, for example, in [6, 7]). The longitudinal and transverse (with respect to the momentum transfer \(q\), which fixes the direction of the \(z\)-axis) response functions \(R_L\) and \(R_T\) are constructed as components of the hadronic tensor \(W^{\mu\nu}\):

\[
R_L(q,\omega) = \left(\frac{q^2}{Q^2}\right)^2 \left[ W^{00} - \frac{\omega}{q} (W^{03} + W^{30}) + \frac{\omega^2}{q^2} W^{33} \right]
\]

\[
R_T(q,\omega) = W^{11} + W^{22}.\]
Fig. 1. – Feynman diagrams contributing to the two-body current with one pion-exchange. The wide line in the correlation diagrams (d)–(g) means a fully-relativistic Dirac propagator for the nucleon.

If gauge invariance is fulfilled, implying that
\[ W_{03} = W_{30} = \left( \frac{\omega}{q} \right) W_{00} = \left( \frac{q}{\omega} \right) W_{33}, \]
then \( R_L \) is simply the time component of the hadronic tensor, namely \( W_{00} \). Hence \( R_L \) and \( R_T \) are determined by the nuclear charge and current distributions, respectively, and they embody the entire dependence upon the nuclear structure.

Working within the framework of the RFG model one can then construct the electromagnetic currents accounting for the effects introduced by pions in first-order perturbation theory (one-pion exchange).

The linked, two-body Feynman diagrams that contribute to electron scattering with one pion-exchange are shown in Fig. 1. The first three correspond to the usual meson-exchange currents (MEC): diagrams (a), (b) refer to the “seagull” current, diagram (c) to the “pion-in-flight” current. The four diagrams (d)–(g) represent the so-called correlation current and are usually not treated as genuine MEC, but as correlation corrections to the nuclear wave function. However, again note that the present approach puts all correlation effects in the current operator and uses an uncorrelated wave function for the initial and final nuclear states.

The general relativistic expressions for the seagull (\( s \)), pion-in-flight (\( p \)) and correlation (\( cor \)) current matrix elements of Fig. 1 are (isospin summations are understood)

\[
(4) \quad j^\mu_s = \frac{f^2}{m_\pi^2} \epsilon_{abc} \vec{\pi}(p_1^\prime) \gamma_5 K_1 u(p_1) \frac{F^V_\mu}{K_1^2 - m_\pi^2} \tau_b \gamma_5 \gamma^\mu u(p_2) + (1 \leftrightarrow 2)
\]

\[
(5) \quad j^\mu_p = \frac{f^2}{m_\pi^2} \epsilon_{abc} \frac{F_\pi (K_1 - K_2)^\mu}{(K_1^2 - m_\pi^2)(K_2^2 - m_\pi^2)} \vec{\pi}(p_1^\prime) \gamma_5 K_1 u(p_1) \vec{\pi}(p_2^\prime) \gamma_5 K_2 u(p_2)
\]

\[
(6) \quad j^\mu_{cor} = \frac{f^2}{m_\pi^2} \vec{\pi}(p_1^\prime) \gamma_5 K_1 u(p_1) \frac{1}{K_1^2 - m_\pi^2}
\]
\[ \propto \pi(p'_2) \left[ \tau_a \gamma_5 K_1 S_F(P_2 + Q)\Gamma^\mu(Q) + \Gamma^\mu(Q)S_F(P'_2 - Q)\tau_a \gamma_5 K_1 \right] u(p_2) + (1 \leftrightarrow 2) . \]

In the above, \(K_1, K_2\) are the four-momenta given to the nucleons 1, 2 by the exchanged pion, and they are defined in Fig. 1, while \(F^V_1\) and \(F_\pi\) are the electromagnetic isovector nucleon and pion form factors, respectively. Furthermore, \(u(p)\) is the free Dirac spinor of a nucleon carrying momentum \(p\), \(S_F(P)\) the nucleon propagator and \(\Gamma^\mu(Q) = F_1 \gamma^\mu + \frac{i}{2m} F_2 \sigma^\mu \nu Q^\nu\) the electromagnetic nucleon vertex (the Galster parametrization is used for the Dirac and Pauli form factors \(F_1\) and \(F_2\)).

A crucial point to be stressed is that the sum of the relativistic seagull, pion-in-flight and correlation currents satisfies current conservation,

\[ Q^\mu (j_s^\mu + j_p^\mu + j_{cor}^\mu) = 0 \, , \]

provided \(F_\pi = F^V_1\). It is also possible \([9, 10]\) to use different phenomenological electromagnetic form factors for the nucleon and pion — even introducing phenomenological form factors at the strong pion-nucleon vertices — without violating current conservation, by appropriate modification in the currents through the generalized Ward-Takahashi identity.

Since the interest here is focused in the one-particle emission induced by the two-body currents introduced above, one needs now to evaluate the matrix element of the above two-body operator between the Fermi gas ground state and a 1p-1h excitations, which are the dominant modes in the quasielastic regime:

\[ \langle ph^{-1} | j^\mu(Q) | F \rangle = \sum_{k < F} \left[ \langle pk | j^\mu(Q) | hh \rangle - \langle pk | j^\mu(Q) | hh \rangle \right] , \]

where the summation runs over all occupied levels in the ground state, and thus includes a sum over spin \((s_k)\) and isospin \((t_k)\) and an integral over the momentum \(k\).

The first and second terms in eq. (8) represent the direct and exchange contribution to the matrix element, respectively. It can be easily verified that in spin-isospin saturated systems the direct term vanishes for the currents (4-6) upon summation over the occupied states. Hence only the exchange term contributes to the p-h matrix elements.

When inserted in (8) the correlation current (6) gives rise to two kinds of diagrams, called “vertex corrections” (VC) and “self-energy” (SE), respectively \([11, 3]\). The self-energy current deserves some special comment, as it diverges: in fact it corresponds to a SE insertion on an external line, which, according to field theory, should not be included in a perturbative expansion. One should then apply a renormalization procedure to dress the external lines by summing up the entire perturbative series of self-energy insertions. In the nuclear context this procedure leads to the relativistic Hartree-Fock approach. A renormalized self-energy current corresponding to one-pion-exchange is constructed in Ref. [3] by renormalizing spinors and energies and by expanding the resulting in-medium one-body current to first order in the square of the pion-nucleon coupling constant.

Once this is done, the electromagnetic inclusive response functions for one-particle emission reactions can be evaluated within the RFG model. The hadronic tensor that
Fig. 2. — Longitudinal (left panels) and transverse (right panels) response functions versus $\omega$ including all first-order contributions (solid) compared with the free result (dashed). The nucleus is $^{40}$Ca, corresponding to a Fermi momentum $k_F = 237 \text{ MeV/c}$. 

arises from the interference of the single-nucleon, one-body (OB) current, $j_{OB}^{\mu}(p, h) = \bar{\pi}(p)\Gamma^\mu u(h)$, with the one-pion-exchange current, $j_p^{\mu} + j_s^{\mu} + j_{cor}^{\mu}$, is for the RFG model with $Z = N$: 

$$W^{\mu \nu} = \frac{3Z}{8\pi k_F q} \sum_{a=p, s, \text{cor}} \int_{k_0}^{k_F} dh(\omega + E_h) \int_0^{2\pi} d\phi_h \sum_{s_p, s_h} \frac{m^2}{E_p E_h} 2\text{Re} \left[ j_{OB}^{\nu}(p, h)^* j_{OB}^{\mu}(p, h) \right].$$  

(9) 

The numerical results obtained for the longitudinal and transverse quasielastic re-
response functions in the 1p-1h sector are shown in Fig. 2. The calculation refers to $Z = N = 20$ and $k_F = 237$ MeV/c, which is representative of nuclei in the vicinity of $^{40}\text{Ca}$. The total responses are displayed in first order of perturbation theory and compared with the zeroth-order ones (free responses) for several momentum transfers. In assessing the impact of the global two-body current contribution to the responses it should be noted that:

1. the overall effect of the two-body currents appears sufficiently modest to justify a posteriori a first-order treatment, their relative contribution ranging from $\sim 5$ to $\sim 15\%$ depending upon the kinematics;

2. the softening at large $q$ appears to be common to both L and T channels, whereas at low $q$ the longitudinal response displays a hardening that is absent in the transverse one;

3. the two-body correlation contribution is nearly vanishing at the peak of the free responses and is roughly symmetrical about the quasielastic peak, implying that their impact on the Coulomb sum rule should be very small.

3. – Parity-violating electron scattering

In this section the parity-violating (PV) effects arising from the weak interaction between the electron and the nucleus are addressed. Such effects, which are negligible in unpolarized electron processes, can be brought to evidence by measuring the asymmetry associated with longitudinally polarized electrons having opposite helicities. In this case the purely electromagnetic cross sections cancel out and one is left with the interference between the electromagnetic and neutral weak currents, corresponding to the exchange of a photon and a $Z^0$, respectively.

An important motivation of parity-violating experiments is the measurement of the single-nucleon form factors, in particular the strange and axial ones: for this reason most experiments are presently being carried out on light nuclei, where the uncertainties associated with the nuclear model are minimized. However, other motivations exist for such studies: specifically, as anticipated in the introduction, the PV response functions display a different sensitivity to nuclear correlations compared with the parity-conserving ones: hence they could not only shed light on the part of the problem concerned with nucleon (and meson) structure, but also are being used as a test of nuclear models. Indeed in [12, 13, 14] a semi-relativistic analysis of the PV responses has been presented, showing the dominance of pionic correlations in the longitudinal channel. In [3] a fully-relativistic calculation has been performed which confirms the above findings, and extend them to higher values of the momentum transfer.

In terms of nuclear response functions the asymmetry reads

$$\mathcal{A} \equiv \frac{d\sigma^+ - d\sigma^-}{d\sigma^+ + d\sigma^-} = A_0 \frac{v_L R^L_{AV} + v_T R^T_{AV} + v_T R'_{AV}}{v_L R^L + v_T R^T},$$

(10)
where \( A_0 = \frac{G Q^2}{2 \sqrt{2} \pi e^2} \), \( G \) being the Fermi constant, and \( v_{L,T,T'} \) are leptonic kinematical factors (see Refs. [7, 8]).

The PV response functions are linked to the interference hadronic tensor \( \tilde{W}_{\mu\nu} \) by the following relations:

\[
R_{AV}^L(q,\omega) = a_A \left( \frac{q^2}{Q^2} \right)^2 \left[ \tilde{W}^{00} - \frac{\omega}{q} (\tilde{W}^{03} + \tilde{W}^{30}) + \frac{\omega^2}{q^2} \tilde{W}^{33} \right]
\]

\[
R_{AV}^T(q,\omega) = a_A \left( \tilde{W}^{11} + \tilde{W}^{22} \right)
\]

\[
R_{V_A}^T(q,\omega) = -i a_V \left( \tilde{W}^{12} - \tilde{W}^{21} \right)
\]

where \( a_A = -1 \) and \( a_V = 4 \sin^2 \theta_W - 1 \). The subscript \( AV \) in the PV responses denotes interferences of axial-vector leptonic currents with vector hadronic currents, and the reverse for the subscript \( VA \).

Within the context of the RFG model the interference hadronic tensor is

\[
(14) \tilde{W}_{\mu\nu} = \frac{3Z}{8\pi k_F q} \int_{k_0}^{k_F} h dh (\omega + E_h) \int_0^{2\pi} d\phi_h \sum_{s_p, s_h} \frac{m^2}{E_p E_h} 2Re \left[ j_{\mu em}^*(p, h) j_{\nu wn}^*(p, h) \right]
\]

where the electromagnetic current \( j_{\mu em}^* \) includes both the single nucleon one-body and the two-body currents discussed in the previous section, i.e. \( j_{\mu em}^* = j_{\mu OB}^* + j_{\mu MEC}^* + j_{\mu cor}^* \). The weak neutral current \( j_{\nu wn}^* \) is instead purely one-body, since the direct coupling of a \( Z^0 \) to the pion is neglected.

As for the parity-conserving sector, the one-body contribution to the three PV responses (11-13) can be evaluated analytically in RFG (see, for example [5] for the explicit expressions of the response functions), whereas the two-body contributions involve multidimensional integrals, to be numerically evaluated.

In Fig. 3 the effect of the pionic physics on the asymmetry in eq. (10) is shown by displaying \( A \) versus \( \omega \) at various values of the momentum transfer \( q \) and of the electron scattering angle \( \theta_e \). The curves represent the free RFG (dashed), and the RFG including MEC and VC (dot-dashed) or the MEC, VC and SE (solid) contributions. It appears that the pionic correlations are mostly felt at low values of \( \theta_e \) and \( q \). Indeed a careful analysis, carried out in Ref. ([1]), shows that the main effect on the asymmetry arises from the vertex corrections, which dominate at low \( q \), in the longitudinal response function. The latter is enhanced at low \( \theta_e \) by the kinematical factor \( v_L \): hence the large modification of the asymmetry at low angles and momentum transfers. On the other hand at large scattering angles the asymmetry is totally insensitive to pions, because the effect of the SE (which gives the main contribution) cancels between the PV and PC responses appearing in the numerator and denominator of eq. (11).

The conclusion can be drawn that the extraction (at large electron angles) of the axial nucleonic form factor \( G_A \) is almost independent of the nuclear model. On the contrary at small angles PV experiments can measure the strange electric content of the nucleon.
Fig. 3. – The PV asymmetry displayed versus $\omega$ for various values of the momentum transfer $q$ and the scattering angle $\theta_e$. Dashed: one-body; dot-dashed: one-body+MEC+VC; solid: total.

only if a good control of the nuclear dynamics is achieved, since the isospin correlations give very large effects. Conversely, interesting insight into the latter can in principle be gained here. The results show that only at very large momentum transfer does the forward-angle asymmetry become insensitive to pionic correlations and hence suitable for assessing the strangeness content of the nucleon.

4. – The $\eta_F$-expansion

In the non-relativistic reductions commonly used in treating the effects of two-body pionic currents in electron scattering reactions, e.g.\cite{15, 16, 17}, not only non-relativistic wave functions have been used, but also non-relativistic current operators derived from a
direct Pauli reduction have been considered. Although these approximations are sufficient in the low-energy regime, they badly fail to hold at high momentum transfer.

On the other hand the exact relativistic calculation outlined in Section 2 is limited, up to now, to the simple pion-exchange potential and to first order perturbation theory. These two assumptions are not too crude in the quasi-elastic region, but they cannot

Fig. 4. – Total transverse response function of $^{40}$Ca including MEC for several values of the momentum transfer, and for $k_F = 237$ MeV/c. Solid: exact relativistic results. The rest of the curves have been computed using the non-relativistic Fermi gas model, with or without relativistic corrections. Dashed: traditional non-relativistic results. Dotted: including relativistic kinematics in the non-relativistic calculations. Dot-dashed: including in addition the new expansion of the OB+MEC currents. The relativistic calculations include a dynamical propagator and $\pi N$ form factor, while the non-relativistic calculations do not include these corrections.
obviously be applied to different kinematical regimes, where higher orders in the perturbative series and different correlations (in particular short-range correlations) are known to be far from negligible.

In order to implement effects of relativity into more realistic nuclear models, new expressions of the one- and two-body electromagnetic currents have been suggested in [18, 19], as well as the $N \rightarrow \Delta$ electromagnetic current [20]: here the current is derived as a non-relativistic expansion in terms of the dimensionless parameter $\eta \equiv p/m$, $p$ being the three-momentum of the struck nucleon. Since the latter is limited by the Fermi momentum, the procedure amounts to an expansion in the parameter $\eta_F = k_F/m \simeq 1/4$.

It should be stressed that such expansion holds valid for any energy and momentum transfer and can therefore be applied even in extreme kinematical conditions.

Although the $\eta_F$-expansion is not needed in the present context, where the calculations are performed without any non-relativistic approximation, it has the merit of yielding recipes to include relativistic corrections in the non-relativistic currents through the simple kinematical factors. For example the relation between the non-relativistic ($NR$) and the new “semi-relativistic” ($SR$) pion-in-flight current so obtained is found to be [21]:

$$j_{p,SR}^\mu = \frac{1}{\sqrt{1 + \tau}} j_{p,NR}^\mu ,$$

where $\tau = |Q^2|/(4m^2)$. Similar relations can be derived for all the components of the one- and two-body nuclear current. By means of this expansion one could in principle extend the validity of existing sophisticated, but non-relativistic, calculations to regimes where they are not at present applicable.

To appreciate the quality of the method in Fig. 4 the total transverse response, including OB+MEC operators, is reported for $q = 500, 1000, 2000$, and $3000$ MeV/c. The solid lines are the exact relativistic result, the dashed lines represent the traditional non-relativistic results, which together with relativistic kinematics give the dotted lines. Finally, the dot-dashed lines correspond to the results using the semi-relativized OB+MEC currents. Whereas the non-relativistic approach is clearly failing already at moderate values of $q$ (of the order $500$ MeV/c), the agreement between the exact model and the $\eta_F$-expanded one is quite good even for very high $q$: these currents are therefore very appropriate and easy to implement in already existing non-relativistic models for the electromagnetic reactions.

5. – Scaling

Before concluding I will shortly illustrate the behavior of pionic correlations in the quasi-elastic peak with respect to scaling and superscaling. A comprehensive introduction to the concept of scaling can be found in Refs. [22, 23, 24, 25]. Here I only remind the basic definitions:

1. Scaling of I kind occurs if the ratio $F^{L,T}(q, \omega)$ between the nuclear response func-
tions $R_{L,T}^{L,T}(q, \omega)$ and the relevant single-nucleon electromagnetic factors becomes, for high values of $q$, function of one single variable, the scaling variable. Several different scaling variables exist in the literature, all of them coalescing into one - or being simply related to each other - for high enough momentum transfers. In the quasi-elastic peak region the natural scaling variable turns out to be $\psi = \pm \sqrt{T_0 / T_F}$,

$$\psi = \pm \sqrt{T_0 / T_F},$$

being

$$T_0 = \frac{1}{2} \left( q \sqrt{1 + \frac{4m^2}{|Q^2|}} - \omega \right)$$

the minimum kinetic energy required to a nucleon to take part in the process and $T_F$ the Fermi kinetic energy. The $+(-)$ sign in (16) refers to the right (left) of the quasi-elastic peak. The analysis of the world data [22, 23, 24] shows that scaling of I kind is reasonably good for $\psi < 0$ and badly violated for $\psi > 0$.

2. Scaling of II kind consists in the independence of the functions $f_{L,T} = k_F \times F_{L,T}$ on the specific nucleus, namely on the Fermi momentum. The analysis of the existing data points to an excellent fulfillment of this scaling in the region $\psi < 0$ and to a not very dramatic breaking of it for $\psi > 0$. When the two kinds of scaling occur the responses are said to “superscale”.

The relativistic Fermi gas model fulfills both kinds of scaling by construction. The observed superscaling behavior of the experimental data [22, 24] offers a clear constraint to the nuclear correlations and can be used as a test of the reliability of the model. It is then natural to ask whether or not the present pionic model superscales.

In [11] the evolution with $q$ of the MEC in the transverse channel has been explored in detail: it has been proven that the relative contribution of the MEC to $R_T$ decreases with $q$, but does not vanish for large values of $q$. In fact, it decreases in going from 0.5 to 1 GeV/c, but then it rapidly saturates at or slightly above $q=1$ GeV/c, where its value stabilizes, typically around 10%. Thus, one can conclude that at momentum transfers above 1 GeV/c scaling of the first kind is satisfied for the MEC contributions considered in this work. Moreover, it is found that for high $q$ the MEC almost vanish for $\omega$ in the vicinity of the QEP.

The evolution with $q$ of the correlation current in the longitudinal and transverse channels has also been discussed at length in [11]. The basic findings are that: 1) the VC do not saturate quite as rapidly as the MEC, although their behavior is rather similar and saturation again occurs somewhere above $q = 1-1.5$ GeV/c: thus, again, scaling of the first kind is achieved at high momentum transfers for these contributions. Moreover, similarly to the MEC case, for high $q$ the VC almost vanish around the QEP.

2) A somehow different behavior is observed in the self-energy, which, due to a delicate
cancellation between the particle and hole dressings, grows - instead of decreasing - with $q$ in the range $q=0.5-2$ GeV/c, and then stabilizes typically at about 30-40% of the free response to the left of the QEP, thus inducing an important softening of the longitudinal and transverse responses.

As far as scaling of second kind is concerned, the two-body current contribution to the response functions is found to grow with $k_F$, in contrast with the free response which decreases as $k_F^{-1}$. More specifically the two-body processes violate the second-kind scaling by roughly three powers of $k_F$. This effect is a rapid function of the Fermi momentum (or equivalently, of the density): for example, if one considers the cases $^2$H/$^4$He/heavy nuclei with Fermi momenta of approximately 55/200/260 MeV/c, respectively, then the 1p-1h MEC contributions amount to 0.1/5/10% of the total transverse response, respectively (normalizing to 10% for the heavy nucleus case).

In summary, in the present pionic model scaling of the first kind is achieved at momentum transfers somewhat below 2 GeV/c, whereas scaling of second kind is badly violated. A similar trend is expected to be followed by the contributions of the heavier mesons, which have been neglected in the present approach: in order to agree with the experimental data, the strength of these contributions cannot be too disruptive.

6. – Conclusions

The impact of pionic correlations and meson-exchange currents on the nuclear electromagnetic response functions, calculated in a fully relativistic context which allows to respect exactly gauge invariance, is found to be modest in the quasi-elastic region. On the other hand the parity-violating asymmetry displays a strong sensitivity to such correlations if the scattering angle and momentum transfers are small. The analysis of the corrections induced by the pion with respect to “superscaling” shows that, whereas they do not disrupt the I-kind scaling of the relativistic Fermi gas, they badly violate scaling of II-kind.

Of course the pion is only one ingredient of the NN force, which, due to its long-range nature, is thought to give the dominant contribution to the responses in the quasi-elastic peak region. In order to deal with different observables and kinematical regions, where the short-range physics is known to play an important role, the present model should be extended to include heavier mesons exchange and currents. A first effort in this direction is performed in Ref. [27], where the modification of the momentum distribution due to the full Bonn potential is analyzed.

* * *

The work presented was carried out in collaboration with J.E. Amaro, J.A. Caballero, T.W. Donnelly and A. Molinari. I would like to thank J.E. Amaro and J.A. Caballero for carefully reading the manuscript and to acknowledge financial support from MEC (Spain) for a sabbatical stay at University of Sevilla (SAB2001-0025).
REFERENCES

[1] Serot B.D. and Walecka J.D., Adv. Nucl. Phys., 16 (1986) 1.
[2] Celenza L.S. and C. M. Shakin C.M., Relativistic Nuclear Physics, (World Scientific, Singapore) 1986.
[3] Amaro J.E., Barbaro M.B., Caballero J.A., Donnelly T.W. and Molinari A., Phys. Rep., 368 (2002) 317.
[4] Donnelly T.W., this volume.
[5] Musolf M.J., Donnelly T.W, Dubach J., Pollock S.J., Kowalski S. and Beise E.J., Nucl. Phys. A, 541 (1992) 525.
[6] Alberico W.M., Molinari A., Donnelly T.W. and Donnelly T.W., Relativistic Nuclear Physics, (World Scientific, Singapore) 1986.
[7] AGRH J.E., Barbaro M.B., Caballero J.A., Donnelly T.W. and Molinari A., Nucl. Phys. A, 643 (1998) 349.
[8] Donnelly T.W. and Sick I., Phys. Rev. Lett., 82 (1999) 3212.
[9] AGRH J.E., Barbaro M.B., Caballero J.A., Donnelly T.W. and Molinari A., Nucl. Phys. A, 643 (1998) 137.
[10] Maieron C., Donnelly T.W. and Sick I., Phys. Rev. C, 65 (2002) 025502.
[11] Alberico W.M., Molinari A., Donnelly T.W., Kronenberg E.L. and Van Orden J.W., Phys. Rev. C, 38 (1988) 1801.