



Scalar and vector form factors of the in-medium nucleon

K. Saito^a, K. Tsushima^b

^a *Tohoku College of Pharmacy, Sendai 981-8558, Japan*

^b *Department of Physics and Astronomy, University of Georgia, Athens, GA 30602, USA*

Received 15 July 2003; accepted 18 September 2003

Editor: W. Haxton

Abstract

Using the quark–meson coupling model, we calculate the form factors at σ - and ω -nucleon strong-interaction vertices in nuclear matter. The Peierls–Yoccoz projection technique is used to take account of center of mass and recoil corrections. We also apply the Lorentz contraction to the internal quark wave function. The form factors are reduced by the nuclear medium relative to those in vacuum. At normal nuclear matter density and $Q^2 = 1 \text{ GeV}^2$, the reduction rate in the scalar form factor is about 15%, which is almost identical to that in the vector one. We parameterize the ratios of the form factors in symmetric nuclear matter to those in vacuum as a function of nuclear density and momentum transfer.

© 2003 Published by Elsevier B.V. Open access under [CC BY license](https://creativecommons.org/licenses/by/4.0/).

PACS: 24.85.+p; 12.39.Ki; 24.10.Jv; 21.30.Fe

Keywords: In-medium form factors; Quark–meson coupling model; Nucleon structure effect; Constituent quark model

The change of hadron properties in a nuclear medium is of fundamental importance in understanding the implication of QCD for nuclear physics. One of the most famous nuclear medium effects may be the nuclear EMC effect [1], and it has stimulated theoretical and experimental efforts to seek nuclear quark–gluon effects for almost two decades.

Recently, the search for modification of the electromagnetic form factors of bound protons has been performed in polarized ($\vec{e}, e' \vec{p}$) scattering experiments on ^{16}O and ^4He nuclei [2]. The experiments measured the ratio of transverse to longitudinal polarization of

the ejected proton, which is proportional to the ratio of electric to magnetic form factors of a proton. However, conventional calculations including free-proton form factors, appropriate optical potentials and bound-state wave functions as well as relativistic corrections, meson-exchange currents (MEC), isobar contributions and final-state interactions, fail to reproduce the observed results in ^4He [2,3]. Indeed, full agreement with the experimental data was only obtained when, in addition to the standard nuclear calculation, a change in the form factors which is caused by the structure modification of bound proton [2,4], was taken into account.

Recent inclusive neutrino experiments on ^{12}C at Los Alamos [5] also suggest that the measured total cross section is about a half of the standard, relativistic shell model calculation including final-state inter-

E-mail addresses: ksaito@tohoku-pharm.ac.jp (K. Saito), tsushima@physast.uga.edu (K. Tsushima).

actions within the distorted wave impulse approximation [6]. In the neutrino reaction, the charged-current vector form factors of bound nucleons are slightly enhanced, while the axial form factors are quenched by the nuclear medium [7]. Finally, the effect of the bound nucleon form factors reduces the total cross section by about 8% relative to that calculated with the free form factors [8]. We stress that this correction is caused by the change of the internal quark wave function at the mean-field level and hence there is no obvious double counting with MEC etc. This is a new effect which should be taken into account additionally to the standard nuclear corrections.

Furthermore, the measurements of polarization transfer observables in exclusive $(\vec{p}, 2\vec{p})$ proton knock-out reactions from various nuclei [9,10] again indicate that it is difficult to account for the measured polarization transfers within the conventional, relativistic distorted wave impulse approximation [11]. To reproduce the measured spin observables, it is necessary to simultaneously reduce the scalar (σ) and vector (ω) coupling constants and the meson masses by about 10–20% [11]. In particular, the analyzing power (A_y), polarization (P) and spin transfer coefficient ($D_{ss'}$) are very sensitive to the change of σ - and ω -nucleon coupling constants and their masses in a nuclear medium. These may again imply the change in the internal structure of bound nucleons.

If the quark substructure of the nucleon is modified depending on the nuclear environment, it would leave traces in a variety of processes and observables, including various form factors. These modifications of bound nucleons can be successfully described within the context of the quark–meson coupling (QMC) model [12]. In the model, the medium effects arise through the self-consistent coupling of σ and ω mesons to confined quarks, rather than to the nucleons. As a result, the internal structure of the bound nucleon is modified by the surrounding nuclear medium.

The electromagnetic form factors of bound nucleons [4] have been studied using an improved cloudy bag model (ICBM) [13,14], together with the QMC model. In the ICBM, a simplified Peierls–Thouless projection technique (the weight function $w(\vec{p})$ appearing in the nucleon wave function is assumed to be unity) is used to account for center of mass (c.m.) and recoil corrections. In addition to it, a Lorentz contraction of the internal quark wave function is included.

The axial form factor in nuclear matter has also been calculated in a similar manner [7]. Furthermore, the form factors at σ - and ω -nucleon strong-interaction vertices in a nuclear medium should also be investigated. The change of these form factors is very significant in understanding how the strong interaction is modified in nuclear matter. It is also expected to play an important role in analyzing the polarization transfer observables in the exclusive $(\vec{p}, 2\vec{p})$ reactions [9,10].

In this Letter, we study the scalar and vector form factors at σ - and ω -nucleon strong-interaction vertices in symmetric nuclear matter. We shall calculate these form factors using a relativistic constituent quark model with a harmonic oscillator (HO) [15] or a linearly rising (LR) confining potential [16] and the Peierls–Yoccoz (PY) projection technique. If we use the “minimax” principle (or the saddle point variational principle) [16,17], it is easy to obtain an approximate solution to the Dirac equation with *any* potential. Since we choose a Gaussian wave function for a confined quark as ansatz, it is possible to calculate the form factors analytically and thus transparent to see how the PY projection and the Lorentz contraction of the quark wave function work in the form factors. Instead, in this exploratory study, we do not include the pion cloud effect which can explicitly be treated in the ICBM. (We will study this effect in a forthcoming paper.)

In the QMC model, the mean-field approximation is applied to the σ and ω meson fields, which couple to confined (u or d) quarks in nuclear matter. Each quark then satisfies the Dirac equation

$$[-i\vec{\alpha} \cdot \vec{\nabla} + \gamma^0 m_q^* + U_{\text{conf}}(r)]\psi(\vec{r}) = E_q \psi(\vec{r}), \quad (1)$$

where $m_q^* = m_q - g_\sigma^q \bar{\sigma}$ and $E_q = \epsilon_q - g_\omega^q \bar{\omega}$ with ϵ_q the quark energy. We take the free quark mass m_q to be 300 MeV. The mean-field values of σ and ω mesons are respectively denoted by $\bar{\sigma}$ and $\bar{\omega}$, and g_σ^q and g_ω^q are the corresponding quark and meson coupling constants. We use a confining potential of HO type, $U_{\text{conf}}(r) = (c/2)(1 + \beta\gamma^0)r^2$, or a LR one, $U_{\text{conf}}(r) = (\lambda/2)(1 + \beta\gamma^0)r$, where β ($0 \leq \beta \leq 1$) controls the strength of the Lorentz vector-type potential. The potential strength is taken to be $c = 0.04 \text{ GeV}^3$ or $\lambda = 0.2 \text{ GeV}^2$ [18].

Although for the LR potential the Dirac equation cannot be solved analytically, the minimax principle allows us to obtain an approximate solution very eas-

ily and accurately [17]. Since the Dirac Hamiltonian does not have a lower bound for the energy spectrum, the usual variational method cannot be applied. The minimax principle amounts to minimizing (maximizing) the energy expectation value of the upper (lower) component of the quark wave function with respect to variational parameters. A trial wave function for the lowest-energy state is usually chosen as

$$\psi(\vec{r}) = N_0 \begin{pmatrix} u(r) \\ i\xi b \vec{\sigma} \cdot \vec{r} u(r) \end{pmatrix} \chi_s, \quad (2)$$

with N_0 a normalization constant, $u(r) = e^{-b^2 r^2/2}$ and b and ξ the variational parameters. These parameters are determined so as to minimize the quark energy ϵ_q with respect to b and maximize it with respect to ξ . Note that for the HO potential with $\beta = 1$ this gives the exact solution [19].

First, we fix the parameters of the model in vacuum. The nucleon mass in vacuum ($\bar{\sigma} = \bar{\omega} = 0$) is given by $M_N = 3\epsilon_q - \epsilon_0$, where ϵ_0 accounts for corrections of c.m. and gluon fluctuations. The parameter ϵ_0 is fitted so as to obtain the free nucleon mass M_N ($= 939$ MeV). The minimax principle then determines the parameters b and ξ . These values are given in Table 1.

In matter, the scalar field couples to the confined quark and hence the quark mass changes depending on the nuclear environment. The nucleon mass in matter M_N^* is then reduced because the σ exchange induces an attractive force between nucleons. In an isosymmetric nuclear matter, the total energy (per nucleon) at nuclear density ρ_B is given by the usual expression in the QMC model [12]

$$E_{\text{tot}} = \frac{4}{\rho_B (2\pi)^3} \int^{k_F} d\vec{k} \sqrt{M_N^{*2} + \vec{k}^2}$$

Table 1

Coupling constants, ϵ_0 , b , ξ , M_N and nuclear incompressibility K . The parameters ϵ_0 , b and ξ are fixed in vacuum, while b^* , ξ^* and M_N^* are calculated at normal nuclear matter density. Here ϵ_0 , b and K are quoted in GeV. The value of β is specified in the first column

	g_σ^2	g_ω^2	ϵ_0	b	ξ	b^*/b	ξ^*/ξ	M_N^*/M_N	K
HO(0)	88.64	120.8	1.08	0.380	0.288	0.941	1.12	0.649	0.392
HO(0.5)	75.38	91.88	1.38	0.425	0.351	0.946	1.14	0.720	0.344
HO(1)	65.12	69.65	1.63	0.464	0.401	0.955	1.15	0.774	0.316
LR(0)	93.95	133.0	1.30	0.364	0.249	0.932	1.11	0.619	0.427
LR(0.5)	85.21	113.5	1.75	0.418	0.304	0.934	1.13	0.667	0.381
LR(1)	76.78	95.16	2.15	0.464	0.349	0.939	1.13	0.712	0.352

$$+ \frac{m_\sigma^2}{2\rho_B} \bar{\sigma}^2 + \frac{g_\omega^2}{2m_\omega^2} \rho_B, \quad (3)$$

where m_σ ($= 550$ MeV) and m_ω ($= 783$ MeV) are respectively the σ and ω meson masses, and g_ω ($= 3g_\omega^q$) is the ω -nucleon coupling constant. The values of the scalar and vector mean fields are, respectively, determined by self-consistency conditions: $(\partial E_{\text{tot}}/\partial \bar{\sigma}) = 0$ and $(\partial E_{\text{tot}}/\partial \bar{\omega}) = 0$. The latter condition ensures baryon number conservation, while the former gives a transcendental equation for the scalar field in matter.

The coupling constants are fitted so as to reproduce the nuclear saturation property ($E_{\text{tot}} - M_N = -15.7$ MeV) at normal nuclear matter density ρ_0 ($= 0.17$ fm $^{-3}$). Note that for each value of ρ_B one has to use the minimax principle to obtain the in-medium parameters b and ξ . The coupling constants and nuclear properties at ρ_0 are listed in Table 1. The σ -nucleon coupling constant g_σ is defined in terms of the quark scalar density S_N : $g_\sigma = 3g_\sigma^q S_N$ ($\bar{\sigma} = 0$), where $S_N(\bar{\sigma}) = \int d\vec{r} \bar{\psi}(r) \psi(r)$.

The wave function for a nucleon moving with momentum \vec{p} can be constructed by the PY projection technique [13,20]:

$$\Psi(\vec{r}_1, \vec{r}_2, \vec{r}_3; \vec{p}) = N(\vec{p}) \int d\vec{x} e^{i\vec{p}\cdot\vec{x}} \Phi(\vec{r}_1, \vec{r}_2, \vec{r}_3; \vec{x}), \quad (4)$$

where $N(\vec{p})$ is a momentum-dependent normalization constant

$$[N(\vec{p})]^{-2} = \int d\vec{r} e^{-i\vec{r}\cdot\vec{p}} [\rho(\vec{r})]^3, \quad (5)$$

with

$$\rho(\vec{r}) = \int \frac{d\vec{k}}{(2\pi)^3} e^{i\vec{k}\cdot\vec{r}} |\phi(\vec{k})|^2. \quad (6)$$

Here ϕ is the quark wave function in momentum space. The localized state Φ is simply given by a

product of the three individual quark wave function

$$\Phi(\vec{r}_1, \vec{r}_2, \vec{r}_3; \vec{x}) = \psi(\vec{r}_1 - \vec{x})\psi(\vec{r}_2 - \vec{x})\psi(\vec{r}_3 - \vec{x}), \quad (7)$$

where \vec{x} refers to the location of the center of the nucleon and \vec{r}_j ($j = 1, 2, 3$) specifies the position of the j th quark.

Because the nucleon consists of three point-like quarks, the expectation value of an operator with respect to the nucleon wave function Eq. (4) may be given by a sum of the individual quark expectation values [20]. In the Breit frame, where the initial (final) momentum of the nucleon is taken to be $-\vec{q}/2$ ($\vec{q}/2$) with \vec{q} the momentum transfer, the scalar and vector form factors are respectively given by

$$\Gamma_{(\nu)}^{(s)}(Q^2) = 3[N(Q^2)]^2 \int d\vec{r} e^{i\vec{q}\cdot\vec{r}} \tilde{\psi}(\vec{r}) \left(\frac{1}{\gamma_0} \right) \tilde{\psi}(\vec{r}, \vec{q}), \quad (8)$$

where $Q^2 \equiv -q_0^2 + \vec{q}^2 = \vec{q}^2$, and we ignore a small tensor term at the ω -nucleon coupling. $\tilde{\psi}$ in Eq. (8) is represented by

$$\tilde{\psi}(\vec{r}, \vec{q}) = \int \frac{d\vec{k}}{(2\pi)^3} e^{i\vec{k}\cdot\vec{r}} \phi(\vec{k}) W(\vec{k}, \vec{q}), \quad (9)$$

where

$$W(\vec{k}, \vec{q}) = \int d\vec{r} e^{-i(\vec{q}/2 + \vec{k})\cdot\vec{r}} [\bar{\rho}(\vec{r})]^2, \quad (10)$$

and

$$\bar{\rho}(\vec{r}) = \int \frac{d\vec{k}}{(2\pi)^3} e^{-i\vec{k}\cdot\vec{r}} \bar{\phi}(\vec{k}) \phi(\vec{k}). \quad (11)$$

Now we can calculate the scalar and vector form factors in nuclear matter analytically:

$$\Gamma_{(\nu)}^{(s)}(Q^2, \rho_B) = \left(\frac{Z_0(\xi^2)}{Y_0^v(\xi^2)} \right) e^{-x^2/6} \frac{\sum_{i=0}^2 (x^2)^i Y_i^{(s)}(\xi^2)}{\sum_{i=0}^3 (x^2)^i Z_i(\xi^2)}, \quad (12)$$

where $x^2 = Q^2/b^2$,

$$Z_0(\xi^2) = 1 + 3\xi^2 + \frac{7}{2}\xi^4 + \frac{25}{18}\xi^6, \quad (13)$$

$$Z_1(\xi^2) = \frac{1}{12}\xi^2 + \frac{1}{9}\xi^4 + \frac{13}{216}\xi^6, \quad (14)$$

$$Z_2(\xi^2) = \frac{1}{432}\xi^4 + \frac{1}{1296}\xi^6, \quad (15)$$

$$Z_3(\xi^2) = \frac{1}{46656}\xi^6, \quad (16)$$

and

$$Y_0^{(s)}(\xi^2) = 1 - \binom{3}{1}\xi^2 + \binom{7}{-7}\xi^4 \mp \frac{25}{18}\xi^6, \quad (17)$$

$$Y_1^{(s)}(\xi^2) = \frac{\binom{9}{-7}}{32}\xi^2 + \frac{\binom{-69}{128}}{128}\xi^4 \pm \frac{335}{1152}\xi^6, \quad (18)$$

$$Y_2^{(s)}(\xi^2) = \pm \frac{1}{128}\xi^4. \quad (19)$$

Recall that the variational parameters ξ and b (thus x^2), which appear in the quark wave function, depend on ρ_B . We have renormalized the vector form factor so that $\Gamma_v = 1$ is maintained at zero momentum transfer. The scalar form factor is also rescaled by the same factor as in the vector case [20].

In contrast, if the c.m. correction is ignored, the form factors are given by

$$\Gamma_{(\nu)}^0(Q^2, \rho_B) = \frac{e^{-x^2/4}}{(1 + \frac{3}{2}\xi^2)} \left[1 \mp \frac{3}{2}\xi^2 \left(1 - \frac{1}{6}x^2 \right) \right]. \quad (20)$$

Because x^2 is small and $\xi \lesssim 0.5$ (for $\rho_B/\rho_0 \leq 2.0$) at small momentum transfer, we can expand the form factors. Up to $\mathcal{O}(x^2)$ or $\mathcal{O}(\xi^2)$, we find that $\Gamma_{(\nu)}^0 = 1 - \binom{3}{0}\xi^2 - x^2/4$, while Eq. (12) gives $\Gamma_{(\nu)}^{(s)} = 1 - \binom{2}{0}\xi^2 - x^2/6$. The c.m. correction thus moderates the reduction of the form factors.

Apart from the c.m. correction, it is also vital to include the Lorentz contraction of the internal quark wave function at moderate or large momentum transfer [13,21]. The full form factors $\tilde{\Gamma}_{(\nu)}^{(s)}$ can be obtained through a simple rescaling [4,13]

$$\Gamma_{(\nu)}^{(s)}(Q^2) \rightarrow \tilde{\Gamma}_{(\nu)}^{(s)}(Q^2) = \eta^* \Gamma_{(\nu)}^{(s)}(\eta^* Q^2), \quad (21)$$

where $\eta^* = (M_N^*/E_N^*)^2$ with $E_N^* = \sqrt{M_N^{*2} + Q^2}/4$. The scaling factor in the argument arises from the coordinate transformation of the struck quark and the prefactor η^* comes from the reduction of the integral measure of two spectator quarks in the Breit frame [4, 13]. Thus, the scaling factor η^* (in vacuum η with M_N) should appear in any nucleon(baryon)–meson form factors if the nucleon (baryon) is assumed to have a three-quark cluster structure.

To illustrate the effects of the c.m. correction and Lorentz contraction on the form factors, we show in Fig. 1 the vector form factor in vacuum. The c.m.

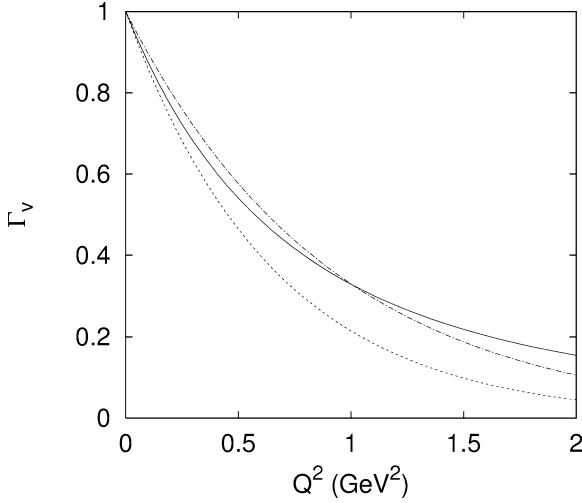


Fig. 1. Vector form factor in vacuum (for the HO case with $\beta = 0.5$). The full result is denoted by the solid curve, while the dot-dashed curve shows the result with the c.m. correction but without the Lorentz contraction. The result without both corrections (Eq. (20)) is denoted by the dotted curve.

correction considerably enhances the form factor in comparison with the result without both effects (see the dotted and dot-dashed curves in the figure). The effect of Lorentz contraction is also important. If the Lorentz contraction is ignored, the form factor drops away like $\sim e^{-x^2/6}$ at large Q^2 . The inclusion of the Lorentz contraction removes this objectionable exponential falloff. Because of the factor η , the form factor is proportional to $1/(1 + Q^2/\Lambda^2)$ and x^2 is modified to $x^2/(1 + Q^2/\Lambda^2)$ with $\Lambda = 2M_N$ (see also Eq. (21)). As a result, the inclusion of the Lorentz contraction enhances the form factor at large Q^2 (see the dot-dashed and solid curves).

Because our aim is to study the density dependence of the form factors in nuclear matter, we consider the ratios of the in-medium form factors to those in vacuum:

$$R_{(v)}^{(s)}(Q^2, \rho_B) = \frac{\tilde{\Gamma}_{(v)}^{(s)}(Q^2, \rho_B)}{\tilde{\Gamma}_{(v)}^{(s)}(Q^2, \rho_B = 0)}. \quad (22)$$

The form factors in symmetric nuclear matter $F_{(v)}^{(s)}$ are thus given by

$$F_{(v)}^{(s)}(Q^2, \rho_B) = R_{(v)}^{(s)}(Q^2, \rho_B) \times F_{(v)}^{\text{emp}}(Q^2), \quad (23)$$

where $F_{(v)}^{\text{emp}}$ are the form factors empirically determined in vacuum [22]. In Fig. 2, the ratio of the in-medium scalar (vector) form factor to that in vacuum is illustrated as a function of Q^2 and ρ_B . (Because the ratios for the LR potential are similar to those for the HO potential, we focus on the HO case for a while.) At $\rho_B/\rho_0 = 1$ and $Q^2 = 1.0 \text{ GeV}^2$, the in-medium scalar (vector) form factor is reduced by 15 (14)% relative to that in vacuum. The reduction rate depends on β very weakly. By contrast, at $\rho_B/\rho_0 = 2$ and $Q^2 = 1.0 \text{ GeV}^2$, the scalar form factor decreases by 35 (29) [24]% for $\beta = 0$ (0.5) [1.0], while the vector form factor diminishes by 28 (26) [22]% for $\beta = 0$ (0.5) [1.0]. At high density the dependence of the reduction on β is thus rather strong, and the reduction rate is correlated with M_N^* (see Table 1).

As in the case of vacuum (see Fig. 1), the effect of Lorentz contraction is again seen at large Q^2 . For example, at $\rho_B/\rho_0 = 2$ and $Q^2 = 1 \text{ GeV}^2$, the vector form factor with the Lorentz contraction is about 7% larger than that without it. We also note that, in the HO case with $\beta = 0.5$, the full vector form factor gives the root-mean-square radius of 0.53 fm. If we neglect the Lorentz contraction effect, it is 0.46 fm.

Finally, we parameterize the ratios for the scalar and vector form factors in Eq. (23). Such parameterizations are very useful in analyzing the experimental results, e.g., for the exclusive $(\vec{p}, 2\vec{p})$ proton knockout reactions [9,10]. With an error less than 0.2%, the ratios can be represented by

$$R_{(v)}^{(s)}(Q^2, \rho_B) = 1 + A_{(v)}^{(s)}(Q^2)(\rho_B/\rho_0) + B_{(v)}^{(s)}(Q^2)(\rho_B/\rho_0)^2, \quad (24)$$

where

$$A_s(y) = -\begin{pmatrix} 0.06829 \\ 0.06323 \end{pmatrix} - \begin{pmatrix} 0.2302 \\ 0.2464 \end{pmatrix} y + \begin{pmatrix} 0.1845 \\ 0.1711 \end{pmatrix} y^2 - \begin{pmatrix} 0.04613 \\ 0.04072 \end{pmatrix} y^3, \quad (25)$$

$$A_v(y) = -\begin{pmatrix} 0.3856 \\ 0.3738 \end{pmatrix} y + \begin{pmatrix} 0.3021 \\ 0.2668 \end{pmatrix} y^2 - \begin{pmatrix} 0.07763 \\ 0.06494 \end{pmatrix} y^3, \quad (26)$$

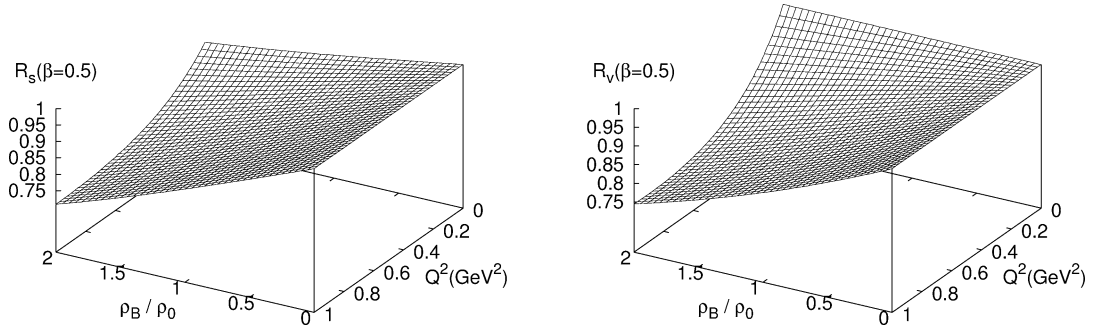


Fig. 2. Ratios for the scalar (left panel) and vector (right panel) form factors in the case of the HO potential with $\beta = 0.5$.

and

$$B_s(y) = \begin{pmatrix} 0.005071 \\ 0.003569 \end{pmatrix} - \begin{pmatrix} 0.02499 \\ 0.03304 \end{pmatrix} y \\ + \begin{pmatrix} 0.09473 \\ 0.1167 \end{pmatrix} y^2 - \begin{pmatrix} 0.1070 \\ 0.1245 \end{pmatrix} y^3 \\ + \begin{pmatrix} 0.03914 \\ 0.04474 \end{pmatrix} y^4, \quad (27)$$

$$B_v(y) = - \begin{pmatrix} 0.04296 \\ 0.05500 \end{pmatrix} y + \begin{pmatrix} 0.2081 \\ 0.2271 \end{pmatrix} y^2 \\ - \begin{pmatrix} 0.2380 \\ 0.2509 \end{pmatrix} y^3 + \begin{pmatrix} 0.08967 \\ 0.09337 \end{pmatrix} y^4. \quad (28)$$

In Eqs. (25)–(28), the upper (lower) numbers are for the case of the HO (LR) potential with $\beta = 0.5$ (1.0), which provides the effective nucleon mass $M_N^*/M_N = 0.71$ – 0.72 at ρ_0 (see Table 1). The in-medium form factors are thus given by Eqs. (23) and (24).

In summary, using the QMC model we have calculated the form factors at σ - and ω -nucleon strong-interaction vertices in symmetric nuclear matter. We have applied both the PY projection technique and the Lorentz contraction of the internal quark wave function. The form factors are reduced by the nuclear medium relative to those in vacuum. The c.m. correction moderates the reduction of the form factors in matter, and the Lorentz contraction is vital at large momentum transfer. We have found that the reduction in the scalar form factor is about 15% at $\rho_B/\rho_0 = 1$ and $Q^2 = 1 \text{ GeV}^2$. This rate is almost identical to that for the vector form factor. In contrast, the scalar and vector form factors are respectively reduced by about 30% and 25% at $\rho_B/\rho_0 = 2$ and $Q^2 = 1 \text{ GeV}^2$. The reduction of the form factors is expected better to reproduce the polarization transfer observables measured at

RCNP and iThemba laboratory [9,10]. We have parameterized the ratios of the form factors in symmetric nuclear matter to those in vacuum. This provides a convenient formula to estimate the in-medium form factors. It is very intriguing to reanalyze the data of polarization transfer observables for exclusive $(\vec{p}, 2\vec{p})$ proton knockout reactions [9,10] including the modification of both the form factors and meson masses in matter [23].

Acknowledgement

K.T. is supported by the Forschungszentrum-Jülich, contract No. 41445282 (COSY-058).

References

- [1] D.F. Geesaman, K. Saito, A.W. Thomas, *Annu. Rev. Nucl. Part. Sci.* 45 (1995) 337.
- [2] S. Malov, et al., *Phys. Rev. C* 62 (2000) 057302; S. Dieterich, et al., *Phys. Lett. B* 500 (2001) 47; R.D. Ransome, *Nucl. Phys. A* 699 (2002) 360c; S. Strauch, et al., *Phys. Rev. Lett.* 91 (2003) 052301; JLab E93-049 Collaboration, S. Strauch, *nucl-ex/0308026*.
- [3] J.J. Kelly, *Phys. Rev. C* 60 (1999) 044609; J.M. Udias, J.R. Vignote, *Phys. Rev. C* 62 (2000) 034302.
- [4] D.H. Lu, A.W. Thomas, K. Tsushima, A.G. Williams, K. Saito, *Phys. Lett. B* 417 (1998) 217; D.H. Lu, K. Tsushima, A.W. Thomas, A.G. Williams, K. Saito, *Phys. Rev. C* 60 (1999) 068201.
- [5] LSND Collaboration, L.B. Auerbach, et al., *Phys. Rev. C* 66 (2002) 015501.
- [6] C. Maieron, M.C. Maartinez, J.A. Caballero, J.M. Udias, *nucl-th/0303075*.
- [7] D.H. Lu, A.W. Thomas, K. Tsushima, *nucl-th/0112001*.
- [8] K. Tsushima, H. Kim, K. Saito, *nucl-th/0307013*.

- [9] T. Noro, et al., *Nucl. Phys. A* 629 (1998) 324c.
- [10] R. Neveling, et al., *Phys. Rev. C* 66 (2002) 034602.
- [11] G.C. Hillhouse, J. Mano, S.M. Wyngaardt, B.I.S. van der Ventel, T. Noro, K. Hatanaka, *Phys. Rev. C* 68 (2003) 034608.
- [12] P.A.M. Guichon, *Phys. Lett. B* 200 (1988) 235;
K. Saito, A.W. Thomas, *Phys. Lett. B* 327 (1994) 9;
P.A.M. Guichon, K. Saito, E. Rodionov, A.W. Thomas, *Nucl. Phys. A* 601 (1996) 349;
K. Saito, K. Tsushima, A.W. Thomas, *Nucl. Phys. A* 609 (1996) 339.
- [13] D.H. Lu, A.W. Thomas, A.G. Williams, *Phys. Rev. C* 57 (1998) 2628.
- [14] A.W. Thomas, *Adv. Nucl. Phys.* 13 (1984) 1;
G.A. Miller, *Int. Rev. Nucl. Phys.* 1 (1984) 190.
- [15] H. Toki, U. Meyer, A. Faessler, R. Brockmann, *Phys. Rev. C* 58 (1998) 3749.
- [16] G. Krein, V.E. Vizcarra, in: V. Guzey, A. Kizilersu, T. Nagae, A.W. Thomas (Eds.), *Proc. of the Workshop on Physics at the Japan Hadron Facility*, World Scientific, Singapore, 2002, nucl-th/0206047.
- [17] J. Franklin, R.L. Intemann, *Phys. Rev. Lett.* 54 (1985) 2068;
J.D. Talman, *Phys. Rev. Lett.* 57 (1986) 1091.
- [18] V.V. Khrushev, V.I. Savrin, S.V. Semenov, *Phys. Lett. B* 525 (2002) 283.
- [19] N. Barik, B.K. Dash, M. Das, *Phys. Rev. D* 31 (1985) 1652;
K. Saito, *Prog. Theor. Phys.* 108 (2002) 609.
- [20] R. Tegen, R. Brockmann, W. Weise, *Z. Phys. A* 307 (1982) 339.
- [21] A.L. Licht, A. Pagnamenta, *Phys. Rev. D* 2 (1970) 1150;
A.L. Licht, A. Pagnamenta, *Phys. Rev. D* 2 (1970) 1156.
- [22] For example, R. Machleidt, *Adv. Nucl. Phys.* 19 (1989) 189.
- [23] K. Saito, K. Tsushima, A.W. Thomas, *Phys. Rev. C* 55 (1997) 2637.