Variation of hadron masses in nuclear matter in the relativistic Hartree approximation

H. Shen\textsuperscript{a} S. Tamenaga\textsuperscript{b} H. Toki\textsuperscript{b}

\textsuperscript{a}Department of physics, Nankai University, Tianjin 300071, China
\textsuperscript{b}Research Center for Nuclear Physics (RCNP), Osaka University, Ibaraki, Osaka 567-0047, Japan

Abstract

We study the modification of hadron masses due to the vacuum polarization using the chiral sigma model, which is extended to generate the $\omega$ meson mass by the sigma condensation in the vacuum in the same way as the nucleon mass. The results obtained in the chiral sigma model are compared with those obtained in the Walecka model which includes $\sigma$ and $\omega$ mesons in a non-chiral fashion. It is shown that both the nucleon mass and the $\omega$ meson mass decrease in nuclear medium, while the $\sigma$ meson mass increases at finite density in the chiral sigma model.

Key words: relativistic Hartree approximation, chiral sigma model

PACS: 21.65.+f, 21.30.Fe, 11.30.Rd

1 Introduction

One of the most interesting topics in nuclear physics is to study how the hadron properties are altered as the environment changes. In particular, the medium modification of hadron masses has attracted a lot of attention both experimentally and theoretically. The observation of enhanced dilepton production from relativistic heavy ion collision experiments [1] could be due to a reduction in the vector meson masses in the medium. Brown and Rho suggested the hypothesis that the vector meson masses drop in nuclear medium according to a simple scaling law [2]. There are many theoretical efforts made to understand the behavior of hadrons in dense matter, including the various

Email addresses: songtc@public.tpt.tj.cn (H. Shen), stame@rcnp.osaka-u.ac.jp (S. Tamenaga), toki@rcnp.osaka-u.ac.jp (H. Toki).
QCD-based methods like QCD sum rules [3] and phenomenological nuclear models such as relativistic mean field approach [4] and quark meson coupling model [5].

The relativistic mean field theory (RMF) has been developed and widely applied to a great variety of problems in nuclear physics [4]. The original version of the RMF theory proposed by Walecka (Walecka model) consists of baryons interacting with each other via the exchange of $\sigma$ and $\omega$ mesons [4]. This model and its later variations can be solved in the mean field approximation, by replacing the meson field operators with their classical expectation values. Since the Walecka model is renormalizable, there is a standard procedure for renormalization of the vacuum polarization contribution carried out by adding required counterterms to the original Lagrangian and subtracting purely vacuum expectation values. It has been discussed that the Walecka model with the vacuum polarization effects taken into account could reproduce reasonably well the saturation properties of nuclear matter and the ground state properties of finite nuclei [4]. However, the Walecka model does not respect the chiral symmetry, which is known to be a very important feature in hadron physics.

The chiral symmetry can be described nicely in the linear sigma model introduced by Gell-Mann and Levy [6], which has been used for various phenomena in hadron physics. It is very natural to use the chiral sigma model for the description of nuclear matter and finite nuclei in the relativistic mean field approximation [4]. It was found that the use of the chiral sigma model in its original form was not satisfactory for the description of nuclear matter. Boguta introduced a dynamical generation of the $\omega$ meson mass in the same way as the nucleon mass, so that a saturating equation of state for nuclear matter could be obtained in the chiral sigma model [7]. This chiral sigma model was applied to study finite nuclei by several groups [8,9].

Many authors have studied the medium modification of hadron masses using some non-chiral models [10,11,12,13]. It has been pointed out that the polarization of the Dirac sea is the most important reason for the reduction of vector meson masses in medium [11]. In this paper, we would like to investigate the variation of hadron masses due to vacuum polarization within the chiral sigma model, and compare with the results obtained in the Walecka model. In section 2, we briefly recapitulate the non-chiral Walecka model and the chiral sigma model, also discuss the procedure for renormalization in these models. In section 3, we explain the model parameters, and show the numerical results. Section 4 is devoted to the summary of this paper.
2 Formalism

In this section, we briefly recapitulate the effective Lagrangian and the renormalization procedure in the Walecka model and in the chiral sigma model. The details regarding the renormalization procedure can be found in earlier references \cite{14,15,16}.

The Lagrangian density of the Walecka model is well known, which involves an explicit description of nucleon and meson degrees of freedom \cite{4}. The Lagrangian density in the Walecka model is given as

\[
\mathcal{L} = \bar{\Psi} \left( i \gamma_\mu \partial^\mu - M - g_\sigma \sigma - g_\omega \gamma_\mu \omega^\mu \right) \Psi \\
+ \frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma - \frac{1}{2} m_\sigma^2 \sigma^2 - \frac{1}{4} W_{\mu\nu} W^{\mu\nu} + \frac{1}{2} m_\omega^2 \omega_\mu \omega^\mu + \delta \mathcal{L},
\]

where \( \Psi, \sigma \) and \( \omega \) are the fields for the nucleon, \( \sigma \) and \( \omega \) mesons with physical masses \( M, m_\sigma \) and \( m_\omega \), respectively, and \( W^{\mu\nu} \equiv \partial^\mu \omega^\nu - \partial^\nu \omega^\mu \). The term \( \delta \mathcal{L} = \delta \mathcal{L}_\sigma + \delta \mathcal{L}_\omega \) contains renormalization counterterms, which are introduced to remove the divergences in the loop calculations within the framework of the relativistic Hartree approximation \cite{4}. The renormalization procedure in the Walecka model has been extensively discussed in Refs. \cite{14,15,16}. Here, we adopt the subtraction scheme given in Ref. \cite{16}. The renormalization counterterms in the Walecka model can be written as

\[
\delta \mathcal{L}_\sigma = \alpha_1 \sigma + \frac{1}{2!} \alpha_2 \sigma^2 + \frac{1}{3!} \alpha_3 \sigma^3 + \frac{1}{4!} \alpha_4 \sigma^4 + \frac{1}{2} \zeta_\sigma \partial_\mu \sigma \partial^\mu \sigma,
\]

\[
\delta \mathcal{L}_\omega = -\frac{1}{4} \zeta_\omega W_{\mu\nu} W^{\mu\nu} - \frac{1}{2} \delta m_\omega^2 \omega_\mu \omega^\mu.
\]

The coefficients are specified by imposing appropriate renormalization conditions. First of all, \( \alpha_1 \) must completely cancel the loop contribution to ensure the stability of the vacuum. The coefficients \( \alpha_2 \) and \( \zeta_\sigma \) can be determined by requiring \( \Pi^R_\sigma|_{M^*=M, q^2=m_\sigma^2} = 0 \) and \( \left. \frac{\partial \Pi^R_\sigma}{\partial q^2} \right|_{M^*=M, q^2=m_\sigma^2} = 0 \). For \( \alpha_3 \) and \( \alpha_4 \), we adopt the usual conditions used in Ref. \cite{4}. The coefficients \( \delta m_\omega^2 \) and \( \zeta_\omega \) can be determined by imposing \( D_\omega|_{M^*=M, q^2=m_\omega^2} = 0 \) and \( \left. \frac{\partial D_\omega}{\partial q^2} \right|_{M^*=M, q^2=m_\omega^2} = 1 \), where \( D_\omega = q^2 - m_\omega^2 + \delta m_\omega^2 - q^2 \Pi^R_\omega \). The explicit expressions for the renormalized meson self-energies in the Walecka model, \( \Pi^R_\sigma \) and \( \Pi^R_\omega \), can be found in Ref. \cite{16}. After carrying out the renormalization procedure, we study the effective masses of \( \sigma \) and \( \omega \) mesons, \( m^*_\sigma \) and \( m^*_\omega \), which can be obtained by searching for the zeros of the inverse propagators,

\[
D_\sigma(M^*, q^2 = m^2_\sigma) = q^2 - m^2_\sigma - \Pi^R_\sigma(M^*, q^2) = 0
\]

\[
D_\omega(M^*, q^2 = m^2_\omega) = q^2 - m^2_\omega + \delta m^2_\omega - q^2 \Pi^R_\omega(M^*, q^2) = 0
\]
We now turn to the chiral sigma model used in Ref. [9]. The Lagrangian density of the chiral sigma model is written as

\[
\mathcal{L} = \bar{\Psi} \left[ i\gamma_\mu \partial^\mu - g_\sigma (\sigma + i\gamma_5 \vec{\tau} \cdot \vec{\pi}) - g_\omega \gamma_\mu \omega^\mu \right] \Psi \\
+ \frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma + \frac{1}{2} \partial_\mu \vec{\pi} \partial^\mu \vec{\pi} - \frac{\mu^2}{2} \left( \sigma^2 + \vec{\pi}^2 \right) - \frac{\lambda}{4} \left( \sigma^2 + \vec{\pi}^2 \right)^2 \\
- \frac{1}{4} W_{\mu\nu} W^{\mu\nu} + \frac{1}{2} \tilde{g}_\omega^2 \left( \sigma^2 + \vec{\pi}^2 \right) \omega_\mu \omega^\mu \\
+ \varepsilon \sigma + \delta \mathcal{L}.
\]  

Here \( \Psi, \pi, \sigma \) and \( \omega \) are the fields for the nucleon, \( \pi, \sigma \) and \( \omega \) mesons. The \( \omega \) meson mass can be generated dynamically by the sigma condensation in the vacuum in the same way as the nucleon mass[7]. In the Lagrangian, an explicit chiral symmetry breaking term \( \varepsilon \sigma \) has been involved, while the term \( \delta \mathcal{L} \) contains the renormalization counterterms. To realize the chiral symmetry in the Nambu-Goldstone mode, a nonzero vacuum expectation value of the \( \sigma \) field, \( \langle \sigma \rangle = \sigma_0 \), is obtained by minimizing the meson effective potential. We now define the new fluctuation field \( \varphi = \sigma - \sigma_0 \), the above Lagrangian is rewritten as

\[
\mathcal{L} = \bar{\Psi} \left[ i\gamma_\mu \partial^\mu - M - g_\sigma \varphi - g_\sigma i\gamma_5 \vec{\tau} \cdot \vec{\pi} - g_\omega \gamma_\mu \omega^\mu \right] \Psi \\
+ \frac{1}{2} \partial_\mu \varphi \partial^\mu \varphi + \frac{1}{2} \partial_\mu \vec{\pi} \partial^\mu \vec{\pi} - \frac{1}{4} W_{\mu\nu} W^{\mu\nu} \\
- \frac{1}{2} m_\pi^2 \vec{\pi}^2 - \frac{1}{2} m_\sigma^2 \varphi^2 + \frac{1}{2} m_\omega^2 \omega_\mu \omega^\mu \\
- \lambda \sigma_0 \varphi^3 - \frac{\lambda}{4} \varphi^4 - \frac{\lambda}{4} \left( 4 \sigma_0 \varphi + 2 \varphi^2 + \vec{\pi}^2 \right) \vec{\pi}^2 \\
+ \frac{1}{2} \tilde{g}_\omega^2 \left( 2 \sigma_0 \varphi + \varphi^2 + \vec{\pi}^2 \right) \omega_\mu \omega^\mu \\
+ \left( \varepsilon - \mu^2 \sigma_0 - \lambda \sigma_0^3 \right) \varphi + \delta \mathcal{L}.
\]  

Here we have dropped a non-essential c-number constant. The energy minimum condition requires the term linear in \( \varphi \) to be zero, \( \varepsilon - \mu^2 \sigma_0 - \lambda \sigma_0^3 = 0 \). The physical masses are related with the parameters in the Lagrangian as

\[
M = g_\sigma \sigma_0, \\
m_\pi^2 = \mu^2 + \lambda \sigma_0^2 = \varepsilon / \sigma_0, \\
m_\sigma^2 = \mu^2 + 3 \lambda \sigma_0^2, \\
m_\omega^2 = \tilde{g}_\omega^2 \sigma_0^2.
\]  

The parameter \( \varepsilon \), which is proportional to \( m_\pi^2 \), represents the order of the explicit chiral symmetry breaking, and the exact chiral limit can be obtained
by setting $\varepsilon = 0$. To perform the renormalization for the chiral sigma model, we need the counterterms of the form

$$\delta L = \delta L_{\sigma \pi} + \delta L_{\omega}.$$  \hfill (12)

Here, $\delta L_{\omega}$ is identical to the one in the Walecka model. The term $\delta L_{\sigma \pi}$ can be written as

$$\delta L_{\sigma \pi} = \alpha_1 \varphi + \frac{1}{2!} \alpha_2 \varphi^2 + \frac{1}{3!} \alpha_3 \varphi^3 + \frac{1}{4!} \alpha_4 \varphi^4$$

$$+ \frac{1}{2!} \beta_2 \pi^2 + \frac{1}{2!} \beta_3 \varphi \pi^2 + \frac{1}{2!} \beta_4 \varphi^2 \pi^2$$

$$+ \frac{1}{2} \zeta_\sigma \partial_\mu \sigma \partial^\mu \sigma + \frac{1}{2} \zeta_\pi \partial_\mu \pi \partial^\mu \pi + \ldots$$  \hfill (13)

In order to respect the chiral symmetry in the renormalization procedure, $\delta L_{\sigma \pi}$ should get back a chiral symmetric form in the limit $\varepsilon \to 0$, as discussed in Ref. [15]. Therefore, much of the arbitrariness in the renormalization procedure is eliminated, and the coefficients $\alpha_i$ and $\beta_i$ in the chiral sigma model are related to each other. We again take the following renormalization conditions,

$$\Pi_{\sigma}^{R}(M^*, q^2 = m^2_\sigma) = 0 \quad \text{and} \quad \frac{\partial \Pi_{\sigma}^{R}}{\partial q^2}(M^*, q^2 = m^2_\sigma) = 0,$$

to specify the independent coefficients in the counterterms. Now we do not need any extra renormalization conditions to specify the coefficients $\alpha_3$ and $\alpha_4$, as done in the Walecka model. After carrying out the renormalization procedure, we can study the effective masses of $\sigma$ and $\omega$ mesons in nuclear medium, $m^*_\sigma$ and $m^*_\omega$. The effective masses are obtained by searching for the zeros of the inverse propagators,

$$D_\sigma(M^*, q^2 = m^2_\sigma) = q^2 - m^2_\sigma - 6\lambda f_\pi \varphi - 3\lambda \varphi^2 + \tilde{g}_\omega^2 \omega^2 - \Pi_{\sigma}^{R}(M^*, q^2) = 0$$  \hfill (14)

$$D_\omega(M^*, q^2 = m^2_\omega) = q^2 - m^2_\omega + \delta m^2_\omega - 2\tilde{g}_\omega^2 f_\pi \varphi - \tilde{g}_\omega^2 \varphi^2 - q^2\Pi_{\omega}^{R}(M^*, q^2) = 0$$  \hfill (15)

where $\Pi_{\omega}^{R}$ is the same as that in the Walecka model, but $\Pi_{\sigma}^{R}$ in the chiral sigma model is different from the one in the Walecka model due to the changes of the coefficients $\alpha_3$ and $\alpha_4$, which can be written as

$$(\Pi_{\sigma}^{R})_{\text{chiral}} = (\Pi_{\sigma}^{R})_{\text{Walecka}} - \Delta \alpha_3 \frac{M^* - M}{g_\sigma} - \frac{\Delta \alpha_4}{2} \left( \frac{M^* - M}{g_\sigma} \right)^2.$$  \hfill (16)

$\Delta \alpha_3$ and $\Delta \alpha_4$ are the differences between the coefficients in the two models, and we find

$$\Delta \alpha_3 = \frac{g^3_\sigma M}{\pi^2} \left\{ 2 - \frac{3m^2_\sigma}{4M^2} - \frac{9}{2} \int_0^1 dx \ln \left[ 1 - \frac{m^2_\sigma}{M^2} x(1 - x) \right] \right\}$$  \hfill (17)
\[ \Delta \alpha_4 = \frac{g^4}{\pi^2} \left\{ 8 - \frac{3m^2_\sigma}{4M^2} - \frac{9}{2} \int_0^1 dx \ln \left[ 1 - \frac{m^2_\sigma}{M^2} x(1 - x) \right] \right\}. \tag{18} \]

3 Numerical results

In Ref. [10], the modification of the \( \omega \) meson mass in nuclear medium due to the vacuum polarization has been studied within the Walecka model. They have taken the values of the masses as, \( M = 939 \) MeV, \( m_\omega = 783 \) MeV, and \( m_\sigma = 520 \) MeV. The coupling constants, \( g^2_\sigma = 66.117 \) and \( g^2_\omega = 79.927 \), were determined by requiring that the renormalized Hartree approximation could reproduce the binding energy \(-15.75\) MeV and the equilibrium Fermi momentum \(1.3\) fm\(^{-1}\) of nuclear matter. Here we use the same parameters to calculate the effective masses of \( \sigma \) and \( \omega \) mesons in nuclear medium, which can be obtained by searching for the zeros of the inverse propagators given in Eqs. (4) and (5). We show in Fig. 1 the effective masses of the nucleon, \( \sigma \) and \( \omega \) mesons as a function of the density. One can see that both \( M^* \) and \( m^*_\omega \) decrease at finite density in the Walecka model, which are in agreement with the results in Ref. [10]. The effective mass of \( \sigma \) meson, \( m^*_\sigma \), decreases at lower densities, and then slightly increases at higher densities. These results will be compared with those obtained in the chiral sigma model.
In the chiral sigma model, we take the vacuum expectation value of the $\sigma$ field as the pion decay constant, $\sigma_0 = f_\pi = 93$ MeV. We adopt the masses $M = 939$ MeV, $m_\pi = 139$ MeV, and $m_\omega = 783$ MeV from their experimental values. Then, the other parameters can be fixed automatically by the following relations, $g_\sigma = \frac{M}{f_\pi} = 10.1$ and $\tilde{g}_\omega = \frac{m_\omega}{f_\pi} = 8.42$. The coupling constants $\mu$ and $\lambda$ depend on $m_\pi$ and $m_\sigma$ through the relations given in Eqs. (9) and (10). In the present model, $m_\sigma$ and $g_\omega$ are taken as free parameters, which can be determined by reproducing the binding energy $-15.75$ MeV and the equilibrium Fermi momentum $1.3$ fm$^{-1}$ of nuclear matter. The fitted values for these two parameters in the renormalized Hartree approximation are $m_\sigma = 715$ MeV and $g_\omega = 4.025$.

We now present the results for hadron masses in nuclear medium using the chiral sigma model, which are obtained by finding the zeros of the inverse propagators given in Eqs. (14) and (15). In Fig. 2, we plot the effective masses of the nucleon, $\sigma$ and $\omega$ mesons as a function of the density. We observe that the reductions of $M^*/M$ and $m_\omega^*/m_\omega$ in the chiral sigma model are slower than those in the Walecka model, but the relationship between these two ratios is kept to be the same in the two models. It is because that the relationship derived in the Walecka model is also valid in the chiral sigma model, which does not depend on the special models and its parameters. In the chiral sigma model, the main feature is that there is not much arbitrariness in the renormalization procedure in order to respect the chiral symmetry. The counterterms

![Fig. 2. The effective masses of the nucleon, $\sigma$ and $\omega$ mesons as a function of the baryon density within the chiral sigma model.](image-url)
with the coefficients $\alpha_3$ and $\alpha_4$ in the chiral sigma model give larger nonlinear contributions, so that it leads to smaller mean field value of $\sigma$ meson, which is equivalent to larger effective nucleon mass. The effective mass of $\sigma$ meson shown in Fig. 2 increases at finite density in contrast to the Walecka model case. It is again due to the large nonlinear $\sigma$ meson interactions.

The existence of $\sigma$ meson has been a controversial subject for many years. Recently, there are a large number of evidences showing its existence [17,18]. However, the nature of $\sigma$ meson as a conventional $q\bar{q}$ state or as a $\pi\pi$ resonant state is still under debate. The density dependence of the $\sigma$ meson properties has been studied both experimentally [19,20] and theoretically [21,22,23,24]. The measurements of the in-medium $\pi\pi$ masses were obtained by the two-pion production experiments induced either by pions (CHAOS collaboration) [19] or by photons (TAPS collaboration) [20] on various nuclei. A significant nuclear-mass dependence of the $\pi\pi$ invariant mass distribution in the $I = J = 0$ channel was observed in the experiments, which could be interpreted as a signature for an in-medium modification of the $\pi\pi$ interaction in the $\sigma$ channel.

On the theoretical side, the density dependence of the $\sigma$ mass has been discussed in several models. Vacas et al. [21] calculated the $\pi\pi$ interaction in the $\sigma$ channel at finite densities in a chiral unitary approach, and found a dropping of the $\sigma$ mass as a function of the density. In Ref. [22], Hatsuda et al. studied the $\sigma$ propagator and found a decrease of the $\sigma$ mass caused by the partial restoration of chiral symmetry. However, the $\sigma$ mass in medium was found to be almost constant in Ref. [23] by using a hybrid model for nuclear matter, in which the nucleon, described as a quark-diquark state using the NJL model, could be moving in self-consistent scalar and vector fields. In the present work, we study the medium modification of hadron masses by using the Walecka model and the chiral sigma model with the vacuum polarization effects taken into account. These models could reproduce reasonably well the saturation properties of nuclear matter and the ground state properties of finite nuclei [4,9]. Therefore, it is very interesting to discuss the density dependence of the $\sigma$ mass in these models. The effective mass of $\sigma$ meson in the Walecka model decreases at lower densities, and then slightly increases at higher densities as shown in Fig. 1. However, we obtain the raise of the $\sigma$ mass in medium, as shown in Fig. 2, in the chiral sigma model using the standard method of the introduction of the counterterms [15]. It is the consequence of the negative energy states. This behavior should be dependent on the renormalization procedure, in particular, how to introduce the counterterms.

4 Summary

We have studied the variation of hadron masses in nuclear matter due to the vacuum polarization using the chiral sigma model, and compared the re-
results with those obtained in the Walecka model. Because of the constraints from chiral symmetry, there is less arbitrariness in the renormalization procedure for the chiral sigma model, as compared with the Walecka model. The renormalized chiral sigma model is able to provide proper binding energy and equilibrium density of nuclear matter, but the $\sigma$ meson mean field value comes out to be too small due to the large nonlinear $\sigma$ meson interactions, and it leads to a small reduction of nucleon mass in medium. The effective mass of $\omega$ meson decreases in nuclear matter. The reduction of $\omega$ meson mass in the chiral sigma model is slower than those in the Walecka model, while the relationship between $M^*/M$ and $m^*_\omega/m_\omega$ is kept to be the same in the two models. For $\sigma$ meson, its effective mass in the chiral sigma model increases at finite density, which is an opposite behavior to the results obtained in the Walecka model. It is very interesting to see that the vacuum polarization do play an important role in the modification of hadron masses in nuclear medium.

Acknowledgments

We acknowledge fruitful discussions with Y. Ogawa and A. Haga. This work was supported in part by the NSFC under contract No. 10135030, and by the EYTP of MOE of China.

References

[1] D. Adamova et al., CERES Collaboration, Phys. Rev. Lett. 91 (2003) 042301.
[2] G. E. Brown and M. Rho, Phys. Rev. Lett. 66 (1991) 2720.
[3] T. Hatsuda and Su H. Lee, Phys. Rev. C 46 (1992) R34.
[4] B. D. Serot and J. D. Walecka, Adv. Nucl. Phys. 16 (1986) 1.
[5] K. Saito and A. W. Thomas, Phys. Rev. C 51 (1995) 2757; K. Saito, K. Tsushima, and A. W. Thomas, Phys. Rev. C 55 (1997) 2637.
[6] M. Gell-Mann and M. Levy, Nuovo Cimento 16 (1960) 705.
[7] J. Boguta, Phys. Lett. B 120 (1983) 34.
[8] V. Fomenko, S. Marcos and L. N. Savushkin, J. Phys. G 19 (1993) 545.
[9] Y. Ogawa, H. Toki, S. Tamenaga, H. Shen, A. Hosaka, S. Sugimoto, and K. Ikeda, Prog. Theor. Phys. 111 (2004) 75.
[10] H. Kurasawa and T. Suzuki, Prog. Theor. Phys. 84 (1990) 1030.
[11] H. Shiomi and T. Hatsuda, Phys. Lett. B 334 (1994) 281.
[12] H.-C. Jean, J. Piekarewicz, and A. G. Williams, Phys. Rev. C 49 (1994) 1981.

[13] K. Saito, K. Tsushima, A. W. Thomas, and A. G. Williams, Phys. Lett. B 433 (1998) 243.

[14] S. A. Chin, Ann. Phys. 108 (1977) 301.

[15] T. Matsui and B. D. Serot, Ann. Phys. 144 (1982) 107.

[16] H. Kurasawa and T. Suzuki, Nucl. Phys. A 490 (1988) 571.

[17] Particle Data Group, S. Eidelman et al., Phys. Lett. B 592 (2004) 1.

[18] E791 collaboration, E. M. Aitala et al., Phys. Rev. Lett. 86 (2001) 770.

[19] P. Camerini et al., Nucl. Phys. A 735 (2004) 89; F. Bonniti et al., Nucl. Phys. A 677 (2000) 213.

[20] J. G. Messchendorp et al., Phys. Rev. Lett. 89 (2002) 222302.

[21] M. J. Vicente Vacas, E. Oset and L. Roca, Nucl. Phys. A 721 (2003) 301c.

[22] T. Hatsuda, T. Kunihiro and H. Shimizu, Phys. Rev. Lett. 82 (1999) 2840.

[23] W. Bentz and A. W. Thomas, Nucl. Phys. A 696 (2001) 138.

[24] G. Chanfray, Nucl. Phys. A 721 (2003) 76c.