

Momentum dependence of the ρ - ω mixing amplitude in a hadronic model

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We calculate the momentum dependence of the ρ - ω mixing amplitude in a purely hadronic model. The basic assumption of the model is that the mixing amplitude is generated by NN loops and thus driven entirely by the neutron-proton mass difference. The value of the amplitude at the ω -meson point is expressed in terms of only the $NN\omega$ and $NN\rho$ coupling constants. Using values for these couplings constrained by empirical two-nucleon data we obtain a value for the mixing amplitude in agreement with experiment. Extending these results to the spacelike region, we find a ρ - ω contribution to the NN interaction that is strongly suppressed and opposite in sign relative to the conventional contribution obtained from using the constant on-shell value for the mixing amplitude.

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The existence of a nonzero mixing matrix element between the isovector ρ meson and the isoscalar ω meson is by now firmly established [1-3]. In addition to the observed (small) branching ratio for the G -parity forbidden decay of the ω meson into two pions, the understanding of the pion form factor at the ω -meson point ($q^2 = m_\omega^2$) necessitates the coherent interplay of two distinct amplitudes; a dominant, G -parity allowed contribution ($\gamma \rightarrow \rho \rightarrow 2\pi$) interfering with a small ($\gamma \rightarrow \omega \rightarrow \rho \rightarrow 2\pi$) amplitude arising from ρ - ω mixing [1, 3].

It has long been recognized that a ρ - ω mixing amplitude would give rise to charge symmetry violation (CSV) in the nucleon-nucleon (NN) force [2, 4, 5]. The ρ - ω contribution to the NN interaction is constructed by employing NN -meson vertices constrained by empirical two-nucleon data. The ρ - ω mixing amplitude, on the other hand, is obtained from a measurement of the pion form factor at the on-shell ω -meson point [1, 3, 6].

In one-boson exchange (OBE) models of the NN force, nucleons interact via the exchange of several mesons possessing different Lorentz and isospin transformation properties [7, 8]. In this paper we are concerned with the mixing between the isoscalar ω meson and the neutral member of the isovector ρ meson. The neutral ω meson couples in a minimal fashion to the conserved baryon current

$$\mathcal{L}_{NN\omega} = g_\omega \bar{\psi} \gamma^\mu \psi \omega_\mu. \quad (1)$$

The isovector ρ meson, on the other hand, has a vector as well as a tensor coupling to the nucleon

$$\mathcal{L}_{NN\rho} = g_\rho \bar{\psi} \gamma^\mu \boldsymbol{\tau} \cdot \boldsymbol{\psi} \boldsymbol{\rho}_\mu + f_\rho \bar{\psi} \sigma^{\mu\nu} \boldsymbol{\tau} \cdot \boldsymbol{\psi} \frac{\partial_\mu}{2M} \boldsymbol{\rho}_\nu. \quad (2)$$

Notice that the above definitions are standard except for possible factors of 2 (some definitions include isospin ver-

tices as $1/2$ or $\tau/2$). In addition, some models include a $NN\omega$ tensor coupling. Since in this paper we choose typical coupling constants determined by the Bonn group, we use their conventions throughout this work [7, 8].

Having constructed an interaction Lagrangian one can then proceed to calculate the contribution from ρ - ω mixing to the NN potential; one obtains (with $\Gamma^\mu \equiv i\sigma^{\mu\nu} q_\nu / 2M$ and $C_\rho \equiv f_\rho / g_\rho$) [2, 4]

$$\hat{V}_{\text{III}}^{\rho\omega}(q) = V_{NN}^{\rho\omega}(q) \gamma^\mu(1) \gamma_\mu(2) [\tau_z(1) + \tau_z(2)], \quad (3a)$$

$$\hat{V}_{\text{IV}}^{\rho\omega}(q) = V_{NN}^{\rho\omega}(q) C_\rho \left[\Gamma^\mu(1) \gamma_\mu(2) \tau_z(1) - \gamma^\mu(1) \Gamma_\mu(2) \tau_z(2) \right], \quad (3b)$$

where

$$V_{NN}^{\rho\omega}(q) = -\frac{g_\rho g_\omega \langle \rho | H | \omega \rangle}{(q^2 - m_\rho^2)(q^2 - m_\omega^2)}. \quad (4)$$

This parameter-free construction of the potential is quite satisfactory since it does not introduce additional parameters beyond those already constrained by charge-symmetry-conserving (CSC) two-nucleon data. More importantly perhaps, most of the differences observed experimentally in the binding energy of mirror nuclei, the Nolen-Schiffer anomaly [9], have been attributed to ρ - ω mixing and explained using the above potential [10]. In addition, ρ - ω mixing plays an important role in explaining the difference between the neutron and proton analyzing power (ΔA) in elastic neutron-proton scattering [11-14]. Indeed, ρ - ω mixing seems to account for half the size of the effect in the Indiana University Cyclotron Facility experiment [3, 12].

One important issue that has been overlooked until very recently, however, is the momentum dependence of the ρ - ω mixing amplitude. To date, most of the theoret-

ical efforts devoted to the understanding of ρ - ω mixing in CSV have assumed the constant on-shell value for the mixing amplitude. Since the relevant momentum transfer carried by the meson exchange between two nucleons is spacelike ($q_\mu^2 < 0$), the use of a mixing amplitude determined at a timelike point is clearly suspect. Recently, Goldman, Henderson, and Thomas have tested the on-shell assumption by constructing a model in which the mixing is driven by the $u - d$ quark mass difference [15]. Since in their model the mixing is generated by $q\bar{q}$ loops, they could examine the momentum dependence of the mixing amplitude and, thus, confront the on-shell assumption. They have concluded that the on-shell assumption may be surprisingly poor.

In this paper we study the momentum dependence of the ρ - ω mixing amplitude in a purely hadronic model. We consider the mixing amplitude to also be generated by fermion loops. The basic assumption of our model, however, is that the mixing amplitude is generated by $N\bar{N}$ loops and thus driven entirely by the small neutron-proton mass difference.

There are several advantages in calculating the ρ - ω mixing amplitude using a purely hadronic description. First, all parameters used in the calculation are known to great accuracy (e.g., nucleon and meson masses) or are constrained by empirical data (e.g., coupling constants) [7, 8, 16, 17]. Furthermore, since field-theoretical models naturally include vacuum corrections, the coupling between nucleons and antinucleons is determined by the underlying theory and, thus, ultimately constrained by the empirical data. These vacuum contributions are an essential part of the relativistic description of a nuclear target. In fact, vacuum corrections are known to be crucial for avoiding the appearance of spurious (center-of-mass) excitations and for maintaining the conservation of the electromagnetic current [18, 19]. Furthermore, hadronic models have been relatively successful in calculating the nuclear response to electromagnetic probes by including, in addition to the traditional particle-hole excitations, vacuum polarization [20, 21]. The self-consistency of these calculations demanded that the same interaction used in the calculation of ground state properties be used for the residual particle-hole and $N\bar{N}$ interaction. Hence, in the calculation of the $N\bar{N}$ loops driving the ρ - ω mixing amplitude [see Eq. (5) below], we will use the same dynamical input as in the calculation of the ρ - ω contribution to the NN potential [Eq. (4)]. Therefore, the ρ - ω mixing amplitude constitutes a parameter-free prediction. While one might question the reliability of a purely hadronic description of the mixing amplitude, we feel that this is balanced by the simplicity and parameter-free nature of the calculation. Second, in the calculation of fermion loops there are no unphysical (e.g., $q\bar{q}$) thresholds appearing near the region of interest [15]. The only thresholds that appear are related to the production of physical ($N\bar{N}$) states and occur far away from the region of interest. Finally, there is no need to introduce *ad hoc* form factors to regularize (infinite) loop integrals (we discuss later the sensitivity of our results to the introduction of form factors). Since the mixing amplitude is sensitive only to the difference between proton and neutron loops,

the difference is finite even though the individual pieces are not (note that the difference between u -quark and d -quark loops should also be finite).

Using standard Feynman rules, the ρ - ω mixing amplitude in the hadronic model described by Eqs. (1) and (2) can be written as

$$\langle \rho | H | \omega \rangle = g_\rho g_\omega q^2 \Pi(q^2), \quad (5)$$

where $\Pi(q^2)$ is the transverse component of the full polarization tensor [see Eqs. (11a) and (11b) below]. Because the ρ meson has, both, a vector and a tensor coupling to the nucleon, one needs to evaluate polarization tensors having Lorentz vector and tensor vertices, i.e.,

$$\Pi^{\mu\nu}(q) = \Pi_{\nu\nu}^{\mu\nu}(q) + C_\rho \Pi_{\nu t}^{\mu\nu}(q), \quad (6)$$

where

$$i\Pi_{\nu\nu}^{\mu\nu}(q) = \int \frac{d^4k}{(2\pi)^4} \text{Tr} [\gamma^\mu G(k+q) \gamma^\nu \tau_z G(k)], \quad (7a)$$

$$i\Pi_{\nu t}^{\mu\nu}(q) = \int \frac{d^4k}{(2\pi)^4} \text{Tr} \left[\gamma^\mu G(k+q) \frac{i\sigma^{\nu\lambda} q_\lambda}{2M} \tau_z G(k) \right]. \quad (7b)$$

The isospin trace can be evaluated by writing isoscalar and isovector components of the nucleon propagator

$$\begin{aligned} G(k) &= \frac{1}{2} G_p(k)(1 + \tau_z) + \frac{1}{2} G_n(k)(1 - \tau_z) \\ &\equiv G_0(k) + G_1(k)\tau_z \end{aligned} \quad (8)$$

in terms of individual proton and neutron contributions

$$G_p(k) = \frac{\not{k} + M_p}{k^2 - M_p^2 + i\epsilon}, \quad G_n(k) = \frac{\not{k} + M_n}{k^2 - M_n^2 + i\epsilon}. \quad (9)$$

After performing the isospin trace one discovers, as expected, that the ρ - ω mixing amplitude is driven by the difference between proton and neutron loops (see Fig. 1)

$$\Pi_{\mu\nu}(q) = \Pi_{\mu\nu}^{(p)}(q) - \Pi_{\mu\nu}^{(n)}(q). \quad (10)$$

The calculation of the vacuum loops is now completely

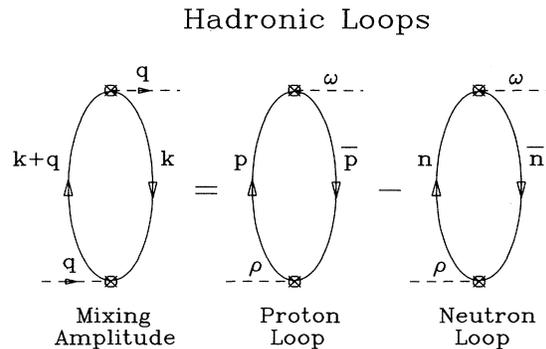


FIG. 1. The ρ - ω mixing amplitude given as the difference between proton and neutron loops. The Feynman diagram on the left labels the external meson momentum q and the loop momentum k .

standard [22]. Since the individual proton and neutron contribution diverge (but not their difference), we isolate the singularities by using dimensional regularization. The Lorentz tensor structure of the polarization can be obtained after performing the appropriate traces; we obtain

$$\Pi_{\nu\nu}^{\mu\nu}(q) = \left(-g^{\mu\nu}q^2 + q^\mu q^\nu \right) \Pi_{\nu\nu}(q^2), \quad (11a)$$

$$\Pi_{\nu t}^{\mu\nu}(q) = \left(-g^{\mu\nu}q^2 + q^\mu q^\nu \right) \Pi_{\nu t}(q^2), \quad (11b)$$

where the unrenormalized polarizations are given (for proton loops) by

$$\Pi_{\nu\nu}^{(p)}(q^2) = -\frac{1}{2\pi^2} \left[\frac{1}{6\epsilon} - \frac{\gamma}{6} - \int_0^1 dx x(1-x) \ln \left(\frac{M_p^2 - x(1-x)q^2}{\Lambda^2} \right) \right], \quad (12a)$$

$$\Pi_{\nu t}^{(p)}(q^2) = -\frac{1}{8\pi^2} \left[\frac{1}{\epsilon} - \gamma - \int_0^1 dx \ln \left(\frac{M_p^2 - x(1-x)q^2}{\Lambda^2} \right) \right]. \quad (12b)$$

In the above equations Λ is an arbitrary renormalization constant, γ is the Euler-Mascheroni constant, and $\epsilon \rightarrow 0$. A finite mixing amplitude is now obtained by taking the difference between proton and neutron contributions

$$\Pi_{\nu\nu}(q^2) = \Pi_{\nu\nu}^{(p)}(q^2) - \Pi_{\nu\nu}^{(n)}(q^2) = \frac{1}{2\pi^2} \int_0^1 dx x(1-x) \ln \left[\frac{M_p^2 - x(1-x)q^2}{M_n^2 - x(1-x)q^2} \right], \quad (13a)$$

$$\Pi_{\nu t}(q^2) = \Pi_{\nu t}^{(p)}(q^2) - \Pi_{\nu t}^{(n)}(q^2) = \frac{1}{8\pi^2} \int_0^1 dx \ln \left[\frac{M_p^2 - x(1-x)q^2}{M_n^2 - x(1-x)q^2} \right]. \quad (13b)$$

Expanding the integrand to first order in the neutron-proton mass difference we obtain the following closed-form expression for the ρ - ω mixing amplitude:

$$\frac{\langle \rho | H | \omega \rangle}{M^2} = \frac{g_\rho g_\omega}{\pi^2} \frac{\Delta M}{M} \times \begin{cases} 1 - \frac{1}{\xi} (1 + \xi^2 + C_\rho) \tan^{-1} \left(\frac{1}{\xi} \right), & \text{for } 0 < q^2 < 4M^2; \\ 1 - \frac{1}{\xi} (1 - \xi^2 + C_\rho) \frac{1}{2} \ln \left(\left| \frac{\xi - 1}{\xi + 1} \right| \right), & \text{otherwise,} \end{cases} \quad (14)$$

where

$$\begin{aligned} M &= \frac{1}{2}(M_n + M_p), \\ \Delta M &= (M_n - M_p), \\ \xi &= \left| 1 - \frac{4M^2}{q^2} \right|^{1/2}. \end{aligned} \quad (15)$$

The above analytic expression is accurate everywhere except in the vicinity of the thresholds ($q^2 = 4M^2$). This equation embodies the central result of the present work. It displays the momentum dependence of the ρ - ω mixing amplitude in terms of three parameters (g_ω , g_ρ , and C_ρ). Having previously constrained these parameters from CSC two-nucleon data, this result provides a parameter-free prediction of the model. In particular, using the above expression at the on-shell ω -meson point, together with the parameters of Table I, we obtain

TABLE I. Meson masses, coupling constants, tensor-to-vector ratio, and cutoff parameters in the Bonn one-boson exchange model (see Table 4.2 of Ref. [7] and Table 4 of Ref. [8]).

Meson	Mass (MeV)	$g^2/4\pi$	$C = f/g$	Λ (MeV)
ρ	770	0.41	6.1	1400
ω	783	10.6	0.0	1500

$$\langle \rho | H | \omega \rangle \Big|_{q^2=m_\omega^2} = -4314 \text{ MeV}^2,$$

which compares well to the experimental value of $(-4520 \pm 600) \text{ MeV}^2$.

The momentum dependence of the ρ - ω mixing amplitude is shown in Fig. 2. Two calculations are displayed. The solid line shows results for the mixing amplitude using Eq. (14) at all values of q^2 . In contrast, the dashed line shows results modified by the introduction of form factors in the spacelike region. These form factors are introduced by modifying the point coupling in the following way:

$$\begin{aligned} g_\rho &\rightarrow g_\rho(q^2) \equiv g_\rho \left(1 - q^2/\Lambda_\rho^2 \right)^{-1}, \\ g_\omega &\rightarrow g_\omega(q^2) \equiv g_\omega \left(1 - q^2/\Lambda_\omega^2 \right)^{-1}. \end{aligned} \quad (16)$$

As before, the numerical values for the cutoffs (Λ) are constrained by empirical two-nucleon data (see Table I).

The topic of form factors, or vertex corrections, now needs some comment. It is clear that the finite size of both nucleons and mesons should modify the naive point coupling at finite q^2 . These vertex corrections, however, need not be included in an *ad hoc* fashion. In fact, vertex corrections can, in principle, be calculated using renormalizable models based on hadronic degrees of freedom. The basic dynamical assumption of these models is that

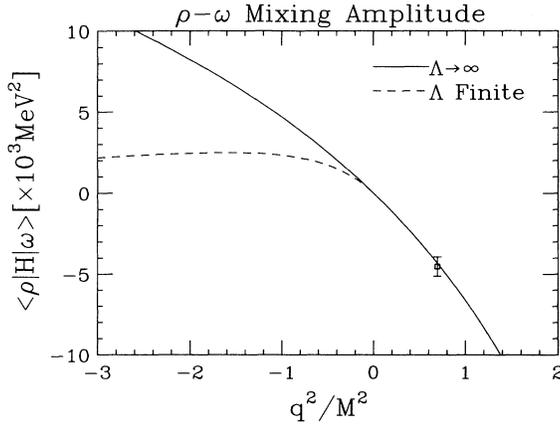


FIG. 2. The ρ - ω mixing amplitude as a function of q^2 with (dashed line) and without (solid line) the inclusion of form factors in the spacelike region. The experimental point was extracted from a measurement of the electromagnetic pion form factor at the ω -meson point [1].

the internal structure of the hadrons can be described in terms of hadronic degrees of freedom alone [17]. Although some progress has recently been made in calculating vertex corrections in hadronic theories [23], much work remains to be done. In particular, very little has been said about vertex corrections in the timelike region. In addition, OBE models of the NN interaction can only constrain the form factors on-shell. Hence, in order to minimize the model assumptions introduced in our calculation, we modify the naive point coupling only in the spacelike region by the introduction of on-shell form factors as prescribed by Eq. (16). The issue of off-shell form factors is clearly an important and open problem.

Because of the smooth behavior of the (dimensionless) transverse polarization over the sampled q^2 region (not

shown), the most important momentum behavior of the mixing amplitude displayed in Fig. 2 is determined by the q^2 factor in Eq. (5). In particular, it reveals that the mixing amplitude has the opposite sign relative to its value at the on-shell point over the entire spacelike region sampled in NN scattering. This behavior is clearly seen in Fig. 3 which shows the ρ - ω contribution to the NN potential. Three calculations are displayed. The solid line shows the NN potential [Eq. (4)] obtained from using the on-shell value for the ρ - ω mixing amplitude. The dashed line uses the same on-shell amplitude but modifies the naive coupling by the inclusion of on-shell form factors at the external nucleon legs. Finally, the dash-dotted line uses the off-shell ρ - ω mixing amplitude and has, both, the off-shell amplitude and the external nucleon legs modified by form factors. In addition to the sign difference displayed over the entire spacelike region, a much suppressed contribution to the NN potential is observed whenever the off-shell mixing amplitude is used.

Finally, in Fig. 4 we show the static contribution to the NN potential in configuration space. These results are obtained from taking the Fourier transform of the three momentum-space potentials displayed in Fig. 3. In particular, with the on-shell mixing amplitude and with no on-shell form factors this r -space potential takes the following form [2, 4]:

$$V_{NN}^{\rho\omega}(r) = -\frac{g_\rho g_\omega}{4\pi} \frac{\langle \rho | H | \omega \rangle}{m_\omega^2 - m_\rho^2} \left(\frac{e^{-m_\rho r}}{r} - \frac{e^{-m_\omega r}}{r} \right). \quad (17)$$

Large differences are clearly seen in the interior for all three potentials. However, due to the strong repulsive nature of the charge-symmetry-conserving NN potential, the two-nucleon wave function will be insensitive to the details of the distorting CSV potential in the interior. More importantly, perhaps, is the occurrence of a node in the potential around $r \sim 0.9$ fm. This is the

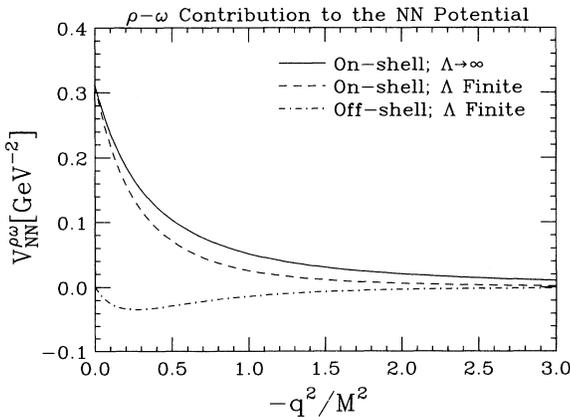


FIG. 3. The contribution from ρ - ω mixing to the NN potential as a function of q^2 using the off-shell value for the mixing amplitude (dash-dotted line), and the on-shell value with (dashed line) and without (solid line) the inclusion of on-shell form factors at the external nucleon legs.

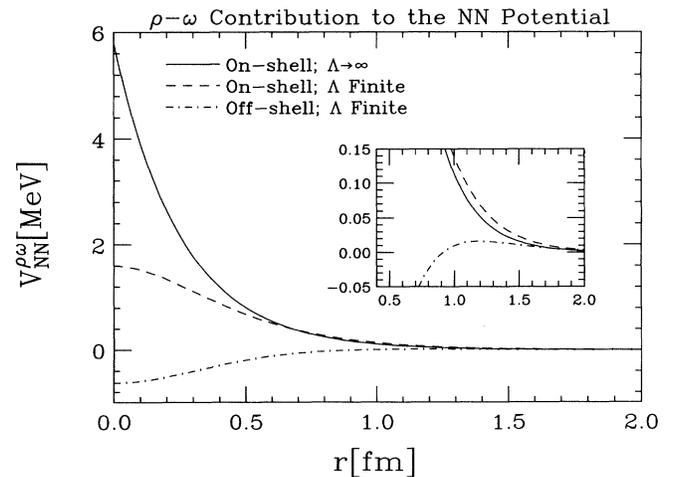


FIG. 4. The contribution from ρ - ω mixing to the NN potential as a function of the NN separation using the off-shell value for the mixing amplitude (dash-dotted line), and the on-shell value with (dashed line) and without (solid line) the inclusion of on-shell form factors at the external nucleon legs.

region where conventional estimates suggest that the ρ - ω contribution, obtained from a competition between the fast falloff of the CSV potential and the suppression of the two-nucleon wave function in the interior, should be larger. Consequently, our findings are in basic agreement with the results obtained in Ref. [15] and are consistent with the view that the ρ - ω contribution to the CSV nucleon-nucleon potential is, effectively, nonexistent.

In summary, we have calculated the momentum dependence of the ρ - ω mixing amplitude in a simple hadronic model. The mixing was assumed to be generated solely by $N\bar{N}$ loops and thus driven by the neutron-proton mass difference. Since the mixing is sensitive only to the difference between proton and neutron loops, the amplitude was found to be finite even though the individual pieces were not. We have presented closed-form analytic expressions for the mixing amplitude in terms of very few parameters. Furthermore, these parameters were obtained from previous fits to two-nucleon data. Hence, our results can be regarded as parameter-free predictions of one-boson exchange models. Using standard values for these parameters we obtained a value for the ρ - ω mixing amplitude at the on-shell ω -meson point in good agreement with experiment. We have extended our results to the spacelike region and have computed the contribution from the off-shell ρ - ω mixing amplitude to the NN potential. These results were compared to a recent calculation of the mixing amplitude in terms of $q\bar{q}$ loops [15]. In spite of the obvious differences between the two models, our findings agree with the main conclusions drawn from that work, namely, that the momentum dependence of the ρ - ω

mixing amplitude is significant and that the occurrence of a node in the NN potential around $r \sim 0.9$ fm severely suppresses the ρ - ω contribution to the CSV potential.

What the impact of these results will be on CSV observables is, at this juncture, hard to predict. For some of them like the Nolen-Schiffer anomaly or the differences in NN scattering lengths one must, at the very least, study the momentum dependence of the π - η mixing amplitude [24, 25]. In addition, one must study medium modifications to the vector-meson propagators. This effect might be important in explaining the Nolen-Schiffer anomaly for medium to heavy nuclei. These issues are currently under investigation.

The neutron-proton analyzing power difference, ΔA , might, however, pose a serious challenge. There, the π - η amplitude does not contribute. At $T_{\text{lab}} = 477$ MeV, ΔA is dominated by the one-pion exchange potential and an absent ρ - ω contribution might not spoil the agreement with experiment. At $T_{\text{lab}} = 183$ MeV, on the other hand, the one-pion exchange contribution is small and ρ - ω mixing dominates [3, 13, 14]. Hence, if as suggested by our findings, and by those of Ref. [15], ρ - ω mixing is indeed severely suppressed, additional CSV mechanisms will have to be found in order to account for the unexplained ΔA at $T_{\text{lab}} = 183$ MeV.

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- [1] L.M. Barkov *et al.*, Nucl. Phys. **B256**, 365 (1985).
 - [2] E.M. Henley and G.A. Miller, in *Mesons in Nuclei*, edited by M. Rho and D.H. Wilkinson (North-Holland, Amsterdam, 1979).
 - [3] G.A. Miller, B.M.K. Nefkens, and I. Šlaus, Phys. Rep. **194**, 1 (1990).
 - [4] Peter C. McNamee, Michael D. Scadron, and Sidney A. Coon, Nucl. Phys. **A249**, 483 (1975).
 - [5] Sidney A. Coon and Michael D. Scadron, Nucl. Phys. **A287**, 381 (1977).
 - [6] Sidney A. Coon and Roger C. Barret, Phys. Rev. C **36**, 2189 (1987).
 - [7] R. Machleidt, in *Proceedings of the Relativistic Dynamics and Quark-Nuclear Physics*, edited by M.B. Johnson and A. Picklesimer (Wiley, New York, 1986).
 - [8] R. Machleidt, K. Holinde, and Ch. Elster, Phys. Rep. **149**, 1 (1987).
 - [9] J.A. Nolen and J.P. Schiffer, Annu. Rev. Nucl. Sci. **19**, 471 (1969).
 - [10] P. Blunden and M.J. Iqbal, Phys. Lett. B **198**, 14 (1987).
 - [11] R. Abegg *et al.*, Phys. Rev. Lett. **56**, 2571 (1986); Phys. Rev. D **39**, 2464 (1989).
 - [12] L.D. Knutson *et al.*, Nucl. Phys. **A508**, 185c (1990); Phys. Rev. Lett. **66**, 1410 (1991); S.E. Vigdor *et al.*, Phys. Rev. C **46**, 410 (1992).
 - [13] G.A. Miller, A.W. Thomas, and A.G. Williams, Phys. Rev. Lett. **56**, 2567 (1986).
 - [14] A.G. Williams, A.W. Thomas, and G.A. Miller, Phys. Rev. C **36**, 1956 (1987).
 - [15] T. Goldman, J.A. Henderson, and A.W. Thomas, Few-Body Syst. **12**, 123 (1992); Mod. Phys. Lett. A (to be published).
 - [16] J.D. Walecka, Ann. Phys. **83**, 491 (1974).
 - [17] B.D. Serot and J.D. Walecka, in *Advances in Nuclear Physics*, edited by J.W. Negele and E. Vogt (Plenum, New York, 1986), Vol. 16.
 - [18] J. Piekarewicz, Nucl. Phys. **A511**, 487 (1990).
 - [19] J.F. Dawson and R.J. Furnstahl, Phys. Rev. C **42**, 2009 (1990).
 - [20] C.J. Horowitz and J. Piekarewicz, Phys. Rev. Lett. **62**, 391 (1989).
 - [21] C.J. Horowitz and J. Piekarewicz, Nucl. Phys. **A511**, 461 (1990).
 - [22] Pierre Ramond, *Field Theory: A Modern Primer* (Benjamin/Cummings, London, 1981).
 - [23] Myriam P. Allendes and Brian D. Serot, Phys. Rev. C **45**, 2975 (1992).
 - [24] Kim Maltman, Los Alamos Report LA-UR-92-2549.
 - [25] J. Piekarewicz, SCRI Report FSU-SCRI-93-38.