M. Poincaré visits Jefferson Lab: Relativistic Models of Few-Nucleon Systems

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I discuss relativistic models of few-nucleon systems, with particular emphasis on calculations of electron-deuteron scattering and the comparison of these calculations with recent data from Jefferson Lab.

1. Introduction

In the “standard model of nuclear physics” nuclear properties are calculated using a Schrödinger equation in which the degrees of freedom are nucleons interacting via energy-independent potentials. This picture can now be solved essentially exactly for light nuclei \([1,2]\). However, it appears it only describes non-relativistic systems, since the Poincaré algebra is satisfied approximately: order-by-order in an expansion in momenta over the nucleon mass \([3,4]\).

Here I review attempts to build a phenomenology of light nuclei which satisfies the Poincaré algebra exactly, and so is manifestly applicable when few-nucleon systems are probed at GeV-scale momentum transfers—the kinematic domain accessed at Jefferson Lab. I will focus on elastic electron-deuteron scattering, since it involves the simplest non-trivial nucleus, and yet also requires the calculation of \(NN\)-system wave functions away from the two-body centre-of-mass frame.

In Section 2 I write down the Poincaré algebra, and sketch some of the different means by which model-builders have attempted to satisfy it. Then in Section 3 I describe the calculation of matrix elements for the interaction of the probe (electron) with a relativistic bound state (the deuteron). I enumerate the places in which “relativistic effects” can occur in the calculation, and show one explicit example of such an effect. I then compare a number of relativistic models of \(ed\) scattering, and address the issue of how we should interpret their widely-varying predictions. I conclude in Section 4.

2. What is a relativistic model?

The laws of quantum mechanics are the same in all inertial frames of reference. In non-relativistic quantum mechanics the generators of the transformations relating one frame of reference to another are the total momentum \(P\) (translations), the total angular momentum \(J\) (rotations), and the boost operator \(K_G\). This last operator can be defined by its action on state vectors: suppose Abigail and Beatrice have their coordinate systems
aligned at $t = 0$, but Abigail moves with velocity $v$ relative to Beatrice:

$$\langle x | \psi(t) \rangle_A = \langle x + vt | \psi(t) \rangle_B \equiv \langle x | e^{iK_G \cdot v} | \psi(t) \rangle_B.$$ (1)

Considering the combined effect on a state vector of all possible pairs of translations, rotations, and boosts then leads to the following commutation relations:

$$[P^i, P^j] = 0; \quad [J^i, J^j] = i \epsilon^{ijk} J^k; \quad [J^i, K^j_G] = i \epsilon^{ijk} K^k_G;$$ (2)

$$[J^i, P^j] = i \epsilon^{ijk} P^k; \quad [K^i_G, K^j_G] = 0; \quad [K^i_G, P^j] = -i \delta^{ij} M.$$ (3)

Meanwhile, the invariance of Schrödinger equation dynamics means that the Hamiltonian, $H$, commutes with $J$ and $P$. $H$ does not commute with $K_G$:

$$[H, P^i] = 0; \quad [H, J^i] = 0; \quad [H, K^i_G] = i P^i.$$ (4)

The relations (2)–(4) define the Galilei algebra, which relates observations in different frames of reference, and must be obeyed in non-relativistic quantum mechanics.

The Poincaré algebra also relates observations in different frames, but now it implements Einsteinian, rather than Galilean, relativity. The changes in the algebra seem simple enough: the only commutators which change are the last two in (3). They become:

$$[K^i_G, K^j_G] = -i \epsilon^{ijk} J^k; \quad [K^i_G, P^j] = -i \delta^{ij} H.$$ (5)

Any “relativistic” quantum theory should obey the Poincaré algebra. The first of the two relations (5) states that relativistic boosts $K$, in contrast to their Galilean cousins, are not commutative. The second is the key commutator which makes relativistic quantum theories difficult to construct: it encodes the absence of absolute time in relativistic theories, since it connects the generator of time translations, $H$, to the generators of spatial boosts and translations. In any interacting theory $H$ is the sum of a free part $H^0$ and a potential energy $V$, so Eq. (5)(b) means that either $K$ or $P$ must depend on $V$.

This is one way to understand Dirac’s famous “three forms” of relativistic quantum mechanics [5]. In instant form (which will be the form I mainly discuss here) $J$ and $P$ depend on $H^0$ alone, and so are the same in the interacting theory as they are for free particles. In contrast, $K$ depends on $V$, and so will be different for systems with different dynamics. We say that the angular and linear momentum are kinematical generators, while boosts are “dynamical”. In point form $P$ is dynamical, and $K$ is kinematic. Finally, in front form linear combinations of the 9 Poincaré algebra commutators are taken, such that only two, rather than three, generators need involve $V$. In front form rotations about the $\hat{x}$ and $\hat{y}$ axes are dynamical. It has, though, the great advantage that boosts in the $\hat{z}$-direction are kinematic.

One example of theories that manifestly obey the Poincaré algebra are relativistic quantum field theories (QFT), such as Quantum Electrodynamics. Field theories can be quantized so that they use instant-form, front-form, or point-form dynamics \(1\). If instant form is adopted then $J$ and $P$ can be written down directly from the stress-energy-momentum tensor of the free theory, and do not connect different Fock-space sectors of the theory \(2\).

\(1\)See Ref. [6] for calculations of electron-deuteron scattering in a front-form QFT-based approach.

\(2\)I do not consider problems of gauge invariance in such definitions here.
In contrast, $K$ involves the particle-number-changing interaction piece of the QFT Hamiltonian, and so connects different sectors of the Fock space. In a weak-coupling theory perturbative forms of the boost are useful, but in nuclear physics, we deal with coupling constants of order 1, and some non-perturbative approximation to the boost, which manages to preserve—at least approximately—the Poincaré algebra, is needed. This means that the application of QFT to relativistic nuclear physics is not straightforward.

Manifest covariance is a sufficient, but not necessary, condition for a theory to be relativistic. The Bakamjian-Thomas (BT) construction is a technique by which a $P$, $J$, and $K$ that obey the Poincaré algebra may be constructed, once a potential energy $V$ is given. Crucially, particle number is conserved in the BT construction—as long as $V$ does not change it. In instant form, we first introduce the intrinsic spin $j$, the operator $X = i \nabla P$, and the invariant mass of the system $M$. Some algebra then shows that defining gives a set of operators that obeys the Poincaré algebra. Note that $K$ depends on $V$, but the BT $P$ and $J$ are the same as in the free theory. BT constructions for front form and point form are also possible. Some results for $e^\pm d$ scattering employing them will be displayed below.

Suppose then that we have a relativistic quantum theory, with dynamics—an $H_0$ and a $V$—which is God-given in the rest frame of the many-body quantum system. Considering, for the time being, only one state in the spectrum, we write:

$$H_0|\psi\rangle_0 = E_0|\psi\rangle_0.$$  

Relativistic covariance imposes constraints on matrix elements in this theory. For instance, if $J$ is an operator corresponding to an observable which is a Lorentz scalar, then we should have:

$$\langle\psi|J_0|\psi\rangle_0 = Q\langle\psi|J_Q|\psi\rangle_Q,$$  

where the subscripts indicate operators and states constructed by observers in frames in which the system is, respectively, at rest, and moving with total momentum $Q$. Quantum-mechanical equivalences of the type imply the presence of unitary transformations relating wave functions and operators:

$$|\psi\rangle_Q = U_Q|\psi\rangle_0; \quad J_Q = U_QJ_0U_Q^\dagger.$$  

$U_Q$ is a unitary representation of the Poincaré group. Of course, in a manifestly covariant theory is satisfied automatically, as long as $|\psi\rangle_Q$ and $J_Q$ are constructed consistently.

3. Electron-deuteron as an example

3.1. Observables

The $e^\pm d$ differential cross section, can (up to corrections of $O(\alpha^2)$) be written in terms of two structure functions, $A$ and $B$:

$$\frac{d\sigma}{d\Omega} = \frac{d\sigma}{d\Omega_{\text{Mott}}} \left[ A(Q^2) + B(Q^2) \tan^2 \left( \frac{\theta_e}{2} \right) \right],$$  

$$H = \sqrt{M^2 + P^2}; \quad J = j + X \times P; \quad K = -\frac{1}{2}\{H, X\} - \frac{P \times J}{H + M};$$  

where $\theta_e$ is the c.m.-frame electron scattering angle, $q^2 = (p'_e - p_e)^2 \equiv -Q^2$ is the virtuality of the photon exchanged between the electron and the nucleus, and $\frac{d\sigma}{d\Omega}_\text{Mott}$ is the Mott ed cross-section. $A$ and $B$ are related to the deuteron form factors $G_C$, $G_Q$, and $G_M$ by:

$$A = G_C^2 + \frac{2}{3} \eta G_M^2 + \frac{8}{9} \eta^2 M_4^2 G_Q^2; \quad B = \frac{4}{3} \eta (1 + \eta) G_M^2; \quad \eta \equiv \sqrt{1 + \frac{Q^2}{4M^2_d}}. \quad (11)$$

We also consider the tensor-polarization observable, $T_{20}$, which is defined from the ratios:

$$x = \frac{2}{3} \eta \frac{G_Q}{G_C}; \quad y = \frac{2}{3} \eta \left( \frac{G_M}{G_C} \right)^2 \left[ \frac{1}{2} + (1 + \eta) \tan^2 \left( \frac{\theta_e}{2} \right) \right]; \quad (12)$$

$$T_{20} = \sqrt{2} \frac{x(x + 2) + y/2}{1 + 2(x^2 + y)}. \quad (13)$$

Measurements of $T_{20}$, $A$, and $B$ at a fixed $Q^2$ enable the extraction of $G_C$, $G_Q$, and $G_M$, as is done, for instance in Ref. [9] out to $Q^2 \sim 2$ GeV$^2$. In the Breit frame (see Fig. 1) the deuteron charge, quadrupole, and magnetic form factors are linear combinations of matrix elements of the $NN$ four-current $J^\mu$ between deuteron magnetic substates, for example:

$$G_C = \frac{1}{3|e|} \left( \langle 1 | J^0 | 1 \rangle + \langle 0 | J^0 | 0 \rangle + \langle -1 | J^0 | -1 \rangle \right). \quad (14)$$

![Figure 1. Three momenta of the deuteron and virtual photon in the Breit frame.](image)

### 3.2. Calculations

Thus to calculate $A$, $B$, and $T_{20}$ we either need an explicit calculation of the wave functions for deuteron magnetic substates $|m_z\rangle_{\pm q/2}$, or we must employ a unitary representation of the Poincaré group to write:

$$q/2 \langle \psi | J_{\text{Breit}}^\mu | \psi \rangle - q/2 = 0 \langle \psi | U_{q/2} \tilde{J}_{\text{Breit}}^\mu U_{-q/2} | \psi \rangle 0. \quad (15)$$

To simplify matters, we now expand $J_{\text{Breit}}^\mu = \tilde{j}_{\text{Breit}}^\mu + J_{2B,\text{Breit}}^\mu$, where $j^\mu$ is the single-nucleon current, and drop $J_{2B}^\mu$. This leads to a “Relativistic Impulse Approximation” (RIA).

If, as we are doing here, we take the view that the Poincaré group relates processes in different frames, and remain agnostic about relativity’s impact on $|\psi\rangle_0$, there are two potential sources of relativistic effects in the RIA. First, the one-body current:

$$j_{\text{Breit}}^\mu = \bar{u}(k'_1) \left[ F_1(Q^2) \gamma^\mu + F_2(Q^2) i \frac{\sigma^\mu \epsilon}{2M} \right] u(k_1) \quad (16)$$

contains effects traditionally called “relativistic”, since they arise in a Dirac equation treatment of the nucleon, e.g. Thomas precession, the Foldy contribution to $\langle r^2 \rangle$. 
Second, there are boost effects arising from the $U$’s. I will mention one simple example of these, that the boost incorporates length contraction. In instant form this has the consequence that the deuteron wave function, computed in the $NN$ rest frame, is evaluated at a value of $p$ which is length-contracted with respect to the value in the Breit frame [10,11,12]:

$$\langle p|U^{-q/2}|\psi\rangle_0 \approx \frac{1}{\sqrt{\eta}}|p_c|\psi_0 \equiv \frac{1}{\sqrt{\eta}}\psi(p_c),$$  \hspace{1cm} (17)

where $p_c$’s components perpendicular to $q$ are the same as $p$’s, but the component parallel to $q$ is $p\cdot\hat{q}/\eta$. Keeping only this length-contraction effect, at present, and approximating $\bar{u}\gamma^0u$ by the first term in its $p/M$ expansion—1—it is easy to show that:

$$q/2\langle m_z|J^0|m_z\rangle_{-q/2} = F_1(Q^2) \int \frac{d^3p_c}{(2\pi)^3} \psi^*_M \left(p_c + \frac{q}{2\eta}\right) \psi_M(p_c).$$  \hspace{1cm} (18)

This is exactly the result for this matrix element in the non-relativistic impulse approximation (NRIA), but with $q$ replaced by $q/\eta$. Expanding $\eta$ in powers of $Q/M$ one might think that the difference between the NRIA and the instant-form RIA will be very small unless $Q^2$ is large. In fact though, this effect is non-negligible, especially near the minimum of $G_C$ at $Q \sim 800$ MeV, since the charge form factor varies rapidly there.

3.3. Survey of relativistic approaches

This then is one relativistic effect, in one approach to relativistic dynamics. What of other approaches? A detailed description of various models which attempt to deal with this regime can be found in the recent reviews [13,14,15]. I will provide no more than a rough sketch here, in each case discussing only the relativistic impulse approximation for each model. My plots for these results are taken from the Gross and Gilman review [14].

Considering calculations where quantum mechanical wave functions generated in the rest frame are boosted to the Breit frame, we can choose instant form, front form, or point form. Examples of all three appear in Fig. 2. For the instant form we have a calculation by Forest, Schiavilla, and Riska (FSR) [12]. It proceeds essentially as described in the previous subsection, although no $p/M$ expansion is made for the nucleon current, and Wigner rotations are included (approximately) in the boost. The input wave function used is the AV18. This result includes some two-body-current contributions.

For the front form we have two representatives: a calculation by Lev, Pacé, and Salme (LPS), which works in a frame where $Q_\perp = 0$. In Ref. [16] several input wave functions were used, but results displayed here are for the NijmII wave function. We also show a light-front calculation by Carbonell and Karmanov (CK), in which an attempt to restore manifest rotational invariance is made [17]. This is done by allowing the orientation of the light front to be an additional vector in the theory, and then eliminating dependence on it. Unfortunately, only approximate calculations in this approach exist as yet.

Moving to the point form, I first note that for Eq. (16) to be correct I should have $j_\mu$ depending on $k'_1 - k_1$. In the point form (and also in some versions of the front form [10]) this need not be the same as $q$ [18]. A consequence of this is that the squared momentum transfer to the struck nucleon is $f(p)Q^2\eta^2$, where the function $f \gtrsim 1$, but depends on the nucleon relative momentum $p$. A calculation by Allen, Klink, and Polyzou (AKP) using this form of dynamics, together with the AV18 wave function is displayed below [18].
In any one of these forms of quantum mechanics one can take the one-body current operator and evaluate its matrix element between deuteron wave functions, defining that as the relativistic impulse approximation. However, in the front form with \( Q^+ = 0 \) (and also in the point-form calculation of AKP) an impulse approximation such as this implicitly includes some two-body current effects. That is because the full consequences of Poincaré covariance and current conservation are not necessarily respected in calculations of individual matrix elements, and must be imposed on the calculation by a procedure equivalent to introducing two-body currents which act specifically to restore these symmetries.

Also shown in Fig. 2 are two calculations in which \(|\psi\rangle_{\pm \mathbf{q}/2}\) is calculated dynamically, that is to say, the dynamical equation of the theory in the moving frame is solved to obtain \(|\psi\rangle\). Obviously in order to do this one needs an \( NN \) model in which \( V \) can be calculated in an arbitrary frame, and this is simplest for meson-exchange models. The calculation of van Orden, Devine and Gross \([19]\) (vOG) uses an interaction developed in the “spectator formalism” and fit to \( NN \) data \([20]\). That of Phillips, Wallace, and Devine (PWD) employs a two-body equation which incorporates the effects of Z-graphs and relativistic kinematics in the \( NN \) system, but only has approximate boost covariance \([21]\). There the Bonn-B interaction is chosen. Both calculations have impulse approximations which respect deuteron four-current conservation.

### 3.4. Results

Turning now to the results shown in Fig. 2 we see that there are a wide range of RIA predictions. All calculations predict a rapid fall in \( A \), but they agree with the experimental data to differing degrees, especially once \( Q^2 \sim 2 \text{ GeV}^2 \). Agreement with the data at the level of the RIA is not essential for success, since explicit two-body currents can, and probably should, be added to all these calculations. However, large two-body currents might cast doubt upon the efficiency of an expansion in hadronic degrees of freedom. For further discrimination between models it is good to divide out the rapid overall trend, as is done in Ref. \([14]\). I will not do a detailed comparison of models with experiment here.

What is clear is that even were the data to be removed from the slide, there would be no definite prediction for the RIA value of \( A \) at, say, \( Q^2 = 2 \text{ GeV}^2 \). This is true even though most of these calculations respect the Poincaré algebra and current conservation, and remains true even if we restrict ourselves to those calculations using the same rest-frame dynamics (e.g. AKP and FSR). Symmetries and rest-frame dynamics alone are not enough to unambiguously predict \( A \) at these high values of \( Q^2 \).

This is perhaps even clearer in \( B \), where the position of the minimum is an extremely sensitive test of the dynamics: both the implementation of the Poincaré algebra and the choice of potential. But again, it is worth stressing that even calculations with the same rest-frame dynamics produce markedly different predictions for the minimum: the position of the minimum in the AKP and FSR calculations differs by about 30%. And front form and point form even predict opposite directions of the “relativistic shift” of this minimum \([16,18,22]\).

Finally, in \( T_{20} \), most models seem to do reasonably well. The JLab Hall C experiment \([25]\), which extended measurements of \( t_{20} \) out to \( Q^2 = 1.7 \text{ GeV}^2 \), shows that relativistic hadronic models of deuterium have some success in predicting this observable. However, detailed comparison with the accurate data at \( Q^2 \lesssim 0.5 \text{ GeV}^2 \) reveals
Figure 2. Comparison of different relativistic calculations for $A, B,$ and $T_{20}$. All calculations use single-nucleon form factors due to Mergell et al. [23]. The calculations, in order of their minima in $B$, are: CK (long dot-dashed line), PWD (dashed double-dotted line), AKP (short dot-dashed line), two versions of the vOG calculation (solid line and long-dashed line), LPS (dotted line), a quark-model calculation not discussed here (widely-spaced dotted line), FSR (medium-dashed line), and a calculation using a $p/M$ expansion by Arenhoevel, Ritz, and Wilbois [24], also not discussed here (short-dashed line). Figure from Ref. [14], thanks to Ron Gilman and Franz Gross.
that models have some problems in precisely reproducing this data.

To further understand the range of predictions for $A$ and $B$ above 1 GeV$^2$ I want to consider a gedanken-calculation, or actually two gedanken-calculation. Calculations 1 and 2 both begin with the same rest-frame $|\psi\rangle$ and the same one-body $j^\mu$, but they have two different implementations of the Poincaré algebra, i.e.

$$|\psi\rangle_Q^{(1)} = U_Q^{(1)}|\psi\rangle_0; \quad |\psi\rangle_Q^{(2)} = U_Q^{(2)}|\psi\rangle_0.$$  \hspace{1cm} (19)

Since the two $U$‘s are both unitary representations of the Poincaré algebra there is a unitary transformation that can make an exact equivalence between matrix elements in formulation 1 and those in formulation 2. If I choose:

$$J_\mu^{(2)} = U_{q/2}^{(2)}U_{q/2}^{(1)\dagger} j_\mu U_{-q/2}^{(1)}U_{-q/2}^{(2)\dagger},$$  \hspace{1cm} (20)

then:

$$\langle q/2 | J_\mu^{(2)} | q/2 \rangle = \langle q/2 | j_\mu | q/2 \rangle.$$  \hspace{1cm} (21)

Crucially though, the unitary transformation (20) results in $J_\mu^{(2)}$ having two-body pieces, even though $j_\mu$ is only a one-body operator. Thus, the equivalence (21) only holds if these two-body currents are included in $J_\mu^{(2)}$. If the relativistic impulse approximation is used in formulation 2 also, then the results of the two different calculations will differ, by an amount given by the size of the two-body piece induced in Eq. (20).

4. Conclusion

Henri Poincaré once said “Les faits ne parlent pas”: facts do not speak. Nevertheless, let me attempt to interpret the results presented here.

I draw two lessons from them. Firstly, there is no unique relativistic impulse approximation prediction for electron-deuteron scattering. The RIA is only defined once a particular implementation of the Poincaré algebra is chosen. Associatedly, the size of the variation in Fig. 2 is indicative of the size of two-body contributions to $A$ and $B$.

Does this mean then, that the situation is hopeless? Not at all: it is just that when looking at an RIA result one must know how the Poincaré algebra was implemented. Thus, the equivalence (21) only holds if these two-body currents are included in $J_\mu^{(2)}$. If the relativistic impulse approximation is used in formulation 2 also, then the results of the two different calculations will differ, by an amount given by the size of the two-body piece induced in Eq. (20).

More data would be useful in order to help different theorists pin down these contributions in their formalisms. A theory-independent statement about the size of two-body currents is impossible at $Q^2 \sim 2 \text{ GeV}^2$, but once a specific theory is chosen data is needed to calibrate the two-body currents which contribute significantly to elastic electron-deuteron scattering at these momentum transfers. Especially useful would be more data around the minimum of $B$, since this is a barometer sensitive to many model elements.

Having worked hard to tune-up our calculations to reproduce elastic electron-deuteron data, an important future step is the extension of these calculations to deuteron electrodisintegration (for steps in this direction see [27,28]). Finally, much of what was said
here also applies to the three-nucleon system. The boosted $NN$ t-matrix is input to the Faddeev equations for $NNN$ bound and scattering states. Different implementations of the Poincaré algebra lead to different boosted $NN$ t-matrices. The unitary equivalence between these implementations is maintained only if the associated three-body forces are kept in the three-nucleon calculation.

**Acknowledgements**

I thank Fritz Coester, Ron Gilman, Franz Gross, Brad Keister, Jerry Miller, Wayne Polyzou, John Tjon, and especially Steve Wallace for stimulating discussions on the topics discussed herein and, in some cases, for sharing their results with me for use in my talk. This work was supported by the U. S. Department of Energy under grants DE-FG02-93ER40756 and DE-FG02-02ER41218.

**REFERENCES**

1. J. Carlson and R. Schiavilla, Rev. Mod. Phys. **70**, 743 (1998).
2. S. Pieper, these proceedings.
3. J. L. Forest, V. R. Pandharipande and J. L. Friar, Phys. Rev. C **52**, 568 (1995).
4. R. A. Krajcik and L. L. Foldy, Phys. Rev. D **10**, 1777 (1974).
5. P. A. M. Dirac, Rev. Mod. Phys. **21**, 392 (1949).
6. J. R. Cooke and G. A. Miller, Phys. Rev. C **66**, 034002 (2002).
7. B. Bakamjian and L. H. Thomas, Phys. Rev. **92**, 1300 (1953).
8. B. D. Keister, in “Few-body Problems in Physics, Williamsburg, 1994”, F. Gross, editor (AIP Press, New York, 1994).
9. D. Abbott et al., Eur. Phys. J. **A7**, 421 (2000).
10. J. Adam and H. Arenhoevel, Nucl. Phys. A **614**, 289 (1997).
11. S. J. Wallace, Phys. Rev. Lett. **87**, 180401 (2002).
12. R. Schiavilla and V. R. Pandharipande, Phys. Rev. C **65**, 064009 (2002).
13. M. Garcon and J. W. Van Orden, Adv. Nucl. Phys. **26**, 293 (2001).
14. R. Gilman and F. Gross, J. Phys. G **28**, R37 (2002).
15. I. Sick, Prog. Part. Nucl. Phys. **47**, 245 (2001).
16. F. M. Lev, E. Pace and G. Salme, Phys. Rev. C **62**, 064004 (2000).
17. J. Carbonell and V. A. Karmanov, Eur. Phys. J. A **6**, 9 (1999).
18. T. W. Allen, W. H. Klink and W. N. Polyzou, Phys. Rev. C **63**, 034002 (2001).
19. J. W. Van Orden, N. Devine and F. Gross, Phys. Rev. Lett. **75**, 4369 (1995).
20. F. Gross, J. W. Van Orden and K. Holinde, Phys. Rev. C **45**, 2094 (1992).
21. D. R. Phillips, S. J. Wallace and N. K. Devine, Phys. Rev. C **58**, 2261 (1998).
22. F. Coester and W. Polyzou, private communication.
23. P. Mergell, U. G. Meissner and D. Drechsel, Nucl. Phys. A **596**, 367 (1996).
24. H. Arenhoevel, F. Ritz and T. Wilbois, Phys. Rev. C **61**, 034002 (2000).
25. D. Abbott et al., Phys. Rev. Lett. **84**, 5053 (2000).
26. D. R. Phillips and S. J. Wallace, in progress.
27. R. Schiavilla and D. O. Riska, Phys. Rev. C **43**, 437 (1991).
28. J. Adam, F. Gross, S. Jeschonnek, P. Ulmer and J. W. Van Orden, Phys. Rev. C **66**, 044003 (2002).