LA-UR-88-1664

LA-UR--88-1664

DE88 010946

TITLE: MESIC INTERACTIONS AND TRINUCLEAR PHYSICS

AUTHOR: A. Picklesimer
Physics Division, Los Alamos National Laboratory,
Los Alamos, New Mexico 87545

INVITED TALK:

Workshop on Relativistic Nuclear Many-Body Physics,
Columbus, Ohio, 1988,
Proceedings to be Published by World Scientific.

Los Alamos Los Alamos National Laboratory
Los Alamos, New Mexico 87545

FORM NO 836 R4
ST NO 2629 5/81

MESIC INTERACTIONS AND TRINUCLEAR PHYSICS

A. Picklesimer†
Physics Division
Los Alamos National Laboratory
Los Alamos, New Mexico 87545
USA

Because some simple technical aspects of the Faddeev equations are needed to fully appreciate the topics to be discussed, and since everyone here is probably not intimately familiar with the nuclear three-body problem or the Faddeev equations, I want to begin with a brief review. The structure of the bound state Faddeev equation can be inferred from some simple manipulations. Consider the nonrelativistic Schrödinger equation:

$$[G_o^{-1} - \sum_i v_i] \sum_j \psi_j = 0 \tag{1}$$

where the potential $V = \sum v_i$, i=1-3, is given as a sum of two-body forces and we have used the "odd man out" notation where, e.g., v_1 denotes the interaction between particles 2 and 3. In (1), $G_o^{-1} = E - H_o$ where H_o is the free three-body Hamiltonian, i.e. the sum of the kinetic energy operators of the three particles $H_o = \sum K_i$, and we have anticipated an expansion of the bound state Ψ in components, $\Psi = \sum \psi_i$. Equation (1) can be rewritten as

$$\sum_i [\, G_o^{-1} - v_i \,] \psi_i = \sum_i \sum_{j \neq i} \psi_j \tag{2}$$

and by equating components as suggested by (2)

$$G_i^{-1} \psi_i = v_i \sum_{j \neq i} \psi_j \tag{3}$$

where $G_i^{-1} = G_o^{-1} - v_i$. Equation (3) characterizes the decomposition $\Psi = \sum \psi_i$ as one in which ψ_i is that part of Ψ in which the i^{th} pair interacts last. Now since $G_i v_i = G_o T_i$, where $T_i = v_i + v_i\, G_o\, T_i$ is a free two-body t-matrix embedded in the three-body space, the bound state equation becomes just

$$\psi_i = G_o(E_t)\, T_i(E_t - K_i)\, P\, \psi_i \tag{4}$$

where we have denoted the bound state (triton) energy as E_t and have made explicit the energy argument of G_o and the two-body t-matrix T_i. In writing (4) we have also made use of the overall antisymmetry of Ψ, and the symmetric nature of the interpretation of the ψ_i, to infer that the ψ_i satisfy $\psi_i = -E_{ij}\, \psi_j$, where E_{ij} is the exchange operator for particles 'i' and 'j', so that P denotes[1] the sum of the three-particle cyclic and acyclic permutation operators (ψ_i is antisymmetric in the pair 'i' since T_i is).

Equation (4) exhibits all of the technical aspects of the problem we need to consider explicitly[1]. First, in the three-body center of mass $\psi_i = \psi_i(\vec{p}, \vec{q})$ is a function of two momenta, the momentum $\vec{q}$ of the spectator particle 'i' and the relative momentum $\vec{p}$ of the interacting pair. The operator T_i is simply a two-particle t-matrix whose "energy" happens to be the operator $(E-K_i)$ rather than just E. Making explicit the dependence on K_i and the pair center of mass energy in momentum space, $T_i(E-K_i) \rightarrow t_i(E-3q^2/4m)$, where t_i is now just the familiar two-body t-matrix evaluated at the energy $(E-3q^2/4m)$. Thus

- Because P is nondiagonal in the spectator momentum $\vec{q}$, the triton equation scans over the two-body t-matrix for all $E < E_t$ as $E-3q^2/4m$ varies over the range of q. Thus, the triton binding energy is determined by the behavior of the two-body t-matrix at energies below about -8.5 MeV, i.e. below the physical range where it is parameterized.
- Because P couples different angular momentum states, the interaction kernel $t_i P$ in (4) couples together a number of basis states in any representation chosen. The triton equation (4) is generally expressed as a coupled channel problem, where a specific channel correponds to the coupling of a specific angular momentum (LSJ) of the interacting pair to a specific

angular momentum state of the spectator (lsj) to form the spin 1/2 triton. Equation (4) is then solved for a finite set of channels.

The choice of which channels to include and which to omit is somewhat arbitrary, but certain choices have become standard. Typically, one specifies a maximum value of J and whether or not to include odd parity (π) pair states. For example, the choices ($J \leq 1$; $\pi = +$), ($J \leq 2$; $\pi = +$), ($J \leq 2$; $\pi = \pm$), and ($J \leq 4$; $\pi = \pm$), involves 5, 9, 18, and 34 channels, respectively. Five channel results are of interest because most of the binding energy effects are already found at this level. Table I lists the channels which comprise the five-channel case.

TABLE I. The three-body five-channel basis.

CHANNEL	L	S	J	ℓ	s	j
1	0	0	0	0	1/2	1/2
2	0	1	1	0	1/2	1/2
3	0	1	1	2	1/2	3/2
4	2	1	1	0	1/2	1/2
5	2	1	1	2	1/2	3/2

Of special interest in this talk will be the two-channel case which consists of the first two entries listed in Table I. This case is of particular interest because most of the binding energy is already present for this coupling, and because it also turns out to contain most of the interesting physics at issue for the triton binding energy defect. For the two channel case, L=0 and the tensor-force-like term of the t-matrix does not contribute to the triton calculation of (4). This is not to say that the NN tensor force does not contribute since the full triplet potential is used to generate the NN t-matrix. However, this is a major simplification which we will later exploit. Now let's assess the status of triton binding energy predictions from nonrelativistic potential models. For nearly two decades, and until very recently,[2,3] attempts to reconcile nonrelativistic potential theory with the experimental value of the triton binding energy consistently met with failure[4-8]. This pattern persisted despite the introduction of a host of independent "realistic" potential models[9-14] which differed in detail, yet provided accurate descriptions of two-body bound state and nucleon-nucleon (NN) scattering data[15]. The basic features common to these attempts included:

- Use of the nonrelativistic Schrödinger equation.
- Unretarded, static (energy-independent) potentials.
- Potentials with (varying) nonlocalities (velocity dependences).
- NN forces only, i.e., no three- or many-body forces.

Thus it appeared that the strictures of this theoretical framework were perhaps incompatible with the triton binding and that strong three-body forces were required to understand the triton[4,5,7,8]. However, it has long been known that an approximately linear relationship exists between predictions for E_t and for the deuteron % D state (P_D) for a given model, and that the physical binding lies roughly on the line formed by plotting the predictions for E_t vs. the prediction for P_D made by various models[6,16]. Just what drives this relationship has been poorly understood, but we'll get to that shortly.

The status quo changed recently when it was found that the energy-independent (static) potential approximations to the full Bonn interaction[17] predict a triton binding energy which is in good agreement with the experimental value of 8.48 MeV, well within a reasonable uncertainty[2,3]. Table II shows representative predictions for E_t from a variety of models, decomposed according to the number of channels included in the calculation.

TABLE II. Channel decomposition of the predictions made for the triton binding energy (in MeV) by various realistic forces.

POTENTIAL	P_D	2	5	9	18	34
RSC	6.47	6.59	7.04	7.21	7.23	7.35
PARIS	5.77		7.30		7.38	7.64
SSC	5.45		7.46	7.52	7.49	7.53
V14	6.08		7.44	7.57	7.57	7.67
TRS	5.92		7.49	7.56	7.52	7.56
BONN	4.38	8.16	8.36	8.44	8.32	8.35

Several remarks should be made in regard to Table II:

- Excepting the Bonn results, the static potentials display a consistent defect in underpredicting the triton binding.
- Excepting the Bonn and V14 potentials which are fitted to np ^1S_o data, all of the other potentials shown in Table II are fitted to pp ^1S_o data. Thus all of the "other" potentials in Table II will gain in binding energy when charge symmetry breaking is taken into account, whereas the Bonn and V14 potentials will lose some binding[18].
- Potential models with larger P_D tend to get more binding energy contribution from higher partial waves than do models with lower P_D.
- The static Bonn prediction breaks the trend of persistent failure to predict the triton binding. Just how this results as a direct consequence of the low P_D of the static Bonn model is considered shortly.
- The essence of the binding energy results for the various potentials is already present at the five channel level, moreover, the essential distinction between the predictions of the other realistic forces and the Bonn potential is already present at the two channel level.

Thus from this point on we can concentrate our attention on the ^1S_o and $^3S_1$ partial waves which comprise the two-channel case. Since the only input to the two channel calculation are the ^1S_o and $^3S_1$ two-body t-matrices, the major differences between the predictions of the static Bonn model and those of the other potentials evidently arises from differences in these S-wave t-matrices. The ^1S_o part of the interaction can be characterized by the scattering length it produces and the $^3S_1$ part by its P_D prediction. It turns out that realistic variations in the singlet scattering length (and effective range), in the context of charge symmetry breaking effects, lead to variations in the prediction of E_t on the order of 200 KeV. Thus, apart from such significant, but relatively minor, variations induced by the ^1S_o channel, we can focus on the $^3S_1$ NN partial wave from here on[3,18]. The principal source of the distinction between the Bonn and other models, apparently the result of differing tensor/central force admixtures, evidently resides therein.

To clarify the mechanism through which this obtains, we note that both the central and tensor forces are attractive in the even parity J=1 state, and that different realistic forces achieve the same ~ 2.2 Mev deuteron binding through different relative strengths of their tensor and central components. A stronger tensor force implies a weaker central force, and vice versa, in order to achieve the observed deuteron binding. In the triton, the tensor force apparently contributes relatively less to the binding than does the central force in order to produce the observed correlation: that a lower P_D yields a more strongly bound triton. In order to examine more closely the mechanism by which, all other things being equal, a stronger tensor force implies a weaker triton binding, we found it useful to generate variants of the static Bonn OBEPQ potential which were as alike as possible, differ mainly in their P_D prediction, and fit all other low-energy parameters with high accuracy. This approach is intended to ensure that extraneous off-shell variations among these models is minimal. Table III lists the P_D and the two- and five-channel E_t predictions associated with these models, as compared to the OBEPQ and the RSC potential[3].

TABLE III. Predictions for E_t vs. potential model.

POTENTIAL	P_D	2-CH	5-CH
A(OBEPQ)	4.38	8.16	8.36
B	5.03	7.87	8.14
C	5.60	7.62	7.94
RSC	6.47	6.59	7.04

Armed with our variable-tensor models we now exploit, in two ways, the fact that we are able to restrict ourselves to the $^3S_1$ channel[3]. First, the two-body t-matrix T(E) satisfies the equation

$$T(E) = V + VG_o(E)T(E) \tag{5}$$

so that, with $V = V_C + V_T$, where V_C denotes the central and V_T the tensor force, we have for the coupled S and D partial waves:

$$T^S(E) = V_C^S + V_C^S G_o(E)T^S(E) + V_T^{SD}G_o(E)T^{DS}(E) \tag{6a}$$

$$T^{DS}(E) = V_T^{DS} + V_T^{DS} G_o(E) T^S(E) + V^D G_o(E) T^{DS}(E) \tag{6b}$$

where $V^{SS} = V^S$, etc. Thus, the $^3S_1$ t-matrix $T^S(E)$ is given by

$$T^S(E) = U^S(E) + U^S(E) G_o(E) T^S(E) \tag{7a}$$

for the energy-dependent effective central potential

$$U^S(E) = V_C^S + V_T^{SD}[G_o^{-1}(E) - V^D]^{-1} V_T^{DS} \tag{7b}$$

which, in first order is just

$$U^S(E)^{(1)} = V^S + V_T^{SD} G_o(E) V_T^{DS} \quad . \tag{7c}$$

Equations (7) express the characteristic mechanism by which tensor force attraction is reduced below threshold, namely by the energy dependence explicit in (7c) and, more precisely, in (7b). Given the solution of (6) for $T^S(E)$, it is a simple matter in momentum-space to turn off the tensor term in (6a), and solve for the potential, thus obtaining the unique effective central potential $U^S(E)$ which yields $T^S(E)$ via (7a). The advantage of $U^S(E)$ over the tensor force is simply that we have traded the latter for a simple parametric dependence.

Given $U^S(E)$, we now exploit the second special advantage of the $^3S_1$ channel. There exists a nuclear bound state at energy β ($\simeq$ -2.2 MeV) in this channel called the deuteron. Because realistic interactions correctly "predict" β, the effective potential $U^S(\beta)$ must also yield this bound state. More generally, we can gauge the strength of the potential $U^S(E)$ for arbitrary values of E simply by treating E as a parameter, fixing its value arbitrarily, and determining what "deuteron" binding energy $|\beta|$ results. The results of such calculations are shown in Figure 1 for the effective potentials generated by the Bonn OBEPQ potential and for the variable tensor models B and C described above. The striking features of Fig. 1 are twofold. First for each of the models, for energies below the deuteron energy, the strength of the effective interaction decreases monotonically with E. Thus $T^S(E)$ corresponds to an ever weakening potential $U^S(E)$ as the parametric energy $E = E_t - 3q^2/4m$ decreases over the range of values relevant to the triton. This implies that as the momentum of the spectator nucleon increases, the contribution to the triton binding weakens. The decrease

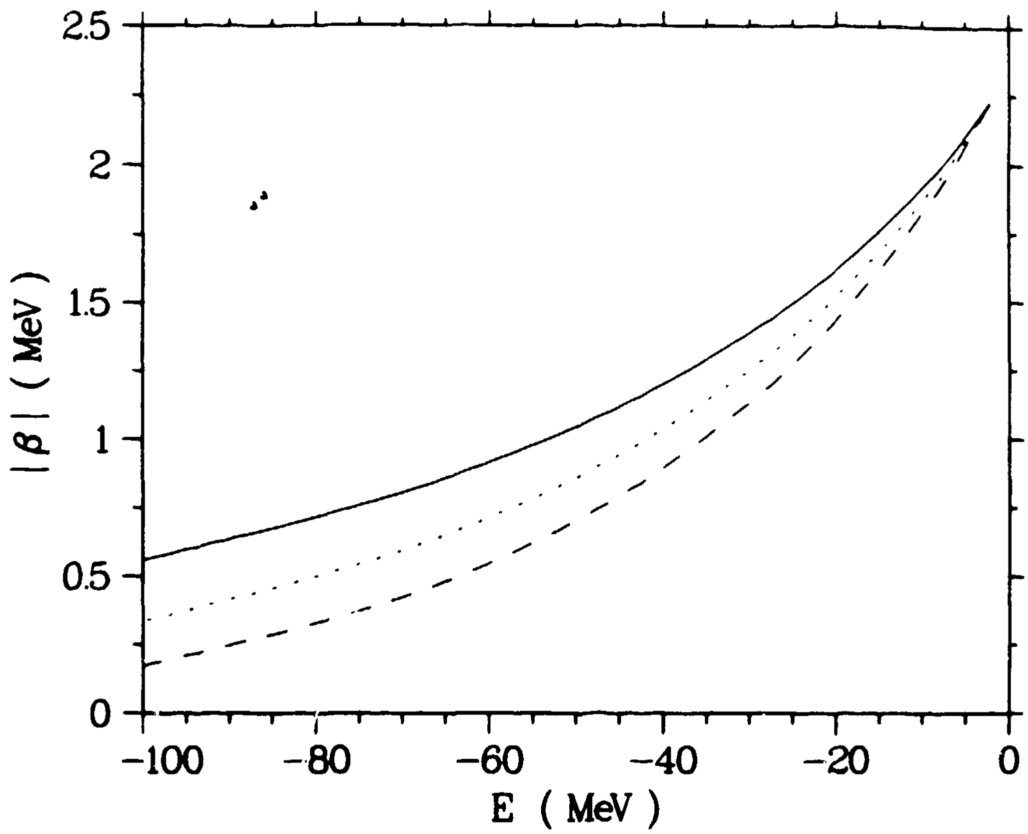

FIGURE 1 The binding energy of the "deuteron" calculated from the effective central potential $U^S(E)$ plotted as a function of E. The solid, dotted, and dashed lines represent, respectively, potentials A (OBEPQ), B, and C [See text].

in potential strength as E moves from β down to about -100 MeV is about 1.5 to 2.0 MeV for the range of models (P_D = 4-7%) considered here. Second, as the P_D associated with the various static potentials increases, the effective strength of the interaction decreases universally for all E below β, i.e., the curves do not cross. We note that these results depend upon the full character of $U^S(E)$ as expressed in (7b), the first order approximation $U^S(E)^{(1)}$ of (7c) displays erratic behavior even at energy β. Thus, the picture which emerges from Fig. 1 is a delicate consequence of the full dynamics entailed in the effective central potential $U^S(E)$.

Thus in the calculation of the triton binding energy, as $T^S(E)$ is sampled over its parametric dependence on E, the stronger the tensor force of the model, the weaker the effective potential associated with $T^S(E)$ *for all* relevant E. Hence, we are justified in making the following statements:

- The (attractive) strength of $U^S(E)$ decreases monotonically with E.
- Stronger tensor forces imply universally weaker $U^S(E)$ *for all* $E < E_t$.
- More generally, the effective potentials associated with stronger tensor forces are weaker; the triton $^3S_1$ channel everywhere "sees" weaker attraction for models with stronger tensor forces.

Thus the "reconciliation" of the triton binding energy and nuclear potential theory observed with the static OBEPQ potential essentially results from the fact that this Bonn potential differs from the other realistic potentials in having a much weaker tensor force, typified by a deuteron D-state percentage (P_D) of only 4.4%, versus a minimum of about 5.5% for the others[2,3].

Although the success of the static Bonn potentials is gratifying from the point of view of nonrelativistic potential theory, the foregoing raises an issue of central importance. We have just seen that the implicit energy-dependence associated with the unique Bonn tensor potential strength is crucial to the success of the static Bonn potentials. But the successful triton predictions derive from these static, potential-model approximations to the full, retarded, *energy-dependent* Bonn interaction rather than from the full interaction itself. Thus, one is led to inquire whether or not the static potential-model results straightforwardly reflect the physical content of the full meson-theoretic Bonn model.

As far as the two-body data is concerned, all of the Bonn interactions agree very well in their predictions in the important low-energy regime, including P_D. One regime where both the momentum-space (k-space) OBEPQ potential and the simpler coordinate-space (r-space) approximations differ from the full interaction is in their predictions for NN $\ell > 0$ scattering phase shifts above $\sim$ 100 MeV[17]. However, the phase shift predictions of the two static potentials deviate as much from one another as they do from the full interaction in this region, and yet the static potentials produce nearly identical predictions for the triton binding energy[2]. Just how the high-energy phase shifts are reflected in the triton binding energy is not clear, but there seems to be no reason to believe, on the basis of the phase shift behavior alone, that the full Bonn interaction would deviate greatly from the two static potentials in its triton binding energy predictions. While the static potentials' close agreement for (positive-energy) scattering observables may reasonably be taken to be indicative of their substantial agreement for energies relevant to the triton binding as well, one cannot similarly infer that the same is true of the full interaction. The reason for this is the triton predictions' dependence on the NN t-matrix for negative values of the parametric energy, $E = E_t - 3q^2/4m$. Simply put, one expects some differences in the negative-energy behavior of the t-matrix corresponding to the energy-dependent interaction relative to the behavior of its energy-independent counterparts.

Other indications of the relevance to the three-body problem of dispersive energy-dependent terms in the two-body interaction come from studies of virtual isobar (Δ-resonance) contributions. For example, an extensive consideration of isobar effects correlated in the two- and three-body systems is detailed in Refs. 19 and 20. In this model only one Δ is allowed at a time, there is no diagonal Δ-nucleon interaction, and the coupled-channel approach employed permits only a subset of the time-ordered diagrams of interest[19]. Within this context energy-dependent effects are studied by adding to a static potential (in this case the Paris potential[14]) the difference between nonstatic single-isobar diagrams and a limiting static value, yielding an energy-dependent interaction. This simplified model differs from a realistic meson-theoretic interaction such as the Bonn interaction in that energy-dependent effects from one boson exchange, double-Δ diagrams, and irreducible multi-meson exchange are not represented and only a small subset of the single-Δ time-ordered diagrams can be considered. Nevertheless, the model of Refs. 19 and 20 allows one to treat not only two-body energy dependence in the three-body problem but also the coupled-channel approach allows a considerable consistency between the model of the two- and three-body forces and the inclusion of nucleon self-energy modifications to the two-body force. One can easily infer the repulsive nature of the energy-dependent effects in the model of Ref. 19 from the structure of second-order perturbation theory; the net repulsive effect from the two-body energy-dependence is found to be large, reducing the predicted triton binding energy by .58 MeV. For later reference, we note that it also turns out in Ref. 19 that three-body force and coupled-channel effects more than compensate this repulsion to yield a net attractive contribution of 300 KeV to the triton binding. Perturbative methods, partly justified on the basis of Ref. 19, were then developed in Ref. 20 to obtain a similar estimate for the role of double-Δ diagrams. The result was a large 1.3 MeV reduction in the triton binding which in this case is not compensated by three-body effects. Although Ref. 20 emphasizes the uncertain, exploratory nature of this result the implication is clear: energy-dependent effects must be taken seriously.

Because its semi-relativistic context[17] is not consistent with the use of nonrelativistic three-body equations, direct use of the full Bonn interaction in Faddeev calculations is inappropriate. However, an energy-dependent one boson exchange potential (OBEP) representation of the full interaction exists and this can be easily (but only approximately) corrected for its use in a nonrelativistic context, as described below. This approximation to the full interaction, the "OBEPT" interaction of Ref. 17, provides a much more accurate representation of the full interaction than its energy-independent counterparts since it retains the exact form of the one-meson exchange diagrams and higher-order (isobar and multi-meson) diagrams are also well-described by such a parameterization[17]. Because no static limit is imposed, W(E) incorporates to some degree the energy-dependence of the full interaction. Like the full interaction, W(E) is constructed for use in a scattering equation of the Lippmann-Schwinger form which employs a relativistic Schrödinger propagator,[21] wherein nonrelativistic kinetic energies are replaced by relativistic ones. It therefore yields a two-body t-matrix which is suited for use in the corresponding relativistic extension of the nonrelativistic three-body equations, obtained by applying the Bonn Lagrangian, time-ordered perturbation theory, and the Bonn truncation schemes[17] to the three-nucleon system. However, the free three-body propagator of such an approach is the relativistic Schrödinger propagator. Thus, we can not immediately apply W(E) in our nonrelativistic study of the triton binding energy.

To adapt the quasipotential OBEPT to our purposes we insert W(E), $E = 2[k^2+m^2]^{1/2}$, into the nonrelativistic ($E = k^2/m$) Lippmann-Schwinger equation in the center of mass, rather than into the corresponding relativistic Schrödinger equation, and then readjust the parameters of W(E) to recover the predictions for two-body observables[22]. Only slight adjustments are needed, and in fact, this correction has little real effect on the predicted triton binding energy. At any rate, the process above yields an energy-dependent quasi-potential, W'(E), adapted for use in nonrelativistic momentum-space three-body calculations. We note that W'(E) is, like W(E), a more faithful representation of the full interaction than are the static Bonn potentials; only a very minor approximation has been introduced in obtaining W'(E) from W(E), as compared to the substantial modifications which result from the additional requirement of complete energy-independence. It would require very delicate interplay between the propagators internal to W(E) and those of the relativistic Schrödinger equation to invalidate this approximation.

One can now use the energy-dependent Schrödinger quasi-potential W'(E) to gauge the extent to which the triton predictions of the full Bonn interaction can be expected to mirror those of its energy-independent approximations. The result is remarkable;[22] the five-channel triton binding energy predicted on the basis of W'(E) is only 6.73 MeV. This is to be compared to a five-channel result based on OBEPQ of 8.36 MeV. Thus the actual repulsive retardation effect found here to characterize a full meson-theoretic interaction, namely the Bonn interaction, is a full three times the effect observed in the model studied in Ref. 19. That this result has nothing to do with the distinction between W(E) and W'(E) is already seen at the two-channel level where W'(E) yields 6.65 MeV, W(E) [uncorrected for its use in nonrelativistic three-body equations] 6.48 MeV, and OBEPQ 8.16 MeV.

To understand this result we proceed as before and study the strength of the effective central potentials $U^S(E)$ associated with W'(E) and W(E). The results of such calculations[22] are shown in Figure 2 along with the corresponding results for the $U^S(E)$ based on OBEPQ and the two other static potentials used in obtaining Fig. 1. Here we see a qualitatively different behavior from that of

Fig. 1. Although W(E) and W'(E) correspond to low P_D (P_D = 4.24% and P_D = 3.95%, respectively), i.e. to weaker tensor forces, the fall-off of their effective

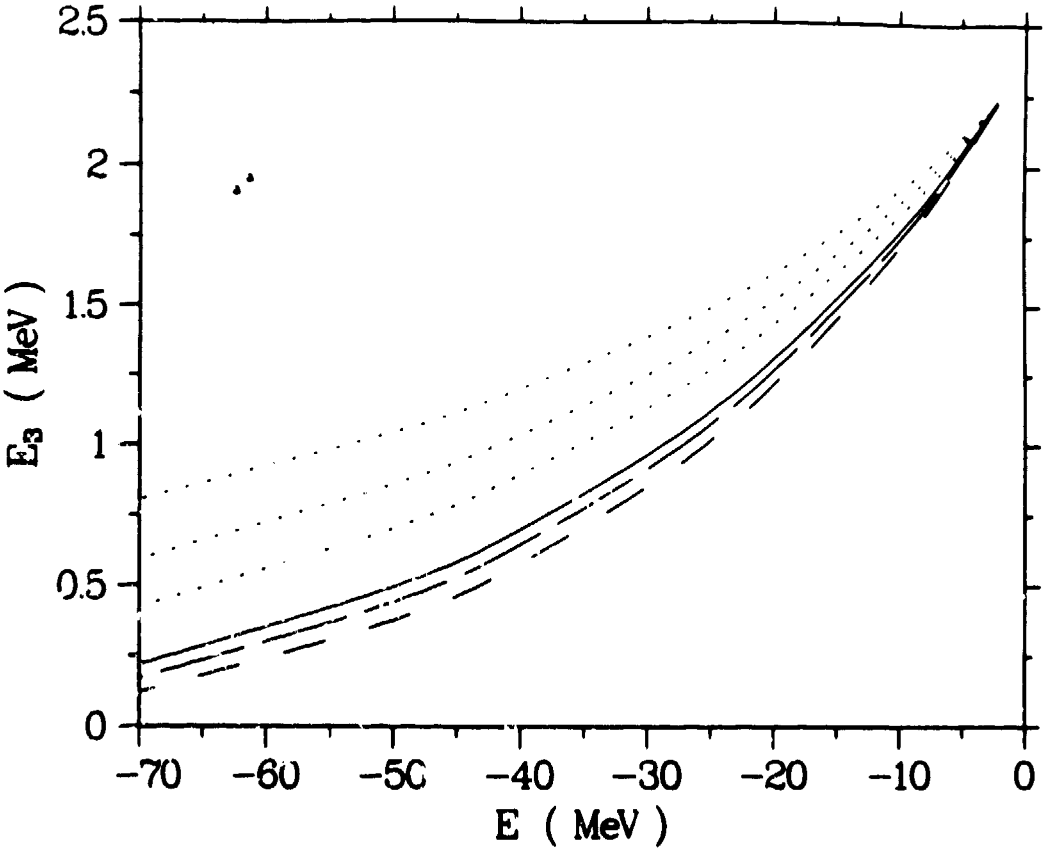

FIGURE 2 The binding energy of the "deuteron", calculated from the effective potentials $U^5(E)$ as a function of E, for a variety of underlying realistic potential models (See text). The dotted curves represent the static potentials shown in Fig. 1; from top to bottom they correspond to P_D = 4.38%, 5.03% and 5.60%, respectively. The solid curve corresponds to the Reid potential, P_D = 6.48%. The dashed curves correspond to the energy-dependent interactions W(E) and W'(E) described in the text. The short-dash curve represents the OBEPT interaction of Ref. 17, P_D = 4.24%, the long-dash curve our adaptation of OBEPT, W'(E), P_D = 3.95%.

attraction in Fig. 2 as (−E) increases breaks the pattern established in Fig. 1 for the static potentials and does not resemble that corresponding to OBEPQ at all. In fact it more nearly resembles that of the Reid potential, also shown in Fig. 2, for which P_D = 6.48%. Evidently the explicit energy dependence of W(E) [or W'(E)] adds to the energy dependence induced by the tensor force so that the net result roughly corresponds to that of the large P_D "realistic" forces. From Fig. 2 it is also evident that the enhanced attractiveness of the static Bonn potentials, which results in their successful triton predictions, is due to the fact that they reflect only the implicit energy dependence (low P_D) of W(E) and not the explicit energy dependence. Thus, the successful triton predictions of the static Bonn interactions rest on the fact that they do not reproduce the negative-energy behavior of the full interaction!

If the triton predictions of the nonrelativistic static Bonn potentials are to retain any significance with regard to their meson-theoretic origins, it is necessary to reconcile the triton prediction made by W'(E) [or W(E)] with the experimental value. The only way to do this is to appeal to the meson theory which prescribed the energy dependence in the first place. Thus, the resolution of this issue must lie with the set of three-nucleon equations which consistently incorporate the energy dependence of the full interaction. The interaction W(E) (and for that matter W'(E)) which enters into the three-body calculation is just (an approximation to) a collection of meson exchange and virtual-Δ diagrams. Although this set of diagrams possesses a high degree of self-consistency in the two-body problem, this is NOT the case when the two nucleons are embedded in a three- or many-body problem.

The reason for this is that the meson-theoretic framework implies the existence of additional diagrams in, e.g., the three-body circumstance, which are required to retain theoretical consistency and which arise from the retarded nat[illegible] of the two-body diagrams. The basic Δ three-body-force diagram was found in the model of Ref. 19 to provide an additional attractive contribution to E_t of .8 MeV. During the time two of the fermions interact, either or both of them may interact, perhaps repeatedly, with the third fermion. Inclusion of such self-energy diagrams can be viewed as effectively modifying the energy dependence of the set of two-body diagrams which comprise W(E) in the three-body space. If this results in a net attraction then this will tend to offset the $(3q^2/4m)$ kinetic term in the potential $W(E_t\text{-}3q^2/4m)$ and thus effectively reduce the energy dependence of W(E) and W'(E) by biasing the range of important energies E in Fig. 2 closer to E_t. This would move the corresponding triton predictions back toward those obtained using the static potentials. A small attractive contribution to E_t associated with such self-energy modifications has been observed in the model of Ref. 19. One can speculate from Fig. 2 that the requisite three-body forces would have to supply a shift in the parametric energy of ~ 5 MeV or more. If this turns out to be the case, then the appropriate two-body static potential theory limit would be understood to result from a cancellation between energy-dependent meson-theoretic two-body effects and the associated three-body forces[19]. In this event, we will have obtained a much deeper understanding of nuclear dynamics. If this picture does not obtain, then the successful triton predictions of the static Bonn potentials are disconnected from the meson-theoretic framework and their fundamental significance is greatly diminished. The issue of the triton binding energy has progressed to the point where two-body and three-body forces need to be treated in a unified and consistent manner. Although one can certainly "fit" the triton using either a three-body force or an energy-dependent phenomenology, this has little theoretical significance. What is required to resolve this issue is a consistent three-body treatment based upon the physical model which underlies meson-theoretic interactions. An understanding of trinuclear binding can never be attained on the basis of ad-hoc static potential models alone, but must include an elaboration and understanding of energy-dependent effects present or neglected in the two-body force and this must be done in conjunction with the relevant associated three-body mechanisms. Thus the successful prediction of the triton binding energy by the static Bonn potentials, which called into question the need for three-body forces in describing the triton, has led full circle back to an appeal to three-body forces.

† *The work upon which this talk is based was performed in collaboration with R. A. Brandenburg, G. S. Chulick, R. Machleidt, and R. M. Thaler.*

REFERENCES

1. W. Glöckle, *Nucl. Phys.* **A381**, 343 (1982).
2. R. A. Brandenburg, G. S. Chulick, R. Machleidt, A. Picklesimer, and R. M. Thaler, LA-UR-86-3700, Los Alamos, (1986); T. Sasakawa, *Nucl. Phys.* **A463**, 327c (1987).

3. R. A. Brandenburg, G. S. Chulick, R. Machleidt, A. Picklesimer, and R. M. Thaler, *Phys. Rev.* **C37**, 1245 (1988).

4. J. L. Friar, B. F. Gibson and G. L. Payne, *Comments Nuc. Part. Phys.* **11**, 51 (1983).

5. G. L. Payne, J. L. Friar, and B. F. Gibson, *Phys. Rev. Lett.* **55**, 374 (1985).

6. A. C. Phillips, *Rep. Prog. Phys.* **40**, 905 (1977).

7. J. L. Friar, et al., in **Progress at LAMPF**, LA-10738-PR, Los Alamos, May, 1986, p.220.

8. C. R. Chen, G. L. Payne, J. L. Friar, and B. F. Gibson, *Phys. Rev.* **C33**, 401 (1986) and *Phys. Rev.* **C31**, 2266 (1985).

9. For example, those of Refs. 10-14.

10. R. V. Reid, *Ann. Phys. (N.Y.)* **50**, 411 (1968).

11. R. B. Wiringa, R. A. Smith, and T. A. Ainsworth, *Phys. Rev.* **C29**, 1207 (1984).

12. R. de Tourreil, and D. W. L. Sprung, *Nucl. Phys.* **A201**, 193 (1973).

13. R. de Toureil, B. Rouben, and D. W. L. Sprung, *Nucl. Phys.* **A242**, 445 (1975).

14. M. Lacombe et al., *Phys. Rev.* **C21**, 861 (1980).

15. See, e.g., Refs. 1 and 3.

16. A. C. Phillips, *Nucl. Phys.* **A107**, 209 (1968).

17. R. Machleidt, K. Holinde, and Ch. Elster, *Phys. Rep.* **149**, 1 (1987).

18. R. A. Brandenburg, G. S. Chulick, R. Machleidt, A. Picklesimer, and R. M. Thaler, *Phys. Rev.* **C**, In Press.

19. Ch. Hajduk, P. U. Sauer and W. Strueve, *Nucl. Phys.* **A405**, 581 (1983).

20. Ch. Hajduk, P. U. Sauer and Shin Nan Yang, *Nucl. Phys.* **A405**, 605 (1983).

21. See, e.g., Appendix A of Ref. 17.

22. Due to their relativistic nature both the full Bonn interaction and W(E) depend on the total two-body center of mass momentum. These non-Galilean frame dependences are neglected in our work, as they also necessarily are whenever static potentials are employed.

23. R. A. Brandenburg, G. S. Chulick, Y. E. Kim, D. J. Klepacki, R. Machleidt, A. Picklesimer, and R. M. Thaler, *Phys. Rev.* **C37**, 781 (1988).