PHYSICAL REVIEW C VOLUME 27, NUMBER 4 APRIL 1983

Weak nuclear interactions in a hybrid baryon-quark model: *p-p* asymmetry

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(Received 12 November 1982)

The hybrid multi-baryon-quark shell model of nuclei is developed for the study of weak interactions in nuclei. Application to the parity violating asymmetry in *p-p* elastic scattering is carried out.

[NUCLEAR REACTIONS Weak asymmetry in elastic *pp* scattering. Energies up to 800 MeV. Quark model of weak interactions in nuclei. Quark structure of nuclei.]

I. INTRODUCTION

For many years boson exchange models have been used to provide a theoretical description of weak interactions in nuclei. Included in these studies is the parity violating asymmetry in *pp* elastic scattering. Experiments on the total asymmetries have been carried out[1] at 15, 45, and 6000 MeV. The theoretical values at the lower energies are not in obvious disagreement with the data, while at 6 GeV the calculations are an order of magnitude too small.[1]

From a fundamental point of view these boson exchange models for weak nuclear interactions are not satisfactory. The physics of short distance plays a crucial role in many of the weak processes, such as the *pp* asymmetry. In the traditional boson exchange model the short-range weak interaction is derived by using ρ and ω exchange from essentially structureless nucleons. This picture has not been justified in light of modern views that nucleons are composite quark structures with a radius of about 1 fm. Even as a phenomenological model this theory is cumbersome, since six parameters are needed in the weak Hamiltonian for ρ, ω exchange.

Recently, there has been a great deal of interest in developing a quark description of nuclear physics at short distance. There is now an extensive body of work on the *N-N* interaction,[2] and a number of studies of quark structure[3] using various reactions. During the past year a hybrid multi-baryon-quark shell model, based on projection operators in coordinate space, has been introduced.[4] This model potentially could provide the basis for a quantitative description of the short-distance (quark) structure of nuclei. However, it has phenomenological aspects which require that a number of processes be studied and compared to experiment in order to complete its development. During the past year, applications have been made to electromagnetic processes in the deuteron,[4] pionic absorption,[5] and the lifetime of Λ's in nuclei.[6] There also has been significant development of the theory.[7]

In the present paper we develop the hybrid model for the treatment of weak interactions in nuclei, and apply this model to the *pp* parity violating asymmetry. In Sec. II the general model for weak processes is given. The theoretical treatment of the *pp* asymmetry is given in Sec. III, and the results for a specific model are discussed in Sec. IV.

II. EFFECTIVE WEAK INTERACTION IN THE HYBRID MODEL

In this section we briefly review the traditional meson exchange model and the hybrid model for weak interactions in nuclei, and discuss their relationship.

A. Traditional model of weak interactions in nuclei

The traditional weak interaction model used in nuclear physics[8] is a boson exchange model. In the version which is generally used the form of the effective weak potential is taken from exchange of π, ρ, and ω mesons as illustrated in Fig. 1(a), with the coupling constants as essentially free parameters for the ρ, ω terms. The potential has the general form of

$$H_{\rm trad} = V_\pi + V_{\rho,\omega} \ , \quad (1a)$$

with V_π arising from pion and $V_{\rho,\omega}$ from ρ and ω exchange, and

$$V_{\pi}^{\rm PV}=V_{\pi}(\Delta I=1,\Delta S=0,S=1)e^{-m_{\pi}r}/r \ ,$$

$$V_{\rho,\omega}^{\rm PV}=\{V_1(\Delta I=0,I=1,\Delta S=1)+V_2(\Delta I=0,\Delta S=1)+V_3(\Delta I=1,\Delta S=0,S=1)\}e^{-m_v r}/r \ , \qquad (1b)$$

where V_π, V_1, V_2, and V_3 are momentum dependent operators. There are also contributions from intermediate N^* excitations, as illustrated in Fig. 1(b), but these terms are usually included as effective weak interactions in the processes of Fig. 1(a). The notation of Eq. (1b), e.g., is that the parity violating potential arising from pion exchange, $V_\pi^{\rm PV}$, changes isospin by one unit, does not change the spin, and operates in the triplet state. Often the ρ and ω masses are taken as $m_\rho=m_\omega=m_v$, as indicated in Eq. (1b).

There are seven coupling constants in this model, one for the pion term and six for the short range ρ,ω terms. The present program of weak interactions in nuclei based on this traditional model involves both the determination of the parameters of the weak Hamiltonian as well as the calculation of the nuclear matrix elements of this Hamiltonian. A systematic program for the determination of the weak Hamiltonian parameters has begun,[9] and considerable effort is underway in the implementation of this systematic program.[10]

In implementing this agenda, one takes advantage of the fact that the coupling constants occur as parameters for terms with different space-spin-isospin structure. The hope is that by studying experimentally and theoretically a variety of nuclear transitions and processes one can make use of the great richness of quantum numbers found in nuclear states to determine the individual coupling constants. However, because of the difficulty of the experiments much of the data is available for transitions that require complicated nuclear structure calculations.

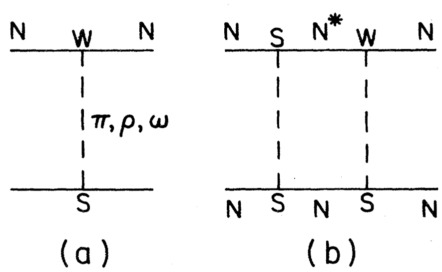

FIG. 1. Boson exchange model of parity violating nucleon-nucleon potential. $S(W)$ means that the boson-nucleon interaction is strong (weak). (a) one boson exchange, (b) two boson exchange.

Recently there has been progress in relating[11,12] this model of weak interaction based on meson exchange to quark quantum chromodynamics models. In Ref. 9 it has been shown explicitly how one can derive ranges of values for the constants needed for the traditional model [Eq. (1)] of Fig. 1(a) from the quark model and the standard model of electroweak interactions. However, only the quark, gauge boson processes involving the weak annihilation of the antiquark in the meson, depicted in Figs. 2(a)—(c), are included. Other important processes, such as that depicted in Fig. 2(d), are not included. Most important from our point of view is that the six-quark processes of Fig. 2(e) are not directly included.

One should also note that the theoretical calculations in the traditional model involve short distances, for which there is a great deal of uncertainty. Moreover, we believe that the physics of this traditional treatment at distances <1 fm is not consistent with present concepts for the structure of hadrons. This takes us to the hybrid model, where we separate the interactions by projections in coordinate space before we begin to analyze the individual problems according to the quantum numbers involved.

B. Hybrid two-baryon-quark shell model

In this subsection we present the hybrid model of Refs. 4 and 7 as applied to weak interactions in nuclei. The method makes use of projection operators in coordinate space to separate the multi-baryon sys-

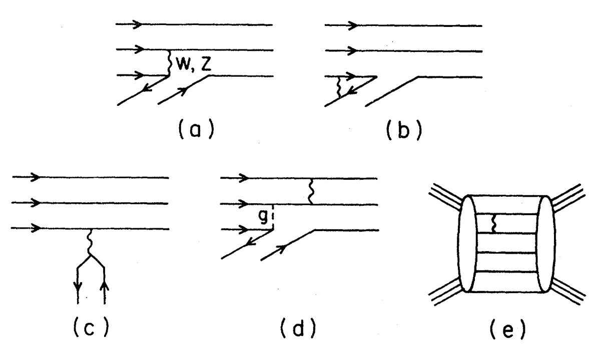

FIG. 2. (a)—(d) Quark models of the weak boson-nucleon interaction. The wiggly line represents a $W(Z)$ exchange, the dashed line a gluon exchange. (e) Six quark weak interaction process.

tem into an interior region where quark coordinates are explicitly used and an exterior region in which the quarks are confined to hadrons. Here we only consider the two-nucleon system, but the generalization to many-baryon systems can easily be done.

1. Effective weak interaction Hamiltonian

Using the projection operators $P_>$ and $P_<$ we project the weak Hamiltonian into the exterior (hadronic) region and interior (quark) region

$$P_> H^{(w)} P_> = H^{(w)}_{\text{hadronic}} , \tag{2a}$$

$$P_< H^{(w)} P_< = H^{w}_{6q} . \tag{2b}$$

The projection operators are defined by the relative separation, r, of the center of the two nucleons, which is a well-defined variable either in terms of a quark model description of the nucleons or a traditional description. The hadronic parts of the weak interaction is to be used for $r > r_0$, and the six-quark part for $r < r_0$, where r_0 is expected to be ≈ 0.8 fm from the earlier work.[4]

The exterior weak interaction must include pion exchange and effects of two pion exchange, illustrated in Figs. 1(a) and (b). There might also be "tails" of heavy meson contributions, but we do not consider this here. Also, we include the two-pion contributions to the weak potentials, which involve N^* intermediate state excitations, as effective boson exchange contributions. Thus our model for the weak interaction is

$$H^w = H^w_\pi, \quad r > r_0$$

$$= H^w_{6q}, \quad r < r_0 , \tag{3}$$

$$H^w_\pi = i(f_\pi/M)(\vec{\tau}_1 \times \vec{\tau}_2)_0(\vec{\sigma}_1 + \vec{\sigma}_2)\cdot[\vec{p}, e^{-m_\pi r}/4\pi r] , \tag{4}$$

and the H^w_{6q} effective Hamiltonian is to be taken from the electroweak theory for quarks. We now discuss this interior Hamiltonian, which is illustrated in Fig. 2(e).

The weak interaction between quarks arising from gauge boson exchange as illustrated in Fig. 3(a) has the form

$$H^{\text{eff}(0)} = -g^2 \int d^4x\, D_F(x^2, M^2) T[J^+_\mu(x) J^\mu(0)] , \tag{5}$$

where $J_\mu(x)$ is the quark current, D_F the gauge boson propagator, and g the electroweak coupling constant. Restricting ourselves to the $\Delta S = 0$ sector, at low and medium energies the effective interaction for u and d quarks, with the neglect of gluonic corrections, is given by the Weinberg-Salam model,[13,14] with a neutral current part,

$$H^{\text{eff}(0)}_{\text{NC}} = \frac{G}{2\sqrt{2}}(1 - \tfrac{8}{3}\sin^2\theta_w)\bar{u}\gamma_\mu u[\bar{u}\gamma^\mu\gamma_5 u - \bar{d}\gamma^\mu\gamma_5 d] + \frac{G}{2\sqrt{2}}(1 - \tfrac{4}{3}\sin^2\theta_w)\bar{d}\gamma_\mu d[\bar{d}\gamma^\mu\gamma_5 d - \bar{u}\gamma^\mu\gamma_5 u] , \tag{6}$$

corresponding to the exchange of Z^0 gauge bosons, and a charge current part

$$H^{\text{eff}(0)}_{\text{CC}} = \frac{G}{\sqrt{2}}\{\bar{d}\gamma_\mu u \bar{u}\gamma^\mu\gamma_5 d + \bar{d}\gamma_\mu\gamma_5 u \bar{u}\gamma^\mu d\} , \tag{7}$$

corresponding to the exchange of W mesons. This effective Hamiltonian is derived by using the W, Z^0 exchange mechanism between quarks, illustrated in Fig. 3(a), and taking the limit as $M_w{}^2 \approx M_z{}^2 \gg$ momentum transfer involved in the calculation. The two parameters, G $(1.03\times 10^{-5}/M^2)$ and $\sin^2\theta_w$ (0.22) are well known.

However, this Hamiltonian, which is taken directly from the Weinberg-Salam model for leptons,[14] is

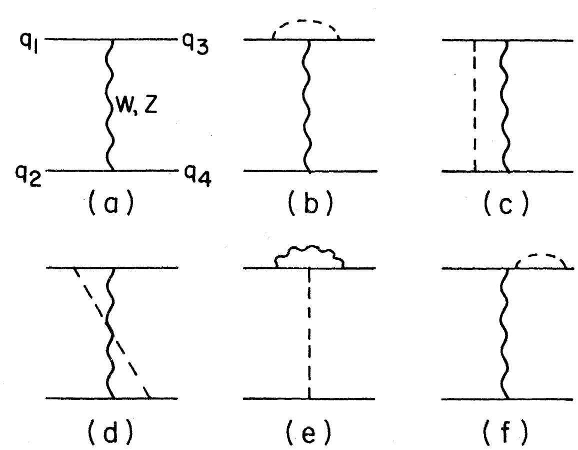

FIG. 3. Weak quark-quark interaction. Wiggly line W,Z; dashed line, gluon.

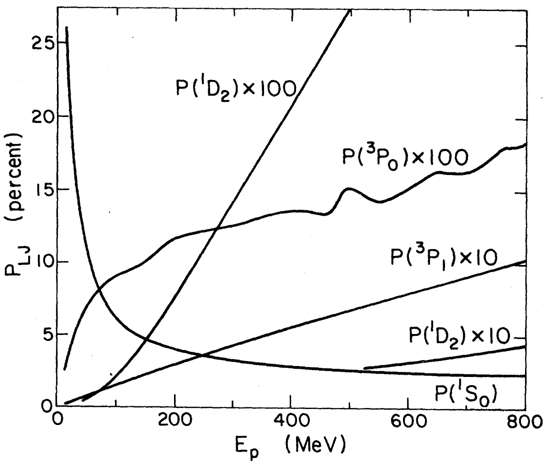

FIG. 4. Quark probabilities, P_{LJ} as a function of proton laboratory kinetic energy, E_p. Here $r_0=0.8$ fm.

not sufficiently accurate for application to hadronic and nuclear weak interactions. This is quite evident from the failure of the model to satisfy the $\Delta I=\frac{1}{2}$ rule in the $\Delta S=1$ processes. The strong interaction (gluon) corrections which are necessary to account for this discrepancy can be carried out using renormalization group methods.[15] Recently there has been a good deal of effort to develop these methods for the $\Delta S=1$ processes,[16] and the resulting weak interaction when used in the hybrid model seems to give an accurate description of Λ decay in nuclear matter.[6] The $\Delta S=0$ processes are more difficult, but several theoretical efforts[9,17] have now provided effective weak Hamiltonians. We briefly review this work to show how it applies to the hybrid model.

Returning to Eq. (5), an operator product expansion is carried out for the time ordered product of currents

$$T[J_\mu^+(x)J^\mu(0)]=\sum_j F_j(X^2)\theta^j \ ,$$

where $F_j(x^2)$ are c-number functions and θ^j are four quark operators. Working at $q^2\approx M_w{}^2$, where QCD perturbation theory is accurate, the effective weak Hamiltonian with all one-gluon loops, some of which are illustrated in Figs. 4(b)—(f), is calculated. The resulting Hamiltonian has the form

$$H^{\rm eff}(M_w)=\frac{G_F}{\sqrt{2}}\sum_j C_j(g(M_w{}^2))\theta^j \ . \quad (8)$$

The coefficients C_j satisfy renormalization group (Callan-Symanzik) equations. These can be readily integrated to obtain the final result:

$$H^{\rm eff}=\frac{G_F}{2}\sum_{j=1}^{16}\left[\frac{\alpha(M_w)}{\alpha(\mu)}\right]^{dj} C_j\bar{\theta}_j \ , \quad (9)$$

where $\alpha(\mu)$ is the QCD running coupling constant, μ is the subtraction point (taken as 1 GeV, characteristic of the masses with which we deal), the d_j are anomalous dimensions, and $\bar{\theta}_j$ are the fully-dressed quark operators. If the quark model which one has available is a correct representation of the system, then the $\bar{\theta}_j$ are just four-quark operators in that representation. Thus to the extent that our hybrid model is a good representation of the two-baryon interior region, Eq. (9) provides an effective weak Hamiltonian.

We use the results of Preston and Goldman[17] for the C_j and d_j and find that the effective weak Hamiltonian which is given in the lowest order given by Eqs. (6) and (7), corrected for gluon processes, becomes

$$H_{qq}^{\rm eff(NC)}=\frac{G}{\sqrt{2}}\{\bar{u}\gamma^\mu u[0.402\bar{u}\gamma_\mu\gamma_5 u-0.185\bar{d}\gamma_\mu\gamma_5 d]+0.333\bar{d}\gamma^\mu d\bar{d}\gamma_\mu\gamma_5 d-0.129\bar{u}\gamma^\mu u\bar{d}\gamma_\mu d\} \ ,$$

$$H_{qq}^{\rm eff(CC)}=\frac{G}{\sqrt{2}}1.562(\bar{u}\gamma^\mu d\bar{d}\gamma_\mu\gamma_5 u+\bar{u}\gamma^\mu\gamma_5 d\bar{d}\gamma_\mu u) \ . \quad (10)$$

In summary, Eqs. (3), (4), and (10) complete our current version for the effective weak nuclear interaction in the hybrid model.

2. *Wave functions*

Using the same projection operators, $P_<$ and $P_>$, as in the previous subsection, the two nucleon wave function in channel α with energy E is given as

$$\Psi_{NN}^\alpha(E)=\Psi_{NN}^{\alpha(E)}(r), \quad r>r_0$$
$$=\Psi_{6q}^{\alpha(E)}(r_i), \quad r<r_0 \ , \quad (11)$$

where

$$\psi_{NN}^{\alpha(E)}(\vec{r})=\phi_{N_1}\phi_{N_2}\Phi_{NN}^{\alpha(E)}(\vec{r}) \ , \quad (12)$$

with ϕ_N being the internal $3q$ internal nucleon wave function and $\Phi_{NN}^{\alpha(E)}(\vec{r})$ the wave function describing the relative motion of the two nucleons. Since the nucleonic polarization is not explicit, $\psi_{NN}^{\alpha(E)}(r)$ is in essence a conventional two-nucleon wave function. The six-quark part of the wave function

$$\Psi_{6q}^{\alpha(E)}\propto\mathscr{A}^\alpha(E)\Psi_{6q}^\alpha \ , \quad (13)$$

with

$$\Psi_{6q}^{\alpha}=\sum C_n^{\alpha}\Psi_{6q}^{n} ,$$

where $\Psi_{6q}^{n}(r_i)$ are six-quark configurations composed of products of single-quark bag wave functions, $\phi_\nu(r_i)$, C_n^{α} are (energy-independent) relative spectroscopic amplitudes for the nth configuration in the channel α, and $\mathscr{A}^{\alpha}(E)$ are the spectroscopic amplitudes which give the probability and phase of the six-quark part of the wave function in channel α. The constant of proportionality is defined in Eq. (21). The single-quark wave functions $\phi_\nu(r_i)$ are taken as MIT bag wave functions,[18]

$$\phi_\nu(r_i)=\eta\begin{bmatrix}\pm i j_l(k_\nu r_i)[Y_l\chi]_{j\mu}\\ j_{l'}(k_\nu r_i)[Y_{l'}\chi]_{j\mu}\end{bmatrix}\equiv\begin{bmatrix}i\mathscr{F}_{t_\nu}[Y_l\chi]_{j\mu}\\ \mathscr{F}_{b_\nu}[Y_{l'}\chi]_{j\mu}\end{bmatrix} , \tag{14}$$

where k are determined by the boundary condition

$$[i\vec{\gamma}\cdot\hat{r}_i\phi_\nu-\phi_\nu]_{\text{surface}}=0$$

giving color confinement. Note that the upper and lower components are wave functions for a particle in a uniform well. The functions ϕ_ν are normalized to unity over the volume of the bag. The bag radius is also taken as r_0.

The bag wave functions of Eq. (14) are simply a convenient representation. The crucial parameters are the C_n and $\mathscr{A}^{\alpha}(E)$. For the low energy applications, such as in Ref. 4, there are only a few coefficients, which can be determined from experiment, allowing one to calculate various low energy results and make predictions. However, for scattering problems over a wide range of energies one must make use of theoretical models, such as that of Ref. 7. This is discussed further in the next section.

III. *p-p* ASYMMETRY

The parity-violating asymmetry in *p-p* scattering associated with the dependence of the cross section on the helicity of the proton beam is an excellent process for the application of the hybrid model defined in the previous section. The asymmetry is defined as

$$A_E(\theta)=\frac{\sigma_E(\theta)^+-\sigma_E(\theta)^-}{\sigma_E(\theta)^++\sigma_E(\theta)^-} , \tag{15}$$

where $\sigma_E(\theta)^{+(-)}$ are the differential cross sections at energy E with the beam proton helicity positive (negative). The total *p-p* asymmetry defined as in Eq. (15) with total cross sections has been measured[1] at energies of 15, 45, and 6000 MeV. Theoretical calculations using the traditional model at the lower energies[19] are not in disagreement[1] with these results; however, the theoretical results[20] at 6 GeV are about an order of magnitude too small within the range of consistent parameters in the traditional model.

The hybrid model of the weak interaction has a simple form for the *p-p* system. This is because the long-range pionic weak contribution vanishes in this application, as can be seen from Eqs. (1b) or (4). Therefore the weak interaction Hamiltonian is

$$H^{w(\mathrm{PV})}(pp)=H_{6q}^{w} ,$$

entirely within r_0, the interior quark region. Thus in comparison to the traditional model[19] in which there appeared unknown ρ and ω coupling constants, the effective Hamiltonian is given in a parameter-free form by Eq. (10). Therefore, we can concentrate on the development of the model for the wave functions. In doing so we make use of the results of Ref. 7.

The asymmetry defined in Eq. (15) is given by[19]

$$A_E(\theta)=\frac{\mathrm{Re}[\langle 00|F^s|00\rangle^*\langle 00|F^w|10\rangle+\sum_{M_{s'}}\langle 1M_{s'}|F^s|10\rangle^*\langle 1M_{s'}|f^w|00\rangle]}{\frac{1}{4}\left[|\langle 00|F^s|00\rangle|^2+\sum_{M_{s'}M_s}|\langle 1M_{s'}|F^s|1M_s\rangle|^2\right]} , \tag{16}$$

where $\langle SM_{s'}|F^s|SM_s\rangle$ is the strong scattering amplitude (parity-conserving) and $\langle S'M_{s'}|f^w|SM_s\rangle$ is the weak (parity-violating) amplitude, with S being the channel spin. Since we wish to carry out the most accurate calculations possible within our model, we avoid the calculation of the strong amplitude. As can be seen from Eq. (1b), $A_E(\theta)$ depends upon an interference between strong and weak amplitudes, and there is a great deal of structure in the E and θ dependence of F^s. A strong qq interaction which is known to comparable accuracy as Eq. (10) for the weak qq interaction is not available. Therefore, we restrict ourselves in the present work to the energy region $E\leq 800$ MeV, and use a phase shift analysis[21] to obtain F^s.

For the determination of the parity violating weak amplitude we must introduce scattering states for the *pp* system, $\psi_{\mathbf{k}}^{(+,-)}$, and calculate the matrix element of the weak Hamiltonian, with

$$f_E^w(\theta)=\langle \psi_{\vec{k}'}^{(-)} | H^w | \psi_{\vec{k}}^{(+)} \rangle \ . \tag{17}$$

We use a δ-function normalization for the plane waves

$$\phi_{\vec{k}}^{\text{PW}}(r)=\frac{1}{(2\pi)^{3/2}}e^{i\vec{k}\cdot\vec{r}} \ , \tag{18}$$

and the standard representation of the scattering states[22]

$$\psi_{\vec{k},s,\nu}^{\pm}(\vec{r})=\sum_{l,l's'JmM} i^{l'}[Y^{l'}(\hat{r})\chi^{s'}]_M^J \mathscr{C}_{m\nu M}^{lsJ} Y_m^{l*}(\hat{k})\phi_{l's'ls;J,k(r)}^{+}\phi_{N_1}\phi_{N_2} \ , \tag{19}$$

where $\mathscr{C}_{m\nu M}^{lsJ}$ is a Clebsch-Gordan coefficient, and $[Y^{l'}(\hat{r})\chi^{s'}]_M^J$ is a vector spherical harmonic with the spin state $\chi^{s'}$ vector coupled to the spherical harmonic $Y^{l'}$.

In order for us to define the scattering states for the inside as well as the outside region, we must normalize the scattering states consistently. Moreover, we would like to introduce the concept of probability, which is generally not required for scattering states. To do this, define the probabilities of the inside and outside parts of the wave function in a given channel $\alpha(E)$ [see Eq. (11)] with respect to the interaction region given by the long-range pion exchange force. Note that in the absence of any long-range interaction, from Eq. (18) the baryon probability in the exterior interaction region $R_\pi \geq r \geq 0$, with $R_\pi = m_\pi^{-1}$, is given by

$$P^{\text{ext}}=V_{R_\pi}/(2\pi)^3, \quad V_{R_\pi}\equiv\frac{4\pi}{3}R_\pi^{\ 3} \ . \tag{20}$$

Therefore we define our inside probability by using scattering wave functions of the form of Eq. (19) but with the normalization of the interior wave function for channel α

$$\int^{r_0}\prod_{i=1}^{6} d^3r_i \, |\Psi_{6q}^{\alpha(E)}(r_i)|^2 = V_{R_\pi}|\mathscr{A}^{\alpha}(E)|^2\sum_n C_n^{\ 2}\equiv V_{R_\pi}P_\alpha \ , \tag{21}$$

where we use the definitions in Eqs. (11) and (13). Thus P_α can be interpreted as a $6q$ probability in the interaction region. The constant of proportionality of Eq. (13) is just V_{R_π}. Then from Eqs. (17) and (18), and this normalization we obtain for the weak matrix elements needed in Eqs. (16) and (17)

$$\langle 00|f^w|10\rangle = V_{R_\pi}\sum_{LL'} i^{L-L'}\mathscr{C}_{000}^{L1L'}\left[\frac{2L+1}{4\pi}\right]^{1/2} Y_0^{L'}(\hat{k}') \times \mathscr{A}^{L'0*}(E)\mathscr{A}^{L1}(E)\langle L'ML'S=0|H_{6q}^w|LMLS=1\ M_s=0\rangle \tag{22}$$

and

$$\langle 1M_{s'}|f^w|00\rangle = V_{R_\pi}\sum_{LL'}(-1)^{1-M_{s'}} i^{L-L'}\mathscr{C}_{M_{s'},-M_{s'},0}^{L'1L}\left[\frac{2L+1}{4\pi}\right]^{1/2} Y_{M_{s'}}(\hat{k}') \times \mathscr{A}^{L'1*}(E)\mathscr{A}^{L0}(E)\langle L'M_{L'}S=1M_{s'}|H^w|LM_LS=0\rangle \ . \tag{23}$$

The six-quark probability in each channel α, P^α, one of the important concepts of the hybrid model, is discussed in detail in Ref. 7. Using the interaction volume as the volume for normalization, one can compare the internal six-quark probability to the exterior two-baryon probability. It is this relative probability which gives a measure of the importance of short-range six-quark phenomena for each channel. These probabilities, of course, are model dependent. The $\mathscr{A}^{LJ}(E)$, taken from Ref. 7, are obtained by probability current conservation and continuity. The

$$P_{LJ}(E)=|\mathscr{A}_{LJ}(E)|^2$$

are taken from Ref. 7, and shown in Fig. 4. The Reid soft core potential[23] is used in these computations.[1] Calculations with the Paris potential[24] give about 50% larger values of the $^1S_0$ quark probability. With this potential we would obtain larger asymmetries (Figs. 5 and 6).

Finally, the matrix elements in Eqs. (22) and (23) are calculated in the quark sector using Eqs. (11), (13), and (14), with H^w given by Eq. (10). These are readily reduced to two-quark matrix elements. The basic two-quark matrix element needed is

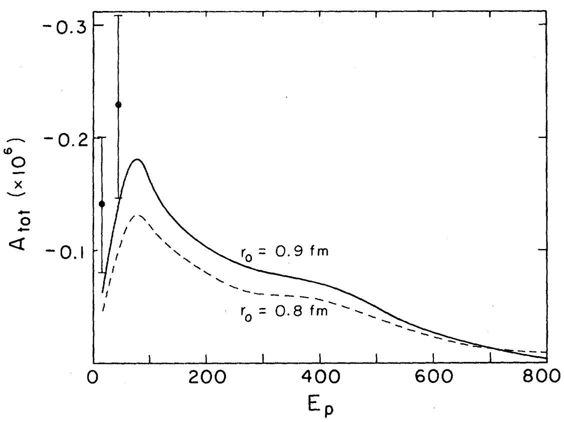

FIG. 5. $A_{tot}\,(E_p)$ as a function of r_0.

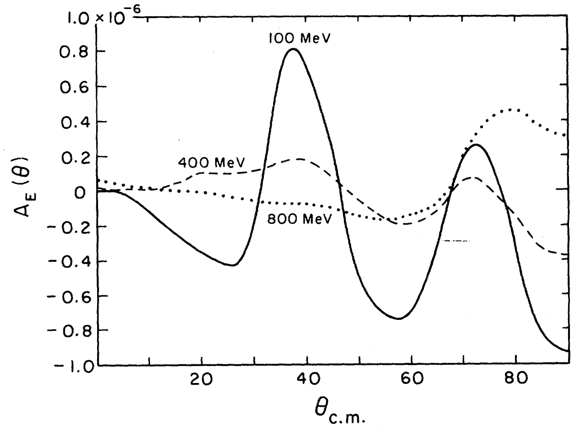

FIG. 6. Angular distributions of $A_E(\theta)$ for E (proton laboratory kinetic energy) of 100, 400, and 800 MeV. $r_0=0.9$ fm.

$$\begin{aligned}\mathscr{M}^w &= \langle q_1',q_2';S'L'J \,|\, \sum_\mu (\gamma_\mu\gamma_5)^{(1)}\gamma^{\mu(2)} + \gamma_\mu^{(1)}(\gamma^\mu\gamma^5)^{(2)}] \,|\, q_1 q_2;SLJ\rangle \\ &= \frac{-i}{\pi}[(2l_1+1)(2l_2+1)(2l_{1'}+1)(2l_{2'}+1)/(2L+1)(2L'+1)]^{1/2}\mathscr{C}^{l_1 l_2 L}_{000}\mathscr{C}^{l_{1'}l_{2'}L}_{000} \\ &\quad\times\{\delta_{S'0}\delta_{S1}\delta_{L'J}\mathscr{C}^{L'1L}_{000}\int dr\, r^2(\mathscr{I}_{t_{1'}}\mathscr{I}_{t_{2'}}+\mathscr{I}_{b_{1'}}\mathscr{I}_{b_{2'}})(\mathscr{I}_{b_1}\mathscr{I}_{t_2}-\mathscr{I}_{t_1}\mathscr{I}_{b_2}) \\ &\quad+\delta_{S'1}\delta_{S0}\delta_{LJ}\mathscr{C}^{L1L1}_{000}\int dr\, r^2(\mathscr{I}_{t_{1'}}\mathscr{I}_{b_{2'}}-\mathscr{I}_{b_{1'}}\mathscr{I}_{t_{2'}})(\mathscr{I}_{t_1}\mathscr{I}_{t_2}+\mathscr{I}_{b_1}\mathscr{I}_{b_2})\ . \end{aligned} \tag{24}$$

In Eq. (24) the l_i are the orbital angular momenta of the upper components of the single quark wave functions, and $\mathscr{I}_b$ and $\mathscr{I}_t$ are radial forms defined in Eq. (14)

This completes the formalism needed for the computation of the asymmetry for $E \le 800$ MeV.

IV. RESULTS AND DISCUSSION

The theoretical calculation of the pp parity violating asymmetry follows from Eq. (16), with the strong amplitude F^s obtained from pp phase shifts and the weak amplitude f^w from Eqs. (22)—(24). All of the quantities needed for this latter calculation have been determined except the configuration choices, C_n, of Eq. (13) and certain phases. These choices complete the model.

In the present work we choose the simplest configurations consistent with current information about bag models. For the $^1S_0$ channel, which is the most important channel in the $E \le 800$ MeV region under consideration, we take an admixture of six quarks in the $s_{1/2}$ single-quark state and p^2s^4 configurations, i.e.,

$$|\psi_{6q}^{^1S_0}(E)\rangle = \sqrt{V_{R_\pi}}\,\mathscr{A}^{^1S_0}(E)[C_0\,|s_{1/2}^6\rangle + C_1\,|p_{3/2}{}^2 s_{1/2}{}^4\rangle + C_2\,|p_{1/2}{}^2 s_{1/2}{}^4\rangle]\ , \tag{25}$$

where $\sum_i C_i^2=1$. We have taken $C_2{}^2/C_1{}^2={}^3P_0$ probability divided by $^3P_1$ probability. The results turned out to be insensitive to the choice of C_2 vs C_1. For the other channels we simply take $(l_j)^1 s_{1/2}{}^5$ configurations. Therefore, the only parameter of the theory is the energy independent value of $C_1{}^2+C_2{}^2$ which we take as 0.3 from the studies of p^2s^4 vs s^6 configurations in quark models for the NN force.[2] The parameters C_0, C_1, and C_2 are chosen to be in phase. With these choices we achieve a wave function with approximate projection of the center of mass.[5] There are two other wave functions which can be constructed from these configurations, one an excited state and the second a state involving mainly spurious center of mass motion.[5] With the values of $\mathscr{A}^{LJ}(E)$ given by the theory of Ref. 7, the only parameter available for the fit of the theory to the data is the choice of the matching radius, r_0. Note that the phase of $\mathscr{A}^{\alpha}(E)$ is just the phase shift $\delta^{\alpha}(E)$ for the channel α and energy E.

The computations involve the evaluation of matrix elements in six-quark states. Although this requires a considerably larger effort than traditional weak interaction N-N calculations, the computations are essentially the same as in the standard nuclear shell model. The use of Dirac single-particle states is only a complication, not an essential difficulty. Having prepared a computer code, subsequent calculations will be relatively routine. Note also that the six-quark part will be no more complicated for complex nuclei than for the two-baryon problem being treated in the present paper. Furthermore, these complications in comparison with the traditional weak model reflect a new level of information being sought—the quark structure of nuclear systems.

The results for the total asymmetry, A_{tot} defined as the angular integral of $A_E(\theta)$, are shown in Fig. 5. The theory is in reasonable agreement in comparison with experiment. There is considerable sensitivity to r_0. For $r_0=0.9$ fm, the best value for the fit to the electrodisintegration of the deuteron in Ref. 4, the theoretical calculation is just within the experimental errors. Note that there is little freedom for us, except to use a different extrapolating NN potential (the Reid soft core is used here), and extend the choice of configurations. The overall results are similar to the old ρ,ω exchange calculations,[19] but our A_{tot} shows more structure. Since the old calculations had great freedom in the choice of parameters, a comparison of the magnitude cannot help us judge the possible differences in these two approaches. Note that at 800 MeV the A_{tot} is changing sign. It would be most interesting to see if the theory could account for the large positive asymmetry ($A_{\text{tot}}\cong 2.65\times 10^{-6}$) at 6 GeV,[25] but we must develop the strong interaction aspects of the model in order to attempt this.

Typical angular distributions are shown in Fig. 6 for $r_0=0.9$ fm. At low energy there is a striking structure in $A(\theta)$ which is not seen in the traditional model.[19] It would be most important to explore this experimentally, as well as theoretically.

We conclude that the hybrid model can give a satisfactory description of the pp parity violating asymmetry, with the weak interaction being given entirely within the quark sector. Systematic experimental studies could be most valuable for the study of the quark structure of nuclei.

The authors would like to thank Dr. John Collins, Dr. Steve Ellis, Dr. Terry Goldman, Dr. Ernest Henley, and Dr. Bruce McKellar for helpful discussions. One of us (L.S.K.) would like to thank the University of Washington Institute for Nuclear Theory for its hospitality during his 1981–82 leave from Carnegie-Mellon University. This work was supported in part by the U.S. Department of Energy and in part by NSF Grant PHY 81-08380.

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