PHYSICAL REVIEW C VOLUME 9, NUMBER 3 MARCH 1974

Parity-violating asymmetry in nucleon-nucleon scattering*

Virginia R. Brown
Lawrence Livermore Laboratory, Livermore, California 94550

Ernest M. Henley and Franz R. Krejs
Physics Department, University of Washington, Seattle, Washington 98195
(Received 20 September 1973)

The weak parity-violating component of the nucleon-nucleon interaction can be studied directly and in detail through the scattering of longitudinally polarized nucleons on nucleons. We use distorted-wave theory to predict the dependence of the differential and total cross sections on the helicity of incident longitudinally polarized protons for elastic proton-nucleon scattering at energies up to ~300 MeV. The weak force used is parametrized by ρ, ω, π, and 2π exchanges, the strong one by a Hamada-Johnston or Bryan-Gersten potential.

[NUCLEAR REACTIONS Weak nucleon-nucleon force. Calculate parity-violating asymmetry for longitudinally polarized nucleon-nucleon scattering.]

I. INTRODUCTION

The theory of weak interactions in its various forms predicts a weak parity-violating (PV) contribution to the nucleon-nucleon interaction.[1,2] In a current-current formulation, this force arises from a weak hadronic current interacting with itself. A study of parity violation in the nucleon-nucleon system thus offers a unique opportunity for studying such a self-interaction. It has been pointed out, for instance,[3] that there is no compelling reason to expect the coupling constant for diagonal processes, i.e., those described by a current interacting with itself, to be identical to that for off-diagonal processes. Broken gauge theories of weak interactions[4] allow such differences. Studies of the weak nucleon-nucleon interaction allow one to investigate these questions, possibly to obtain the sign of the weak coupling constant,[5] and quite generally to determine the nature of the weak interaction of hadrons. Without this knowledge, it is difficult to attempt to understand nuclear PV experiments.

In the conventional Cabibbo theory of the weak interactions, the currents consist of leptonic and hadronic components. The low-energy properties of the leptonic current are well understood; the structure of the hadronic weak currents is obtained from studies of semileptonic processes. The Cabibbo theory has been eminently successful in explaining the experimental results of these leptonic and semileptonic processes. However, the theory has had to be extended or modified to understand nonleptonic decays and reactions.[6] For instance, the empirical $\Delta I = \frac{1}{2}$ rule, which is found to hold to high accuracy in the weak nonleptonic decays of strange baryons and mesons, does not follow from the simple Cabibbo theory. It is possible that this rule will be obtained directly from a theory of weak interactions.[4] In the meantime, the Cabibbo theory is assumed to be generalized by either (a) the addition of (neutral) currents or (b) dynamical octet enhancement.[1,2] The first one of these modifications adds sufficient currents that the $\Delta I = \frac{1}{2}$ rule is obtained, whereas the second modification assumes that the hadronic forces enhance the octet representation of the current-current interaction. Because the $\Delta I = \frac{3}{2}$ is contained in the $\{27\}$ representation, octet enhancement suffices to give the desired result; alternatively, one may assume suppression of the $\{27\}$. There are differences between (a) and (b) which can be tested experimentally, particularly in the weak force between nucleons. Thus, the long-range weak one-pion exchange force is more than an order of magnitude larger for modification (a) than for (b).[1,2] Regardless of the origin of the $\Delta I = \frac{1}{2}$ rule for strangeness-changing decays, related modifications of the Cabibbo theory are expected for the weak nucleon-nucleon force, e.g. suppression of the $\Delta I = 2$ PV force; clearly direct investigations of this force are crucial in understanding the weak interactions of hadrons.

The existence of a parity-violating component of the nuclear force has been ascertained from numerous experiments performed in nuclei. Examples are the α decay[7] of a 2^- state in ^{16}O and observations of the nonvanishing expectation values of pseudoscalar quantities in electromagnetic nuclear transitions,[2,8] particularly the circular polarization of photons. With the exception[9] of the α decay from ^{16}O, published theoretical calculations[2,8] predict effects which are generally an order of magnitude or more too small and of the wrong sign. Although corrections to the nuclear theory-dependent part of the calculations have been suggested[10] and may rectify the discrepancy, it will be difficult, if not impossible, to isolate the nuclear-structure factors from the weak-nuclear-force factors. For this purpose, it is necessary to go directly to the two-nucleon system.[11] Indeed, the circular polarization of photons emitted in the capture of thermal neutrons by protons has been measured.[12] The results and theory disagree,[13] but it is difficult to discern exchange current and parity-violating electromagnetic corrections from the normal effects in this capture process.[14]

The motivation for the study of PV effects in N-N (nucleon-nucleon) scattering should be clear from the above discussion. Measurements of the total cross section of longitudinally polarized projectiles from an unpolarized target (or vice versa) have been suggested[15] and are being attempted.[16] In this experiment, the spin direction is reversed periodically, and a nonvanishing expectation value of $\langle\vec{\sigma}\cdot\hat{p}\rangle$ is sought. Here $\langle\vec{\sigma}\rangle$ represents the initial polarization and $\hat{p}$ a unit vector along the incident momentum. Such an experiment has the advantage of no background from the much larger parity-allowed terms such as $\langle\vec{\sigma}\cdot\vec{p}\times\vec{p}'\rangle$ where $\vec{p}'$ is the scattered momentum. This term is present in differential cross section searches of parity violation where one may seek nonvanishing expectations of $\langle\vec{\sigma}\cdot\hat{p}\rangle$ or $\langle\vec{\sigma}\cdot\hat{p}'\rangle$. Although the magnitude of the PV effect is small (of order 1 ppm), the resonance technique pioneered by Lobashov *et al.*[17] can be exploited. A detector tuned to the frequency of the polarization flipping rate is used to remove parity-conserving effects, such as those due to bremsstrahlung emitted during the scattering. The experimental geometry should, of course, be chosen to mitigate such possible spurious effects. The removal of systematic errors is crucial, since there is no comparison experiment easily available for which a null result would be predicted.

II. WEAK PV POTENTIAL

Since there still is no definitive theory for the weak interactions of hadrons, we consider the simplest and most basic framework, namely, the Cabibbo theory. However, we introduce those modifications of the theory required by experimental studies of the nonleptonic decays of strange particles, that is octet enhancement or $\{27\}$ suppression.

In the Cabibbo theory the Hamiltonian density which gives rise to the weak nucleon-nucleon interaction is

$$H_{\rm wk}=2^{-1/2}\,G\,J_\mu J^{\mu\dagger}\,, \tag{1}$$

with

$$J_\mu=J_\mu^{(1-i2)}\cos\theta+J_\mu'^{(4-i5)}\sin\theta\,, \tag{2}$$

where θ is the Cabibbo angle, $\tan\theta\approx 0.22$. The current J_μ conserves strangeness, whereas J_μ' changes it by -1. Both currents consist of vector and axial-vector components in terms of which the parity-violating strangeness-conserving part of $H_{\rm wk}$ is

$$H_{\rm wk}^{\rm PV}(\Delta S=0)=-2^{-1/2}G[(V_\mu A^{\mu\dagger}+A_\mu V^{\mu\dagger})\cos^2\theta +(V_\mu' A'^{\mu\dagger}+A_\mu' V'^{\mu\dagger})\sin^2\theta]\,. \tag{3}$$

Because J_μ is an isovector, the first term in Eq. (3) transforms like a spurion of isospin $\Delta I=0$ and 2, whereas the second term which arises from an isospinor current behaves like a spurion of isospin $\Delta I=1$. As discussed in the Introduction, the $\Delta I=\frac{1}{2}$ transition matrices are found to dominate those with $\Delta I=\frac{3}{2}$ ($\Delta I=\frac{1}{2}$ rule) for the strangeness-changing decays of baryons. This suggests that the octet representation of the product of two weak currents is enhanced or that the $\{27\}$ representation which supplies the $\Delta I=\frac{3}{2}$ transition matrices is suppressed. We assume that this empirical finding also applies to the strangeness-conserving currents, so that the $\Delta I=2$ part of the first term in Eq. (3), which belongs to the $\{27\}$ representation, is strongly suppressed and thus only retains the isoscalar part of that term. As a consequence, PV forces occur not only between neutrons and protons (n-p) but between protons (p-p) and between neutrons (n-n). This contrasts with older calculations of the PV force,[1,2] where only the exchange of charged mesons is included. In that case one obtains a pure exchange nucleon-nucleon interaction proportional to

$$\tau_1^{(+)}\tau_2^{(-)}\pm\tau_2^{(+)}\tau_1^{(-)}\,.$$

On the other hand, a $\Delta I=0$ weak interaction gives a force which does not depend on isospin or is proportional to $\vec{\tau}_1\cdot\vec{\tau}_2$.

It is assumed that at energies $\lesssim 300$ MeV both the strong and the weak PV nucleon-nucleon interaction can be described by a potential. Because of the strong hadronic repulsion at small distances, a direct contact term between the nucleons is ineffective and the weak force is mediated by the meson cloud surrounding the nucleons. If we

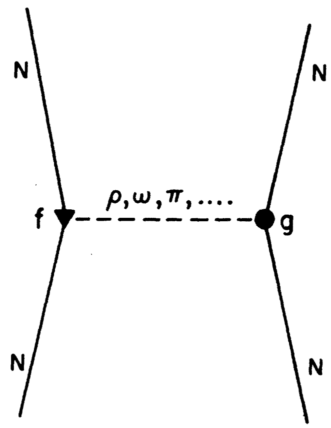

FIG. 1. One-boson-exchange contributions to the weak PV potential. The weak PV vertex is shown as a triangle, the strong one as a circle.

only include single-boson exchange contributions, then the weak PV potential arises from diagrams such as that shown in Fig. 1.

The potential obtained from $J_\mu J^{\mu\dagger}$ has been derived by Michel in a factorization approximation[18] and has also been obtained by other means.[1,2,8] In the factorization approximation, the weak ρ-nucleon vertex is given by

$$2^{-1/2}G\{\langle N|J^\mu|N\rangle\langle\rho|J_\mu^\dagger|0\rangle+\text{h.c.}\}\,. \tag{4}$$

In order to estimate dynamical effects, such as octet enhancement, a reasonable theory of baryonic structure with SU(3) symmetry is required. A quark model could be used, but divergences occur which make the result subject to serious questions. There are models which allow one to carry out suitable renormalizations due to the presence of hadronic forces. One such model is the σ model, which has been generalized to SU(3).[19] However, even without SU(3) symmetry, the calculations required are complex and give results which suggest that octet enhancement may not be large.[20] For this reason, we assume that the $\{27\}$, or $\Delta I=2$ part of Eq. (4), is suppressed rather than that the $\{8\}$ or $\Delta I=0$ part is enhanced. We thus take the ρ-N coupling to be purely isoscalar and write

$$\langle N\rho|H_{\rm PV}|N\rangle=f_\rho\,\bar u_N\gamma^\mu\gamma^5\vec\tau u_N\cdot\vec\rho_\mu\,, \tag{5}$$

with $\vec\rho_\mu$ an isovector spin-1 operator. Although f_ρ is a function of the momentum transfer, this dependence is unimportant at the momentum transfers considered by us, and we identify f_ρ with that obtained in the factorization approximation for the charged ρ-N vertex, Eq. (4),

$$f_\rho\equiv2^{-1}g_{\rho N}{}^{-1}Gm_\rho{}^2\cos^2\theta g_A\approx1.4\times10^{-6}\,, \tag{6}$$

where G is the weak-coupling constant, m_ρ is the mass of the ρ meson, g_A is the axial-vector renormalization constant $g_A\approx1.24$[21] and $g_{\rho N}$ is the strong ρ-N coupling constant, $g_{\rho N}{}^2/4\pi\approx0.62$.

In addition to a weak PV ρ-N coupling, other vector mesons may contribute. With the assumption of isoscalar dominance, ω emission becomes likely. Indeed, a large matrix element for weak $N\to N\omega$ is found with SU(6)W or with a quark model.[22] We write the vertex as

$$\langle N\omega|H_{\rm PV}|N\rangle=f_\omega\bar u\gamma^\mu\gamma^5u_N\omega_\mu\,. \tag{7}$$

With SU(6)W and strong octet dominance, McKellar and Pick find $f_\omega\approx-5f_\rho$, but with other assumptions they obtain $f_\omega\approx\frac14 f_\rho$ etc. For simplicity, we assume $f_\omega=f_\rho$, but also investigate $f_\omega=0$ and $f_\omega=-f_\rho$, since the relative sign of f_ω and f_ρ is unknown.

The potentials which arise from ρ and ω exchanges are readily found from Eqs. (4) to (7) to be

$$\begin{pmatrix}V_{\rm PV}^\rho\\ V_{\rm PV}^\omega\end{pmatrix}=-\frac1M\begin{pmatrix}f_\rho g_{\rho N}\vec\tau_1\cdot\vec\tau_2\\ f_\omega g_{\omega N}\end{pmatrix}\left(\begin{pmatrix}1+\mu_V\\ 1\end{pmatrix}(i\vec\sigma_1\times\vec\sigma_2)\cdot[\vec{\rm p},v(r)]+(\vec\sigma_1-\vec\sigma_2)\cdot\{\vec{\rm p},v(r)\}_+\right)\,, \tag{8}$$

where M is the nucleon mass, $\mu_V=3.70$ is the isovector anomalous magnetic moment, $\vec r=\vec r_1-\vec r_2$, $\vec{\rm p}=\frac12(\vec{\rm p}_1-\vec{\rm p}_2)$ are the relative position and momentum of the nucleons, and $v(r)=\exp(-mr)/4\pi r$ with $m\approx775$ MeV, the average mass of the ρ and ω. In Eq. (8) there is no anomalous magnetic moment term for $V_{\rm PV}^\omega$, similar to μ_V, because of the absence of a magnetic coupling $\propto\bar u(p')i\sigma_{\mu\nu}q^\nu u(p)\omega^\mu$ of the ω to the nucleon.[23] In most of the following results we take $g_{\omega N}=2^{1/2}g_{\rho N}$, a choice made by some nuclear-force theories[23]; however, we also investigate the effect of other ratios $g_{\omega N}/g_{\rho N}$ because the above estimation is a low one.

In addition to the $\cos^2\theta$ or $J_\mu J^{\mu\dagger}$ contribution to the N-N potential, which we represent by ρ and ω exchanges, there is the $\sin^2\theta$ or $J'_\mu J'^{\mu\dagger}$ contribution. The longest-range part of this interaction is the one-pion exchange potential. Here CP conservation forbids π^0 exchange,[1,2] and one finds

$$V_{\rm PV}^{1\pi}=\frac{gfi}{2^{3/2}M}\,(\vec\sigma_1+\vec\sigma_2)\cdot[\vec{\rm p},v_\pi(r)]\,(\vec\tau_1\times\vec\tau_2)^{(3)}\,, \tag{9}$$

where g is the strong π-nucleon coupling constant

$g^2/4\pi \simeq 14.4$, f is the weak π-nucleon coupling constant $f \simeq 4.3\times 10^{-8}$, and $v_\pi(r) = \exp(-m_\pi r)/4\pi r$ with $m_\pi \simeq 136.6$ MeV the charged-pion mass.

As for the strong interactions, in addition to one-boson exchange potentials (OBEP), two-particle exchanges should probably be included. Of these, the two-pion exchange potential (TPEP) is undoubtedly the most important. There are contributions to $V_{\rm PV}^{2\pi}$ from both terms of Eq. (2). Those from $J_\mu J^{\mu\dagger}$ have been estimated by Henley *et al.*,[20] whereas the $\sin^2\theta$ terms have been calculated by Desplanques[24] and by Pirner and Riska.[25] Because of double-counting problems with the ρ and ω exchange contributions and because of the uncertainties in evaluating off-mass-shell contributions, which are required to obtain $V_{\rm PV}^{2\pi}$ proportional to $\cos^2\theta$, we assume that the major part of this term is already included in $V_{\rm PV}^{\rho} + V_{\rm PV}^{\omega}$. On the other hand, because standard applications[1,2] of the Cabibbo theory have heretofore not given any p-p and n-n forces (no neutral meson exchanges and thus only exchange forces), we included TPEP from $J'_\mu J'^{\mu\dagger}$. We will see that this contribution is very small indeed compared to that from ρ and ω exchanges. The potential of Desplanques[24] was used; it differs insignificantly from that of Pirner and Riska,[25] and is

$$V_{\rm PV}^{2\pi} = 2f(\tau_1^{(3)} + \tau_2^{(3)})\Big(i(\vec\sigma_1\times\vec\sigma_2)\cdot[\vec{\rm p}, f_C(r)] + (\vec\sigma_1 - \vec\sigma_2)\cdot\{\vec{\rm p}, f_B(r)\}\Big)\,, \tag{10}$$

where $f_C(r)$ and $f_B(r)$ are continuous superpositions of Yukawa functions whose detailed form is given in Ref. 24.

III. PV ASYMMETRY

With the potential $V_{\rm PV}$ given by the sum

$$V_{\rm PV} = V_{\rm PV}^{\rho} + V_{\rm PV}^{\omega} + V_{\rm PV}^{\pi} + V_{\rm PV}^{2\pi}\,, \tag{11}$$

we can calculate PV effects. As discussed in the Introduction, the simplest experiment is the measurement of an asymmetry in the total cross sections of longitudinally polarized nucleons scattered from a hydrogen target. We have also computed the differential asymmetry $A(\theta)$:

$$A(\theta) = \frac{d\sigma_+/d\Omega - d\sigma_-/d\Omega}{d\sigma/d\Omega}\,, \tag{12a}$$

where $d\sigma_+/d\Omega$ and $d\sigma_-/d\Omega$ are the differential cross sections for positive and negative helicities of the incoming nucleon, respectively, and $d\sigma/d\Omega$ is the unpolarized differential cross section. The total asymmetry is

$$A = (\sigma_+ - \sigma_-)/\sigma\,. \tag{12b}$$

The PV scattering amplitude is calculated in distorted-wave Born approximation. The total scattering amplitude can be written as

$$\mathfrak{F}_{s's} = F_{s's} + f_{s's}\,, \tag{13}$$

where $F_{s's}$ is the normal hadronic amplitude for total spins s and s' and

$$f_{s's} \propto \langle \psi_{s'}^{(-)} | V_{\rm PV} | \psi_s^{(+)}\rangle\,, \tag{14}$$

with distorted waves which were computed for both a Hamada-Johnston[26] and a Bryan-Gersten[27] potential.

Since we are particularly interested in the weak PV force, we compute the asymmetry separately for $V_{\rm PV}^{\pi}$, $V_{\rm PV}^{2\pi}$, and $V_{\rm PV}^{\rho} + V_{\rm PV}^{\omega}$. The one-pion potential of Eq. (9) does not contribute for like nucleons. In this case, it follows from the identity of projectile and target that the asymmetry $A(\theta)$, Eq. (12a), is symmetric about 90°. We actually consider neutron-neutron scattering because we do not take Coulomb effects into account. In the notation of Eq. (13) we find

$$A(\theta) = \frac{2\,{\rm Re}[F_{00}^* f_{01} + F_{11}^* f_{10}]}{\frac{1}{2}(|F_{00}|^2 + F_{11}|^2)}\,. \tag{15}$$

For n-p scattering we do not include $V^{2\pi}$, since

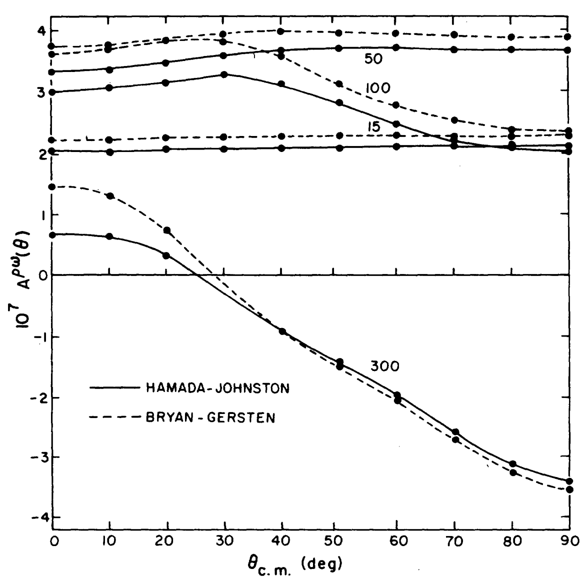

FIG. 2. The asymmetry $A^{\rho\omega}(\theta)$ due to ρ and ω exchanges with $f_\omega g_{\omega N} = 2^{1/2} f_\rho g_{\rho N}$ for like-nucleon scattering as a function of center-of-mass angle, at laboratory energies of 15, 50, 100, and 300 MeV. The solid curves have been calculated with the Hamada-Johnston, the dashed ones with the Bryan-Gersten potential.

this parity-violating interaction gives a negligible contribution to the asymmetry of the like nucleon system. On the other hand, we include the one-pion exchange PV potential, Eq. (9), which contributes to the asymmetry only for n-p scattering.

We supplement the notation used earlier and express the weak scattering amplitude in terms of $f^{I'I}_{s'm',sm}$, where I' and I now denote the isospin of the final and initial states, respectively, and m, m' are magnetic quantum numbers. We find for the numerator in Eq. (12a)

$$\frac{d\sigma_+}{d\Omega}-\frac{d\sigma_-}{d\Omega}=\mathrm{Re}\left\{\sum_{m'}\left[2(F^{11*}_{1m',11}+F^{00*}_{1m',11})(f^{10}_{1m',11}+f^{01}_{1m',11})+(F^{11*}_{1m',10}+F^{00*}_{1m',10})(f^{11}_{1m',00}+f^{00}_{1m',00})\right]\right.$$
$$\left.+(F^{11*}_{00,00}+F^{00*}_{00,00})(f^{11}_{00,10}+f^{00}_{00,10})\right\}, \tag{16a}$$

and for the denominator

$$\frac{d\sigma}{d\Omega}=\frac{1}{4}\left\{\sum_{mm'}\left[|F^{11}_{1m',1m}|^2+|F^{00}_{1m',1m}|^2+2\,\mathrm{Re}(F^{11*}_{1m',1m}F^{00}_{1m',1m})\right]+|F^{11}_{00,00}|^2+|F^{00}_{00,00}|^2+2\,\mathrm{Re}(F^{11*}_{00,00}F^{00}_{00,00})\right\}. \tag{16b}$$

IV. RESULTS AND DISCUSSION

The dominant part of the PV nucleon-nucleon potential is due to ρ and ω exchanges; it is thus of short range. Since the weak amplitude is calculated with waves which are distorted by the strong forces, the amplitude is influenced greatly by the short-range behavior of these distorted waves. It is for this reason that we investigate the PV asymmetry for two different hadronic forces, namely, the hard-core Hamada-Johnston[26] potential and the velocity-dependent Bryan-Gersten potential with no hard core.[27] As might be expected, the asymmetry calculated with the Bryan-Gersten potential is larger than that found with waves distorted by the Hamada-Johnston potential. The enhancement is about twice as large for n-p as for p-p (n-n) scattering. This result can be understood as follows. The low-angular-momentum components are the ones affected most by the hard core. Although $J=0$ contributes for both n-p and p-p, $J=1$ contributes only to the PV n-p scattering amplitude. Hence the enhancement is larger in that case.

We first treat the case of like nucleons, i.e., n-n or p-p. The asymmetry due to ρ and ω exchanges with $f_\omega g_{\omega N}=2^{1/2}f_\rho g_{\rho N}$, is plotted as a function of scattering angle at various energies in Fig. 2 for both the Hamada-Johnston and Bryan-Gersten potentials. In order to discuss the physical interpretation of these results, it is useful to write the total scattering amplitude as

$$\mathfrak{F}=\sum_J\left[f(J)+F(J)\right], \tag{17}$$

where J is the total angular momentum. In the present calculation, angular momenta up to $J=4$ are included. For 15 MeV, the lowest energy considered here, the asymmetry is almost independent of the scattering angle, which indicates that the $J=0$ contribution dominates. As the energy increases, the angular dependence becomes more pronounced since higher angular momenta become more and more important.

The total asymmetry is shown as a function of energy in Fig. 3. We display the asymmetry due to the two-pion potential, multiplied by 10, only as a function of energy because its effect turns

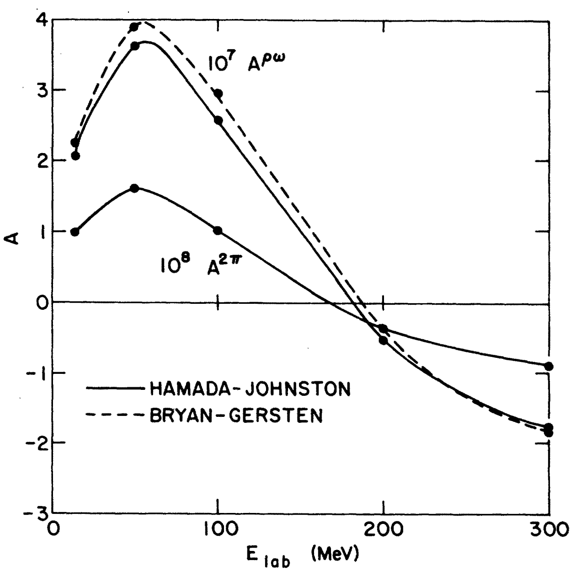

FIG. 3. The total asymmetry for A for like nucleons as a function of laboratory energy. The points correspond to calculated energies. The ρ-ω contributions are shown with $f_\omega g_{\omega N}=2^{1/2}f_\rho g_{\rho N}$. The solid curves are calculated with the Hamada-Johnston, the dashed one with the Bryan-Gersten potential. The 2π contribution for the Hamada-Johnston potential is shown separately; the difference in scale should be noted.

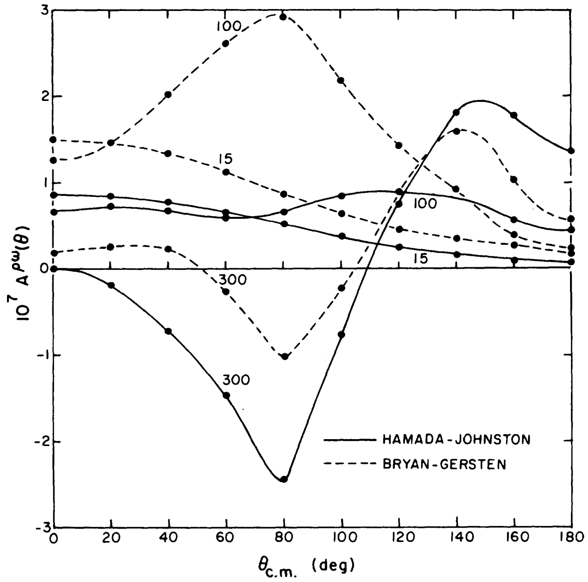

FIG. 4. The sum of the ρ and ω contributions ($f_\omega g_{\omega N} = 2^{1/2} f_\rho g_{\rho N}$) to $A(\theta)$ for n-p scattering as a function of angle at various laboratory energies.

out to be negligible despite the longer range of the two-pion exchange force. We therefore do not include it for the n-p system. To gain some understanding of the behavior shown in Fig. 3, we consider the $J=0$ part of the decomposition given by Eq. (17). For sufficiently small energies both $f_{01}(0)$ and $f_{10}(0)$ are monotonically increasing functions of energy, whereas $F_{00}(0)$ is proportional to $\sin[\delta(^1S_0)]$ and $F_{11}(0)$ to $\sin[\delta(^3P_0)]$. Both phase shifts, here denoted by their spectroscopic nomenclature, $^1S_0$ and $^3P_0$, pass through zero somewhat above 200 MeV. Thus the $J=0$ part of the numerator of Eq. (15) first increases as a function of energy, reaches a maximum, and passes through zero at about 215 MeV. These

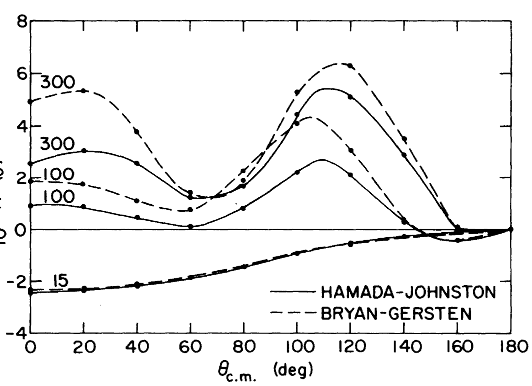

FIG. 5. The 1π contribution to $A(\theta)$ for n-p scattering as a function of angle. Laboratory kinetic energies are given.

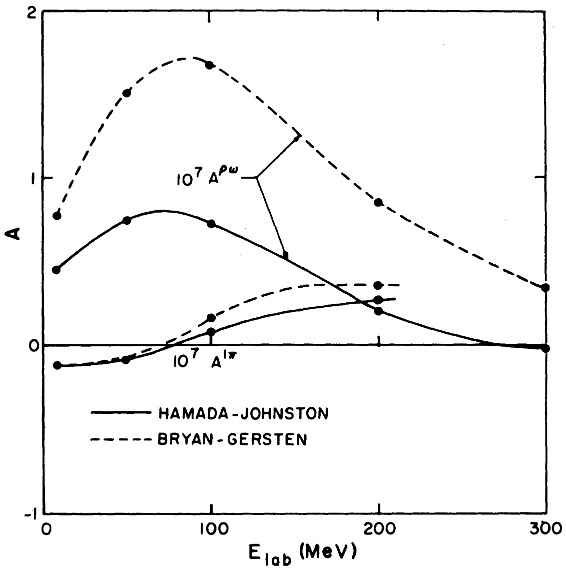

FIG. 6. The ρ-ω ($f_\omega g_{\omega N} = 2^{1/2} f_\rho g_{\rho N}$) and 1π contributions to the total asymmetry A for n-p scattering as a function of energy.

features also can be obtained from the optical theorem. The denominator does not vanish but becomes small, also in the same energy region. The asymmetry due to the $J=0$ part, alone, therefore varies rapidly with energy close to 200 MeV.[28] The higher-angular-momentum contributions to the asymmetry, notably those for $J=2$ and 4, become more important as the energy increases. They have the opposite sign to the $J=0$ part and thus partially cancel it; this results in a

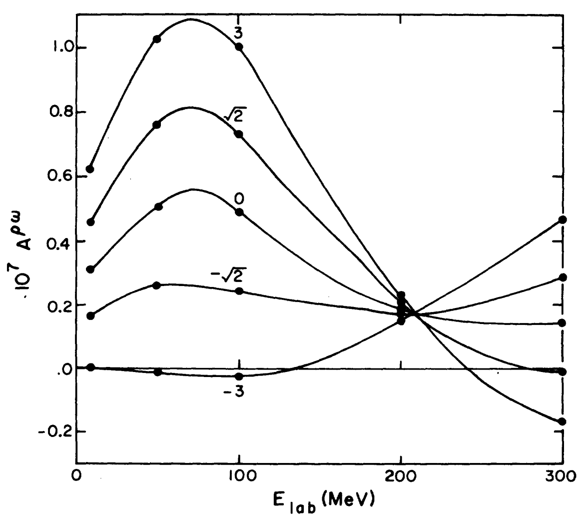

FIG. 7. The total asymmetry for n-p scattering with $f_\omega g_{\omega N}/f_\rho g_{\rho N} = 0$, $\pm 2^{1/2}$, and ± 3 for the Hamada-Johnston potential.

shift of the maximum of A to about 50 MeV and of the zero to approximately 170 MeV, as seen in Fig. 3.[28]

Turning to n-p scattering, we plot the ρ-ω contributions to the differential asymmetry $A(\theta)$ at various energies in Fig. 4. The differential asymmetry due to the one-pion PV exchange potential is shown in Fig. 5. The more-pronounced angular dependence of $A(\theta)$ for n-p scattering than for p-p scattering arises because odd as well as even J values contribute to the PV scattering amplitude in the n-p case. The pion contribution to the differential asymmetry vanishes at 180° at all energies due to the form of the PV potential, Eq. (9).

Figure 6 shows separately the ρ-ω and the one-pion contribution to the total asymmetry as a function of energy. The asymmetry due to ρ-ω exchanges follows the same pattern as for the like nucleon case. The one-pion exchange contribution is only shown up to somewhat above 200 MeV because we did not include $J>4$, and these angular momenta could be important at higher energies due to the long range of the pion exchange force.

We had hoped that A would be larger for the n-p system than for like nucleons, since all J values contribute to the former asymmetry whereas only even ones do for two like particles. This hope was not borne out by the calculation, as can be seen from Fig. 6. One reason is that the partial cancellation of the $J=0$ contribution to the asymmetry by higher angular momenta occurs in the n-p system as well as in the p-p one. Nevertheless, the numerator of Eq. (12b) is indeed larger for n-p than for p-p, but the increase in the denominator is even greater, so that the asymmetry remains small.

In all of the above work we have chosen $f_\omega = f_\rho$ and $g_\omega = 2^{1/2} g_\rho$. However, both the relative sign and the magnitude of f_ω/f_ρ are not known, although quark model and strong $SU(6)_W$ predict $f_\omega/f_\rho < 0$.[22] Furthermore, even the ratio $g_{\omega N}/g_{\rho N}$ is not well determined. Nuclear force fits with various meson exchanges or use of SU(3) together with ϕ-ω mixing give $0.5 \lesssim g_{\omega N}/g_{\rho N} \lesssim 4$.[23,29] On the other hand, there is general agreement that the anomalous magnetic moment coupling $\propto \bar{u}\, i\sigma^{\mu\nu} q_\nu \times u\omega_\mu$ is absent or small for ω exchange.[23,29] This is reflected in Eq. (8). The normal vector coupling, $\propto \bar{u}\gamma^\mu u\omega_\mu$, enters in the PV ω- or ρ-exchange potential only through the combination $f_\omega g_\omega$ or $f_\rho g_\rho$. Because of the uncertainties referred to above, we show the sensitivity of the asymmetry to the ratio $f_\omega g_\omega/f_\rho g_\rho$ in Fig. 7. There, we plot the total n-p asymmetry as a function of energy for $f_\omega g_\omega/f_\rho g_\rho = 0$, $\pm 2^{1/2}$, and ± 3 for the Hamada-Johnston potential with $f_\rho g_{\rho N}$ fixed.

In conclusion, our calculation, which is based on the Cabibbo theory supplemented by $\{27\}$ suppression, predicts an asymmetry which is barely within reach of present-day experiments. Furthermore, we predict that the best energy at which to carry out searches for this asymmetry below appreciable meson production threshold is close to 50 MeV. Although many theoretical assumptions on which the calculation is based do not stand on firm ground, the energy dependence is primarily determined by hadronic interactions and should therefore be trustworthy. Experimental measurements of the asymmetry are of great significance and a study of their energy and angular dependence would shed considerable light on the theory of the weak interaction of hadrons. For instance, in contrast to many previously published papers,[2,8] we predict a PV force between protons, which may even be larger than that between neutrons and protons.

We are grateful to Professor H. Frauenfelder, Dr. R. Mischke, and Dr. D. Nagle for fruitful discussions of the experimental aspects of the problem, and for keeping us informed on the progress of their experimental search for the asymmetry A.

*Work performed in part under the auspices of the U.S. Atomic Energy Commission.

[1]E. M. Henley, Annu. Rev. Nucl. Sci. 19, 367 (1969).
[2]E. Fischbach and D. Tadić, Phys. Rep. 6C, 124 (1973).
[3]M. Gell-Mann, M. L. Goldberger, N. M. Kroll, and F. E. Low, Phys. Rev. 179, 1518 (1969).
[4]S. Weinberg, Phys. Rev. Lett. 19, 1264 (1967); Phys. Rev. D 5, 1412 (1972); H. Georgi and S. L. Glashow, Phys. Rev. Lett. 28, 1494 (1972); Phys. Lett. 44B, 191 (1973); B. W. Lee, Phys. Rev. D 6, 1188 (1972); J. Prentki and B. Zumino, Nucl. Phys. B47, 99 (1972).
[5]T. D. Lee, Phys. Rev. Lett. 26, 801 (1971).
[6]R. E. Marshak, Riazuddin, and C. P. Ryan, *Theory of Weak Interactions* (Wiley Interscience, New York, 1969), Chap. 6.
[7]H. Hättig, K. Hünchen, P. Roth, and H. Wäffler, Nucl. Phys. A137, 144 (1969); H. Hättig, K. Hünchen, and H. Wäffler, Phys. Rev. Lett. 25, 941 (1970).
[8]An up-to-date account of the extensive experimental references is given in M. Gari, Phys. Rep. 6C, 318 (1973).
[9]M. Gari and H. Kümmel, Phys. Rev. Lett. 23, 26 (1969); M. Gari, Phys. Lett. 31B, 627 (1970); E. M. Henley, T. E. Keliher, and D. U. L. Yu, Phys. Rev. Lett. 23, 941 (1969).
[10]G. E. Brown and M. Gari, private communication. See

also E. M. Henley, in The Fifth International Conference on High Energy Physics and Nuclear Structure, Uppsala, 1973 (to be published).

[11]V. R. Brown, E. M. Henley, and F. R. Krejs, Phys. Rev. Lett. 30, 770 (1973).

[12]V. M. Lobashow *et al.*, Nucl. Phys. A197, 241 (1972).

[13]E. Hadjimichael and E. Fishbach, Phys. Rev. D 3, 755 (1971); G. S. Danilov, Yad. Fiz. 147, 88 (1971) [transl.: Sov. J. Nucl. Phys. 14, 443 (1972)].

[14]E. M. Henley, A. H. Huffman, and D. U. L. Yu, Phys. Rev. D 7, 943 (1973); M. Gari and A. H. Huffman, Phys. Rev. C 7, 531 (1973); P. Herczeg and P. Singer, to be published; P. Herczeg and L. Wolfenstein, to be published.

[15]M. Simonius, Phys. Lett. 41B, 415 (1972); to be published; T. D. Lee, private communication.

[16]H. Frauenfelder, D. E. Nagle, and R. E. Mischke, private communication.

[17]V. M. Lobashov, V. A. Nazarenko, L. F. Saenko, L. M. Smotritskii, and G. I. Kharkevich, Zh. Eksp. Teor. Fiz. Pis'ma Red. 3, 76 (1966) [transl.: JETP Lett. 3, 47 (1966)]; Zh. Eksp. Teor. Fiz. Pis'ma Red. 3, 268 (1966) [transl.: JETP Lett. 3, 173 (1966)]; Zh. Eksp. Teor. Fiz. Pis'ma Red. 5, 73 (1967) [transl.: JETP Lett. 5, 59 (1967)].

[18]F. C. Michel, Phys. Rev. 133, B329 (1964).

[19]M. Gell-Mann and M. Levy, Nuovo Cimento 16, 53 (1960); M. Levy, *ibid.* 52A, 23 (1967).

[20]E. M. Henley, T. E. Keliher, W. J. Pardee, and D. U. L. Yu, Phys. Rev. D 8, 1503 (1973).

[21]C. Christenson *et al.*, Phys. Rev. D 5, 1628 (1972).

[22]B. H. J. McKellar and P. Pick, Phys. Rev. D 7, 260 (1973); J. G. Körner, Phys. Lett. 44B, 361 (1973).

[23]See, e.g. G. Kramer, *Springer Tracts in Modern Physics* (Springer, New York, 1970), Vol. 55, p. 152. This feature was not included in our previous treatment of *p-p* scattering, Ref. 11, where, apart from isospin, the same potential was used for ρ and ω exchanges.

[24]B. Desplanques, Phys. Lett. 41B, 461 (1972).

[25]H. J. Pirner and D. O. Riska, Phys. Lett. 44B, 151 (1973).

[26]T. Hamada and I. D. Johnston, Nucl. Phys. 34, 382 (1962).

[27]R. A. Bryan and A. Gersten, Phys. Rev. D 6, 341 (1972); 7, 2802(E) (1973). Fit *D* of this reference has been used in our calculation.

[28]Simonius, Ref. 15, has recently calculated the PV asymmetry for *p-p* scattering. He only includes $J=0$ for the PV matrix element, but $J=0$, 1, and 2 for the parity-conserving matrices. These differences explain why his calculated result compares more closely with our total asymmetry *A* than with that calculated for $J=0$ alone. Simonius also uses Watson's theorem [K. M. Watson, Phys. Rev. 95, 228 (1954)], which is inapplicable for higher values of *J* because the *S* matrix connects more than two states.

[29]M. M. Nagels, T. A. Rijken, and J. J. de Swart, Phys. Rev. Lett. 31, 569 (1973).