be of some theoretical interest.

In general, since the CP nonconservation occurs in the weak-interaction Hamiltonian, there are no large effects to be expected such as an asymmetry in $\eta \rightarrow \pi^+ + \pi^0 + \pi^-$ or $\eta \rightarrow \pi^+ + \pi^- + \gamma$, time-reversal noninvariance in nuclear γ emission, or the occurrence of $\eta \rightarrow \pi^0 + e^+ + e^-$.

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PARITY-NONCONSERVING NUCLEAR FORCES*

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The effective single-particle weak parity-nonconserving potential is recalculated including parity-nonconserving one-pion exchange forces and the two-body correlations induced by the hard cores in the nucleon-nucleon potential. The circular polarization of the γ ray from the 482-keV transition in ^{181}Ta is calculated to be $-(0.2 \pm 0.1) \times 10^{-4}$. The measured value is $-(0.06 \pm 0.01) \times 10^{-4}$, showing that the observed parity nonconservation in nonleptonic $\Delta S = 0$ transitions is in agreement with the Cabibbo theory of weak interactions.

The current-current theory of weak interactions predicts a weak force between nuclei which can in principle be detected through its parity-nonconserving effects.[1] The attempts to calculate such effects and the attempts to observe them have been recently reviewed by Okun.[2] The circular polarization measurements of Lobashov et al.[3] and Boehm and Kankeleit[4] are shown in Table I, and the asymmetry measurements of Abov et al.[5] and Warming et al.[6] are shown in Table II. Also shown in the tables are the results of calculations reported in the literature.[7-11] It is noticeable that the calculations overestimate the effect. It was pointed out by Adams[12] that the hard core of the nucleon-nucleon potential, in keeping the nuclei in the tail of the weak potential, may provide an explanation of this overestimate.

In the present Letter we show that this is in fact the case, and in taking the hard-core correlations into account we significantly improve the agreement with experiment. We also include in the calculation the strangeness-changing weak currents which induce a one-pion exchange, parity-nonconserving potential between nuclei.[13] Enhanced by its long range, and suppressed by the factor $\sin^2\theta_C$, the net effect is 25% of that of strangeness-conserving currents. The relative sign of the two contributions is not fixed by experiment, nor by any reliable theoretical arguments.[14] (The sign of G is taken as

Table I. Circular polarization in γ transitions, in units 10^{-4}.

Transition	Measurements Value	Ref.	Previous calculations Value	Ref.	Present calculations $fg>0$	$fg<0$
^{175}Lu 343 keV	$+(0.2+0.3)$	4	$+(0.3\pm0.2)$ -0.7	8 10	$-(0.1\pm0.05)$	$-(0.15\pm0.1)$
^{175}Lu 396 keV	$+(0.4\pm0.1)$	3	$\pm(0.9\pm0.6)$	7	$\pm(0.3\pm0.2)$	$\pm(0.45\pm0.3)$
^{181}Ta 482 keV	$-(0.06\pm0.01)$ $-(0.1\ \pm0.4)$	3 4	$-(0.6\pm0.3)$ -0.7	8 10	$-(0.2\pm0.1)$	$-(0.3\pm0.1)$
^{203}Tl 273 keV	$-(0.2\ \pm0.3)$	4	$-(0.9\pm0.3)$	9	$-(0.3\pm0.1)$	$-(0.45\pm0.2)$

Table II. Asymmetry coefficient of γ transition after capture of polarized neutrons, in units 10^{-4}.

Transition	Measurements Value	Ref.	Previous calculations Value	Ref.	Present calculations $gf>0$	$gf<0$
^{114}Cd 9.04 MeV	$-(3.7\pm0.9)$ $-(2.5\pm2.2)$	5 6	±6	11	±2	±3

positive in accordance with the intermediate-boson theory, as is customary.) Our final results are shown in Tables I and II. If the sign of gf is taken as positive the agreement with experiment is good. In fact it is probably too good in view of the tenuous link between the weak Hamiltonian and the final results.

Nevertheless it is possible to conclude that the observed parity-nonconserving effects in nuclei do not contradict the presently accepted theory of weak interactions, and that, unfortunately for the theorist, it is necessary to use a correlated nuclear wave function in a quantitative calculation.

The calculation proceeds in three stages: A weak nucleon-nucleon potential is derived, it is then averaged to obtain an effective single-particle potential, and this potential is used to calculate polarizations, asymmetries, etc. Each stage introduces errors which are difficult to estimate.

The weak nucleon-nucleon potential is calculated in different approximations from the strangeness-conserving and the strangeness-changing currents.

Michel[7] calculated the contribution of the strangeness-conserving currents by looking at $\langle NN|H_{\text{p.v.}}^{w}\times|NN\rangle$ in the crossed channel, with $H_{\text{p.v.}}^{w}$ the parity-nonconserving interaction, and writing

$$\langle N\bar{N}|J_V^{1+i2}J_A^{1-i2}|N\bar{N}\rangle\approx\langle N\bar{N}|J_V^{1+i2}|0\rangle\langle 0|J_A^{1-i2}|N\bar{N}\rangle. \tag{1}$$

With ρ dominance of the vector form factor, this can be regarded as a $\rho^{\pm}$-exchange potential, V_ρ. In the nonrelativistic limit,

$$V_\rho=\frac{Gm_\rho^{\;2}}{4\pi\sqrt{2}m_N}\Big\{(\vec{\sigma}^{(1)}-\vec{\sigma}^{(2)})\cdot\{\tfrac{1}{2}(\vec{p}_1-\vec{p}_2),\exp(-m_\rho r_{12})/r_{12}\}$$
$$+(\mu^v+1)(i\vec{\sigma}^{(1)}\times\vec{\sigma}^{(2)})\cdot[\tfrac{1}{2}(\vec{p}_1-\vec{p}_2),\exp(-m_\rho r_{12})/r_{12}]\Big\}T_{12}^{(+)}, \tag{2}$$

where

$$T_{12}^{(\pm)}=\tau_-^{(1)}\tau_+^{(2)}\pm\tau_+^{(1)}\tau_-^{(2)}, \tag{3}$$

and G is the Fermi constant ($Gm_N^{\;2}=1.02\times10^{-5}$) and $\mu^v=3.70$ the isovector anomalous magnetic moment of the nucleon.

The one-pion exchange contribution is calculated by relating $\langle N|H_w|N\pi\rangle$ to the matrix elements for parity-nonconserving hyperon decays to $N\pi$.[13] The potential is

$$V_\pi = \frac{gf}{4\pi\sqrt{2}m_N}(\vec{\sigma}^{(1)}+\vec{\sigma}^{(2)})\cdot[\tfrac{1}{2}\vec{p}_1-\vec{p}_2), \exp(-m_\pi r_{12})/r_{12}]T_{12}^{(-)} \tag{4}$$

with

$$|f| = 5.2\times 10^{-8}, \tag{5}$$

and g is the pseudoscalar pion-nucleon coupling constant ($g^2/4\pi = 14.4$).

One may further show that CP conservation forbids contributions to the parity-nonconserving potential from exchange of neutral spin-zero mesons.[15] Exchange of neutral vector mesons would appear as the next term in the expansion of $\langle N\bar{N}|J_V^{1+i2}J_A^{1-i2}|N\bar{N}\rangle$ in Eq. (1). Thus in taking

$$V(1,2) = V_\rho + V_\pi \tag{6}$$

as the parity-nonconserving nucleon-nucleon potential we have included the dominant terms.

An effective single-particle potential W may be defined by

$$\langle\alpha|W|\beta\rangle = \langle\Psi_\alpha|\sum_{i<j}V(i,j)|\Psi_\beta\rangle, \tag{7}$$

where $|\alpha\rangle$ are one-particle states and Ψ_α are many-particle states with the last particle in the state α, and the other particles in a standard state $\{\gamma_1,\cdots,\gamma_{N-1}\}$. If the wave function Ψ_α is a Slater determinant of the states $|\gamma_i\rangle$ then

$$\langle\alpha|W|\beta\rangle = \sum_i\langle\alpha\gamma_i|V|\beta\gamma_i-\gamma_i\beta\rangle. \tag{8}$$

Michel proposed that the effective potential be defined in nuclear matter with plane-wave states and regarded as the momentum-space representation of a potential which is then used to calculate the nuclear matrix elements. Using V_ρ only, and taking m_ρ infinite, he obtained the potential W_M,

$$\langle\vec{k}'s_z't_z'|W_M|\vec{k}s_zs_t\rangle = \delta_{s_zs_z'}\delta_{t_zt_z'}\delta(\vec{k}-\vec{k}')\frac{(\mu^v+1)G}{\sqrt{2}m_N}\rho\left[\tfrac{1}{2}+t_z\frac{N-Z}{A}\right]\langle s_z'|\vec{\sigma}|s_z\rangle\cdot\vec{k}, \tag{9}$$

which was used in the calculations of Refs. 7-9.[16] ρ is the density of nuclear matter.

To take the hard cores into account we replace the two-particle plane-wave state $|\alpha\gamma_i\rangle$ by the corresponding Bethe-Goldstone state. Using the cannonical transformation from the plane-wave Slater determinent state to the correlated state,[17] this may be shown to take two-particle correlations into account correctly but to neglect higher order correlations.

In the numerical calculation, hard-core correlations were included only in relative s states, using Gomes's approximate wave function[18]

$$u(q,r) = \left\{\frac{\sin qr}{qr}-\frac{\sin qc}{qr}\frac{\mathrm{si}(\beta r)}{\mathrm{si}(\beta c)}\right\}\theta(r-c), \tag{10}$$

where c is the core radius 0.4 fm and $\beta = 1.633\ \mathrm{fm}^{-1}$ corresponding to a healing distance of 1.18 fm. The distinction between singlet and triplet strong nucleon-nucleon potentials was ignored. The Gomes wave function was checked by evaluating the matrix elements using a Moskowski-Scott wave function.[19] The largest variation was 5%.

The result for the matrix elements of the effective potential is

$$\langle\vec{k}'s_zt_z'|W|\vec{k}s_zt_z\rangle = \left\{w_\rho(k)-\frac{gf}{(\mu^v+1)Gm_\pi^{\,2}}2t_zw_\pi(k)\right\}\langle\vec{k}'s_z't_zt|W_M|\vec{k}s_zt_z\rangle, \tag{11}$$

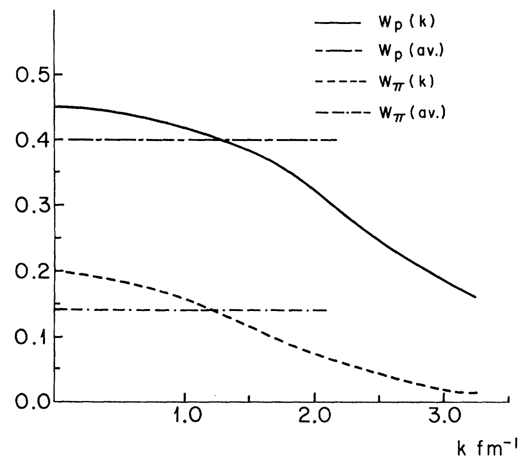

FIG. 1. Variation of effective potential with momentum.

where $w_\rho(k)$ and $w_\pi(k)$ are plotted in Fig. 1.[20]

In the shell-model matrix elements $\langle n'l'j'm'|W \times|nljm\rangle$, the Gaussian cutoff in the wave function effectively limits the k-space integration to a sphere of radius 2.0 fm^{-1} for the nuclei of interest. As both $w_\rho(k)$ and $w_\pi(k)$ vary slowly in this region, we replace them by the average values

$$w_\rho(\mathrm{av})=0.4,\quad w_\pi(\mathrm{av})=0.14$$

The errors introduced, about 10% in w_ρ and 30% in w_π, are lost in the background of the approximations we have made. We can tolerate the large error in the pion term as it is itself a small correction to the rho term.

Polarizations and asymmetries may then be calculated from results obtained from Michel's potential simply by multiplying by the factor $0.4 \mp 0.1(2t_z)$, where the minus sign corresponds to the choice of a positive value for gf.

The results obtained in this way are tabulated in Tables I and II, and are in quite good agreement with the observations.

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[14]One can estimate the sign by calculating the lowest order graphs in an intermediate-boson theory, retaining only that part of the weak current which induces transitions within the baryon octet. When a cutoff is introduced the result is that the pion-exchange "coupling constant," gf, is positive. However, as a similar simple-minded calculation gives the wrong sign for the n-p mass difference, we should not take this model too seriously.

[15]G. Barton, Nuovo Cimento 19, 512 (1961), has shown that a neutral pseudoscalar particle cannot couple to the $\bar{N}N$ system in two CP-conserving couplings, one of which conserves while the other does not conserve P. It is an immediate extension that the same result holds for a neutral scalar particle.

[16]The particle-physics convention is used for the isospin, so that the proton has isospin $+\frac{1}{2}$.

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[20]One may worry that transitions with $\Delta T=1$, which could have been induced by V_π, are missing from the effective potential. However, in heavy nuclei such transitions will be suppressed by a factor $(N-Z+2)^{-1}$, which is about 1/40 in these cases, just as isospin mixing by Coulomb forces is suppressed. [Cf. A. M. Lane and J. Soper, Phys. Letters 1, 28 (1962).]