PHYSICAL REVIEW C

NUCLEAR PHYSICS

THIRD SERIES, VOLUME 36, NUMBER 6 DECEMBER 1987

ρ-ω mixing in nuclear charge asymmetry

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(Received 29 July 1987)

The particle physics input to the NN charge asymmetric potential due to ρ-ω mixing is updated. The calculated changes in scattering lengths $|a_{nn}|-|a_{pp}|\sim 1.3\pm 0.2$ fm are in good agreement with experiment ($\sim 1.4\pm 0.8$ fm), as are the contributions to the binding energy difference in the $A=3$ mirror nuclei.

I. INTRODUCTION

The ρ-ω mixing prediction[1,2] of a positive value of ~ 1 fm for the difference $\Delta a = |a_{nn}| - |a_{pp}|$ of the singlet $^1S_0$ scattering lengths [after correction for the direct electromagnetic (em) effects] was soon confirmed by a remeasurement of the neutron-neutron scattering length a_{nn} by the reaction $\pi^- d \rightarrow \gamma nn$.[3] The theoretical uncertainties in the analysis of this experiment have been steadily minimized[4] and the extracted value of a_{nn} confirmed by another kinematically complete measurement of the same reaction.[5] The major source of that prediction of charge asymmetry is isospin mixing of vector mesons (ρ-ω mixing) in single meson-exchange models of the two-nucleon force. This charge asymmetric potential $\Delta V^{\rho\omega}$ then depends on the electromagnetic[6] transition matrix element $\langle \rho \mid H_{em} \mid \omega \rangle$ between the vector mesons ρ^0 and ω. The magnitude and sign of $\langle \rho \mid H_{em} \mid \omega \rangle$ were just beginning to be determined reliably a decade ago—the magnitudes in Refs. 1 and 2 differ by almost a factor of two because they quoted inconsistent experiments.

Because ρ-ω mixing is a dominant component of some (but not all[7]) experimental tests of nuclear charge asymmetry and because a_{nn} is now reliably known, it is appropriate to update the particle physics input to the potential $\Delta V^{\rho\omega}$ which breaks charge symmetry.[1,2] In this note, we take advantage of a new fit to the world data from the reaction $e^+e^- \rightarrow \pi^+\pi^-$, which is now very accurate in the ρ-ω mass range, to rederive $\langle \rho \mid H_{em} \mid \omega \rangle$ and $\Delta V^{\rho\omega}$. The effects of ρ-ω mixing are then calculated for two simple mirror systems which are not eigenstates of the charge symmetry operator (a rotation by π about the 2-axis of isospin space).[8] These systems are the low-energy two-nucleon scattering parameters already mentioned and the binding energy difference ΔE of the mirror nuclei ^{3}He-^{3}H. We find a modest increase in the nuclear charge asymmetry due to ρ-ω mixing over that of Ref. 1. The formalism of Ref. 1 (and Ref. 2) is retained in the calculation. This means that, if that theoretical framework continues to stand the test of time, this source of charge symmetry breaking has much less uncertainty than ever before because of the greater precision of the particle data in the last decade.

II. THE ρ-ω MIXING

From the experimental point of view, ρ-ω interference has always been observed through the G-parity forbidden $\omega \rightarrow \pi^+\pi^-$ decay of ω mesons produced in different reactions. These include[9] strong interactions such as proton-antiproton annihilations and pion and kaon collisions on nucleons, and electromagnetic processes such as photoproduction on nuclei and e^+e^- annihilations into $\pi^+\pi^-$ pairs. The model assumptions used in data analyses are difficult to verify in the fixed target experiments in which besides the pion pair other hadrons were involved. Because of a "pure" initial state in e^+e^- collider measurements of the pion form factor, model assumptions are required only at the last stage of the analysis, so the colliding ring experiments provide the cleanest way to determine ρ-ω mixing.

The data on the electromagnetic form factor of the pion are concisely summarized in Ref. 10. The timelike region near the ρ (and ω) mass probed by the colliding ring e^+e^- experiments is successfully described by a ρ-meson intermediate state with ρ-ω interference taken into account. Three discussions[11] of vector meson mixing parameters, appearing since our analysis,[1] were written too early to include the recent Novosibirsk data[12] which summed to a number of detected events exceeding by two orders of magnitude the statistics of earlier colliding ring experiments.[13] The important datum from the fit in Ref. 12 to all the timelike data with $\sqrt{s} > 2m_\pi$ is the branching ratio

$$B_{\omega\pi\pi} = \Gamma(\omega \to 2\pi)/\Gamma(\omega \to \text{all})$$

which was determined to be

$$B_{\omega\pi\pi} = (2.3 \pm 0.4 \pm 0.2)\% \ , \qquad (1)$$

where the first error includes the statistical and systematic inaccuracy in the experimental data, while the second is an error due to the different models considered for the data at energies higher than $s \approx m_\rho^2$. Applying the Feynman rules to the amplitude for the G-parity forbidden decay $\omega \to 2\pi$ leads immediately to[1]

$$\Gamma(\omega \to 2\pi) \approx \left| \frac{\langle \rho \mid H_{\text{em}} \mid \omega \rangle}{i m_\rho \Gamma_\rho} \right|^2 \Gamma(\rho \to 2\pi) \ , \qquad (2)$$

and a transition matrix element

$$|\langle \rho \mid H_{\text{em}} \mid \omega \rangle| \approx 0.004\,52 \pm 0.0006 \ \text{GeV}^2 \ , \qquad (3)$$

where we have used the values $m_\rho \sim 775.9 \pm 1.1$ MeV and $\Gamma_\rho \sim 150.5 \pm 3.0$ MeV also from Ref. 12. The result is unchanged if one adopts the Particle Data Group (PDG) values[9] $m_\rho \sim 770 \pm 3$ MeV, $\Gamma_\rho \sim 153 \pm 2$ MeV.

The branching ratio $B_{\omega\pi\pi}$ of Eq. (1) is consistent with the earlier Saclay storage ring result[13] of $(1.6 \pm ^{0.9}_{0.7})\%$ but the latter branching ratio had a larger statistical and systematic error and no estimate of model error. Equation (1) is to be preferred over the PDG world average of $B_{\omega\pi\pi} = (1.7 \pm 0.2)\%$ which is heavily weighted by two precise measurements of photoproduction of vector mesons from complex nuclei. The extraction of $B_{\omega\pi\pi}$ from these latter experiments relies on model assumptions about diffractive scattering of ρ and ω from a nucleus, optical potentials, etc. which are nearly uncontrollable. We note that the very large value of 0.006 GeV^2 for $|\langle \rho \mid H_{\text{em}} \mid \omega \rangle|$ obtained in Ref. 2 was based upon parameters from a 1972 Saclay e^+e^- measurement which had very large statistical and systematic errors and was discarded from the data base of Ref. 12. A final argument in favor of the fit of Ref. 12 is that it extrapolates to the pion charge radius at $s = 0$ in good agreement with the radius obtained from spacelike values obtained by colliding pions with an electron target.

The sign of $\langle \rho \mid H_{\text{em}} \mid \omega \rangle$ has been discussed at some length in Ref. 1. It is determined from the relative phase of the ω to the ρ amplitudes in $e^+e^- \to \pi^+\pi^-$ near m_ρ and m_ω. It is negative. Therefore

$$\langle \rho \mid H_{\text{em}} \mid \omega \rangle = -0.004\,52 \pm 0.0006 \ \text{GeV}^2 \qquad (4)$$

is preferred to the often used value of $\langle \rho \mid H_{\text{em}} \mid \omega \rangle \approx -0.0034 \ \text{GeV}^2$ derived in Ref. 1 from contemporary "world averages."

Theoretical descriptions of meson mixing have evolved from the old SU(3) tadpole parametrization of Ref. 1. Firstly, meson mixing has been attributed to mass ratios of current quarks in a particular scheme of chiral symmetry breaking. Each new experimental datum such as Eq. (4) then provides an opportunity for a new parametrization of the up-down current quark mass ratio; for details of this approach see Ref. 10. Another approach, which can avoid assumptions about the mass ratios, attributes meson mixing to isoscalar quark-annihilation diagrams mediated by gluon exchanges.[14] In the current quark picture, these diagrams are not calculated but reexpressed in terms of measured electromagnetic mass differences. This approach yields a theoretical expectation[15] of $\langle \rho \mid H_{\text{em}} \mid \omega \rangle \approx -0.0029 \pm 0.0012 \ \text{GeV}^2$, just barely consistent in magnitude but significantly with the same sign as experiment. This theoretical treatment of ρ-ω mixing was made by analogy to the annihilation graph driven analysis[15] of $\pi^0\eta\eta'$ mixing which also turns out to be about 30% lower than experiment.[16]

III. THE CHARGE ASYMMETRIC POTENTIAL $\Delta V^{\rho\omega}$

The Feynman rules for the ρ-ω force diagram give[1]

$$[S_{fi} = 1_{fi} - iT_{fi}(2\pi)^4\delta(P_{fi})]$$

$$T^{\rho\omega}_{\text{NN}} = -\frac{H_\mu(\omega \text{N}_1\text{N}_1)H^\mu(\rho \text{N}_2\text{N}_2)\langle \rho \mid H_{\text{em}} \mid \omega \rangle}{(m_\rho^2 - t)(m_\omega^2 - t)} + (1 \leftrightarrow 2) \ , \qquad (5)$$

where we have taken a narrow ρ width to simplify the calculation. We use the specific forms for the VNN couplings

$$H_\mu(\omega \text{NN}) = \tfrac{1}{2} g_\omega \overline{\text{N}}(\gamma_\mu^+ \kappa_S i\sigma_{\mu\nu}\Delta^\nu/2M_{\text{N}})\text{N} \ , \qquad (6a)$$

$$H_\mu(\rho \text{NN}) = \tfrac{1}{2} g_\rho \overline{\text{N}}\tau^3(\gamma_\mu + \kappa_V i\sigma_{\mu\nu}\Delta^\nu/2M_{\text{N}})\text{N} \ , \qquad (6b)$$

specialize to the $T=1$, ${}^1S_0$ state, and make a nonrelativistic reduction of (5) to $O(k^2)$ where $\mathbf{k}$ is the momentum transferred to the vector meson $\mathbf{k} = \mathbf{p}_{1f} - \mathbf{p}_{1i} = \mathbf{p}_{2i} - \mathbf{p}_{2f}$:

$$\Delta T^{\rho\omega} \equiv T^{\rho\omega}_{\text{nn}}({}^1S_0) - T^{\rho\omega}_{\text{pp}}({}^1S_0) = g_\rho g_\omega \langle \rho \mid H_{\text{em}} \mid \omega \rangle \frac{1 + \beta \mathbf{k}^2/M_{\text{N}}^2}{(\mathbf{k}^2 + m_\rho^2)(\mathbf{k}^2 + m_\omega^2)} \ , \qquad (7)$$

so that β is the leading correction to the dominant dipole form ($m_\rho \approx m_\omega$). The ${}^1S_0$ potential is defined in Born approximation as the Fourier transform of (7) in the limit $m_\rho = m_\omega = m_V$,

$$\Delta V^{\rho\omega} = \frac{g_\rho g_\omega}{4\pi} \frac{\langle \rho \mid H_{\text{em}} \mid \omega \rangle}{2m_V} e^{-m_V r} \left[1 + \beta \left[\frac{2}{m_V r} - 1 \right] \frac{m_V^2}{M_{\text{N}}^2} \right] . \qquad (8)$$

This form of $\Delta V^{\rho\omega}$ is an exceptionally accurate representation of the Fourier transform of the exact form (7) as can be seen from a power series expansion in

$$[(m_\omega - m_\rho)/(m_\omega + m_\rho)]^2 \sim 10^{-5} .$$

The form (8) has the advantage of (i) displaying the insensitivity of $\Delta V^{\rho\omega}$ to the rho mass which is still imperfectly known, (ii) making clear that the dominant term of the potential is exponential rather than Yukawa, and (iii) isolating the relativistic corrections of order M_N^{-2} which correct the dominant $O(1)$ term by about 10% in a consistent nonrelativistic reduction.[2] Note that the second term in brackets is of opposite sign to the first for all values of $r \gtrsim 2/mV \approx \frac{1}{2}$ fm.

The constant β in (7) and (8) is (determined by independent methods in Refs. 1 and 2)

$$\beta = \tfrac{1}{2}[\kappa_S \kappa_V + \tfrac{1}{2}(\kappa_S + \kappa_V)] , \tag{9}$$

and represents the local relativistic correction to the dominant exponential. Friar and Gibson[2] have identified all corrections of this order (M_N^{-2}) including a nonlocal momentum-dependent term and a retardation term. Fortunately their calculations show that the $\sim 20\%$ correction from (9) is partially cancelled by the nonlocal terms and the sum of all (M_N^{-2}) terms is about a 10% reduction of the leading exponential of (8). Rather than attempt to calculate with momentum-dependent potentials, we will give numerical results for both $O(1)$ and $O(M_N^{-2})$ terms in (8) and let the reader estimate the full nonrelativistic correction.

To proceed one must know the coupling constants g_ρ, g_ω, κ_S, and κ_V. According to the vector dominance model of the electromagnetic form factor for hadrons, κ_V and κ_S are the isovector and isoscalar anomalous magnetic moments of the nucleon: $\kappa_V = 3.70$ and $\kappa_S = -0.12$. As the vector dominance hypothesis[17] is the basis of the extraction of ρ-ω mixing from the pion form factor and universality ($g_\rho \approx g_{\rho\pi\pi} \approx f_\rho$) appears to be in good agreement with the data it appears most consistent to extract g_ρ and g_ω from the e^+e^- scattering data. Utilizing the most recent data on $\Gamma(\rho \to e^+e^-)$ (Ref. 12) and $\Gamma(\omega \to e^+e^-)$[9] one obtains $g_\rho^2/4\pi \approx 2.4$ as expected from universality but a slight increase over Ref. 1 in the ωNN coupling constant $g_\omega^2/4\pi \approx 21.0 \pm 1.3$. This estimate of g_ρ (at $t = m_\rho^2$) is in agreement with the narrow width approximation to ρ exchange in πN scattering of $g_\rho \approx 2.2$.[18] The present estimate of g_ω is rather conservative compared to those of the compilations[19] which range from $g_\omega^2/4\pi \approx 32 \pm 6$ extracted from a dispersion relation analysis of NN scattering[20] to $g_\omega^2/4\pi = 96 \pm 48$ according to the electromagnetic form factors of the nucleon.[21] [The definitions of VNN coupling constants vary according to authors. Without a definition such as the Hamiltonian densities (6), which are lacking in Refs. 20 and 21, one cannot be sure of the factors of 2 in a published g_ρ or g_ω—we think our interpretation is correct.]

The mixing parameter $\langle \rho | H_{em} | \omega \rangle$ and the coupling constants are all on-mass-shell values ($t \approx m_V^2$) with the exception of κ_V and κ_S which, by vector dominance, apply at the $t=0$ (the photon mass shell). The ρ-mass-shell[18] value of κ_V is 6.6, nearly twice as large as the vector dominance prediction of $\kappa_V = 3.7$. A prescription for this rapid variation of κ_V with t which has the effect of weakening the magnetic ρ coupling in the region of interest is given in Ref. 22. It appears an unnecessary refinement for our purposes since $\beta \sim \kappa_V/4$ (neglecting the small κ_S) increases from ~ 0.9 to ~ 1.6; well within the range of values already found[1] to be small corrections to the dominant term of $\Delta V^{\rho\omega}$. In particular, we cannot verify the suggestion that the magnetic couplings of the ρ and ω can cancel the dominant charge coupling term so that even the sign of the scattering length difference is uncertain.[23] Only with the very large $\kappa_S \sim 0.7$ chosen in Ref. 23 can β of (8) be large enough to change the sign of charge asymmetry effects. There is no evidence, however, for deviation from the vector dominance value of $\kappa_S \approx 0$. Empirical values are $\kappa_S \leq 0.2$ from form factor data[21] and $\kappa_S = 0.14 \pm 0.20$ from NN scattering.[20]

Finally we repeat the important observation that the product

$$g_\rho g_\omega \langle \rho | H_{em} | \omega \rangle \langle 0 , \tag{10}$$

independent of the ϕ-ω mixing angle[1] and that the magnetic couplings do not change the overall sign of $\Delta V_{\rho\omega}$.

It appears then that one could simply scale the earlier results for $\Delta a/a$ ($\Delta a = |a_{nn}| - |a_{pp}|$) by the ratio ~ 1.4 of the strengths (10) of the present and earlier[1] $\Delta V^{\rho\omega}$. While this is true for the scattering lengths, new experimental data require a fresh calculation of the contribution of $\Delta V^{\rho\omega}$ to ΔE for the trinucleon bound state.

IV. SCATTERING LENGTH AND EFFECTIVE RANGE

The empirical value of a_{nn} does not suffer from the model dependence[24] of the subtraction of the direct electromagnetic (em) interaction needed to extract the true nuclear scattering length from the experimental a_{nn}. The experimental a_{nn} found from $\pi^- d \to \gamma nn$ in which only the photon was detected is[3]

$$a_{nn} = -18.5 \pm 0.4 \pm 0.2 \text{ fm} ,$$

in excellent agreement with the kinematically complete determination from the same reaction[5]

$$a_{nn} = -18.7 \pm 0.6 \text{ fm} ,$$

which folds the negligible theoretical error into the quoted error. The value of a_{nn} from the photon spectrum is presumably more accurate because of the very high number of events recorded in that experiment.

The nuclear a_{pp} is difficult to extract because of nonlocalities in the electromagnetic interaction between the two protons. If we accept the usual subtraction of the latter from the measured $a_{pp}^{exp} = -7.828 \pm 0.008$ fm, the result is[8]

$$a_{pp} = -17.1 \pm 0.2 \text{ fm} .$$

Another extracted value of $a_{pp} \sim -17.9$ fm is quoted in

TABLE I. Changes in scattering length $\Delta a = |a_{nn}| - |a_{pp}|$ and effective range $\Delta r = r_{nn} - r_{pp}$ as $\Delta V^{\rho\omega}$ is added to the charge symmetric Reid and de Tourreil–Rouben–Sprung potentials. The local nonrelativistic term of Eq. (8), labeled NR ($\beta=0$) and the local relativistic correction are shown separately for vector dominance magnetic (VDM) couplings ($\beta=0.637$) and the stronger on-mass-shell rho coupling $\kappa_V = 6.6$ labeled by ($\beta=1.22$).

	Reid		dTRS	
	Δa (fm)	Δr (fm)	Δa (fm)	Δr (fm)
NR ($\beta=0$)	1.39	−0.026	1.64	−0.032
Local ($\beta=0.637$)	−0.17	+0.005	−0.16	+0.003
Total VDM	1.12	−0.021	1.48	−0.029
Local ($\beta=1.22$)	−.49	+0.008	−0.29	+0.005
Total ($\kappa_V=6.6$)	0.90	−0.018	1.35	−0.027

a 1982 compilation[19] but the calculation has yet to be published. Charge asymmetry in these mirror states is then characterized by[3,8]

$$(|a_{nn}| - |a_{pp}|)_{exp} \approx 1.4 \pm 0.8 \text{ fm} , \quad (11)$$

and

$$(r_{nn} - r_{pp})_{exp} \approx 0.04 \pm 0.25 \text{ fm} .$$

We present in Table I shifts in a and the effective range r obtained by the variable phase method for the potential $\Delta V^{\rho\omega}$ added to two local "realistic" charge symmetric potentials. We chose the Reid soft-core potential[25] which has a large repulsion at short distances and the de Tourreil–Rouben–Sprung (dTRS) potential[26] which has a "supersoft core." It can be shown analytically[1] that the importance of the second order terms in (8) is increased if the charge symmetric force does not allow the nucleons to be close to one another. The exact numerical results also display this trend. Only the local second order terms displayed in (8) have been calculated here; they are partially cancelled by nonlocal terms of the same order (M_N^{-2}) as explained in Ref. 2. The results are not overly sensitive to the rho mass; a 2% increase in Δa follows from a reduction of $m_\rho = 776$ MeV of Table I to the PDG figure of 770 MeV. A similar 3–5 % increase in Δa is observed if one imposes a monopole form factor on the charge couplings g_ρ and g_ω according to the prescription of Ref. 22 or a smaller cutoff mass of 1400 MeV suggested in Ref. 7.

We conclude from Table I the ρ-ω mixing can account for most of the measured charge asymmetry of NN scattering lengths. Other contributions are $\lesssim 0.3$ fm and are discussed in Refs. 4 and 27. The early predictions[1] of dominance of ρ-ω mixing have been borne out by scattering length measurements[3-5] and, in the present paper, refined by contemporary and precise mixing parameters.

V. BINDING ENERGY DIFFERENCE OF ^{3}He-^{3}H NUCLEI

A perturbative estimate of the Coulomb contribution to the ^{3}He-^{3}H binding energy difference can be made directly from the experimental elastic electron scattering form factors for ^{3}He and ^{3}H, so that model-dependent nuclear wave functions are bypassed.[28] This method was extended to estimate all the direct (but local) electromagnetic (em) contributions and supplemented by model-dependent wave function estimates of the contributions from the small nonlocal terms of the em interaction and the proton-neutron mass difference.[29] Of the complete 764 keV binding energy difference all but 81±29 keV was attributed to these direct em contributions.[29] This "nearly model-independent" number then characterizes the nuclear charge asymmetry in the $A=3$ system. It is model independent in the sense that some knowledge of the three-nucleon wave function in the nuclear interior is available in the form of the em form factors of electron scattering in contrast[24] to the interior of the two-nucleon wave function which depends on models (potentials, quark models, etc.).

The "hyperspherical" formula used in these estimates

$$\langle {}^3\text{He} | U | {}^3\text{He} \rangle - \langle {}^3\text{H} | U | {}^3\text{H} \rangle = \frac{1}{2\pi^2[(3)^{1/2}]^3} \int_0^\infty q^2 dq\, 2a_S(q^2/3)a_V(q^2/3)U(q^2/3)[F_S(q^2)+F_V(q)] . \quad (12)$$

employs the isoscalar (S) and isovector (V) combinations of the $A=3$ form factors. It is most often used to estimate the static Coulomb contribution with $a_S = \frac{1}{2}(G_E^p + G_E^n)$ and $a_V = \frac{1}{2}(G_E^p - G_E^n)$ and $U(q^2) = 4\pi e^2/q^2$, but can be used with any local spin-independent potential with appropriate a_S and a_V.[29] If one neglects the small neutron Sachs form factor G_E^n then[30]

$$F_S + F_V \approx [4F({}^3\text{He}) - F({}^3\text{H})]/3G_E^p , \quad (13)$$

which displays the dependence of the formula on the individual form factors.

The $A=3$ form factors were known only up to $q^2 \approx 8$ fm^{-2} for ^{3}H at the time of the estimate of Ref. 29. This was sufficient for near convergence of the integral for the overwhelmingly dominant static Coulomb contribution. A plausible extrapolation of the ^{3}H charge form factor was made to estimate contributions of shorter range potentials of the direct em type and charge asymmetric potentials such as $\Delta V^{\rho\omega}$.

The second great advance in experimental knowledge which prompts this reappraisal of ρ-ω mixing is the new availability[31] of high quality data on the charge form factors of both ^{3}He and ^{3}H up to $q^2 \approx 50(21)$ fm^{-2}, respectively. The newly measured ^{3}H form factor differs slightly at high q^2 from the extrapolation assumed in Ref. 29. This has little effect on the static Coulomb contribution which depends mainly on the low q^2 data. If, however, the interaction M in (12) is a short range Yukawa corresponding to vector meson exchange, the value of the integral decreases by nearly a factor of 2 as the measured ^{3}H form factor is substituted for the extrapolation assumed in Ref. 29. Thus a recalculation is in order.

The total direct em contributions to the binding energy difference yield $693 \pm 19 \pm 5$ keV so that the measure of charge asymmetry in the $A=3$ mirror system is the nonzero

$$\Delta E = 80 \pm 19 \pm 5 \text{ keV} , \qquad (14)$$

where the first error is due to the error envelope on the measured charge form factors[32] of Eq. (12) and the second is an educated guess at the shortcomings[33] of (12) and the model dependence of the meson-exchange corrections to the measured form factors. For $F(^3\text{He})$ and $F(^3\text{H})$ we used the current parametrization of the world data[32] and chose the nucleon form factors G_E^p and G_E^n of Höhler *et al.*[21] because of their attention to normalization problems in the data. The experimental charge form factors include both the one-body quantities expected in (12) and meson exchange current (MEC) contributions which are model dependent. We have subtracted two calculations[34,35] of the dominant pair diagram (in pseudoscalar πNN coupling[36]) from the experimental data and found that the MEC corrections to the static Coulomb contribution are small ($\sim 3\%$ increase) and differ by about 2 keV in the range of q^2 in which the MEC calculations can be compared. For our purposes, the measure of charge asymmetry in (14) is adequate, although a reanalysis of the ^{3}H data (weighted by a factor of $\frac{1}{4}$ to the ^{3}He data [Eq. (13)])may modify the final extraction by 1 or 2 %.[37]

Next, we turn to estimates of the contribution

$$\langle {}^3\text{He} \mid V^{\rho\omega} \mid {}^3\text{He} \rangle - \langle {}^3\text{H} \mid V^{\rho\omega} \mid {}^3\text{H} \rangle = -\langle {}^3\text{He} \mid \Delta V^{\rho\omega} \mid {}^3\text{He} \rangle ,$$

of ρ-ω mixing to the $A=3$ binding energy difference. The $^1S_0$ state comprises 90% of the two-body $T=1$ component of the trinucleon so that a good estimate of this contribution can be made by taking expectation values of the $^1S_0$ form (8) of ΔV. The results are shown in Table II for model wave functions[38] obtained by solving the Faddeev equations with the dTRS and Reid soft core charge symmetric potentials. Because the estimate is perturbative, one may simply scale the second order local relativistic corrections of line 2 of Table II by the value of β desired.

TABLE II. Contribution of $\Delta V^{\rho\omega}$ to the binding energy difference between ^{3}He and ^{3}H (in keV). The "experimental" extraction is 80 ± 24 keV [Eq. (14) of text].

	"Model free"[a]	dTRS[b]	Reid[c]
NR ($\beta=0$)	92 ± 12	78	58
Local ($\beta=0.637$)	-3 ± 2	-8	-11
Total ($\kappa_V=3.7$)	89 ± 14	70	47
$\langle r^2 \rangle_{\text{He}}^{1/2}$	1.77[d]	1.83[e]	1.86[e]

[a]From charge form factor data corrected for meson exchange currents.
[b]From a trinucleon wave function evaluated with dTRS potential (Ref. 38).
[c]From a trinucleon wave function evaluated with Reid potential (Ref. 38).
[d]Experimental rms radius of ^{3}He in fm for point nucleons.
[e]Model point rms radii of the two trinucleon wave functions (Ref. 38).

These trinucleon wave functions are, however, unable to account simultaneously for the total binding energy and for the em form factors of the nuclei.[39] In a model nucleus as binding decreases, the system swells, and the short-ranged $\Delta V^{\rho\omega}$ contribution will decrease from its value in the physical nucleus. This "scaling" is illustrated by the bottom two rows of Table II. It has been suggested[29] that the model-independent method of Ref. 28 can be used to estimate the binding energy difference of ^{3}He and ^{3}H due to any spin-independent local interaction. Because the hyperspherical approximation to the two-body correlation function lacks the hole present at short distances in the true correlation function, the expectation values obtained from (12) are larger than from model wave functions corresponding to the same charge form factor. For the long ranged Coulomb potential this overestimate is about 1%;[33] it is presumably greater for short ranged potentials. The reliability of the "model-independent" estimate furnished by (12) for short-ranged potentials has not been tested as thoroughly as it has been for the Coulomb potential,[29,33] but at least the size of the nuclear system is correct. Replacing the full potential by its $^1S_0$ part we find $\Delta V^{\rho\omega}$ contributes $\sim 89 \pm 14$ keV. The model-dependent MEC contributions also affect this result. If they are not subtracted from the experimental charge from factors the contribution of $\Delta V^{\rho\omega}$ is $\sim 49 \pm 14$ keV. The non-negligible contribution of $\pi\eta\eta'$ mixing is even more strongly affected by MEC and will be discussed in another paper. The new ^{3}H data at higher q^2 does, however, lend confidence in the "model-independent" estimate if Eq. (12) and MEC models are indeed under control for short range potentials.

VI. CONCLUSIONS

We have shown that the charge asymmetric potential due to ρ-ω mixing is about 140% stronger than previous estimates. The input to the calculation was a new measurement of the G-parity forbidden decay $\omega \rightarrow 2\pi$ which has extremely low statistical errors and an inherently clean interpretation. The effect of this source of charge asymmetry on the scattering length difference $\Delta a = |a_{nn}| - |a_{pp}|$ is $\Delta a \sim 1$ fm in good agreement with the experimental findings of $\Delta a \sim 1.4 \pm 0.8$ fm. Its contribution to the binding energy differences in the $A = 3$ system is also in agreement with experiment.

ACKNOWLEDGMENTS

This research was supported in part by National Science Foundation Grant PHY 86-0636801 and by a visiting fellowship for S.A.C. from the Science and Engineering Research Council. We would like to acknowledge a useful conversation with Ernest Henley on the presentation of these results. Granddon Yen supplied us with the code for the nucleon electromagnetic form factors.

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