

Charge-Symmetry-Breaking Potentials from Isospin-Violating Meson-Baryon Coupling Constants

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We consider charge-symmetry violations in the nucleon-nucleon force that result from isospin-violating meson-baryon coupling constants. The isospin violations arise from the difference in the up and down constituent quark masses. We show that a class IV charge-symmetry-breaking potential exists in the resulting ω exchange contribution to the NN force. The magnitude of this contribution is consistent with that phenomenologically required by the measured difference of n and p analyzing powers in elastic $\vec{n}-\vec{p}$ scattering at 183 MeV.

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Isospin violation manifests itself in a number of hadronic and nuclear observables. The scattering length differences of the pp and nn systems [1], the binding energy differences of mirror nuclei [2,3], and the n and p analyzing power difference in elastic $\vec{n}-\vec{p}$ scattering [4–6] are all examples. These effects presumably originate in the differing mass and electromagnetic interactions of the up and down quarks. Consequently, the confrontation of theoretical calculations of isospin-violating observables with experiment is of continuing interest, as it potentially grants us new insight into hadronic structure and offers constraints on phenomenological models of QCD.

Our focus will be on charge-symmetry-breaking (CSB) in the np system, which is generated by the so-called class IV CSB potentials [7]. In contrast, the class III CSB potentials respect isospin in the np system, but distinguish pp from nn systems. There have been several calculations of these potentials in the context of meson exchange models [8–11]. In such a picture, three distinct CSB contributions to the NN force exist. CSB contributions can arise from (i) isovector-isoscalar mixing in the meson propagators, (ii) isospin breaking in the meson-nucleon coupling constants, and (iii) isospin breaking in the nucleon wave function. In addition, there are electromagnetic contributions, such as the photon's coupling to the neutron's anomalous magnetic moment. In principle, all these effects contribute to CSB observables; one wishes to combine them in a dynamical model. Several different sources of CSB contribute to the nonzero analyzing power difference $\Delta A \equiv A_n - A_p$ measured in polarized, elastic $n-p$ scattering. Those studied so far include the exchange of charged pions and rhos, the photon's coupling to the neutron's anomalous magnetic moment, and $\rho-\omega$ mixing. The last is large because of the small mass difference between the ρ and ω ; the exchanged ρ can convert into an ω . This mixing is clearly seen in $e^+e^- \rightarrow \pi^+\pi^-$ cross section measurements at the ω production point [12]. The $\rho-\omega$ mixing amplitude that fits the e^+e^- data also ex-

plains the ΔA measurement at 183 MeV [4] and accounts for a large fraction of the binding energy difference seen in the $A = 3$ systems [2]. However, it has been suggested that the $\rho-\omega$ mixing amplitude depends on the momentum transfer q [13]. Indeed, several authors argue that the q^2 dependence is large and that the resulting isospin-violating potential is small at the spacelike momentum transfers relevant for CSB experiments [14–19]. The issue continues to be controversial [20–22].

If the $\rho-\omega$ mixing potential is, in fact, small, then the CSB contributions discussed so far no longer suffice to fit the data [23]. Yet other sources of isospin violation could well exist. Indeed, one ought to consider isospin violation arising from the nucleon's intrinsic wave function as well as from the vector-meson-nucleon coupling constants. These are sources of additional isospin violation; they deserve examination regardless of the q^2 dependence of the $\rho-\omega$ mixing amplitude.

In this paper we focus on isospin violation in the vector meson couplings to the nucleon. As most CSB studies have focused on the role of mechanisms (i), that is, $\rho-\omega$ mixing, and (iii)—through the sensitivity of charged pion and rho exchange to the nucleon mass difference—discussed above, we study (ii) exclusively, as we wish to understand its impact. We assume that isospin violation in the nucleon's internal wave function, while undoubtedly nonzero, is negligibly small due to the large mass difference between the nucleon and the $\Delta(1910)$ —the first P_{31} baryon. The $\rho-\omega$ mass difference, in contrast, is a mere 12 MeV. There is no argument, however, that protects the isospin symmetry of the vector-meson-nucleon coupling constants. In the following we examine the isospin violation arising from the mass difference of the up and down quarks. Electromagnetic radiative corrections have been estimated earlier [24].

Dmitrašinović and Pollock have studied the isospin-violating electroweak form factors of the nucleon in a simple constituent quark model [25]. These are poten-

tially important for interpreting parity-violating electron-nucleon scattering in terms of the nucleon's strange quark content, as the Z^0 coupling is sensitive to isospin violation. They find that the isoscalar quark current $\bar{u}\gamma_\mu u + \bar{d}\gamma_\mu d$ has a larger matrix element in the proton than in the neutron as the up quark has a larger magnetic moment than a down quark. The size of the violation, which is about 1%, is set by the ratio of the difference in up and down constituent quark masses to their average mass.

Henley and Zhang have calculated the isospin dependence of the vector-meson-nucleon couplings in a constituent quark model, through explicit calculation of the quark model wave function overlaps [26]. Our results for the isospin-violating couplings are very similar to theirs.

Two assumptions define our model. First, the vector mesons are assumed to couple to the appropriate isospin components of the nucleon's electromagnetic current. This assumption is in the spirit of the vector meson dominance model. Second, we assume this current can be estimated—at low q^2 —in a nonrelativistic, constituent quark model. Such models give good descriptions of the nucleon magnetic moments.

Perhaps the simplest way to realize these assumptions is in a hybrid quark-meson model, in which the mesons couple directly to the quarks. However, this picture is not required. A model with composite vector mesons can still satisfy our assumptions.

In our model the vector mesons couple to the nucleon's electromagnetic current, which we decompose into isoscalar and isovector quark components, appropriate for the coupling of the ω and ρ , respectively, to the nucleon. In the quark model, the isoscalar electromagnetic charge of the up and down quarks is $e_i^{(0)} = 1/3$, whereas the isovector electromagnetic charge of the up quark is $e_u^{(1)} = 1$ and that of the down quark is $e_d^{(1)} = -1$. The vector quark current is

$$J^\mu = e_u \bar{u}\gamma^\mu u + e_d \bar{d}\gamma^\mu d; \quad (1)$$

the constituent quarks are assumed elementary. We are interested in computing the vector coupling g_N^V and the tensor coupling f_N^V of the nucleon to the vector mesons ρ and ω . That is,

$$\begin{aligned} \langle N(p', s') | J_V^\mu(q) | N(p, s) \rangle = \\ \bar{U}(p', s') \left[g_N^V \gamma^\mu + i f_N^V \sigma^{\mu\nu} \frac{(p' - p)_\nu}{2M_N} \right] U(p, s). \end{aligned} \quad (2)$$

Note that $U(p, s)$ denotes an on-shell nucleon spinor of mass M_N , momentum p , and spin s . The couplings g_N^V and f_N^V are functions of the four-momenta at the vertex, here $g_N^V(q^2)$ and $f_N^V(q^2)$ ($q \equiv p' - p$), though we presume the couplings constant in our region of interest. We compute the couplings at $q^2 = 0$, as the nonrelativistic quark model is best suited to an estimate in the static limit. For low-energy scattering experiments, such as the 183 MeV \bar{n} - \vec{p} analyzing power measurement [4], this limit should be

reasonable. We obtain g_N^V and f_N^V by examining the nonrelativistic reduction of Eq. (2) and then computing the matrix elements of the resulting operators in the quark model. Thus, we evaluate

$$\frac{g_N^V}{\bar{g}^V} = \sum_{i=1}^3 \langle N \uparrow | e_i^{(\tau)} | N \uparrow \rangle = \sum_{i=1}^3 e_i^{(\tau)}, \quad (3a)$$

$$\frac{g_N^V + f_N^V}{2M_N \bar{g}^V} = \sum_{i=1}^3 \langle N \uparrow | \frac{e_i^{(\tau)}}{2m_i} \sigma_i^z | N \uparrow \rangle \quad (3b)$$

in the nucleon rest frame. One sums over the charges and magnetic moments of the quark i . The symbol $e_i^{(\tau)}$ denotes the appropriate isospin component of the electromagnetic charge of the quarks; the ω -nucleon coupling, for example, is determined by the isoscalar quark charge. The couplings g_N^V and f_N^V are written explicitly in units of \bar{g}^V , the isospin-averaged vector coupling of the vector mesons to the nucleon. Note that \bar{g}^ω and \bar{g}^ρ are known from fits to NN scattering and to the properties of the deuteron [27,28]. The ket $|N \uparrow\rangle$ denotes a nucleon state with spin up. We are interested in evaluating the isospin-violating contributions to g_N^V and f_N^V and use the full SU(6) wave function for the nucleon [29]. In this limit, the magnetic moments are independent of the spatial distribution of the wave function, so that they follow immediately from the spin structure of the nucleon. Note that g_N^V , in contrast, depends only on the nucleon's flavor structure. Consequently, the calculation of the vector and tensor couplings proceeds straightforwardly. The difference between the up and down quark masses can generate isospin violations in the meson-baryon couplings. Introducing

$$m \equiv \frac{1}{2}(m_d + m_u), \quad \Delta m \equiv (m_d - m_u), \quad (4)$$

Eq. (3) implies that

$$g_N^\omega = \bar{g}^\omega, \quad g_N^\rho = \bar{g}^\rho \quad (5)$$

and, defining

$$\frac{f_N^V}{2M_N} \equiv \frac{f_{(0)}^V + f_{(1)}^V \tau_z}{2M}, \quad (6a)$$

that

$$f_{(0)}^\omega = 0, \quad f_{(1)}^\omega = \frac{5}{6} \frac{\Delta m}{m} \bar{g}^\omega, \quad (6b)$$

$$f_{(0)}^\rho = \frac{3}{2} \frac{\Delta m}{m} \bar{g}^\rho, \quad f_{(1)}^\rho = 4\bar{g}^\rho, \quad (6c)$$

where $M = (M_n + M_p)/2$ denotes the mean nucleon mass—the isospin breaking we compute includes the effect of the neutron-proton mass difference. We have chosen $m = M/3 = 313$ MeV in Eqs. (6b) and (6c). We adopt this choice throughout the paper. Note that τ_z acts at the hadronic level, so that $\tau_z |p\rangle = +|p\rangle$ and so on. The isospin-breaking corrections contribute to the tensor couplings exclusively—the vector couplings are unchanged. Moreover, these corrections are *isovector* for

the ω coupling, and are *isoscalar* for the ρ coupling. Thus, their appearance simulates ρ - ω mixing; this will become explicit when we discuss the resulting CSB potentials. Note that $\Delta m > 0$ in the constituent quark model [30]; the up quark, which is lighter, generates a larger anomalous magnetic moment for the proton.

Before discussing the isospin-breaking corrections in detail, let us consider the isospin-symmetric results for g_N^V and f_N^V . That is,

$$\frac{f_{(0)}^\omega}{g_N^\omega} = 0, \quad \frac{f_{(1)}^\rho}{g_N^\rho} = 4. \quad (7)$$

These nonrelativistic quark model (NRQM) results are qualitatively consistent with the f_N^V/g_N^V ratios that emerge from phenomenological fits to the NN interaction [27,28]—recall that the Bonn B potential parameters [28], for example, are $f_N^\omega/g_N^\omega = 0$ and $f_N^\rho/g_N^\rho = 6.1$. These successes are intimately connected to the NRQM's ability to describe the nucleon magnetic moments. In the above model, the anomalous magnetic moment is purely isovector: $\kappa_N = 2\tau_z$. Note that $\kappa_p^{\text{exp}} = 1.79$ and $\kappa_n^{\text{exp}} = -1.91$. The above successes encourage us to use the NRQM to compute the isospin-violating corrections to these coupling constants as well. These corrections are given in Eq. (6).

Henley and Zhang have also examined the impact of quark mass difference effects on the vector-meson-nucleon coupling constants [26]. They adopt an “effective perturbative QCD model”: they calculate the nucleon-nucleon-meson vertex in terms of nonrelativistic, constituent quarks and connect the produced $q\bar{q}$ pair with the other quarks

$$\begin{aligned} V_{\text{IV}}^{\rho+\omega}(q^2 \rightarrow 0) &= \left(\frac{5}{6} \frac{\bar{g}^{\omega^2}}{m_\omega^2} - \frac{3}{2} \frac{\bar{g}^{\rho^2}}{m_\rho^2} \right) \left(\frac{\Delta m}{m} \right) \frac{i[\vec{\sigma}(1) - \vec{\sigma}(2)] \cdot \vec{q} \times \vec{P}}{4M^2} [\tau_z(1) - \tau_z(2)] \\ &\equiv C(q^2 = 0) \frac{i[\vec{\sigma}(1) - \vec{\sigma}(2)] \cdot \vec{q} \times \vec{P}}{4M^2} [\tau_z(1) - \tau_z(2)], \end{aligned} \quad (10)$$

where $\vec{P} = \vec{p}_1' + \vec{p}_1$. This expression sums the contributions from single ρ and ω exchanges. Let us compare the strength of the CSB potentials given in Eqs. (9) and (10). The $C(q^2 = 0)$ of Eq. (10), in terms of the Bonn B potential parameters [$\bar{g}_\omega^2(q^2 = 0)/4\pi = 11.13$; $\bar{g}_\rho^2(q^2 = 0)/4\pi = 0.42$] [28], is

$$\begin{aligned} C(q^2 = 0) &= 1.77 \times 10^2 \frac{\Delta m}{m} \text{ GeV}^{-2} \\ &\approx 2.32 \text{ GeV}^{-2}. \end{aligned} \quad (11)$$

Note that the sign of Eq. (11) is determined by the ω contribution—in the Bonn model $\bar{g}^{\omega^2}/\bar{g}^{\rho^2} \approx 27$. The last estimate for C results when one uses the “lower bound” of Δm , $\Delta m = 4.1$ MeV, of Lichtenberg [30]. The strength of the class IV ρ - ω mixing potential in Eq. (9) at $q^2 = 0$, on the other hand, is

$$\begin{aligned} C_{\rho-\omega}^{\text{on-shell}}(q^2 = 0) &= -\frac{\bar{g}^\omega f_{(1)}^\rho}{m_\rho^2 m_\omega^2} \langle \rho | H | \omega \rangle |_{q^2 = m_\omega^2} \\ &\approx 2.07 \text{ GeV}^{-2}, \end{aligned} \quad (12)$$

via perturbative one-gluon exchange. We are able to reproduce the isospin breaking they compute in the vector-meson-nucleon coupling constants. The isospin breaking of their model can apparently be generated on rather general grounds.

We shall now compute the CSB potentials that arise from the isospin-violating couplings in Eq. (6). In a one boson exchange approximation, we obtain the following CSB potentials for ω and ρ exchanges:

$$V_{\text{CSB}}^\omega = -\left(\frac{f_{(1)}^\omega \bar{g}^\omega}{q^2 - m_\omega^2} \right) \hat{V}(1,2), \quad (8a)$$

$$V_{\text{CSB}}^\rho = -\left(\frac{f_{(0)}^\rho \bar{g}^\rho}{q^2 - m_\rho^2} \right) \hat{V}'(1,2), \quad (8b)$$

where $\hat{V}(1,2) = \Gamma^\mu(1)\gamma_\mu(2)\tau_z(1) - \gamma^\mu(1)\Gamma_\mu(2)\tau_z(2)$ and $\Gamma^\mu = i\sigma^{\mu\nu}q_\nu/2M$ with $q = p_1' - p_1$. Note that $\hat{V}'(1,2)$ is of the form of $\hat{V}(1,2)$ with the exchange $\tau_z(1) \leftrightarrow \tau_z(2)$. The above are identical in form to the CSB potential from ρ - ω mixing. That is,

$$V_{\text{CSB}}^{\rho-\omega} = -\frac{\bar{g}^\omega f_{(1)}^\rho \langle \rho | H | \omega \rangle}{(q^2 - m_\rho^2)(q^2 - m_\omega^2)} \hat{V}(1,2). \quad (9)$$

In the case of ρ exchange, the couplings of Eq. (6) also generate a class III CSB potential. The complete nonrelativistic reduction of the potentials of Eq. (8) can be found in Ref. [7]. Extracting the class IV pieces of Eq. (8) and combining them in the nonrelativistic limit yields, in the center-of-mass frame, the class IV potential [7]

where we have used the *on-shell* value of the ρ - ω mixing matrix element, $\langle \rho | H | \omega \rangle = -4520 \pm 600 \text{ MeV}^2$ [12], for purposes of comparison. We have also used the Bonn B value for $f_{(1)}^\rho$, $f_{(1)}^\rho = 6.1\bar{g}^\rho$. A potential of the magnitude of Eq. (12) is needed for a successful description of the 183 MeV ΔA data [4]. Thus, isospin violation in the meson-baryon coupling constants suffices *alone* to generate the qualitative magnitude of the phenomenologically required class IV CSB potential. Note that the additional propagator in Eq. (9) implies that its q^2 dependence will likely differ from that of Eq. (8); however, the magnitude of ΔA at 183 MeV is controlled by the $q^2 = 0$ potential [31].

Here we have focused on the class IV CSB potential that arises from isospin violations in the vector-meson-nucleon coupling constants. The resulting class IV CSB potential is identical in structure to that which arises from ρ - ω mixing. Moreover, its magnitude is commensurate in size with that phenomenologically required to explain the Indiana University Cyclotron Facility ΔA measure-

ment [4]. If the ρ - ω mixing amplitude is q^2 dependent and the isospin-violating potential small for the space-like momentum transfers relevant to the above experiment [13–19], then we have found a source of isospin violation that can fill the role demanded by the data. If the ρ - ω mixing amplitude is not q^2 dependent [20], then the total CSB potential is probably too large to fit the data.

The isospin breaking we compute in the vector-meson-nucleon coupling constants at $q^2 = 0$ arises on rather general grounds. We assume that the vector mesons couple to the appropriate isospin components of the electromagnetic current; we compute these components of the current in the nonrelativistic constituent quark model. The magnitude of the isospin breaking we predict depends numerically on only \bar{g}_ω and $\Delta m/m$. The results we obtain do not depend on the details of the nucleon's structure in the NRQM. Indeed, our results depend merely on the manifest spin and flavor structure of the nucleon in the SU(6) limit. Consequently, we believe our estimate to have little model dependence. This is why we reproduce the isospin breaking of Henley and Zhang's more complicated quark model [26]. The isospin breaking we predict could have a nontrivial q^2 dependence. This requires a detailed model calculation beyond the scope of our present approach.

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