Axial vector current and coupling of the quark in the instanton model

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We compute the axial vector current, form factor, and coupling for quarks in the instanton liquid model with two light flavors. Non-local current corrections are derived, as required by the effective 't Hooft interaction. We obtain a pion-mediated axial form factor and an axial vector coupling which, when simply applied in the non-relativistic limit for constituent quarks, matches the experimental value to within a few percent, both in and out of the chiral limit.

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I. INTRODUCTION

The axial vector coupling constant of the nucleon, deduced from measurements of neutron beta decay, is known to be $|G_A/G_V| = 1.2670 \pm 0.0035$ [1] in vacuum. Strict chiral invariance would have this ratio be unity, and the substantial difference suggests a deviation in excess of current quark mass effects. Thus it has long been argued that the increase in G_A is a result of the spontaneous chiral symmetry breaking that characterizes the QCD vacuum.

Although additional chiral-invariant meson-nucleon couplings can remedy this problem in the linear sigma model [2], a more fundamental explanation is warranted. The issue is complicated by the difficulties experienced by lattice QCD practitioners [3], in that such calculations tend to yield values significantly lower than experiment. Possible explanations have recently centered on finite-size effects [4], although the specific physics is unresolved [5].

One might also try to deduce the nucleon's axial vector coupling by studying that of a constituent quark, denoted in this paper as g_A. In the non-relativistic limit,

$$G_A = \frac{5}{3} g_A, \tag{1}$$

but incomplete knowledge of the quark wave function leaves the relativistic corrections unknown. However, there has been considerable success in treating the nucleon as massive constituent quarks [6], and if the deviation of G_A from unity is in fact due to chiral forces the constituent approach should encompass the relevant physics. At the quark level chiral symmetry also demands an axial vector coupling of unity, which generates an excess of nucleon axial vector charge when Eq. (1) is directly applied. Studies addressing this predict negative corrections of $\mathcal{O}(1)$ to g_A using the MIT bag model [7] in the large-N_c limit [8,9] or via perturbative corrections [10].

In this paper we consider corrections to g_A from instantons. The instanton liquid model (ILM), in which the QCD vacuum is populated by classical gauge configurations [11,12], dynamically generates the constituent quark mass and has consistently led to phenomenological success when applied to problems related to spontaneous chiral symmetry breaking [13,14]. As detailed below, we find a 25% reduction in the axial vector coupling from instantons, with only slight sensitivity to the chiral limit. This is in excellent agreement with experiment if corrections to Eq. (1) are small.

To determine the axial vector coupling we will first carefully analyze the axial vector current. Since instanton-induced interactions between quarks are non-local, the currents associated with conserved quantities are non-trivial. Next, given the necessity of pions in axial vector phenomena, meson correlation functions are briefly reviewed. We then derive the axial vector current and verify the importance of the pion pole in the axial vector form factor, $h_A(q^2)$. Finally, the leading corrections to the coupling g_A are computed, taking into account multi-instanton effects and substantial non-local vertex corrections. In the chiral limit we find $g_A = 0.77$ and estimate that with a physical pion mass $g_A = 0.80$. With the naive application of Eq. (1), these correspond to a nucleon axial vector couping of 1.28 and 1.34, respectively.

II. NON-LOCALITY AND CONSERVED CURRENTS

The effects of an instanton liquid background are most economically encoded in an effective quark action [15]. Taking the exact zero-mode solutions to the Dirac equation and isolating them as the dominant low-energy effects leads to a $2N_f$-quark vertex function [16] similar to the Nambu–Jona-Lasinio model of QCD [17–19]. A well-studied simplification of this model, the chiral random matrix theory, has shown that chiral symmetry breaking via random overlap integrals reproduces Dirac eigenvalue correlations as computed on the lattice [20]. Thus such an approach seems to indeed contain the essence of non-perturbative QCD. Furthermore, the ILM relies on only two parameters: the instanton density and the average instanton size. Both of these were phenomenologically fixed long ago with the vacuum energy density and chiral condensate [21,15] and measured on the lattice [22]. These parameters determine the diluteness of the instanton liquid, the ratio of the average size to inter-instanton spacing, to be about 1/3, a somewhat small parameter which allows for perturbative treatment of the instanton-induced vertex.

The Eulcidean effective action for N_f flavors in the chiral limit is written [14]

$$S=-\int d^4x\,\psi^\dagger(x)i\partial\!\!\!/\psi(x)+i\lambda\int dU\,d^4z \times\prod_f^{N_f}\left[-d^4x_f\,d^4y_f\,\psi_f^\dagger(x_f)i\partial\!\!\!/\Phi(x_f-z,U)\times\tilde{\Phi}(y_f-z,U)i\partial\!\!\!/\psi^f(y_f)\right], \tag{2}$$

where U is the instanton's $2\times N_c$ color/spin orientation matrix and z is its position. The $\Phi(x)$ is the zero mode solution for fermions in the field of one instanton, and its Fourier transform is the form factor

$$f(p)=2x\left[I_1(x)K_0(x)-I_0(x)K_1(x)+\frac{1}{x}I_1(x)K_1(x)\right]_{x=p\rho/2}, \tag{3}$$

which provides a natural momentum cutoff of scale ρ^{-1}. We have left out additional (quarkless) λ terms which relate this coupling constant, which is in fact a Lagrange multiplier, to the instanton density, N/V.

Because of non-locality Eq. (2) is not invariant under symmetry transformations. Since bulk observables are often insensitive to this problem it is usually ignored; however, when addressing currents the dependence is crucial. Literature exists in which non-local interactions are modified to be invariant [23–29] and we follow the same procedure here. Taking the non-local four-fermion interaction as the starting point, it then becomes a matter of multiplying each quark operator by a path-ordered exponential in the background of a source gauge field a_μ, replacing the Euclidean quark fields as

$$\psi(x)\rightarrow W(z,x)\psi(x),$$

$$\psi^\dagger(x)\rightarrow\psi^\dagger(x)\gamma_0W(x,z)\gamma_0,$$

with the operator [23]

$$W(x,y)=\mathcal{P}\exp\left(-i\int_x^y ds_\mu a_\mu^a T^a\right).$$

As has been pointed out in the cited works, the operators $W(x,y)$ must be path ordered and the choice of path integrated over is not unique. However, we are primarily concerned with the longitudinal currents, for which results are independent of any particular choice of path [27].

Thus we write the modified action as

$$S=-\int d^4x\,\psi^\dagger(x)i\partial\!\!\!/\psi(x)+i\lambda\int dU\,d^4z \times\prod_f^{N_f}\left[-d^4x_fd^4y_f\,\psi_f^\dagger(x_f)\gamma_0W(x_f,z)\gamma_0 i\partial\!\!\!/\Phi(x_f-z)\times\Phi^\dagger(y_f-z)i\partial\!\!\!/W(z,y_f)\psi^f(y_f)\right]. \tag{4}$$

We now concentrate on the case of two light quarks, $N_f=2$.

The Noether currents can be evaluated directly from the action. In momentum space,

$$j_\mu^a(q)=-\left.\frac{\delta S}{\delta a_\mu^a(q)}\right|_{a_\mu^a=0}.$$

We thus have, after Fourier transforming the fields and averaging over colors,

$$j_\mu(q)=\int\frac{d^4p}{(2\pi)^4}\psi^\dagger(p)\gamma_\mu T\psi(p+q)+j_\mu^L(q)+j_\mu^R(q), \tag{5}$$

where the left-handed non-local term is

$$j_\mu^L(q)=\frac{\lambda(2\pi\rho)^4}{N_c^2-1}\epsilon^{f_1f_2}\epsilon_{g_1g_2}\left(\delta_{j_1}^{i_1}\delta_{j_2}^{i_2}-\frac{1}{N_c}\delta_{j_2}^{i_1}\delta_{j_1}^{i_2}\right)\int d^4z\,d^4x\frac{d^4p_1}{(2\pi)^4}\frac{d^4p_2}{(2\pi)^4}\frac{d^4k_1}{(2\pi)^4}\frac{d^4k_2}{(2\pi)^4}\frac{d^4p'}{(2\pi)^4}$$
$$\times\int_x^z ds_\mu e^{-iq\cdot s}\{e^{-ix\cdot(p_1-p')-iz\cdot(p'+p_2-k_1-k_2)}[\psi_L^\dagger(p_1)\gamma_0T\gamma_0]_{f_1i_1}\psi_L^{g_1j_1}(k_1)f(p')f(k_1)$$
$$-e^{ix\cdot(k_1-p')-iz\cdot(p_1+p_2-p'-k_2)}\psi_{Lf_1i_1}^\dagger(p_1)[T\psi_L(k_1)]^{g_1j_1}f(p_1)f(p')\}\psi_{Lf_2i_2}^\dagger(p_2)\psi_L^{g_2j_2}(k_2)f(p_2)f(k_2), \tag{6}$$

and the right-handed $j_\mu^R(q)$ is similarly defined. Group indices have been suppressed; note that these are carried by the operators T. The instanton or anti-instanton at the core of each vertex naturally splits the quark fields into right- or left-handed spinors, $\psi_{R/L}=\frac{1}{2}(1\pm\gamma_5)\psi$. Since we are only concerned here with spin structures of the identity matrix or γ_5, the decomposition is trivial. Carrying out the spatial integrations, we have

$$j_\mu^L(q)=\frac{q_\mu}{q^2}\frac{i\lambda(2\pi\rho)^4}{N_c^2-1}\epsilon^{f_1f_2}\epsilon_{g_1g_2}\left(\delta_{j_1}^{i_1}\delta_{j_2}^{i_2}-\frac{1}{N_c}\delta_{j_2}^{i_1}\delta_{j_1}^{i_2}\right)\int\frac{d^4p_1}{(2\pi)^4}\frac{d^4p_2}{(2\pi)^4}\frac{d^4k_1}{(2\pi)^4}\frac{d^4k_2}{(2\pi)^4}\delta(p_1+p_2-k_1-k_2+q)$$
$$\times\{[\psi_L^\dagger(p_1)\gamma_0T\gamma_0]_{f_1i_1}\psi_L^{g_1j_1}(k_1)[f(p_1)-f(p_1+q)]f(k_1)-\psi_{Lf_1i_1}^\dagger(p_1)[T\psi_L(k_1)]^{g_1j_1}$$
$$\times f(p_1)[f(k_1)-f(k_1-q)]\}\psi_{Lf_2i_2}^\dagger(p_2)\psi_L^{g_2j_2}(k_2)f(p_2)f(k_2). \tag{7}$$

In the event of spontaneous symmetry breaking, quark operators which overlap with the condensate can be paired into loops [26,30]. When chiral symmetry is broken an effective quark mass is generated, dependent on four-momentum $p=\sqrt{p_\mu p_\mu}$ as

$$M(p)=Mf(p)^2.$$

The constant M is the solution to a gap equation obtained by minimization with respect to the coupling λ [31],

$$\frac{N}{V}=4N_c\int\frac{d^4p}{(2\pi)^4}\frac{M(p)^2}{p^2+M(p)^2}, \tag{8}$$

and the constant $N/V=(200\text{ MeV})^4$ is the instanton density. With the phenomenological $\rho=1/3$ fm, one finds $M=345$ MeV. In this case the current, Eq. (7), reduces to

$$j_\mu^L(q)=-i\frac{q_\mu}{q^2}\int\frac{d^4p}{(2\pi)^4}\psi_L^\dagger(p)\Bigg\{M(p+q)\gamma_0T\gamma_0-M(p)T+\frac{\lambda(2\pi\rho)^4}{N_c}f(p)f(p+q)$$
$$\times\int\frac{d^4k}{(2\pi)^4}f(k)\frac{M(k)}{k^2+M(k)^2}\mathrm{Tr}\left[f(k+q)\frac{1-\gamma_5}{2}\gamma_0T\gamma_0-f(k-q)T\frac{1-\gamma_5}{2}\right]\Bigg\}\psi_L(p+q), \tag{9}$$

where the trace is over spin only.

Further reduction of this still-unwieldy expression requires specifying a particular current and associated group elements T. The simplest example is the quark number current, B_μ, where $T=\mathbb{1}$:

$$B_\mu(q)=\int\frac{d^4p}{(2\pi)^4}\psi^\dagger(p)\left\{\gamma_\mu-i\frac{q_\mu}{q^2}[M(p+q)-M(p)]\right\}$$
$$\times\psi(p+q). \tag{10}$$

As $q\to0$, this becomes [28,32]

$$B_\mu(0)=\int\frac{d^4p}{(2\pi)^4}\psi^\dagger(p)\left(\gamma_\mu-i\frac{dM(p)}{dp_\mu}\right)\psi(p).$$

It is easy now to see that the q-dependent term in Eq. (10), a result of the non-locality of the interaction, is necessary for current conservation. With Eq. (10) and the Dirac equation,

$$[\not p-iM(p)]\psi(p)=0, \tag{11}$$

we see the standard cancellation ensues and

$$q_\mu B_\mu(q)=0. \tag{12}$$

This is the simplest illustration of the more general rule, which becomes more complicated when the axial isovector current is considered. But before analyzing this, we review the pion's derivation in the ILM.

III. MESON CORRELATION FUNCTIONS

The pion's role in axial vector charge currents has been long established. To be made explicit in the ILM, one must sum all ladder diagrams using the 't Hooft vertex, a series shown in Fig. 1. This was originally done in Ref. [33], although the advent of the effective action (4) has since made this procedure simpler.

In the chiral limit, one finds the expected pole in the pion propagator as $q^2\to0$. Furthermore, the use of non-local corrections to the axial currents has been shown to produce the required transversality in the mesonic current and properly relate the pion decay and renormalization constants [26]. Following the results of these cited works, we have the pion correlation function of Fig. 1:

$$\Pi_\pi^{AB}(q^2)=\int\frac{d^4p}{(2\pi)^4}\langle[\psi^\dagger(p)\tau^A\gamma_5\psi(p+q)][\psi^\dagger(p+q)\tau^B\gamma_5\psi(p)]\rangle$$
$$=-\frac{\delta^{AB}}{4N_c^2}\left(\frac{1}{i\lambda}-\frac{2}{N_cM^2}\int\frac{d^4p}{(2\pi)^4}\frac{M(p)M(p+q)[p\cdot(p+q)+M(p)M(p+q)]}{[p^2+M(p)^2][(p+q)^2+M(p+q)^2]}\right)^{-1}. \tag{13}$$

The gap equation (8) can be written [32]

$$\frac{1}{i\lambda}=\frac{2}{N_cM^2}\int\frac{d^4p}{(2\pi)^4}\frac{M(p)^2}{p^2+M(p)^2}, \tag{14}$$

and one finds, upon expanding in small q^2,

$$\Pi_\pi^{AB}(q^2) = -\frac{\delta^{AB}}{4N_c q^2}\left(\frac{1}{M^2}\int\frac{d^4p}{(2\pi)^4}\frac{M(p)^2 - \frac{1}{2}pM(p)M'(p) + \frac{1}{4}p^2M'(p)^2}{[p^2+M(p)^2]^2}\right)^{-1} \equiv -\frac{g_{\pi qq}^2}{q^2}\delta^{AB}. \tag{15}$$

The coupling constant is determined thusly by the compositeness condition [34] and, with the standard instanton parameters listed above, we find $g_{\pi qq}=3.78$. The pion and other channels were analyzed in a Nambu–Jona-Lasinio (NJL) model with alternative form factors by Plant and Birse, who found $g_{\pi qq}=3.44$ for the pion and the sigma parameters $g_{\sigma qq}=3.51$ and $m_\sigma=443$ MeV [27]. This required extracting the sigma pole, and instead of repeating this procedure or fitting the sigma correlation function [35] we will simply scale our pion-quark coupling in a similar manner to obtain $g_{\sigma qq}=3.86$. These two channels will be sufficient in the following analysis of the axial vector properties of the quark.

IV. AXIAL VECTOR CURRENT AND FORM FACTOR

We now address the axial isovector current. A key experimental quantity is the matrix element of the current, $A_\mu^a(p_2-p_1)$, taken with initial and final state nucleons of momenta p_1 and p_2. Here we compute the analogous quantity with quark fields:

$$\langle\psi(p_2)|A_\mu^a(q)|\psi(p_1)\rangle = u^\dagger(p_2)\frac{\tau^a}{2}[\gamma_\mu\gamma_5 g_A(q^2) + q_\mu\gamma_5 h_A(q^2)]u(p_1), \tag{16}$$

where the $u^\dagger(p_2)$ and $u(p_1)$ are the Euclidean spinor solutions of the Dirac equation for free quarks, and $q=p_2-p_1$. We are ultimately interested in the limit of $p_1=p_2=0$, or on-shell quarks at zero momentum transfer, with the constant $g_A=g_A(0)$ the axial vector coupling. First we consider $h_A(q^2)$, the axial vector form factor. With a nonzero g_A and current conservation, it is clear in the chiral limit that $h_A(q^2)$ has a pole at $q^2=0$, identified as the pion. Through explicit construction, we will show this pole is manifest in the ILM.

Inserting $T=\gamma_5\tau/2$ into Eq. (9), we have bare and single instanton (and anti-instanton) contributions in

$$A_\mu^a(q)_{2a} = \int\frac{d^4p}{(2\pi)^4}\psi^\dagger(p)\frac{\tau^a}{2}\gamma_\mu\gamma_5\psi(p+q) + i\frac{q_\mu}{q^2}\int\frac{d^4p}{(2\pi)^4}\psi^\dagger(p)\frac{\tau^a}{2}\gamma_5\Bigg\{M(p+q)+M(p) - \frac{2i\lambda(2\pi\rho)^4}{N_c}f(p)f(p+q)\int\frac{d^4k}{(2\pi)^4}f(k) \times[f(k+q)+f(k-q)]\frac{M(k)}{k^2+M(k)^2}\Bigg\}\psi(p+q). \tag{17}$$

Graphically, this is Fig. 2(a).

Clearly, no amount of labor will reveal this to be a conserved current; as the appearance of the chiral Goldstone boson is a multi-instanton effect, so too is axial vector current conservation. Pion tadpoles, as rendered in Figs. 2(b) and 2(c), are necessary to remove the finite divergence of A_μ^a. The first is a bare current insertion, the second the nonlocal piece. Their sum is

$$A_\mu^a(q)_{2b+2c} = -\frac{q_\mu}{q^2}\frac{2\lambda(2\pi\rho)^4}{N_c}\int\frac{d^4p}{(2\pi)^4}f(p)f(p+q) \times\psi^\dagger(p)\frac{\tau^a}{2}\gamma_5\psi(p+q)\int\frac{d^4k}{(2\pi)^4}f(k) \times[f(k+q)+f(k-q)]\frac{M(k)}{k^2+M(k)^2}. \tag{18}$$

This cancels the final term in Eq. (17) and the full current is

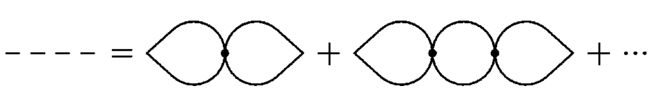

FIG. 1. The series of instanton ladder diagrams summed to obtain meson correlation functions.

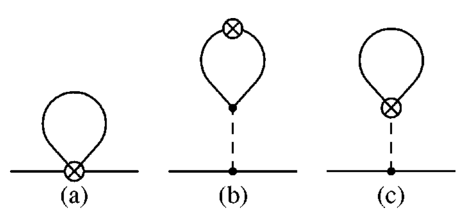

FIG. 2. Diagrams contributing to the axial vector form factor, $h_A(q^2)$. Circled crosses are current insertions and the dashed line denotes a pion.

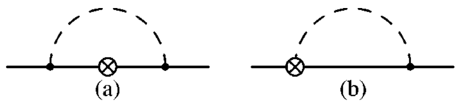

FIG. 3. Diagrams contributing to the axial vector coupling constant, g_A. Circled crosses are current insertions and dashed lines denote either a pion or sigma meson propagator.

$$A^a_\mu(q)=\int \frac{d^4p}{(2\pi)^4}\psi^\dagger(p)\left\{\gamma_\mu+\frac{q_\mu}{q^2}[iM(p)+iM(p+q)]\right\}\times\gamma_5\frac{\tau^a}{2}\psi(p+q). \tag{19}$$

The first term is the bare axial vector coupling, unity with the definition of Eq. (16), and the second is the pion pole carried by the axial vector form factor,

$$h_A(q^2)=-i\frac{M(p+q)+M(p)}{q^2}. \tag{20}$$

Note that at long wavelength the leading divergent part of the form factor is the non-local, quark level version of the classic nucleon result, $H_A(q^2)=-2iM_N/q^2$. The pion pole similarly couples separately to each constituent quark.

It is now clear, with the Dirac equation (11), that

$$q_\mu A^a_\mu(q)=0. \tag{21}$$

Equivalently, with the propagator,

$$S(p)=\frac{\not p+iM(p)}{p^2+M(p)^2},$$

the Ward-Takahashi identity is satisfied:

$$q_\mu\Gamma^a_{\mu5}=-\frac{\tau^a}{2}[\gamma_5S(p+q)^{-1}+S(p)^{-1}\gamma_5]. \tag{22}$$

V. AXIAL VECTOR COUPLING CONSTANT

The axial vector form factor involved the effects of multiple instantons, with pions transferring *t*-channel momentum. Similar propagation in the *s*-channel contributes to the axial vector coupling constant. While still dominant, the pion is not the only relevant resonance in this channel. Since the sigma's coupling constant is of comparable size, and its mass is only slightly greater than that of a constituent quark, its effects will also be taken into account. Higher resonances, such as the ρ and A_1 vectors, are both substantially heavier and have quark-meson couplings a mere third of the scalar and pseudoscalar [27,35]. With the leading contributions of order g^2_{iqq}, these are safely ignored.

In Eq. (19) we have seen that the bare g_A is unity. To this we first add the pion contribution of Fig. 3(a), which is

$$g_A^{3a,\pi}=-\frac{g^2_{\pi qq}}{2M}\int\frac{d^4p}{(2\pi)^4}\frac{M(p)}{p^2}\frac{-p^2+2M(p)^2}{[p^2+M(p)^2]^2}. \tag{23}$$

After evaluation, the identical sigma contribution differs only by the coupling constant and massive propagator in the integrand,

$$g_A^{3a,\sigma}=-\frac{g^2_{\sigma qq}}{2M}\int\frac{d^4p}{(2\pi)^4}\frac{M(p)}{p^2+m^2_\sigma}\frac{-p^2+2M(p)^2}{[p^2+M(p)^2]^2}. \tag{24}$$

These expressions evaluate to 9.1×10^{-4} and 0.019, respectively, when we take the parameters described in Sec. III.

The non-local contributions couple the axial current directly to the four-quark vertex, as shown in Fig. 3(c). With an internal pion, we find

$$g_A^{3b,\pi}=\frac{3g^2_{\pi qq}}{2M}\int\frac{d^4p}{(2\pi)^4}\frac{M'(p)}{p[p^2+M(p)^2]}, \tag{25}$$

and for the sigma,

$$g_A^{3b,\sigma}=\frac{g^2_{\sigma qq}}{2M}\int\frac{d^4p}{(2\pi)^4}\frac{pM'(p)}{(p^2+m^2_\sigma)[p^2+M(p)^2]}. \tag{26}$$

Numerically, we obtain -0.20 and -0.047 for these corrections. Note that both numbers are significantly larger than those of Fig. 3(a), i.e. the non-local corrections are not suppressed in any systematic way.

In the chiral limit, we therefore find

$$g_A=1+9.1\times10^{-4}+0.019-0.20-0.047=0.77,$$

in excellent agreement with the nonrelativistically deduced value.

The vector current, to which the axial vector coupling is compared experimentally, receives no contributions from the non-local vertices. Furthermore, a critical sign difference in the numerator as compared to Eqs. (23) and (24) leads to a finite but very small contribution from each diagram, summing to only 2% of the bare coupling. Thus we can ignore such instanton effects in the coupling g_V and take it to be unity.

In order to estimate the correction from finite current quark masses, we assume larger masses for the internal meson propagators. With the physical pion mass, $m_\pi=138$ MeV, the instanton calculations of Hutter [35] generate a sigma mass of $m_\sigma=540$ MeV. Reevaluating, we naturally find a deduction in these subtractions,

$$g_A=1+0.014+0.017-0.19-0.042=0.80.$$

However, this is a mere 5% increase compared to the chiral limit, and the result remains close to the pseudo-experimental value.

We are ultimately interested in the nucleon axial vector coupling, and in the nonrelativistic limit we find

$$G_A = \frac{5}{3} g_A = \begin{cases} 1.28 & \text{for } m_\pi = 0, \\ 1.34 & \text{for } m_\pi = 140 \text{ MeV}. \end{cases} \tag{27}$$

The physical value, $G_A = 1.2670$ [1],[1] lies a few percent below.

VI. CONCLUSIONS

Through an analysis of current conservation in the instanton liquid model, we found that terms arising from the nonlocality of the 't Hooft vertex were central to maintaining the symmetries of the theory. This can be thought of as a correction for the truncation of quark basis states, for we have isolated the zero modes and removed them from the dynamics. Reintroducing the symmetries to the volume of this extended interaction vertex via path-ordered exponentials recovers the symmetries, as has been noted by other authors some time ago [24,25]. A general form of current corrections has been computed here, followed by specialization to the axial isovector case.

After these modifications the pion pole appears as a *t*-channel quark-antiquark correlation, governing each constituent quark's axial vector form factor and maintaining current conservation. Similar multi-instanton corrections, when applied with *s*-channel chiral mesons, lead to an axial vector coupling constant of the quark of $g_A = 0.77$. In the nonrelativistic limit this corresponds to a nucleon coupling of $G_A = 1.28$, within 2% of the experimental value. Away from the chiral limit, when the pion is set to its physical mass, agreement is within 5%. These numbers are obtained with the two standard inputs of the instanton vacuum, the number density and average size, and no additional parameters. The nonlocal current terms lead to the dominant contributions and cannot be disregarded. While there are additional corrections from vector mesons, the significantly lower coupling constants and higher masses render them negligible.

Our results suggest that the ILM contains the primary effects which lead to the observed vacuum axial vector coupling of 1.267. Given that this variance from unity is a likely result of spontaneous chiral symmetry breaking, a central result of the ILM, it is not surprising that instantons contribute favorably. It is, however, noteworthy in its implication that relativistic wave function corrections to the constituent quark are nearly negligible.

[1]Our convention differs from that reported by the Particle Data Group by an overall sign.

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