Semileptonic decay constants of octet baryons in the chiral quark-soliton model

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Based on the recent study of the magnetic moments and axial-vector constants within the framework of the chiral quark-soliton model, we investigate the baryon semileptonic decay constants (f_1, f_2) and (g_1, g_2). Employing the relations between the diagonal transition matrix elements and off-diagonal ones in the vector and axial-vector channels, we obtain the ratios of baryon semileptonic decay constants f_2/f_1 and g_1/f_1. The F/D ratio is also discussed and found that the value predicted by the present model naturally lies between that of the Skyrme model and that of the nonrelativistic quark model. The singlet axial-vector constant $g_A^{(0)}$ can be expressed in terms of the F/D ratio and $g_A^{(3)}$ in the present model and turns out to be small. The results are compared with available experimental data and found to be in good agreement with them. In addition, the induced pseudotensor coupling constants g_2/f_1 are calculated, the SU(3) symmetry breaking being considered. The results indicate that the effect of SU(3) symmetry breaking might play an important role for some decay modes in hyperon semileptonic decay. [S0556-2821(97)02923-8]

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I. INTRODUCTION

Baryon semileptonic decays have played an important role in various facets to the understanding of the structure of baryons. For example, they provide information on the Cabibbo-Kobayashi-Maskawa (CKM) angles $|V_{ud}|$ and $|V_{us}|$ as well as the F/D ratio. Recently, baryon semileptonic decays have gained new interest in respect of a series of experiments measuring the first moment of the spin-dependent structure function $g_1(x)$ [1–4], since it is related to SU(3) invariant matrix elements of hyperon semileptonic processes F and D.

A recent high-precision measurement of $\Sigma^- \to n + e^- + \overline{\nu}$ [5] shows a hint that the effect of SU(3) symmetry breaking might be important to describe baryon semileptonic decays. Another experiment measuring g_1/f_1 in $\Lambda \to p + e^- + \overline{\nu}$ was conducted by Dworkin *et al.* with high statistics [6]. The result of Ref. [6] prefers the hypothesis that the weak magnetism is less than the CVC (conserved vector current) prediction $(f_2/f_1 = 0.97)$. In fact, Ref. [6] obtained $f_2/f_1 = 0.15 \pm 0.30$ at which the fit yields the minimum of χ^2. This value is really far from the CVC hypothesis. Also from the theoretical point of view there were already serious doubts about the strong postulate of exact SU(3) symmetry in the Cabibbo theory [7–9].

The effects of SU(3) symmetry breaking in baryon semileptonic decays have been extensively studied from various points of view [9–13]. Donoghue *et al.* [9] made a careful analysis of hyperon semileptonic decays, considering the pattern of symmetry breaking based on the quark model. It was asserted in Ref. [9] that SU(3) breaking comes from two sources: the mismatch of the wave functions for the quarks and the recoil effect. However, Roos reanalyzed hyperon semileptonic decays [10] including recent data and showed that the scheme of symmetry breaking by Donoghue *et al.* fails to fit correctly the g_1/f_1 ratio for $\Lambda \to p + e^- + \overline{\nu}$. The mismatch of strange-quark wave functions worsens the fit. Only by introducing a substantial second-class axial-vector coupling g_2 in the $\Lambda \to p + e^- + \overline{\nu}$, one could fit the data [10]. Avenarius [11] studied also SU(3) symmetry breaking in semileptonic hyperon decays, based on the *Ansatz* that SU(3) symmetry in the polarization at the current quark level is kept while SU(3) symmetry at the constituent quark level is broken. The result was that with SU(3) symmetry broken the F/D ratio (0.73 ± 0.09) turned out to be larger than that of the case of SU(3) symmetry (0.59 ± 0.02). Ehrnsperger and Schäfer [12] came to a rather different conclusion. They showed that the effects of SU(3) symmetry breaking lead to a reduction of the F/D ratio (0.49 ± 0.08). Quite recently, Ratcliffe [13] reexamined SU(3) symmetry-breaking effects in hyperon semileptonic decays. What he obtained is that $F/D = 0.582$ for an SU(3) symmetric fit and $F/D = 0.570$ for an SU(3) breaking fit. The result of Ref. [13] indicates that the effects of SU(3) symmetry breaking is rather tiny. There have been also similar discussions related to the validity of SU(3) symmetry in the context of the spin structure of the proton. In particular, Lipkin strongly criticized the use of SU(3) symmetry in studying the spin structure of the proton [14]. The topic of SU(3) symmetry breaking in hyperon semileptonic decays seems to be evidently far from the settlement yet and very difficult to be analyzed without relying on particular *Ansätze*.

Recently, we investigated the magnetic moments of the baryon octet [15] and the axial-vector constants of the nucleon [16] within the framework of the chiral quark-soliton model (χQSM) [17], taking into account the $1/N_c$ rotational corrections and linear m_s corrections which furnish the effect of SU(3) symmetry breaking. The magnetic moments and axial-vector constants which are given in terms of diagonal matrix elements can be, however, related to the off-

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diagonal transition form factors, i.e., semileptonic weak form factors f_1, f_2, and g_1. Therefore, we can easily evaluate them, using former calculations. The presence of the m_s corrections allows the nonvanishing values of induced pseudotensor coupling constants g_2.

The large N_c limit ($N_c \to \infty$) provides a useful guideline in understanding the low-energy properties of the baryon systematically [18], though in reality N_c is equal to 3. In the large N_c, the nucleon can be viewed as a classical soliton of the pion field. An example of the dynamical realization of this idea is given by the Skyrme model [19]. However, the χQSM presents a more realistic picture than the Skyrme model. In the light of chiral perturbation theory, the effective chiral action on which the χQSM is based contains automatically the four-derivative Gasser-Leutwyler terms [20] and the venerable Wess-Zumino term [21] with correct coefficients [22–24]. Moreover, the χQSM interpolates between the Skyrme model and the nonrelativistic model (NRQM) [25,26], because the χQSM is ideologically close to the Skyrme model in the limit of large soliton size while as the size of the soliton approaches zero the χQSM reproduces the results of the NRQM.

The aim of the present paper is to investigate the semileptonic decay constants (f_1, f_2) and (g_1, g_2) within the framework of the χQSM. The outline of the paper is as follows. In the next section, we sketch briefly the basic formalism of the χQSM. In Sec. III, we first discuss the semileptonic decay constants in SU(3) flavor symmetry. As it should be, it is shown that the induced pseudotensor coupling constant $g_2(0)$ vanishes in SU(3) symmetry within the framework of the χQSM. The F/D ratio is also discussed. The singlet axial-vector constant $g_A^{(0)}$ is shown to be related to the F/D ratio and the axial-vector constant $g_A^{(3)}$. Considering the strange quark mass m_s, we discuss the effect of SU(3) symmetry breaking on the coupling constants f_1, f_2, and g_1. The g_2 is also evaluated. We compare the results with those obtained in the case of SU(3) symmetry. The deviations from the Cabibbo theory are discussed in detail. In Sec. IV, we summarize the present work and draw the conclusion.

II. GENERAL FORMALISM

The transition matrix element $\mathcal{M}_{B_1\to B_2 l \overline{\nu_l}}$ for the process $B_1 \to B_2 l \overline{\nu_l}$ can be written as

$$\mathcal{M}_{B_1\to B_2 l \overline{\nu_l}} = \frac{G}{\sqrt{2}} \begin{pmatrix} V_{ud} & (\text{for } \Delta S=0) \\ V_{us} & (\text{for } |\Delta S|=1) \end{pmatrix} \times \langle B_2|J_\mu^W|B_1\rangle \overline{u_l}(p_l)\gamma^\mu(1-\gamma_5)u_{\overline{\nu_l}}(p_\nu), \tag{1}$$

where G denotes the effective Fermi coupling constant and V_{ud}, V_{us} stand for the Cabibbo-Kobayashi-Maskawa angles. The leptonic current $\overline{u_l}(p_l)\gamma^\mu(1-\gamma_5)u_{\overline{\nu_l}}(p_\nu)$ is the known part. The hadronic weak current J_μ^W has following spin and flavor structures:

$$J_\mu^W = \begin{cases} \overline{\psi(x)}\gamma_\mu(1-\gamma_5)\frac{1}{2}(\lambda_1 \pm i\lambda_2)\psi(x) & \text{for } \Delta S=0, \\ \overline{\psi(x)}\gamma_\mu(1-\gamma_5)\frac{1}{2}(\lambda_4 \pm i\lambda_5)\psi(x) & \text{for } |\Delta S|=1. \end{cases} \tag{2}$$

The transition matrix element of the hadronic weak current $\langle B_2|J_\mu^W|B_1\rangle$ can be expressed in terms of six independent form factors:

$$\langle B_2|J_\mu^W|B_1\rangle = \overline{u}_{B_2}(p_2)\left[\left\{f_1(q^2)\gamma_\mu - \frac{if_2(q^2)}{M_1}\sigma_{\mu\nu}q^\nu + \frac{f_3(q^2)}{M_1}q_\mu\right\} - \left\{g_1(q^2)\gamma_\mu - \frac{ig_2(q^2)}{M_1}\sigma_{\mu\nu}q^\nu + \frac{g_3(q^2)}{M_1}q_\mu\right\}\gamma_5\right]u_{B_1}(p_1), \tag{3}$$

with the momentum transfer $q = p_2 - p_1$. The form factors f_i and g_i are real quantities depending only on the square of the momentum transfer in the case of CP-invariant processes. We can safely neglect f_3 and g_3 for the reason that on account of q_μ their contribution to the decay rate is proportional to the ratio $m_l^2/M_1^2 \ll 1$, where m_l represents the mass of the lepton (e or μ) in the final state and M_1 that of the baryon in the initial state.

It is already well known how to deal with the hadronic matrix elements such as $\langle B_2|J_\mu^W|B_1\rangle$ in the χQSM (for a review see [27]). Hence, we shall briefly explain how to calculate them with regard to the semileptonic processes. The χQSM is characterized by a low-momenta QCD partition function [28] which is given by a functional integral over eight pseudoscalar and quark fields:

$$\mathcal{Z} = \int \mathcal{D}\psi\mathcal{D}\psi^\dagger\mathcal{D}\pi^a \exp\left(\int d^4x\,\psi^\dagger\beta(-i\partial\!\!\!/ + \hat{m} + MU^{\gamma_5})\psi\right). \tag{4}$$

Integrating out quark fields leaves us with the effective chiral action

$$\mathcal{Z} = \int \mathcal{D}\pi^a \exp(-S_{\text{eff}}[\pi]), \tag{5}$$

where

$$S_{\text{eff}} = -\text{Sp}\ln\beta(-i\partial\!\!\!/ + \hat{m} + MU^{\gamma_5}) \tag{6}$$

with the pseudoscalar chiral field

$$U^{\gamma_5} = \exp(i\pi^a\lambda^a\gamma_5) = \frac{1+\gamma_5}{2}U + \frac{1-\gamma_5}{2}U^\dagger. \tag{7}$$

$\hat{m}$ is the matrix of the current quark mass given by

$$\hat{m} = \text{diag}(m_u, m_d, m_s) = m_0\mathbf{1} + m_8\lambda_8. \tag{8}$$

λ^a represent the usual Gell-Mann matrices normalized as $\text{tr}(\lambda^a\lambda^b) = 2\delta^{ab}$. Here, we have assumed isospin symmetry ($m_u = m_d = \overline{m}$). M stands for the dynamical quark mass arising from the spontaneous chiral symmetry breaking, which is in general momentum-dependent [28]. We regard M as a constant and employ the proper-time regularization for convenience. The m_0 and m_8 in Eq. (8) are respectively defined by

$$m_0=\frac{2\overline{m}+m_s}{3}\simeq\frac{m_s}{3}, \quad m_8=\frac{\overline{m}-m_s}{\sqrt{3}}\simeq\frac{-m_s}{\sqrt{3}}. \tag{9}$$

The operator iD is expressed in Euclidean space in terms of the Euclidean time derivative ∂_τ and the Dirac one–particle Hamiltonian $H(U^{\gamma_5})$:

$$iD=\partial_\tau+H(U^{\gamma_5})+\beta\hat{m}, \tag{10}$$

with

$$H(U^{\gamma_5})=\frac{\vec{\alpha}\cdot\nabla}{i}+\beta M U^{\gamma_5}. \tag{11}$$

β and $\vec{\alpha}$ are the well-known Dirac Hermitian matrices. The U is assumed to have a structure corresponding to the embedding of the SU(2) hedgehog into SU(3):

$$U_c=\begin{pmatrix} U_0 & 0 \\ 0 & 1 \end{pmatrix}, \tag{12}$$

with

$$U_0=\exp[i\hat{r}\cdot\vec{\tau}P(r)]. \tag{13}$$

$P(r)$ is called profile function. The partition function of Eq. (4) can be simplified by the saddle point approximation which is exact in the large N_c limit. One ends up with a stationary profile function $P(r)$ which is evaluated by solving the Euler-Lagrange equation corresponding to $\delta S_{\mathrm{eff}}/\delta P(r)=0$. This gives a static classical field U_0. The soliton is quantized by introducing collective coordinates corresponding to $\mathrm{SU(3)_f}$ rotations of the soliton in flavor space [and simultaneously $\mathrm{SU(2)_{spin}}$ in spin space]

$$U(t,\vec{x})=R(t)U_c(\vec{x})R^\dagger(t), \tag{14}$$

where $R(t)$ is a time-dependent SU(3) matrix. The quantum states from this quantization are identified with the SU(3) baryons according to their quantum numbers. In the large N_c limit, the angular velocity of the soliton $\Omega=R^\dagger(t)\dot{R}(t)$ can be regarded as a small parameter, so that we can use it as an expansion parameter. After the rotation, the Dirac differential operator Eq. (10) can be expressed as

$$i\widetilde{D}=[\partial_\tau+H(U^{\gamma_5})+\Omega(t)+\gamma_4 R^\dagger(t)\hat{m}R(t)]. \tag{15}$$

Then the propagator $(i\widetilde{D})^{-1}$ can be expanded with regard to the angular velocity Ω and the strange quark mass m_s:

$$\frac{1}{i\widetilde{D}}\simeq\frac{1}{\omega+iH}-\frac{1}{\omega+iH}\Omega\frac{1}{\omega+iH} -\frac{1}{\omega+iH}\gamma_4 R^\dagger(t)\hat{m}R(t)\frac{1}{\omega+iH}. \tag{16}$$

The transition matrix element of the hadronic weak current given by Eq. (3) can be related to the correlation function

$$\left\langle 0\left|J_{B_1}\left(\vec{x},\frac{T}{2}\right)\overline{\psi}\hat{\Gamma}\hat{O}\psi J^\dagger_{B_2}\left(\vec{y},-\frac{T}{2}\right)\right|0\right\rangle \tag{17}$$

at large Euclidean time T. $\hat{\Gamma}$ and $\hat{O}$ are abbreviations for the corresponding spin and flavor operators. The J_B denotes the baryon current which is constructed from N_c quark fields

$$J_B=\frac{1}{N_c!}\varepsilon^{i_1\cdots i_{N_c}}\Gamma^{\alpha_1\cdots\alpha_{N_c}}_{SS_3II_3Y}\psi_{\alpha_1 i_1}\cdots\psi_{\alpha_{N_c}i_{N_c}}. \tag{18}$$

$\alpha_1\cdots\alpha_{N_c}$ are spin-isospin indices, $i_1\cdots i_{N_c}$ are color indices, and the matrices $\Gamma^{\alpha_1\cdots\alpha_{N_c}}_{SS_3II_3Y}$ are taken to endow the corresponding current with the quantum numbers SS_3II_3Y. The $J_B(J_B^\dagger)$ plays a role of annihilating (creating) the baryon state at large T. The rotational $1/N_c$ and linear m_s corrections being taken into account as shown in Eq. (16), the relevant transition matrix elements can be written as

$$\begin{aligned}(f_1+f_2)^{(B_1\to B_2)}=&\,w_1\langle B_2|D^{(8)}_{X3}|B_1\rangle+w_2 d_{pq3}\langle B_2|D^{(8)}_{Xp}\,\hat{S}_q|B_1\rangle\\&+\frac{w_3}{\sqrt{3}}\langle B_2|D^{(8)}_{X8}\,\hat{S}_3|B_1\rangle\\&+m_s\Bigg[\frac{w_4}{\sqrt{3}}d_{pq3}\langle B_2|D^{(8)}_{Xp}\,D^{(8)}_{8q}|B_1\rangle\\&+w_5\langle B_2|(D^{(8)}_{X3}\,D^{(8)}_{88}+D^{(8)}_{X8}\,D^{(8)}_{83})|B_1\rangle\\&+w_6\langle B_2|(D^{(8)}_{X3}\,D^{(8)}_{88}-D^{(8)}_{X8}\,D^{(8)}_{83})|B_1\rangle\Bigg]\end{aligned} \tag{19}$$

for the transition magnetic moments $[f_1(0)+f_2(0)]^{(B_1\to B_2)}$ and

$$\begin{aligned}g_1^{(B_1\to B_2)}=&\,a_1\langle B_2|D^{(8)}_{X3}|B_1\rangle+a_2 d_{pq3}\langle B_2|D^{(8)}_{Xp}\,\hat{S}_q|B_1\rangle\\&+\frac{a_3}{\sqrt{3}}\langle B_2|D^{(8)}_{X8}\,\hat{S}_3|B_1\rangle\\&+m_s\Bigg[\frac{a_4}{\sqrt{3}}d_{pq3}\langle B_2|D^{(8)}_{Xp}\,D^{(8)}_{8q}|B_1\rangle\\&+a_5\langle B_2|(D^{(8)}_{X3}\,D^{(8)}_{88}+D^{(8)}_{X8}\,D^{(8)}_{83})|B_1\rangle\\&+a_6\langle B_2|(D^{(8)}_{X3}\,D^{(8)}_{88}-D^{(8)}_{X8}\,D^{(8)}_{83})|B_1\rangle\Bigg],\end{aligned} \tag{20}$$

for the transition axial-vector constants $g_1^{(B_1\to B_2)}(0)$. The induced pseudotensor coupling constants $g_2^{(B_1\to B_2)}$ are expressed by

$$\frac{g_2^{(B_1\to B_2)}}{M_{B_1}} = 4m_s(\beta_1 i f_{ab3}+\beta_2 i\varepsilon_{ab3})\cdot\langle B_2|D_{Xa}D_{8b}|B_1\rangle. \tag{21}$$

The parameters w_i, a_i, and β_j depend on dynamics of the chiral soliton. As for the expressions for w_i and a_i, one can find them in Ref. [29] and Ref. [16], respectively. β_1 and β_2 can be written explicitly as [30,31]

$$\beta_1+\beta_2=\frac{iN_c}{24\sqrt{3}}\int\frac{d\omega}{2\pi}\,\mathrm{Sp}\left(\frac{1}{\omega+iH}\gamma_4\tau^i\frac{1}{\omega+iH}i\varepsilon_{ijk}\tau^j x^k\gamma_5\right), \tag{22}$$

$$\beta_1=-\frac{iN_c}{12\sqrt{3}}\int\frac{d\omega}{2\pi}\,\mathrm{Sp}\left(\left\{\frac{1}{\omega+iH}\gamma_4,\frac{1}{\omega+iH_0}\right\}\gamma_5\vec{x}\cdot\vec{\tau}\right). \tag{23}$$

The remarkable feature of the soliton picture of the baryons is that the singlet axial-vector charge of the nucleon $g_A^{(0)}$ is expressed in terms of the *same* parameters a_i as in Eq. (20):

$$g_A^{(0)}=\frac{1}{2}a_3+\sqrt{3}m_s(a_5-a_6). \tag{24}$$

Hence, in this picture the value of $g_A^{(0)}$ can be extracted by fitting the data on semileptonic decays *without* resorting to those on polarized deep inelastic scattering [see the next section for the analysis of the SU(3)-symmetric case]. With SU(3) symmetry explicitly broken by m_s, the collective Hamiltonian is no more SU(3) symmetric. The pure octet states are mixed with the higher representations such as antidecuplet states. Therefore, the baryon wave function with spin $S=1/2$ requires the modification due to the strange quark mass (m_s). Since we treat m_s perturbatively up to the first order, the collective wave function of the baryon octet can be written as

$$\Psi_B(R)=\Psi_B^{(8)}(R)+m_s\,c_{\overline{10}}^B\;\Psi_B^{(\overline{10})}(R)+m_s c_{27}^B\Psi_B^{(27)}(R), \tag{25}$$

where

$$c_{\overline{10}}^B=\frac{\sqrt{5}}{15}(\sigma-r_1)\begin{bmatrix}1\\0\\1\\0\end{bmatrix}I_2,\quad c_{27}^B=\frac{1}{75}(3\sigma+r_1-4r_2)\begin{bmatrix}\sqrt{6}\\3\\2\\\sqrt{6}\end{bmatrix}I_2 \tag{26}$$

in the basis of $[N,\ \Lambda,\ \Sigma,\ \Xi]$. Here, B denotes the SU(3) octet baryons with the spin 1/2. The constant σ is related to the nucleon sigma term $\Sigma=\overline{m}\langle N|\overline{u}u+\overline{d}d|N\rangle=3/2\overline{m}\sigma$ and r_i designates K_i/I_i, where K_i stands for the anomalous moments of inertia defined in Ref. [32]. The collective wave function can be explicitly written in terms of the SU(3) Wigner $D^{(\mu)}$ function

$$\Psi_B^{(\mu)}=(-)^{S_3-1/2}\sqrt{\dim(\mu)}[D^{(\mu)}_{(YTT_3)(-1SS_3)}]^*. \tag{27}$$

TABLE I. The expressions of $f_1(0)$, $f_2(0)$, and $g_1(0)$ in exact SU(3) symmetry. The κ_p and the κ_n denote the anomalous magnetic moments of the proton and the neutron, respectively.

Decay mode	$f_1(0)$	$f_2(0)$	$g_1(0)$
$n\to p$	1	$\frac{1}{2}(\kappa_p-\kappa_n)$	$F+D$
$\Sigma^-\to\Sigma^0$	$\sqrt{2}$	$\frac{1}{\sqrt{2}}(\kappa_p+\frac{1}{2}\kappa_n)$	$\sqrt{2}F$
$\Sigma^\pm\to\Lambda$	0	$\pm\sqrt{\frac{3}{8}}\kappa_n$	$\mp\sqrt{\frac{2}{3}}D$
$\Lambda\to p$	$\sqrt{\frac{3}{2}}$	$\frac{1}{2}\sqrt{\frac{3}{2}}\kappa_p$	$\sqrt{\frac{3}{2}}(F+D/3)$
$\Sigma^-\to n$	1	$\frac{1}{2}(\kappa_p+2\kappa_n)$	$F-D$
$\Xi^-\to\Lambda$	$\sqrt{\frac{3}{2}}$	$\frac{1}{2}\sqrt{\frac{3}{2}}(\kappa_p+\kappa_n)$	$\sqrt{\frac{3}{2}}(F-D/3)$
$\Xi^-\to\Sigma^0$	$\frac{1}{\sqrt{2}}$	$\frac{1}{2\sqrt{2}}(\kappa_p-\kappa_n)$	$\sqrt{\frac{1}{2}}(F+D)$

Hence, we have two different contributions of SU(3) symmetry breaking: One from the effective Lagrangian and the other from the wave function corrections. All contributions of SU(3) symmetry breaking are kept in linear order of m_s. Apart from these two contributions, we shall see in the next section that in the case of $f_2(0)$ the mass differences between octet baryons come into play. Hence, on the whole, we have three different sources for the SU(3) symmetry breaking in the present model.

III. SEMILEPTONIC DECAY CONSTANTS IN THE χQSM

A. Exact SU(3) symmetry

In exact SU(3) symmetry, the vector coupling constants $f_1(0)$ and $f_2(0)$ can be simply expressed in terms of the anomalous magnetic moments of the proton and neutron. Similarly, $g_1(0)$ in all decay modes can be parametrized in terms of two SU(3)-invariant constants F and D [33]:

$$g_1^{(B_1\to B_2)}=FC_F^{B_1\to B_2}+DC_D^{B_1\to B_2}, \tag{28}$$

where $C_F^{B_1\to B_2}$ and $C_D^{B_1\to B_2}$ are SU(3) Clebsch-Gordan coefficients that appear when an octet operator is sandwiched between octet states. The superscript refers to the hadrons involved and subscripts F and D denote the antisymmetric and symmetric parts. We list the expressions for $f_1(0)$, $f_2(0)$, and $g_1(0)$ in Table I.[1]

The pseudotensor coupling constants $g_2(0)$ are all predicted to be zero in exact SU(3) symmetry because of G parity. In fact, it can be shown that $g_2(0)$ vanish in the present model. To do so, it is of great use to introduce a transformation

$$G_5=\tau_2 C\gamma_5, \tag{29}$$

[1]Note that the signs in the $\Lambda\to p$, and $\Sigma^-\to n$ modes are different from Ref. [33]. However, the ratios $f_2(0)/f_1(0)$ and $g_1(0)/f_1(0)$ are not affected. We have employed the phase convention in the manner of De Swart [34,35].

where C is the operator of charge conjugation: $C^{-1}\gamma_\mu C=-\gamma_\mu^T$. Under this transformation, the one-body Dirac Hamiltonian, Dirac, and Pauli matrices are respectively changed as follows:

$$G_5^{-1}HG_5=H^T,$$
$$G_5^{-1}\gamma_\alpha G_5=\gamma_\alpha^T \quad \text{for} \quad \alpha=1,\ldots,5,$$
$$G_5^{-1}\tau_a G_5=-\tau_a^T. \tag{30}$$

The pertinent trace for the leading order contribution to $g_2(0)$ can be written as

$$\mathrm{tr}\left[\left\langle \vec{x}\left|\frac{1}{\omega+iH}\gamma_5\tau_a x_i\right|\vec{x}\right\rangle\right]. \tag{31}$$

Utilizing the G_5 transformation and the properties of the trace $\mathrm{tr}(M^T)=\mathrm{tr}(M)$ and $\mathrm{tr}(WMW^{-1})=\mathrm{tr}(M)$, we can show that the trace of the leading contribution vanishes:

$$\begin{aligned}\mathrm{tr}&\left[\left\langle \vec{x}\left|\frac{1}{\omega+iH}\gamma_5\tau_a x_i\right|\vec{x}\right\rangle\right]\\ &=\mathrm{tr}\left[\left\langle \vec{x}\left|G_5^{-1}\frac{1}{\omega+iH}\gamma_5\tau_a x_i G_5\right|\vec{x}\right\rangle\right]^T\\ &=-\mathrm{tr}\left[\left\langle \vec{x}\left|\frac{1}{\omega+iH}\gamma_5\tau_a x_i\right|\vec{x}\right\rangle\right].\end{aligned} \tag{32}$$

Similarly we can prove that the $1/N_c$ rotational corrections to $g_2(0)$ also disappear.

In exact SU(3) symmetry, $g_1^{(B_1\to B_2)}$ can be written in terms of three independent dynamic quantities a_i calculable in the present model:

$$g_1^{(B_1\to B_2)}=a_1\langle B_2|D_{X3}^{(8)}|B_1\rangle+a_2 d_{pq3}\langle B_2|D_{Xp}^{(8)}\hat{S}_q|B_1\rangle + \frac{a_3}{\sqrt{3}}\langle B_2|D_{X8}^{(8)}\hat{S}_3|B_1\rangle. \tag{33}$$

As for the process $n\to pe^-\overline{\nu}$, $g_1^{(n\to p)}$ becomes just the axial-vector coupling constant $g_A^{(3)}$. After some straightforward manipulation [16], we end up with the expression for $g_A^{(3)}$ [36,37]:

$$g_A^{(3)}=\frac{7}{30}\left(-a_1+\frac{1}{2}a_2+\frac{1}{14}a_3\right). \tag{34}$$

We can also obtain the singlet axial-vector coupling constant $g_A^{(0)}$:

$$g_A^{(0)}=\frac{1}{2}a_3. \tag{35}$$

It is of great interest to see how the χQSM plays an interpolating role between the Skyrme model and the NRQM. In the limit of the NRQM the dynamic quantities a_i in Eq.(33) are, respectively [38,39],

$$a_1=-5, \quad a_2=4, \quad a_3=2, \tag{36}$$

which give the correct values of $g_A^{(3)}$ and $g_A^{(0)}$ in the NRQM, i.e., $g_A^{(3)}=5/3$ and $g_A^{(0)}=1$. The ratio F/D then can be written in terms of a_i:

$$\frac{F}{D}=\frac{5}{9}\left(\frac{-a_1+(1/2)a_2+(1/2)a_3}{-a_1+(1/2)a_2-(1/6)a_3}\right). \tag{37}$$

In the limit of the Skyrme model, i.e., when the size of the soliton is very large [25], $g_A^{(0)}$ approaches zero or a_3 vanishes. Hence, the F/D ratio becomes obviously 5/9 in this limit. This result is exactly the same as what was obtained by Bijnens *et al.* [40] and Chemtob [41]. On the other hand, in the limit of the NRQM, i.e., in the limit of zero soliton size, the present model gives the value of $F/D=2/3$ which is exactly the same value as that of the SU(6) NRQM. The present model predicts F/D to be 0.61—correspondingly, we have $g_A^{(0)}=0.36$ [37]—which lies between the value from the Skyrme model (5/9) and that from the NRQM (2/3). The χQSM shows here again interpolation between the Skyrme model and the NRQM [25,26]. Notably, the smallness of the singlet axial-vector charge $g_A^{(0)}$ is directly related to the fact that F/D does not deviate much from that of the Skyrme model (5/9), where $g_A^{(0)}$ is known to be zero [42]. Using Eqs. (34), (35), (37), we can express the singlet axial-vector constant $g_A^{(0)}$ in terms of the F/D ratio and $g_A^{(3)}$:

$$g_A^{(0)}=\frac{9g_A^{(3)}}{1+F/D}\left(\frac{F}{D}-\frac{5}{9}\right). \tag{38}$$

Substituting the value of $F/D=0.582$ obtained in a recent analysis [13] and $g_A^{(3)}=1.26$, one gets $g_A^{(0)}=0.19$.

The g_1 is normally determined in experiments with g_2 assumed to be zero. However, Hsueh *et al.* [5] extracted for the first time the induced pseudotensor coupling constant g_2 in $\Sigma^-\to ne^-\overline{\nu}$ decay. This new experimental result gives a reduced value for $g_1(0)$: Instead of $g_1(0)=0.328\pm0.019$ and $g_2=0$, $g_1(0)=0.20\pm0.08$ and $g_2(0)=-0.56\pm0.37$ are obtained. These results are remarkable, since they indicate that the effect of SU(3) symmetry breaking might play an important role in baryon semileptonic decays. In fact, these new experimental results have triggered discussions about the effect of SU(3) symmetry breaking in hyperon semileptonic decays [11,13]. Hence, we need to consider the explicit SU(3) symmetry breaking.

B. SU(3) symmetry-breaking effects

As we have shown in Sec. II, the strange quark mass m_s provides the effect of SU(3) symmetry breaking in two different forms: One from the effective action and the other from the wave function corrections. Apart from these two contributions, the mass difference between baryon states must be considered in the case of $f_2(0)$.

By switching on SU(3) symmetry breaking to the first order in m_s, we obtain the expressions for $f_2(0)$ and $g_1(0)$ deviating from those listed in Table I. It should be noted that by the Ademollo-Gatto theorem [43] $f_1(0)$ do not get any contribution from linear m_s corrections. We choose the combination of the magnetic moments in which all corrections from $1/N_c$ are canceled to avoid ambiguity arising from the

TABLE II. The results of $f_2(0)/f_1(0)$. The column under $SU(3)_{sym}$ lists the case of exact SU(3) symmetry with the experimental data of κ_p and κ_n while the next column shows the case of broken SU(3) symmetry by m_s with the experimental data of κ_B [44]. The third column presents the numerical results of $f_2(0)$ in the χQSM with the constituent quark mass $M=420$ MeV, while the fourth one shows those of $f_2(0)$ in the χQSM with the same constituent quark mass. The experimental data are taken from Ref. [6] for $\Lambda\to pe^-\overline{\nu}$ and from Ref. [5] for $\Sigma^-\to ne^-\overline{\nu}$.

Decay mode	$SU(3)_{sym}$	$SU(3)_{br}$	χQSM_{sym}	χQSM_{br}	Expt.
$n\to p$	1.853	1.853	1.41	1.58	
$\Sigma^-\to\Sigma^0$	0.418	0.516 ± 0.012	0.25	0.43	
$\Sigma^-\to\Lambda$	1.435 [a]	1.625 ± 0.011 [a]	0.95 [a]	1.33 [a]	
$\Lambda\to p$	0.896	0.787 ± 0.004	0.64	0.74	0.15 ± 0.30
$\Sigma^-\to n$	-1.017	-1.010 ± 0.016	-0.92	-1.18	$-0.96\pm0.07\pm0.13$
$\Xi^-\to\Lambda$	-0.060	-0.093 ± 0.006	-0.14	-0.18	
$\Xi^-\to\Sigma^0$	1.853	1.725 ± 0.011	1.41	2.06	

[a] Instead of f_2/f_1, we list $\sqrt{\frac{3}{2}}f_2$.

relations between the hyperon magnetic moments [15]. Hence we get the unambiguous relations between hyperon magnetic moments and off-diagonal matrix elements:

$$f_2^{(n\to p)}(0)=\frac{1}{2}(\kappa_p-\kappa_n),$$

$$f_2^{(\Sigma^-\to\Sigma^0)}(0)=\frac{M_\Sigma}{2M_N}\frac{1}{\sqrt{2}}(\kappa_{\Sigma^+}-\kappa_{\Sigma^-}),$$

$$f_2^{(\Sigma^\pm\to\Lambda)}(0)=\mp\frac{M_\Sigma}{\sqrt{2}M_N}\mu_{\Sigma^0\Lambda},$$

$$\mu_{\Sigma^0\Lambda}=\frac{1}{\sqrt{3}}\left(-\kappa_n+\frac{1}{4}(\kappa_{\Sigma^+}+\kappa_{\Sigma^-})-\kappa_{\Xi^0}+\frac{3}{2}\kappa_\Lambda\right),$$

$$f_2^{(\Lambda\to p)}(0)=\frac{M_\Lambda}{2\sqrt{6}M_N}\left(\kappa_n+\frac{5}{2}\kappa_p+\frac{1}{2}\kappa_{\Sigma^-}-3\kappa_\Lambda-\kappa_{\Xi^-}\right),$$

$$f_2^{(\Sigma^-\to n)}(0)=\frac{M_\Sigma}{2M_N}\left(\kappa_n+\frac{1}{2}\kappa_p-\kappa_{\Sigma^-}-\frac{1}{2}\kappa_{\Sigma^+}\right),$$

$$f_2^{(\Xi^-\to\Lambda)}(0)=\frac{M_\Xi}{2\sqrt{6}M_N}\left(\kappa_p-\frac{1}{2}\kappa_{\Sigma^+}+3\kappa_\Lambda-\kappa_{\Xi^0}-\frac{5}{2}\kappa_{\Xi^-}\right),$$

$$f_2^{(\Xi^-\to\Sigma^0)}(0)=\frac{M_\Xi}{2\sqrt{2}M_N}\left(\kappa_{\Sigma^+}+\frac{1}{2}\kappa_{\Sigma^-}-\kappa_{\Xi^0}-\frac{1}{2}\kappa_{\Xi^-}\right), \tag{39}$$

where κ_B is the anomalous magnetic moment corresponding to the baryon B. These relations have corrections of order $O(m_s^2)$ and $O(m_s/N_c)$ which are assumed to be small. In exact SU(3) symmetry these relations reduce again to the SU(3) relations shown in Table I. In Table II we compare the results with SU(3) symmetry breaking to those in SU(3) symmetry. The experimental data for $f_2(0)$ are taken from Ref. [6] for $\Lambda\to pe^-\overline{\nu}$ and from Ref. [5] for $\Sigma^-\to ne^-\overline{\nu}$. Incorporating experimental values for magnetic moments taken from [44] into formulas in Table I and those in Eq. (39), respectively, we obtain the $f_2(0)/f_1(0)$ ratios for seven different channels.[2] Let us first compare the first two columns. Apart from the $|\Delta S|=0$ modes for which the effect of SU(3) symmetry breaking is observed in around 10%, we can see the comparably large effect of the SU(3) symmetry breaking. Considering the SU(3) symmetry breaking inherent already in experimental magnetic moments, one can say that the effect of SU(3) symmetry breaking is even larger. From the comparison of the third and fourth columns, we can find the effects of SU(3) symmetry breaking. The effects on $f_2(0)/f_1(0)$ are noticeably large in almost every decay mode. In particular, the deviation from SU(3) symmetry appearing in the $\Sigma^-\to ne^-\overline{\nu}$ mode is remarkable. Indeed, the effects of SU(3) symmetry breaking pull $f_2(0)/f_1(0)(\Sigma^-\to n)$ drastically down from its SU(3) symmetric value, so that it turns out to be in good agreement with the data [5].

The expressions of $g_1(0)$ for the case of SU(3) symmetry breaking can be obtained similarly to the SU(3) symmetric case (see Table I). For convenience, we define the baryonic axial-vector constants as

$$g_A^B=\left\langle B\left|\left(g_A^{(3)}+\frac{1}{\sqrt{3}}g_A^{(8)}\right)\right|B\right\rangle. \tag{40}$$

Then we can write the expressions for $g_1(0)(B_1\to B_2)$ similar to Eq. (39):

$$g_1^{(n\to p)}(0)=\frac{1}{2}(g_A^p-g_A^n),$$

$$g_1^{(\Sigma^-\to\Sigma^0)}(0)=\frac{1}{2\sqrt{2}}(g_A^{\Sigma^+}-g_A^{\Sigma^-}),$$

[2] Note that though we use the SU(3) symmetric expressions to obtain the first column in Table II the results nevertheless include a part of the SU(3) symmetry breaking through the experimental data. However, by doing that we can see at least the effect of SU(3) symmetry breaking within the framework of the χQSM.

$$g_1^{(\Sigma^\pm\to\Lambda)}(0)=\mp\frac{1}{2}\sqrt{\frac{2}{3}}\left(-g_A^n+\frac{1}{2}g_A^{\Sigma^0}-g_A^{\Xi^0}+\frac{3}{2}g_A^{\Lambda}\right),$$

$$g_1^{(\Lambda\to p)}(0)=\frac{1}{2\sqrt{6}}\left(g_A^n+\frac{5}{2}g_A^p+\frac{1}{2}g_A^{\Sigma^-}-3g_A^{\Lambda}-g_A^{\Xi^-}\right),$$

$$g_1^{(\Sigma^-\to n)}(0)=\frac{1}{2}\left(g_A^n+\frac{1}{2}g_A^p-g_A^{\Sigma^-}-\frac{1}{2}g_A^{\Sigma^+}\right),$$

$$g_1^{(\Xi^-\to\Lambda)}(0)=\frac{1}{2\sqrt{6}}\left(g_A^p-\frac{1}{2}g_A^{\Sigma^+}+3g_A^{\Lambda}-g_A^{\Xi^0}-\frac{5}{2}g_A^{\Xi^-}\right),$$

$$g_1^{(\Xi^-\to\Sigma^0)}(0)=\frac{1}{2\sqrt{2}}\left(g_A^{\Sigma^+}+\frac{1}{2}g_A^{\Sigma^-}-g_A^{\Xi^0}-\frac{1}{2}g_A^{\Xi^-}\right). \tag{41}$$

Making use of Eq. (20), we obtain the sum rule between six different decay modes

$$g_1^{(n\to p)}=\frac{1}{2}(\sqrt{2}g_1^{(\Sigma^-\to\Sigma^0)}+\sqrt{6}g_1^{(\Sigma^+\to\Lambda)}-\sqrt{6}g_1^{(\Lambda\to p)}$$
$$+g_1^{(\Sigma^-\to n)}-2\sqrt{2}g_1^{(\Xi^-\to\Sigma^0)}). \tag{42}$$

In order to verify Eq. (42), more accurate experimental data are required than presently available. In the SU(3) limit the right-hand side of Eq. (42) becomes $D+F$ in accordance with the Cabibbo theory [47].

In Table III, the results of $g_1(0)/f_1(0)$ with SU(3) symmetry breaking are compared to those in SU(3) symmetry. The effect of SU(3) symmetry breaking is measured in 5–10 %, which is not that strong. Compared to the case of $f_2(0)/f_1(0)$, the effect of SU(3) symmetry breaking is rather soft. This is partly due to the fact that in the axial-vector channel the mass differences do not come into play, which is somewhat in line with the argument of Ref. [13]. The results predicted by the χQSM are in good agreement with the experimenta data [44] within about 15% which is a typical predictive power of the model. In particular, the results agree with the data remarkably in $|\Delta S|=1$ channels.

TABLE III. The results of $g_1(0)/f_1(0)$. The column under $\chi\text{QSM}_{\text{sym}}$ lists the numerical results in exact SU(3) symmetry while in the next column shows the case of broken SU(3) symmetry by m_s. The constituent quark mass $M=420$ MeV is used. Most experimental data are taken from the Particle Data Group [44]. The data for the $\Sigma^-\to\Lambda$ mode is taken from [45] while that for the $\Xi^-\to\Sigma^0$ mode is from [46].

Decay mode	$\chi\text{QSM}_{\text{sym}}$	$\chi\text{QSM}_{\text{br}}$	Expt.
$n\to p$	1.33	1.42	1.2573 ± 0.0028
$\Sigma^-\to\Sigma^0$	0.50	0.55	
$\Sigma^-\to\Lambda$	0.83 [a]	0.91 [a]	0.720 ± 0.020 [a]
$\Lambda\to p$	0.78	0.73	0.718 ± 0.015
$\Sigma^-\to n$	-0.33	-0.31	-0.340 ± 0.017
$\Xi^-\to\Lambda$	0.23	0.22	0.25 ± 0.05
$\Xi^-\to\Sigma^0$	1.33	1.29	$1.25^{+0.14}_{-0.16}$

[a]Instead of g_1/f_1, we list $\sqrt{\frac{3}{2}}g_1$.

It is also interesting to see that in the present model the ratio of g_1/f_1 between $\Lambda\to pe^-\overline{\nu}$ decay and $\Sigma^-\to ne^-\overline{\nu}$ decay is well reproduced

$$\frac{g_1/f_1(\Lambda\to pe^-\overline{\nu})}{g_1/f_1(\Sigma^-\to ne^-\overline{\nu})}=-2.28 \quad (\text{Expt:} \ -2.11\pm0.15). \tag{43}$$

It was pointed out that this ratio is *a priori* constrained to -3 in quark models with SU(6) symmetry [48,49], which is noticeably larger than the experimental value.

It is also interesting to compare the present results with those from the Skyrme model with vector mesons [50]. Ref-

TABLE IV. The induced pseudotensor coupling constants ratio $g_2(0)/f_1(0)$. The results are compared with Refs. [49,51].

Decay mode	Expression	χQSM	Schlumpf	Kellett
$n\to p$	0	0	0	0.29
$\Sigma^-\to\Sigma^0$	0	0	0	
$\Sigma^-\to\Lambda$	$-\frac{4m_s}{5\sqrt{3}}M_{\Sigma^-}\beta_2$ [a]	-0.029 [a]	0 [a]	0.18 [a]
$\Lambda\to p$	$\frac{8m_s}{5\sqrt{3}}M_{\Lambda}(\beta_1+\frac{1}{2}\beta_2)$	0.046	0.023	0.25
$\Sigma^-\to n$	$-\frac{4m_s}{15\sqrt{3}}M_{\Sigma^-}(3\beta_1+\beta_2)$	-0.020	-0.007	-0.09
$\Xi^-\to\Lambda$	$\frac{2m_s}{5\sqrt{3}}M_{\Xi^-}\beta_1$	0.006	0.008	
$\Xi^-\to\Sigma^0$	$\frac{2m_s}{15\sqrt{3}}M_{\Xi^-}(21\beta_1+16\beta_2)$	0.125	0.04	

[a]Instead of g_2/f_1, we list $\sqrt{\frac{3}{2}}g_2$.

erence [50] presented the g_1/f_1 ratio in five different channels. Except for the $\Lambda \to p$ mode, the present model seems to be far better than Ref. [50]. For example, we obtain the $|g_1/f_1| = 0.31$ for the $\Sigma^- \to n$ mode comparable to the experimental data 0.34 ± 0.017, while Ref. [50] yields 0.24.

From Eq. (21), we can obtain the ratio $g_2(0)/f_1(0)$ in terms of β_1 and β_2. In Table IV, the expressions and numerical results for them are listed. Numerically, β_2 is much larger than β_1, which explains why the g_2/f_1 for the $\Xi^- \to \Sigma^0$ mode turns out to be much greater than those for the other modes. Our results are compared with results of an SU(6) relativistic quark model [51] as well as with those of a light-front relativistic quark model [49]. We see that our results are close—within a factor of 2—to the results of Ref. [49], whereas they differ by almost an order of magnitude from those of Ref. [51].

IV. SUMMARY AND CONCLUSION

The aim of the present work has been to investigate baryon semileptonic decays within the framework of the chiral quark-soliton model (χQSM). In particular, the role of SU(3) symmetry breaking in baryon semileptonic decays have been discussed. Based on the recent studies on the octet magnetic moments and axial-vector constants, we have obtained the ratios f_2/f_1 and g_1/f_1 without and with SU(3) symmetry breaking, respectively. In exact SU(3) symmetry, we have shown that g_2 vanishes in the present model, as it should be. We have discussed also that for the F/D ratio the χQSM (0.61) interpolates the Skyrme model (5/9) and the NRQM (2/3).

It was found that the effect of SU(3) symmetry breaking differs in different channels and modes. In general, SU(3) symmetry breaking contributes strongly to the ratio of the vector coupling constants $f_2(0)/f_1(0)$ while it does not much to the ratio $g_1(0)/f_1(0)$. In addition, we have evaluated the ratio g_2/f_1. This is the first calculation of the g_2/f_1 in soliton models. The results come out to be small except for the $\Xi^- \to \Sigma^0$ mode. Due to lack of experimental data, the values of $g_2(0)$ we calculated are predictions.

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