Chiral corrections to the baryon axial currents

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Chiral logarithmic corrections to the baryon axial vector currents in semileptonic hyperon decay are computed including intermediate octet and decuplet states, and it is shown that the spin-$\frac{3}{2}$ decuplet partially cancels the large spin-$\frac{1}{2}$ octet contribution. The corrections to the quark currents are also computed.

In a recent work [1], a new formulation of the low-energy chiral effective lagrangian for pions and baryons was presented [#1]. By treating baryons in the chiral effective theory as heavy static fields, an improved derivative expansion for the baryon fields was obtained. In this reformulation, higher dimension terms with extra derivatives acting on baryon fields are suppressed to the same extent as higher dimension terms in which the derivatives instead act upon pion fields. The derivative expansion for both pions and baryons becomes an expansion in powers of (k/Λ_χ), where k is a momentum of the order of the pion mass and $\Lambda_\chi \sim 1$ GeV is the chiral symmetry breaking scale which suppresses non-renormalizable terms in the chiral effective theory.

Using this chiral effective lagrangian for baryons, ref. [1] calculated SU(3) breaking chiral logarithmic corrections to axial currents of the baryon octet due to meson loops [#2]. These corrections were shown to be order one. Given that the radiative corrections are so large, it is rather surprising that SU(3) is a useful symmetry for the description of hyperon semileptonic decays. However, it turns out that the experimental data is in reasonable agreement, not only with the lowest order SU(3) symmetric predictions, but also with the predictions including the large SU(3) breaking chiral logarithmic correction [#3]. A fit including the chiral logarithmic corrections leads to very different values for D and F than the SU(3) symmetric fit without logarithms. It is non-trivial that the large one-loop correction can be compensated by a shift in the values of D and F, since the radiative corrections do not satisfy the original SU(3) relations. Thus, it appears to be purely accidental that values for D and F consistent with experimental data can be obtained in both cases. Although experimental agreement can be achieved with and without the large radiative corrections, the better fit to the data occurs without the radiative corrections included. In addition, large radiation corrections to SU(3) symmetry makes it difficult to rationalize the agreement of SU(3) symmetry with the data. For these reasons, we wish to investigate whether SU(3) symmetry is restored in some manner.

There are several possible ways in which the radiative corrections might be smaller than the values computed in ref. [1]. The baryon decuplet–octet splitting is comparable to the pseudoscalar meson masses, e.g. $\Sigma^* - p \approx 450$ MeV, which is comparable to the K mass. Thus one possibility is that the decuplet corrections cancel the octet corrections. In this paper, we calculate the radiative corrections to the

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[#1] The term pion will be used to refer to the entire Goldstone boson octet.

[#2] These corrections were computed earlier by Bijnens, Sonoda, and Wise [2] who inadvertently omitted wave function renormalization.

[#3] For both fits, errors on experimental measurements of g_A were increased by a nominal 0.2, which accounts for the effects of higher order terms in the chiral lagrangian involving the symmetry breaking mass matrix. If errors are not enlarged, neither fit leads to a reasonable χ^2.

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octet axial currents due to intermediate decuplet states, and show that it partially cancels the octet contribution. In the limit that the baryon binding is weak compared with the meson masses, the radiative corrections are computed by first calculating corrections to the quark currents, and then computing the matrix element of the renormalized quark currents in the baryon. This calculation is also presented.

We first calculate the chiral corrections to the baryon currents due to intermediate decuplet fields. These corrections will be computed neglecting the decuplet–octet mass splitting. As with the spin-$\frac{1}{2}$ baryon octet fields, the decuplet fields will be treated as heavy fermion fields in the effective chiral lagrangian. In order to include the spin-$\frac{3}{2}$ decuplet, it is necessary to introduce a Rarita–Schwinger field T^{μ}_{abc} which contains both a Lorentz index and a spinor index, and obeys the constraint $\gamma_\mu T^{\mu}_{abc}=0$. The indices abc represent SU(3) tensor indices which are completely symmetrized to form an SU(3) decuplet. Since we are interested in renormalization of axial baryon octet currents, decuplet fields will only appear as internal lines of Feynman diagrams. Although it is easy to write down the most general lowest order lagrangian involving decuplet fields, we will introduce only those terms which are needed for the present calculation. The π, ξ and Σ fields are defined by

$$\pi=\frac{1}{\sqrt{2}}\times\begin{pmatrix} \frac{1}{\sqrt{2}}\pi^0+\frac{1}{\sqrt{6}}\eta & \pi^+ & K^+ \\ \pi^- & -\frac{1}{\sqrt{2}}\pi^0+\frac{1}{\sqrt{6}}\eta & K^0 \\ K^- & \bar{K}^0 & -\frac{2}{\sqrt{6}}\eta \end{pmatrix}, \tag{1}$$

$$\xi=\exp(\mathrm{i}\pi/f)\,,\quad \Sigma=\xi^2=\exp(2\mathrm{i}\pi/f)\,, \tag{2}$$

where $f\approx 93$ MeV. The combinations of pion fields which transform as a vector and an axial vector,

$$V^\mu=\tfrac{1}{2}(\xi\,\partial^\mu\xi^\dagger+\xi^\dagger\,\partial^\mu\xi)\,,$$
$$A^\mu=\tfrac{1}{2}\mathrm{i}(\xi\partial^\mu\xi^\dagger-\xi^\dagger\,\partial^\mu\xi)\,, \tag{3}$$

are useful for constructing the lagrangian. The lagrangian for the decuplet field is given by

$$L^{10}=\mathrm{i}\bar{T}^\mu \not{\mathscr{D}} T_\mu - m_T\bar{T}^\mu T_\mu+\mathscr{C}(\bar{T}^\mu A_\mu B+\bar{B}A_\mu T^\mu)$$
$$+\mathscr{H}\bar{T}^\mu\gamma_\nu\gamma_5 A^\nu T_\mu\,, \tag{4}$$

where T^μ is the spin-$\frac{3}{2}$ decuplet containing the Δ, Σ^*, Ξ^* and Ω^- isospin multiplets and B is the spin-$\frac{1}{2}$ baryon octet. SU(3) indices have been suppressed. Note that the $\mathscr{C}$ term is parity invariant because T^μ is an axial vector under parity. The covariant derivative acting on T^μ is

$$\mathscr{D}^\nu T^\mu_{abc}=\partial^\nu T^\mu_{abc}+(V^\nu)^d_a T^\mu_{dbc}+(V^\nu)^d_b T^\mu_{adc}$$
$$+(V^\nu)^d_c TA^\mu_{abd}\,. \tag{5}$$

Under an $\mathrm{SU(3)_L\times SU(3)_R}$ transformation, the fields in (4) transform as follows:

$$\Sigma\to L\Sigma R^\dagger\,,\quad \xi\to L\xi U^\dagger=U\xi R^\dagger\,,$$
$$B\to UBU^\dagger\,,\quad T^\mu_{abc}\to U^d_a U^e_b U^f_c T^\mu_{def}\,, \tag{6}$$

where U is defined by the transformation law of ξ above.

In lagrangian (4), the baryon decuplet and octet have not been treated as heavy static fields. The lagrangian can be rewritten in terms of new fields T^μ_v and B_v with definite velocity v^μ which are related to the original fields T^μ and B by [3]

$$T^\mu_v(x)=\exp(\mathrm{i}m_T\not{v}v_\mu x^\mu)\;T^\mu(x)\,,$$
$$B_v(x)=\exp(\mathrm{i}m_B\not{v}v_\mu x^\mu)\;B(x)\,. \tag{7}$$

In terms of these velocity-dependent fields, lagrangian (4) becomes (neglecting the baryon–octet mass difference m_T-m_B)

$$L^{10}_v=\mathrm{i}\bar{T}^\mu_v\not{\mathscr{D}}T_{v\mu}+\mathscr{C}(\bar{T}^\mu_v A_\mu B_v+\bar{B}_v A_\mu T^\mu_v)$$
$$+\mathscr{H}\bar{T}^\mu_v\gamma_\nu\gamma_5A^\nu T_{v\mu}\,. \tag{8}$$

As in the case of the baryon field in ref. [1], the Dirac gamma matrix structure in lagrangian L^{10}_v can be eliminated using the velocity four-vector v^μ and the spin operators S^μ_v. The analysis is identical to that found in ref. [1] and will not be repeated here. It is important to note, however, that for the spin-$\frac{3}{2}$ fields, S^μ_v is not the total spin operator of the field. Rather

$$S^\alpha_v T^\mu_v=\tfrac{1}{2}[\sigma^\alpha-(\sigma\cdot v)v^\alpha]T^\mu_v\,,$$
$$S^2_v T^\mu_v=-\tfrac{3}{4}T^\mu_v\,. \tag{9}$$

Thus, the spin operators act only on the spinor por-

Fig. 1. Wavefunction renormalization graph. The solid triangle represents vertices proportional to $\mathscr{C}$ in L_v^{10}. Octet and decuplet fields are represented by single and triple lines respectively.

tion of the Rarita–Schwinger field. With this one caveat, the new lagrangian without gamma matrices follows directly from the previous analysis,

$$L_v^{10} = \mathrm{i}\bar{T}_v^\mu (v\cdot\mathscr{D}) T_{v\mu} + \mathscr{C}(\bar{T}_v^\mu A_\mu B_v + \bar{B}_v A_\mu T_v^\mu) + 2\mathscr{H}\bar{T}_v^\mu S_{v\nu} A^\nu T_{v\mu} . \tag{10}$$

The decuplet contribution to the axial current can be derived from L_v^{10} using the Noether procedure,

$$\begin{aligned} J_\mu^A &= \tfrac{1}{2} v^\mu \bar{T}_v^\nu (\xi T^A \xi^\dagger - \xi^\dagger T^A \xi) T_{v\nu} \\ &+ \mathscr{H}\bar{T}_v^\nu S_v^\mu (\xi T^A \xi^\dagger + \xi^\dagger T^A \xi) T_{v\nu} \\ &+ \tfrac{1}{2}\mathscr{C}\bar{T}_{v\mu} (\xi T^A \xi^\dagger + \xi^\dagger T^A \xi) B_v \\ &+ \tfrac{1}{2}\mathscr{C}\bar{B}_v (\xi T^A \xi^\dagger + \xi^\dagger T^A \xi) T_{v\mu} . \end{aligned} \tag{11}$$

The propagator for the Rarita–Schwinger field involves a polarization projector due to it extra Lorentz index. There are four positive energy spinor solutions $\mathscr{U}_i^\mu$, $i=1, ..., 4$, to the Dirac equation of motion consistent with the constraint $\gamma_\mu T^\mu = 0$. For the computation of Feynman graphs, we need the polarization sum

$$P_v^{\mu\nu} = \sum_i \mathscr{U}_i^\mu \bar{\mathscr{U}}_i^\nu = (v^\mu v^\nu - g^{\mu\nu}) - \tfrac{4}{3} S_v^\mu S_v^\nu . \tag{12}$$

The decuplet propagator is then given by $\mathrm{i}P_v^{\mu\nu}/(v\cdot k)$. Using the fact that $v\cdot S_v = 0$, the product of two projectors can be easily computed, $P_v^{\mu\nu} P_{v\nu}^{\lambda} = -P_v^{\mu\lambda}$. In addition, the identities

$$\begin{aligned} &P_v^{\mu\nu} v_\nu = P_v^{\mu\nu} v_\mu = 0 , \quad P_v^{\mu\nu} S_{v\nu} = S_{v\mu} P_v^{\mu\nu} = 0 , \\ &P_v^{\mu\nu} S_{v\mu} = -\tfrac{4}{3} S_v^\nu , \quad S_{v\nu} P^{\mu\nu} = -\tfrac{4}{3} S_v^\mu , \end{aligned} \tag{13}$$

are useful for the computation of Feynman diagrams.

We are now ready to compute the effects of the baryon decuplet on the renormalization of baryon octet axial vector currents. The additional Feynman graphs involving decuplet fields which contribute to the renormalization of the baryon octet axial vector currents are shown in figs. 1, 2. The calculation was performed in the isospin symmetry limit $m_u = m_d = 0$, and M_η^2 was rewritten as $\frac{4}{3} M_K^2$ using the Gell-Mann–Okubo formula. The renormalization of the axial current can be written in the form

$$\begin{aligned} &\langle B_i | J_\mu^A | B_j \rangle \\ &= \left[\alpha_{ij}^A + (\bar{\beta}_{ij}^A - \bar{\lambda}_{ij}\alpha_{ij}^A) \frac{M_K^2}{16\pi^2 f^2} \ln\left(\frac{M_K^2}{\mu^2}\right)\right] \\ &\times \bar{u}_{B_i} \gamma_\mu \gamma_5 u_{B_j} , \end{aligned} \tag{14}$$

where α_{ij}^A is the lowest order result, $\bar{\lambda}_{ij} = \lambda_{ij} + \lambda'_{ij}$ is the one-loop correction due to wavefunction renormalization,

$$\sqrt{Z_i Z_j} \doteq 1 + \bar{\lambda}_{ij} \frac{M_K^2}{16\pi^2 f^2} \ln\left(\frac{M_K^2}{\mu^2}\right) ,$$

$\bar{\beta}_{ij}^A = \beta_{ij}^A + \beta_{ij}'^A$ is the correction due to all other graphs, and u is a spinor. The coefficients α_{ij}^A, λ_{ij}, and β_{ij}^A due to intermediate octet states are given in ref. [1]. The coefficients λ'_{ij} and $\beta_{ij}'^A$ are the new contributions from the graphs in figs. 1, 2. The coefficients λ'_{ij} are given by

$$\begin{aligned} &\lambda'_{pn} = \tfrac{1}{2}\mathscr{C}^2 , \quad \lambda'_{\Lambda\Xi^-} = \tfrac{19}{12}\mathscr{C}^2 , \\ &\lambda'_{\Lambda\Sigma^-} = \tfrac{5}{3}\mathscr{C}^2 , \quad \lambda'_{n\Sigma^-} = \tfrac{17}{12}\mathscr{C}^2 , \\ &\lambda'_{\Xi^0\Xi^-} = \tfrac{13}{6}\mathscr{C}^2 , \quad \lambda'_{\Sigma^0\Xi^-} = \tfrac{27}{12}\mathscr{C}^2 = \lambda'_{\Sigma^+\Xi^0} , \\ &\lambda'_{p\Lambda} = \tfrac{3}{4}\mathscr{C}^2 , \quad \lambda'_{pp} = \tfrac{1}{2}\mathscr{C}^2 . \end{aligned} \tag{15}$$

Fig. 2 yields the coefficients $\beta_{ij}'^A$:

$$\begin{aligned} &\beta_{pn}'^{1+i2} = -\tfrac{10}{81}\mathscr{H}\mathscr{C}^2 + (\tfrac{2}{3}D + \tfrac{2}{9}F)\mathscr{C}^2 , \\ &\beta_{\Lambda\Sigma^-}'^{1+i2} = -\frac{5}{27\sqrt{6}}\mathscr{H}\mathscr{C}^2 + \frac{8}{3\sqrt{6}}(D+F)\mathscr{C}^2 , \\ &\beta_{\Xi^0\Xi^-}'^{1+i2} = \tfrac{20}{81}\mathscr{H}\mathscr{C}^2 + (\tfrac{14}{27}D + 2F)\mathscr{C}^2 , \end{aligned} \tag{16}$$

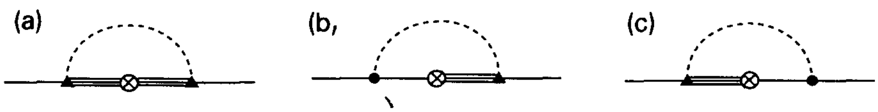

Fig. 2. The Feynman graphs which contribute to the renormalization of the axial currents of baryon octet fields. The solid triangle represents vertices proportional to $\mathscr{C}$ in L_v^{10} and the circled cross represents an insertion of the axial current (11). Octet and decuplet fields are represented by single and triple lines respectively. Graph (a) is proportional to $\mathscr{H}\mathscr{C}^2$ and graphs (b) are proportional to $\mathscr{C}^2$.

$$\beta'^{4+i5}_{p\Lambda}=\frac{5}{9\sqrt{6}}\mathscr{H}\mathscr{C}^2-\frac{1}{\sqrt{6}}(D+F)\mathscr{C}^2\,,$$

$$\beta'^{4+i5}_{\Lambda\Xi^-}=-\frac{5}{9\sqrt{6}}\mathscr{H}\mathscr{C}^2+\frac{1}{\sqrt{6}}(\tfrac{34}{18}D-F)\mathscr{C}^2\,,$$

$$\beta'^{4+i5}_{n\Sigma^-}=\tfrac{5}{81}\mathscr{H}\mathscr{C}^2+(-\tfrac{1}{27}D+F)\mathscr{C}^2\,,$$

$$\beta'^{4+i5}_{\Sigma^0\Xi^-}=-\frac{55}{81\sqrt{2}}\mathscr{H}\mathscr{C}^2+\frac{1}{\sqrt{2}}(\tfrac{19}{9}D+\tfrac{17}{9}F)\mathscr{C}^2$$

$$=\frac{1}{\sqrt{2}}\beta'^{4+i5}_{\Sigma^+\Xi^0}\,,$$

$$\beta'^{8}_{pp}=\frac{1}{\sqrt{3}}(D-F)\mathscr{C}^2\,. \qquad (16\text{ cont'd})$$

For completeness, eq. (15) and (16) include the values of λ'_{pp} and β'^{8}_{pp} which are relevant for analysis of the strangeness content of the proton.

A lowest order fit to the hyperon semileptonic decays gives $F\approx 0.47$ and $D\approx 0.81$ [4]. Chiral logarithmic corrections involving only baryon octet fields prefer the values $F\approx 0.33$ and $D\approx 0.56$ [1]. When intermediate decuplet states are included, there are two additional parameters, $\mathscr{C}^2$ and $\mathscr{H}$. The parameter $\mathscr{C}^2$ can be determined by a fit to the $\Delta\to N\pi$ decay rate, $|\mathscr{C}|\sim 1.6$. The parameter $\mathscr{H}$ is not determined experimentally, because the π–Δ coupling is unknown. Thus, the fit to semileptonic hyperon decays involves one additional parameter when intermediate decuplet states are allowed. For $|\mathscr{C}|=1.6$, the three parameter fit yields $F=0.40\pm 0.03$, $D=0.61\pm 0.04$ and $\mathscr{H}=-1.9\pm 0.7$ with $\chi^2=2.1$ for 11 degrees of freedom. The total chiral correction to the axial currents is significantly reduced by including intermediate decuplet states. The decays have corrections which are now less than 30%, with the exception of $\Sigma\to n$ (67%) and $\Lambda\to p$ (46%).

The radiative corrections to the baryon currents were computed above by including intermediate decuplet and octet states. This is a good approximation in the limit in which the decuplet–octet splitting is small compared with the K mass, but all the radially excited baryons have much higher energy. There is another way of approaching the problem of radiative corrections. One can consider the baryon as a weakly bound state of quarks. In this limit, the corrections are computed by *first* computing the radiative corrections to the quark currents, and then computing the matrix elements of the currents between baryon states. The graphs contributing to the renormalization of the baryon currents in this approach are shown in fig. 3. The pion couplings are proportional to the quark spin operator, and so can flip the quark's spin, though they do not affect the quark's momentum. Thus the intermediate states in the quark picture can be thought of as baryon states with all possible combinations of the quark spin, which is equivalent to including intermediate spin-$\frac{1}{2}$ and spin-$\frac{3}{2}$ states.

The quark calculation is not identical to the baryon calculation, however, the difference occurs in graphs such as those in fig. 4. In the quark picture, these graphs do not have a logarithmic divergence, and so do not contribute to the chiral log corrections. In the baryon picture, the intermediate state is to be considered as a baryon, and the graph is divergent. Which picture is appropriate depends on the momentum of the meson. If the momentum is small compared with the *binding energy*, then the baryon picture is correct; if it is large, then the quark picture is correct. The important scale for a composite object such as a baryon is the binding energy, not the mass of the object. If the binding is weak, then the pion is emitted from a quark, and is absorbed by another quark before the quarks can interact. This affects the quark–quark potential, but does not contribute to the renormalization of the current. If the binding is strong, then the quarks interact with each other strongly between the times the pion is emitted and absorbed. Thus the in-

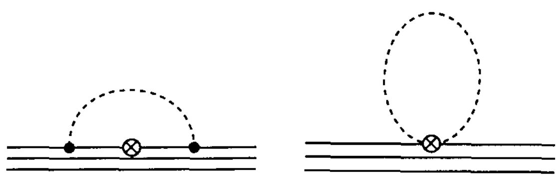

Fig. 3. Graphs contributing to chiral corrections to the baryon currents in the quark picture. The solid lines represent quark fields.

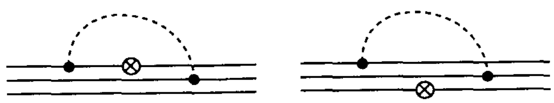

Fig. 4. Some of the graphs which contribute to chiral logarithms in the baryon picture but not in the quark picture.

termediate state is better thought of as a baryon, rather than as three quarks [#4].

The quark–pion lagrangian needed to compute corrections in the quark picture [5] is

$$L=\mathrm{i}\bar{\psi}(\partial\!\!\!/+V\!\!\!/)\psi+g_{\mathrm{A}}\bar{\psi}A\!\!\!/\gamma_5\psi + \tfrac{1}{4}f^2\,\mathrm{Tr}(\partial_\mu\Sigma\,\partial^\mu\Sigma^\dagger)+a\,\mathrm{Tr}\,M(\Sigma+\Sigma^\dagger)\,, \tag{17}$$

whence the quark axial currents are

$$J_\mu^{\mathrm{A}}=\tfrac{1}{2}v^\mu\bar{\psi}(\xi T^{\mathrm{A}}\xi^\dagger-\xi^\dagger T^{\mathrm{A}}\xi)\psi + g_{\mathrm{A}}\bar{\psi}S_v^\mu(\xi T^{\mathrm{A}}\xi^\dagger+\xi^\dagger T^{\mathrm{A}}\xi)\psi + \tfrac{1}{2}\mathrm{i}f^2\,\mathrm{Tr}(T^{\mathrm{A}}\,\partial_\mu\Sigma^\dagger\Sigma-T^{\mathrm{A}}\,\partial_\mu\Sigma\Sigma^\dagger)\,. \tag{18}$$

The renormalized quark current can be written as

$$\langle\psi_i|J_\mu^{\mathrm{A}}|\psi_j\rangle = \left[\alpha_{ij}^{\mathrm{A}}+(\beta_{ij}^{\mathrm{A}}-\lambda_{ij}\alpha_{ij}^{\mathrm{A}})\,\frac{M_K^2}{16\pi^2f^2}\ln\left(\frac{M_K^2}{\mu^2}\right)\right] \times\bar{u}_i\gamma_\mu\gamma_5u_j\,, \tag{19}$$

$$\alpha_{ud}^{1+i2}=\alpha_{us}^{4+i5}=g_{\mathrm{A}}\,, \tag{20}$$

$$\beta_{ud}^{1+i2}=-\tfrac{1}{9}g_{\mathrm{A}}^3-\tfrac{1}{2}g_{\mathrm{A}}\,,\quad \beta_{us}^{4+i5}=\tfrac{2}{9}g_{\mathrm{A}}^3-\tfrac{5}{4}g_{\mathrm{A}}\,, \tag{21}$$

$$\lambda_{ud}=\tfrac{11}{6}g_A^2\,,\quad \lambda_{us}=\tfrac{37}{12}g_{\mathrm{A}}^2\,. \tag{22}$$

The renormalization of the flavor diagonal quark currents is more complicated, because the singlet and octet currents can mix.

$$\begin{pmatrix}J_0^\mu\\ J_8^\mu\end{pmatrix}\to\begin{pmatrix}1-x\frac{32}{9}g_{\mathrm{A}}^2 & x\frac{10}{9}g_{\mathrm{A}}^2\\ x\frac{20}{9}g_{\mathrm{A}}^2 & 1-x(\frac{19}{6}g_{\mathrm{A}}^2-\frac{3}{2})\end{pmatrix}\begin{pmatrix}J_0^\mu\\ J_8^\mu\end{pmatrix}, \tag{23}$$

where

$$\begin{pmatrix}J_0^\mu\\ J_8^\mu\end{pmatrix}=\begin{pmatrix}\bar{u}\gamma^\mu\gamma_5u+\bar{d}\gamma^\mu\gamma_5d+\bar{s}\gamma^\mu\gamma_5s\\ \bar{u}\gamma^\mu\gamma_5u+\bar{d}\gamma^\mu\gamma_5d-2\bar{s}\gamma^\mu\gamma_5s\end{pmatrix}, \tag{24}$$

and

$$x=\frac{M_K^2}{16\pi^2f^2}\ln\left(\frac{M_K^2}{\mu^2}\right). \tag{25}$$

A fit without the chiral corrections gives $g_{\mathrm{A}}\approx0.8$. A fit including chiral corrections gives $g_{\mathrm{A}}\approx0.56$, with $\xi^2=2$ for 13 degrees of freedom. The strangeness changing decays have 50% corrections, and the strangeness conserving decays have 30% corrections. Note that the quark calculation, which has only one free parameter, fits the data as well as the three parameter fit using baryons.

The quark calculation can be compared with the baryon calculation, by computing the matrix elements of the quark currents in the baryon using nonrelativistic quark model wavefunctions [#5]. F, D, $\mathscr{C}$, and $\mathscr{H}$ are determined in terms of g_{A},

$$F=\tfrac{2}{3}g_{\mathrm{A}}\,,\quad D=g_{\mathrm{A}}\,,\quad \mathscr{C}=-2g_{\mathrm{A}}\,,\quad \mathscr{H}=-3g_{\mathrm{A}}\,, \tag{26}$$

which agrees with the parameters determined by the baryon fit. However, one can see that taking baryon matrix elements of quark currents (19) is not equivalent to substituting the parameter values (26) into the baryon currents (14) because the graphs of fig. 4 are included in the baryon calculation but not in the quark calculation.

In summary, we have computed the chiral logarithmic corrections to the baryon axial currents in two different ways, using intermediate octets and decuplets, and using matrix elements of renormalized quark currents. In either case, the radiative corrections are significantly smaller than a calculation including intermediate octet states alone. This provides indirect evidence that intermediate states which differ from the initial/final baryon by a quark spin flip are important. However, it is not possible to decide from our calculation whether using octets and decuplets is better than using quarks. (The one advantage of the quark calculation is that it has fewer free parameters.) In either of the two calculations, some of the radiative corrections are still substantial (~50%).

[#4] This remark is clearer in the case of a weakly bound state such as positronium. The QED corrections due to momenta higher than the binding energy (e.g. due to electron vacuum polarization diagrams) should be calculated in terms of electron operators. The resulting theory is then combined with the positronium wave function to compute radiative corrections to positronium spectra. It would not be useful to compute the loop effects using intermediate positronium states.

[#5] The complete quark model calculation must also include SU(3) breaking in the quark wavefunctions (which does not have a chiral log) which has been computed by Donoghue et al. [6], but will not be included here.

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