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Factorization contributions and the breaking of the $\Delta I = \frac{1}{2}$ rule in weak $\Lambda N \rho$ and $\Sigma N \rho$ couplings

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We compute the modified factorization contributions to the $\Lambda \to N\rho$ and $\Sigma \to N\rho$ couplings and demonstrate that these contributions naturally include $\Delta I = \frac{3}{2}$ terms which are comparable ($\simeq 0.4$ to -0.8 times) in magnitude to the corresponding $\Delta I = \frac{1}{2}$ terms. As a consequence, we conclude models which treat vector meson exchange contributions to the weak conversion process $\Lambda N \to NN$ assuming such weak couplings to satisfy the $\Delta I = \frac{1}{2}$ rule are unlikely to be reliable.

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The $\Delta I = \frac{1}{2}$ rule is a prominent feature of observed $\Delta S = 1$ nonleptonic weak interactions (K decay and hyperon decay). Not only is the ratio of $\Delta I = \frac{1}{2}$ to $\Delta I = \frac{3}{2}$ amplitudes considerably enhanced over that of the corresponding un-QCD-modified operator strengths, but also the nonleptonic decays completely dominate semileptonic decay modes, indicating a significant enhancement of the $\Delta I = \frac{1}{2}$ amplitudes. As a consequence of this observation, it has become conventional, in the absence of other evidence, to assume the validity of the $\Delta I = \frac{1}{2}$ rule for all $\Delta S = 1$ nonleptonic weak interactions. In particular, in the meson-exchange treatment of $\Lambda N \to NN$ it has been assumed that the relevant weak baryon-meson couplings satisfy the rule. In the case of the π couplings, this is known empirically, from hyperon decay, to be a valid assumption, but no similar experimental support exists for the assumption that vector meson couplings satisfy the rule. In this Brief Report we argue that, for the latter couplings, one may indeed expect significant violations of the $\Delta I = \frac{1}{2}$ rule. We base this statement on an evaluation of factorization contributions to the couplings and show below how, for such contributions, the structure of QCD modifications to the weak interactions are such as to distinguish naturally the pseudoscalar and vector cases.

As is well known, the effects of QCD on the $\Delta S = 1$ nonleptonic interactions can be taken into account perturbatively, down to a scale $\simeq 1$ GeV where the strong interactions begin to become truly strong [1–4] . One obtains, for the effective $\Delta S = 1$ nonleptonic Hamiltonian

$$\mathcal{H}_{\text{eff}} = -\sqrt{2} G \sin\theta_C \cos\theta_C \sum_{i=1}^{6} c_i O_i \ , \tag{1}$$

where the operators, O_i, have the form

$$\begin{aligned}
O_1 &= \bar{d}_L\gamma_\mu s_L\, \bar{u}_L\gamma^\mu u_L \,-\, \bar{u}_L\gamma_\mu s_L\, \bar{d}_L\gamma^\mu u_L \ , \\
O_2 &= \bar{d}_L\gamma_\mu s_L\, \bar{u}_L\gamma^\mu u_L \,+\, \bar{u}_L\gamma_\mu s_L\, \bar{d}_L\gamma^\mu u_L \\
&\quad + 2\, \bar{d}_L\gamma_\mu s_L\, \bar{d}_L\gamma^\mu d_L + 2\, \bar{d}_L\gamma_\mu s_L\, \bar{s}_L\gamma^\mu s_L \ , \\
O_3 &= \bar{d}_L\gamma_\mu s_L\, \bar{u}_L\gamma^\mu u_L \,+\, \bar{u}_L\gamma_\mu s_L\, \bar{d}_L\gamma^\mu u_L \\
&\quad + 2\, \bar{d}_L\gamma_\mu s_L\, \bar{d}_L\gamma^\mu d_L - 3\, \bar{d}_L\gamma_\mu s_L\, \bar{s}_L\gamma^\mu s_L \ , \\
O_4 &= \bar{d}_L\gamma_\mu s_L\, \bar{u}_L\gamma^\mu u_L \,+\, \bar{u}_L\gamma_\mu s_L\, \bar{d}_L\gamma^\mu u_L \\
&\quad - \bar{d}_L\gamma_\mu s_L\, \bar{d}_L\gamma^\mu d_L \ , \\
O_5 &= \bar{d}_L\gamma_\mu \lambda^a s_L (\bar{u}_R\gamma^\mu \lambda^a u_R \,+\, \bar{d}_R\gamma^\mu\lambda^a d_R \\
&\quad + \bar{s}_R\gamma^\mu\lambda^a s_R) \ , \\
O_6 &= \bar{d}_L\gamma_\mu s_L (\bar{u}_R\gamma^\mu u_R \,+\, \bar{d}_R\gamma^\mu d_R \,+\, \bar{s}_R\gamma^\mu s_R) \ ,
\end{aligned} \tag{2}$$

and the coefficients, c_i, are scale dependent and calculable perturbatively. The operators $O_1, \ldots, O_6$ in Eq. (2) have the specific (*flavor*, *isospin*) quantum numbers $(8, \frac{1}{2})$, $(8, \frac{1}{2})$, $(27, \frac{1}{2})$, $(27, \frac{3}{2})$, $(8, \frac{1}{2})$, and $(8, \frac{1}{2})$, respectively. The operators $O_{5,6}$, with left-right (LR) chiral structure are due to penguin graphs. The leading (O_1) term typically has a coefficient, $c_1 \simeq 4\, c_4$ at a scale $\simeq 1$ GeV, O_4 being the only $\Delta I = \frac{3}{2}$ operator, indicating that a portion of the observed experimental enhancement results from QCD modifications of the relative operator strengths. The additional factor of $\simeq$ 4–5 in the observed amplitude ratios not accounted for by this modification, however, must be associated with specific dynamics in the matrix elements of the operators. In the case of K decay, it seems likely that a significant portion of this dynamical enhancement is associated with final state interactions (FSI), the $\Delta I = \frac{1}{2}$ operators leading to the attractive $I = 0$ $\pi\pi$ s-wave final state, the $\Delta I = \frac{3}{2}$ operator to the replusive $I = 2$ state [5–7]. A similar explanation is not, however, tenable for hyperon decays, since, at least for Λ and Σ , the final state phases are known to be small. An old idea [8], which provides an attractive (if qualitative) alternative for these decays, is based on the observation that there are large enhancements of the penguin operator matrix elements in what is usually called the factorization approximation, these enhancements resulting from the different, LR, chiral structure of these operators. We briefly describe this approximation below.

In the approximation that FSI may be neglected, one may separate the graphs contributing to baryon-meson, $B' \to BM$, weak transitions into two classes: "external," describing those graphs in which both a quark and antiquark line from the effective quark-level weak vertex end up in the final state meson, M, and "internal," describing all other graphs. The advantage of this classification is that the external contributions are effectively

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"factorized" into a product of the matrix elements of two currents, one connecting B' to B, and one connecting M to the vacuum. These current matrix elements are completely known in terms of baryon semileptonic decay form factors and meson decay constants, so that the external or factorization contributions may be reliably calculated. This is not the case for the internal contributions. (Note that it is essential to use the modified form of factorization, in which Fierz rearrangements of the O_i are also taken into account, or one will fail to satisfy the correct isospin relations between factorization matrix elements of the operators.) As mentioned above, when M is a pion, the Fierz-rearranged contributions of $O_{5,6}$ (the nonrearranged expectation is zero because the $\Delta S = 1$ current portion of the nonrearranged form is a flavor singlet) contain a large enhancement relative to the matrix elements of the LL operators. Since these operators are pure $\Delta I = \frac{1}{2}$ this provides an attractive qualitative explanation of the $\Delta I = \frac{1}{2}$ rule, especially when combined with the observation that the $\Delta I = \frac{3}{2}$ pieces of the internal contributions would vanish in the naive quark model limit (in which the baryons contain only the leading three-quark color-singlet Fock space component) owing to the color symmetry of the $\Delta I = \frac{3}{2}$ operator, O_4. The prescription of simply ignoring the internal contributions is called the "factorization approximation." The predictions of this approximation are actually rather ill defined, since the values of the Wilson coefficients of the penguin operators, which arise from the evolution below a scale $\simeq m_c$, are quite sensitive to the precise scale chosen, making amplitudes where these terms occur multiplied by a large enhancement factor also quite sensitive to the scale choice. What can more safely be determined are factorization contributions to quantities which do not involve the (enhanced) penguin operators, for example, the $\Delta I = \frac{3}{2}$ contributions to hyperon decay. Here, if one takes the coefficient c_4 to be evaluated at a scale of 1 GeV and uses the true $\Delta I = \frac{3}{2}$ amplitudes obtained after making corrections for Σ-Λ and π_3-π_8 mixing in the physical amplitudes [9] (the p-wave Λ and Ξ amplitudes are increased by $\simeq 400\%$ and $\simeq 100\%$ by these corrections), one finds (1) good fits to the s- and p-wave $\Delta I = \frac{3}{2}$ Ξ amplitudes and s-wave Σ triangle discrepancy (2) that the p-wave Σ triangle discrepancy is underestimated by a factor of 4, and (3) that the s- and p-wave Λ amplitudes are overestimated by a factor of $3-4$ (the p-wave factorization contribution also being opposite in sign to the experimental value). Although one should bear in mind that the experimental errors on the $\Delta I = \frac{3}{2}$ amplitudes are rather large (apart from the Λ s wave, the factorization predictions fall within $\simeq 2\sigma$ of the central experimental value), it seems safe to conclude that, while the factorization contributions are of roughly the correct magnitude, there are additional non-negligible contributions from the internal graphs (even for $\Delta I = \frac{3}{2}$). In what follows we will be treating weak vector meson couplings for which, as we will see below, the factorization contributions of the penguin operators vanish. We will then see that the remaining factorization contributions involve large violations of the $\Delta I = \frac{1}{2}$ rule and, in light of the above discussion, argue that one should, therefore, expect some portion of this violation to survive in the total couplings.

Let us turn to the evaluation of the factorization contributions to the weak $\Lambda N\rho$ and $\Sigma N\rho$ couplings (the corresponding contributions to both $\Lambda N\omega$ and $\Sigma N\omega$ couplings are small and, even for $\Sigma N\omega$, satisfy the $\Delta I = \frac{1}{2}$ rule, so we will not discuss them further). We employ the effective weak Hamiltonian of Eqs. (1) and (2), with coefficients c_i evaluated at a scale 1 GeV [4,10] : $c_1 = -1.90$, $c_2 = 0.14$, $c_3 = 0.10$, $c_4 = 0.49$. The coefficients $c_{5,6}$ are not needed since the factorization contributions to the weak ρ couplings of the operators $O_{5,6}$ vanish. This fact follows from the observation that these operators contain color-singlet flavor-octet nonstrange currents only in Fierz-rearranged form, in which form, owing to the original LR chiral structure, only scalar and pseudoscalar currents are involved, the vacuum-to-ρ matrix elements of which automatically vanish. The remaining contributions are straightforward to work out.

We define the effective weak couplings via

$$\langle N(p')\rho|H_{\text{eff}}|Y(p)\rangle = \epsilon_\mu^{(\rho)*}\bar{u}_N(p')\left[f_1^w\gamma^\mu - i\frac{\sigma^{\mu\nu}q_\nu}{2m_N}f_2^w + g_1^w\gamma^\mu\gamma_5 - i\frac{\sigma^{\mu\nu}q_\nu}{2m_N}g_2^w\gamma_5\right]u_Y(p)\ , \tag{3}$$

where $q = p-p'$ and $\epsilon_\mu^{(\rho)}$ is the ρ polarization vector, and the baryon transition form factors via

$$\langle B'(p')|\ V_\mu - A_\mu\ |B(p)\rangle = \bar{u}_{B'}(p')\left[f_1\gamma_\mu - i\frac{\sigma_{\mu\nu}q^\nu}{2m_N}f_2 + g_1\gamma_\mu\gamma_5 + i\frac{q_\mu\gamma_5}{2m_N}g_3\right]u_B(p) \tag{4}$$

where we have dropped the 2nd class current form factors, f_3 and g_2 in Eq. (4). In using Eq. (4) below, we will assume that f_1 and f_2 are given by their conserved vector current (CVC) values, and take g_1/f_1 from hyperon semileptonic decay data [11]. g_3 does not enter the expressions for the factorization contributions to the couplings due to the transversality of the ρ polarization vector. Defining the ρ decay constant, f_ρ, by

$$\langle O|V_\mu^3|\rho^0(\vec{q})\rangle = f_\rho m_\rho^2\epsilon_\mu^{(\rho)}(\vec{q}), \tag{5}$$

we then obtain for the factorization contributions to the $\Lambda \to p\rho^-$ couplings

$$f_i^w = \sqrt{2}K\left[-\tfrac{2}{3}c_1 + \tfrac{4}{3}c_2 + \tfrac{4}{3}c_3 + \tfrac{4}{3}c_4\right]f_i^{\Lambda p} \quad (i=1,2)$$
$$g_1^w = f_1^w\,[g_1^{\Lambda p}/f_1^{\Lambda p}], \tag{6}$$

with

$$K = \frac{G_F}{4\sqrt{2}}\sin(2\theta_C)f_\rho m_\rho^2\ , \tag{7}$$

where G_F, θ_C are the Fermi constant and Cabibbo angle, respectively, and $f_i^{\Lambda p}$, $g_i^{\Lambda p}$ the form factors relevant to $\langle p|\bar{u}(\gamma_\mu - \gamma_\mu\gamma_5)s|\Lambda\rangle$. We have, from CVC, $f_1^{\Lambda p}(0) = -\sqrt{\frac{3}{2}}$, $f_2^{\Lambda p}(0)/f_1^{\Lambda p}(0) = 1.63$ and, from Λ semileptonic decay data, $g_1^{\Lambda p}(0)/f_1^{\Lambda p}(0) = -0.72$. As mentioned above, $g_2^{B'B}$ is 2nd class, and assumed to be zero. Since $O_1, \ldots, O_3$ are $\Delta I = \frac{1}{2}$ and O_4 $\Delta I = \frac{3}{2}$

the corresponding $\Lambda \to n\rho^0$ contributions follow from isospin Clebsch-Gordan coefficients, leading to

$$f_i^w = -K\,[-\tfrac{2}{3}c_1 + \tfrac{4}{3}c_2 + \tfrac{4}{3}c_3 - \tfrac{8}{3}c_4]\,f_i^{\Lambda p} \quad (i = 1,2)$$
$$g_1^w = f_1^w\,[g_1^{\Lambda p}/f_1^{\Lambda p}]\,. \tag{8}$$

One finds similarly, for the factorization contributions to the $\Sigma^- \to n\rho^-$ couplings,

$$f_i^w = \sqrt{2}K\,[-\tfrac{2}{3}c_1 + \tfrac{4}{3}c_2 + \tfrac{4}{3}c_3 + \tfrac{4}{3}c_4]\,f_i^{\Sigma^- n} \quad (i = 1,2)$$
$$g_1^w = f_1^w\,[g_1^{\Sigma^- n}/f_1^{\Sigma^- n}]\,, \tag{9}$$

where, from CVC, $f_1^{\Sigma^- n}(0) = -1$, $f_2^{\Sigma^- n}(0)/f_1^{\Sigma^- n}(0) = -1.86$ (compatible with experiment [11]) and, from experiment, $g_1^{\Sigma^- n}(0)/f_1^{\Sigma^- n}(0) = 0.34$ [11] . From the form of the operators $O_1, \ldots, O_4$, one sees that there are no terms containing simultaneously both a $\bar{u}$ and d field, and hence both the $\Delta I = \frac{1}{2}$ and $\Delta I = \frac{3}{2}$ factorization contributions to the $\Sigma^+ \to n\rho^+$ couplings vanish. This reduces the number of independent reduced matrix elements for both the $\Delta I = \frac{1}{2}$ and $\Delta I = \frac{3}{2}$ operators from two to one, and one may then show that the $\Delta I = \frac{1}{2}$ factorization contributions to the $\Sigma N\rho$ couplings are in the ratios $1 : \frac{1}{\sqrt{2}} : \frac{1}{\sqrt{2}} : \frac{1}{2}$ and the $\Delta I = \frac{3}{2}$ contributions in the ratios $1 : -\sqrt{2} : \frac{1}{\sqrt{2}} : -1$, for $\Sigma^- \to n\rho^-$, $\Sigma^+ \to p\rho^0$, $\Sigma^0 \to p\rho^-$, and $\Sigma^0 \to n\rho^0$, respectively. Expressions for the weak $\Sigma^+ \to p\rho^0$, $\Sigma^0 \to p\rho^-$, and $\Sigma^0 \to n\rho^0$ couplings are then readily obtained from those of Eq. (9).

From expressions (6), (8), (9), and the discussion below Eq. (9), we see that the relative strength of the $\Delta I = \frac{1}{2}$ to $\Delta I = \frac{3}{2}$ contributions to the weak couplings is determined, in all cases, by the factor $[-\frac{2}{3}c_1 + \frac{4}{3}c_2 + \frac{4}{3}c_3 + \frac{4}{3}c_4]$ for the ρ^- couplings, and $[-\frac{2}{3}c_1 + \frac{4}{3}c_2 + \frac{4}{3}c_3 - \frac{8}{3}c_4]$ for the ρ^0 couplings. For the ρ^- couplings, using the c_i values quoted above (corresponding to a scale 1 GeV), the contributions of the 27_F c_3 and c_4 terms are 0.54 times those of the 8_F c_1 and c_2 terms. The leading $\Delta I = \frac{3}{2}$ term is 0.52 times the leading $\Delta I = \frac{1}{2}$ term and 0.41 times the net $\Delta I = \frac{1}{2}$ contribution. Similarly, for the ρ^0 couplings, the 27_F contributions are -0.81 times the 8_F contributions, while the leading $\Delta I = \frac{3}{2}$ contribution is -1.03 times the leading $\Delta I = \frac{1}{2}$ contribution and -0.82 times the net $\Delta I = \frac{1}{2}$ contribution. The factorization contributions to the weak $\Lambda N\rho$ and $\Sigma N\rho$ couplings thus badly violate the $\Delta I = \frac{1}{2}$ rule. The basic reason for this is the complete absence of the penguin contributions, which had large enhancements in the π coupling case.

We present the numerical results for the factorization contributions to the $\Lambda N\rho$ and $\Sigma N\rho$ weak couplings f_1^w, f_2^w, and g_1^w in Table I. The ratios reflect the strong violation of the $\Delta I = \frac{1}{2}$ rule discussed above. This violation will, however, be reflected in the full couplings only if the factorization contributions represent a moderate to sizable fraction of the full couplings. To see whether or not this is likely to be the case, we consider two existing models which have made predictions for the $\Lambda N\rho$ (though not the $\Sigma N\rho$) couplings [12,13]. In Table II we compare the factorization contributions to the $\Lambda \to p\rho^-$ couplings obtained above with the values obtained in the models of Refs. [12,13] (adjusted to Particle Data Group conventions for γ_5). We will return to a brief discussion of the models below, but for the moment, two features of the table are of note. First, the model predictions differ considerably, most significantly for the parity-violating g_1^w coupling. Second, the factorization contributions are $\simeq \frac{1}{3}$ of the full predictions for f_1^w, f_2^w and between $\simeq \frac{1}{4}$ and 1 times the full prediction for g_1^w. Also, as we will discuss below, there are significant uncertainties in the model predictions. From Table II it thus appears to us extremely unlikely that one can ignore the $\Delta I = \frac{3}{2}$ components of the $\Lambda N\rho$ and $\Sigma N\rho$ couplings.

TABLE I. Factorization contributions to the weak $\Lambda N\rho$, $\Sigma N\rho$ couplings. All entries in units of 10^{-7}.

Process	f_1^w	f_2^w	g_1^w
$\Lambda \to p\rho^-$	-4.1	-6.6	2.9
$\Lambda \to n\rho^0$	0.36	0.58	-0.26
$\Sigma^+ \to p\rho^0$	-0.29	0.54	-0.10
$\Sigma^+ \to n\rho^+$	0	0	0
$\Sigma^0 \to p\rho^-$	-2.3	4.3	-0.79
$\Sigma^0 \to n\rho^0$	-0.21	0.38	-0.07
$\Sigma^- \to n\rho^-$	-3.3	6.1	-1.1

A few words are in order concerning the models of Refs. [12,13] which we have used to gauge the potential importance of the factorization contributions. The model of Ref. [12] provides the framework for the weak couplings of the meson-exchange treatment of $\Lambda N \to NN$ employed by Dubach *et al.* [14]. Here the parity-violating (PV) g_1^w coupling is obtained from the known PV $\Lambda p\pi^-$ and $\Sigma^+ p\pi^o$ couplings via the $SU(6)_w$ treatment of Desplanques, Donoghue, and Holstein [15], where factorization estimates have been used to provide values for the two $SU(6)_w$ reduced matrix elements, a_V and a_T, present in the $\Lambda N\rho$ couplings but not in the $\Lambda N\pi$ and $\Sigma N\pi$ couplings [15]. The parity-conserving (PC) f_1^w, f_2^w couplings are obtained via a pole model analysis which includes ground state baryon pole terms and K^* pole terms, the strong vector meson couplings being obtained from an $SU(3)_F$ vector dominance model (VDM) treatment together with the weak baryon-baryon transition matrix elements from an analogous treatment of the PC $\Lambda \to N\pi$ and $\Sigma \to N\pi$ amplitudes. In Ref. [13], the PV g_1^w coupling is obtained from a pole model treatment which keeps only baryon poles belonging to the $(70, 1^-)$ multiplet of (ordinary) SU(6). The required weak baryon transition matrix elements between ground state and negative parity excited state baryons are taken from a treatment [16] of hyperon s-wave π decays which includes the leading commutator terms as well as the negative parity baryon poles, and which fixes the PV baryon-baryon matrix elements by assuming (1) that the D/F ratio for the

TABLE II. Comparison of $q^2 = 0$ factorization and model values for $\Lambda p\rho^-$ couplings. All entries in units of 10^{-7}. The models are as discussed in the text.

Coupling	Factorization	Ref. [12]	Ref. [13]
f_1^w	-4.1	-15.0	-10.6
f_2^w	-6.6	-22.6	-24.3
g_1^w	2.9	3.4	12.0

weak baryon transitions is -1 and (2) that experimentally observed deviations from a modified Lee-Sugawara sum rule are due entirely to the negative parity baryon poles. The required strong couplings are obtained using VDM arguments, together with information on the scalar multipoles in S_{11}, S_{11}' electroproduction. The PC f_1^w, f_2^w couplings are obtained using a pole model treatment which includes $\frac{1}{2}^{+*}$ baryon resonance pole contributions, in addition to the ground state baryon pole and K^* pole terms of Ref. [12]. The strong couplings required are obtained again using VDM arguments, together with data on P_{11} radiative decays and assumptions about the scalar multipoles in P_{11} electroproduction. The weak baryon-baryon couplings for the $\frac{1}{2}^{+*}$ pole terms are taken from a pole term analysis of hyperon p-wave π decays [16] which (1) assumes an F/D ratio of -1 for the weak baryon transitions and (2) fixes the overall strength by optimizing the full fit to the experimental p-wave amplitudes. This fit, however, employs a K-π weak transition strength in its K-pole graphs an order of magnitude greater than that extracted from $K \to \pi\pi$ [17], which makes the whole procedure appear somewhat dubious. The $\frac{1}{2}^{+*}$ pole contributions to f_1^w in Ref. [13] are negligible, but this is not true of the corresponding contributions to f_2^w. Finally, the K^* pole contributions are obtained using VDM plus $SU(3)_F$ arguments for the strong K^* couplings and a factorization treatment, which keeps only the O_1, O_2 terms of $\mathcal{H}_{\rm eff}$ and drops the Fierz-rearranged contributions, for the $K^* - \rho$ weak transition. This is, in fact, a rather suspect way to treat factorization contributions, even if they were expected to represent well the full coupling. Indeed, if one keeps all terms in $\mathcal{H}_{\rm eff}$, one finds the same linear combinations of the c_i occurring for the charged and neutral K^*-ρ mixing terms as occur for the charged and neutral ρ weak couplings above; i.e., there is very significant breaking of the $\Delta I = \frac{1}{2}$ rule for the K^* pole terms.

As can be seen from the discussion above, there are many assumptions and approximations which enter the models of Refs. [12,13]. As such, the model values for the weak couplings quoted in Table II will involve significant uncertainties whose sizes are difficult to quantify. We feel, however, that the arguments leading to the ground state baryon pole contributions, especially to f_1^w, are likely to be the most reliable, so these contributions provide a useful benchmark. For $\Lambda \to p\rho^-$, these are -10.5×10^{-7} and -11.3×10^{-7} for the models of Refs. [12,13], respectively. Similarly, for the f_2^w $\Lambda \to p\rho^-$ coupling, these contributions are -14.8×10^{-7} and -9.3×10^{-7}, respectively. Note that only the n pole term contributes to f_1^w, but that both n and Σ^+ pole terms contribute to f_2^w. In the latter case, there is considerable cancellation between the two terms, which makes the actual result rather sensitive to possible $SU(3)_F$ breaking in the relative strengths of the strong $pn\rho^-$ and $\Sigma^+\Lambda\rho^-$ couplings, a $\pm30\%$ variation of the the $\Sigma^+\Lambda\rho^-$ strength from its $SU(3)_F$ value, for example, producing a variation of $\pm7 \times 10^{-7}$ ($\pm10 \times 10^{-7}$) in the corresponding ground state baryon pole contribution to f_2^w for the parametrizations of Refs. [12,13], respectively. One should also note that the close agreement of the f_2^w values in the two models is actually a numerical accident, since the $\frac{1}{2}^{+*}$ baryon pole terms of Ref. [13], completely absent in Ref. [12], contribute $\simeq \frac{1}{3}$ of the quoted f_2^w value.

To summarize, we have evaluated the factorization contributions to the weak $\Lambda N\rho$ and $\Sigma N\rho$ couplings and find that they involve large violations of the $\Delta I = \frac{1}{2}$ rule. Since the size of these contributions is not small on the scale of values to be expected for the full couplings, it appears very unlikely that it is safe to assume the validity of the $\Delta I = \frac{1}{2}$ rule as an input when determining the weak $\Lambda N\rho$ couplings to be used in treating $\Lambda N \to NN$ in the meson-exchange framework.

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