Challenges in Hyperon Decays

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Abstract

We give a personal overview of some of the unsolved problems related to hyperon decays. We cover nonleptonic decays, radiative decays and magnetic moments. Some of the theoretical issues are also touched upon.

1 Introduction

While the frontier of high energy physics marches on to higher and higher energies, there are many unsolved puzzles left over in low-energy physics. Hyperon decays account for many of them. These puzzles are generally not considered serious enough to deserve major attention. Our inability to calculate accurately with the strong interaction is usually the reason invoked to explain why these puzzles remain, and why they may not be directly relevant to our understanding of the fundamental physics.

Nevertheless, one should be reminded that, even if their solutions do not require changes in the fundamental physics, it is still of utmost importance to understand the dynamics of low-energy hadronic physics through the thicket of the strong interaction. In addition, one should keep in mind that surprises can potentially occur that alter our understanding at a fundamental level. One example is the potential new source of CP violation that may reveal itself in hyperon decays. There are of course many other corners in which new physics may have already revealed itself in the data, but is not yet recognized by us because of our inability to quantify the strong interaction. One example that stands out is the $\Delta I = 1/2$ rule that had been repeatedly used in the literature to motivate new physics in the past.

Since hyperon decays is a vast and difficult subject [1], it is beyond my ability to review the field in detail. My goal is only to recount some studies with some personal perspective so as to help motivate more future experiments on this subject.

2 Problems and Tools

Various tools or models have been invented to understand hadronic dynamics without facing the strong-coupling issues head on. One can start by taking advantage of only the symmetry property which presumably is respected by the strong dynamics, except for the potential spontaneous symmetry breaking by vacuum condensates. One first identifies, or assumes, the low-energy degrees of freedom and their associated symmetry, and then proceeds to write down the most general low-energy effective theory describing these degrees of freedom consistent with the required symmetry. This is the approach taken by chiral perturbation theory (ChPT) [2] and the QCD sum rule [3]. These low-energy effective theories typically involve infinite numbers of terms, or couplings. To make such theories useful some expansion parameters and an associated cutoff scheme are necessary to derive predictions. For example in the simplest ChPT, the octet of Goldstone bosons associated with chiral-symmetry breaking are identified as the low-energy degrees of freedom, and using energy and momentum as expansion parameters, an infinite series of Lagrangian terms can be written down according to the chiral symmetry. The higher-dimensional operators are suppressed by powers of p/Λ_χ where Λ_χ is the chiral-symmetry-breaking scale. Similar expansion series are employed in the QCD-sum-rule approach.

Unfortunately, there is no guarantee either that the expansion is convergent, or that the degrees of freedom included are sufficient. This is especially the case when the energy involved in the hadronic process is not limited to very small p/Λ_χ and can sometimes be quite close to where the high resonance occurs. In such a case it is often impossible to justify the approximation or to account for the data without extra inputs. However, there are so far no reliable principles about how such extra inputs should be invoked or deployed. As a result, the theory has degenerated into models that are mostly invoked to account for particular special subsets of phenomenology. For example, it is well known that the lower-order chiral-perturbation theory cannot account for the data on hyperon radiative decays [4, 5] even after the heavy-baryon chiral-perturbation-theory formalism is employed [6, 7, 8]. In an attempt to resolve this puzzle, some additional resonances have been added in the analysis recently [9]. While the resulting "model" can account for the data (in fact it can account for the data in more than one way), it is not clear how much fundamental understanding is actually gained in the process. With these general comments on the theoretical difficulties, we can go over the specific challenges.

The challenges facing us in hyperon decays can be summarized in the following categories:

1. *CP* violation;
2. Semileptonic decays;
3. Nonleptonic decays;
4. Radiative decays;
5. Magnetic moments.

I shall leave the topic of CP violation to German Valencia and the topic of semileptonic decays to Earl C. Swallow in this workshop.

3 Nonleptonic Decays

The nonleptonic decays, $B \to B'\pi$, include:

$$\Lambda \to p\pi^- \ ; \ \ \Lambda \to n\pi^0 \ ;$$

$$\Sigma^+ \to n\pi^+ \ ; \ \ \Sigma^+ \to p\pi^0 \ ; \ \ \Sigma^- \to n\pi^- \ ;$$

$$\Xi^0 \to \Lambda\pi^0 \ ; \ \ \Xi^- \to \Lambda\pi^- \ .$$

The amplitude can be written in matrix form as

$$\mathbf{M} = \bar{u}_{B'}(p')(\mathbf{A} + \mathbf{B}\gamma_5)u_B(p)$$

where the **A** amplitudes are parity violating and S-wave, while the **B** amplitudes are parity conserving and P-wave.

3.1 $\Delta I = 1/2$ Rule

Each amplitude can be further decomposed into two isospin channels. For example, the A amplitude for $\Lambda \to p\pi^-$ can be decomposed into $A(\Lambda \to p\pi^-) = \sqrt{2}A^{1/2}_{\Lambda_-} - A^{3/2}_{\Lambda_-}$ and similarly for the B amplitudes. Experimentally [10], the ratio $A^{3/2}/A^{1/2}$ for each S-wave (A) or P-wave (B) amplitude is found to be small and of the order of that in the nonleptonic kaon decays. For example,

$$\frac{A^{3/2}}{A^{1/2}} = \begin{matrix} \\ \Lambda \\ \Sigma \\ \Xi \end{matrix} \begin{matrix} \begin{matrix} S & \quad P \end{matrix} \\ \begin{pmatrix} 0.014 & 0.006 \\ -0.017 & -0.047 \\ 0.034 & 0.023 \end{pmatrix} \end{matrix}$$

The fact that these ratios, as well as the corresponding one in kaon nonleptonic decays, are all of the order of 0.02 in absolute value indicates that the $\Delta I = 1/2$ rule may reflect something general about the strangeness-changing nonleptonic decays. However, so far we have not been able to understand this selection rule in general.

For example, in kaon decays, the factor of 20 in $\Delta I = 1/2$ enhancement is partially accounted for (by a factor of 3 to 4) by the renormalization group (RG) effect of QCD between the weak scale and the low-energy kaon scale. However, numerically this enhancement is still insufficient and this actually leads to many proposals in the literature of new physics to account for it.

For hyperons, the understanding improves a little bit. The similar RG enhancement effect is still operating. Pati and Woo, and Miura and Minamikawa [11], discovered another potential enhancement effect. They observed that, using a Fierz transformation, the $O^{\Delta I=3/2}$

operator that can mediate hyperon decay can be written in such a way that it is symmetric under the exchange of color indices of either the (u, s) pair or the (d, s) pair. If the baryons are taken naively as consisting of three valence quarks only, then the antisymmetry of the color wavefunction as well as this color property of the $\Delta I = 3/2$ operator imply that only the $\Delta I = 1/2$ operator can mediate the transition $< B' \,|\, H_W |\, B >$, where H_W is the weak Hamiltonian, that is, $< B' |\, O^{\Delta I = 3/2} |\, B > = 0$. Of course, $B \to B'$ matrix elements are not the only mechanism for $B \to B'\pi$ decays. And the baryon wavefunction is not exactly three-quark. However, it is a good indication that the $\Delta I = 1/2$ amplitudes are dynamically favored. The most unsatisfactory aspect of this explanation is probably that the same mechanism clearly does not apply to kaons. Therefore it does not provide a general understanding of the strangeness-changing nonleptonic decays. Nevertheless, with such a qualitative mechanism at hand, one can conclude that, overall, the $\Delta I = 1/2$ rule is probably not as serious a problem theoretically as in the case of kaons.

3.2 Nonleptonic enhancement

The data clearly show that the typical nonleptonic branching ratios, say, $\Gamma(\Lambda \to p\pi^-)$, are about 10^3 larger than the typical semileptonic ones, say, $\Gamma(\Lambda \to pe\bar{\nu}_e)$. Why? We have not been able to account for this based on fundamental principles. There were various attempts to resolve this difference using the quark model plus some QCD inputs many years ago. However, none can be considered compelling. This problem has not received very much recent attention. Nevertheless, I am wondering whether one can derive some better qualitative understanding by learning from the recent work on related issues in the charm and especially the B systems, such as heavy quark perturbation theory [12] or QCD-inspired factorization models [13].

3.3 The S-wave/P-wave problem

In chiral models or the current-algebra approach, S-wave ($\mathbf{A}$) is due to a contact term or, in ChPT language, the lowest order term in the chiral Lagrangian. It is typically represented as in Fig. 1a. Note that the same vertex does not contribute to the P-wave amplitudes, $\mathbf{B}$. To get P-wave amplitudes, one needs to invoke the pole diagrams in Figs. 1b and 1c. The lowest order weak Lagrangian gives

$$L_W^0 = h_D \mathrm{Tr}(\bar{\mathbf{B}}\{\lambda, \mathbf{B}\}) + h_F \mathrm{Tr}(\bar{\mathbf{B}}[\lambda, \mathbf{B}])$$

where $\lambda = \xi^\dagger \lambda_6 \xi$, $\xi^2 = \Sigma$ is the usual Goldstone boson in matrix form. L_W^0 contributes to S-wave amplitudes, A, directly, but not to B. The two parameters h_D and h_F can be used to fit the S-wave amplitudes quite well at this order. In particular, $A(\Sigma^+ \to n\pi^+)$ is found to be zero at this order while experimentally, indeed, it is much smaller than the others.

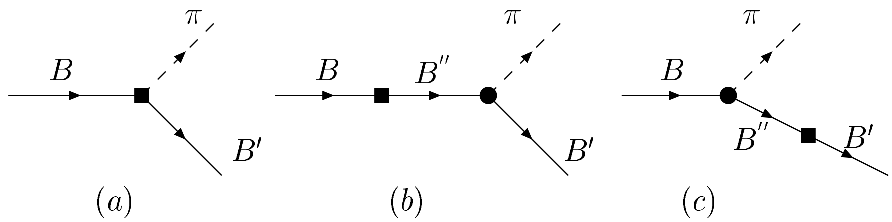

Figure 1. Diagrams for $B \to B'\pi$. Here circles and squares represent strong and weak vertices respectively.

Using the fitted h_D and h_F, together with strong-interaction parameters, one can calculate the pole diagrams and predict the P-wave amplitudes, **B**. The result turns out to be a poor fit to many of the **B** amplitudes extracted from the data. Actually, alternatively, one can also choose to use h_D and h_F to fit the P-wave amplitudes, and use them to predict S-wave amplitudes. The result is an equally poor fit. The is the famous S-wave/P-wave problem.

There are a few potential ways out of the puzzle. One can see if the inclusion of a higher-order effective Lagrangian can rectify the situation. However, there are just too many parameters in the next-order Lagrangian and there is no unique, convincing fit to the data. Alternatively, one can adopt a model that includes new resonances, such as spin-1/2, parity-odd resonances or spin-1/2, parity-even Roper resonances, as the intermediate states B'' in the pole diagrams [9]. The result can fit both S- and P-wave amplitudes, however the fit is clearly model dependent.

4 Radiative Decays

The radiative decays, with branching fractions as reported in [10], are:

$\Sigma^+(uus) \to p(uud)\gamma$	BF $\approx (1.23 \pm 0.05) \times 10^{-3}$	**
$\Sigma^0 \to n\gamma$	BF too small to see due to electromagnetic decay	
$\Sigma^0 \to \Lambda\gamma$	BF $\approx$ 100% electromagnetic decay	
$\Xi^0 \to \Sigma^0\gamma$	BF $\approx (3.5 \pm 0.4) \times 10^{-3}$	
$\Xi^0 \to \Lambda\gamma$	BF $\approx (1.18 \pm 0.30) \times 10^{-3}$	
$\Lambda \to n\gamma$	BF $\approx (1.75 \pm 0.15) \times 10^{-3}$	
$\Xi^- \to \Sigma^-\gamma$	BF $\approx (1.27 \pm 0.23) \times 10^{-4}$	**

(Note that ** are charged modes between U-spin doublets.)

Generally, the amplitudes can be rewritten as

$$A = -\frac{e}{M_B + M_{B'}} i\epsilon^{*\mu} q^\nu \bar{u}(p')\sigma_{\mu\nu}(C + D\gamma_5)u(p)$$

Table 1: Present status of decay rates and asymmetry parameters. The numbers are the combined weighted mean from Ref. [4]. Neither the decay rate nor the asymmetry parameter for $\Sigma^0 \to \Lambda + \gamma$ have been measured.

$B_i \to B_f + \gamma$	Γ [10^{-18} GeV]	α	Ref.
$\Lambda \to n + \gamma$	4.07 ± 0.35	–	[15]
$\Xi^0 \to \Lambda + \gamma$	2.4 ± 0.36	0.43 ± 0.44	[16]
$\Xi^0 \to \Sigma^0 + \gamma$	8.1 ± 1.0	0.20 ± 0.32	[17]
$\Sigma^+ \to p + \gamma$	10.1 ± 0.5	-0.76 ± 0.08	[18, 19]
$\Xi^- \to \Sigma^- + \gamma$	0.51 ± 0.092	1.0 ± 1.3	[20]

where C is the parity-conserving magnetic dipole transition (M1) and D is the parity-violating electric dipole transition (E1). The decay rate is $\Gamma \propto |C|^2 + |D|^2$. The asymmetry is $A_\gamma = \dfrac{2Re(C^*D)}{|C|^2 + |D|^2}$, that is, one needs both nonzero C and D amplitudes to get nonzero asymmetry. There is an old theorem [14] by Hara regarding the vanishing of D amplitudes.

4.1 Hara's Theorem

Hara's Theorem: Assuming that the amplitudes are not singular in the $SU(3)_F$ limit, parity violating D amplitudes must vanish for decays between states of a U-spin doublet in the $SU(3)_F$ limit.

The theorem implies that

$$D(\Sigma^+ \to p\gamma) = D(\Xi^- \to \Sigma^-\gamma) = 0$$

and as a result the asymmetries $A_\gamma(\Sigma^+ \to p\gamma) = A_\gamma(\Xi^- \to \Sigma^-\gamma) = 0$.

Accepting the assumption of Hara's theorem, even when $SU(3)$ breaking is taken into account, the asymmetry A_γ should remain small. Unfortunately, experimentally $A_\gamma(\Sigma^+ \to p\gamma)$, which is the only one measured to be nonzero, was found to be -0.76 ± 0.08, indeed large and negative. The others listed in Table 1 have much larger errors.

The hadronic models did not have a great deal of success in explaining the details of the data. All the models of this type (except those that include vector mesons) preserve Hara's theorem in their formulations. General analyses which include $SU(3)$ breaking [21] actually predict a small and positive asymmetry for the Σ^+ decay while the experimental result is negative and relatively large. Models that assume vector-meson dominance [22] can introduce effects that violate Hara's theorem due to mixing of vector mesons with the photon. In models using quarks, it was pointed out [23] that the diagrams in which a W boson is exchanged between two constituent quarks can give rise to violation of Hara's theorem. In addition, models which include vector-meson dominance are in better agreement with the data, though the situation is still not satisfactory. Among hadronic models, the observed negative asymmetry parameter for Σ^+ decay is best accounted for using QCD sum rules [3]. Other approaches can be found in Refs. [24, 25, 26]. Detailed overviews on both experimental and theoretical aspects of weak radiative decays of hyperons are given in Refs. [18, 4, 5].

Many papers [27, 28], including some recent ones [28], questioned the assumption of the Hara Theorem on general grounds. The point is that it is possible that the parity-violating amplitude D becomes singular when one takes the $SU(3)$ limit. This points out that how one treats the $SU(3)$ breaking can have a crucial influence on the outcome of the analysis. As elaborated in the magnetic moment section below, we believe that more careful analyses are still wanting.

One may try to employ ChPT, useful in describing low-energy hadronic processes involving only mesons, to tackle hyperons. For application to processes involving baryons, it is most consistently formulated in the heavy-baryon formulation [12], in which the $SU(3)$-invariant baryon mass, $\dot{m}$, is removed by a field transformation (see Ref. [6] for details). In this approach an amplitude for a given process is expanded in external pion four-momenta, q, baryon *residual* four-momenta, k, and the quark mass, m_s. If one neglects the up and down quark masses, one can assume that q, k, and m_s are of the same order and represent their value as E (and adopt the convention that k and m_s are of the same order in the chiral expansion). Perturbation theory is reliable only when E is smaller than the chiral-symmetry-breaking scale Λ_χ. In the heavy-baryon formulation there is an additional expansion in $1/\dot{m}$. However, all these terms can be effectively absorbed in counterterms of the theory [6, 7, 8].

Weak radiative decays of hyperons have been studied before in the context of ChPT by Jenkins, Luke, Manohar and Savage [26] and Neufeld [25]. Jenkins *et al.* and Neufeld calculated the amplitude up to the one-loop level. These loop diagrams give contributions to the amplitudes which are at least of $\mathcal{O}(\mathcal{E}^{\in})$ in the chiral expansion. However, tree-level direct-emission diagrams from the next-to-leading-order weak Lagrangian, which give contributions of $\mathcal{O}(\mathcal{E})$ to the amplitudes, were not considered [6, 7]. In a series of papers [6, 7, 8] we consistently calculated the leading-order amplitude of weak radiative decays of hyperons in ChPT. At this order, no loop contributions need to be considered. However, one does need to take into account the higher-order terms in the weak chiral Lagrangian. We showed that these terms give rise to a violation of Hara's theorem. As a consequence the decay rates for the charged decays $\Sigma^+ \to p + \gamma$ and $\Xi^- \to \Sigma^-\gamma$ can be accounted for consistently. We showed that, in leading order, ChPT predicts the ratios of the decay amplitudes of all the neutral channels as functions of the baryon masses only. We compared these predictions with the data. Furthermore, the asymmetry parameters still vanish in this leading-order calculation. However, this is not necessarily inconsistent with the data in the expansion scheme of ChPT.

In particular, the diagrams contributing to the leading-order amplitude are the tree diagrams given in Fig. 2. There are two kinds of diagrams: the direct-emission diagrams, Fig. 2a, and the baryon-pole diagrams, Fig. 2b and Fig. 2c. Loop diagrams give rise to contributions of higher order. The pole diagrams contribute only to the parity-conserving form factor C, in accordance with the Lee-Swift theorem [29]. Unfortunately as concluded in Ref. [8], within the context of the simplest heavy-baryon chiral-perturbation theory without inclusion of any additional resonances, the understanding of many of the experimental results in radiative hyperon decays, including the asymmetry, is still beyond reach. However, if one is willing to introduce additional resonances, together with the additional parameters that come with them, it is possible to fit the data with these additional parameters [9].

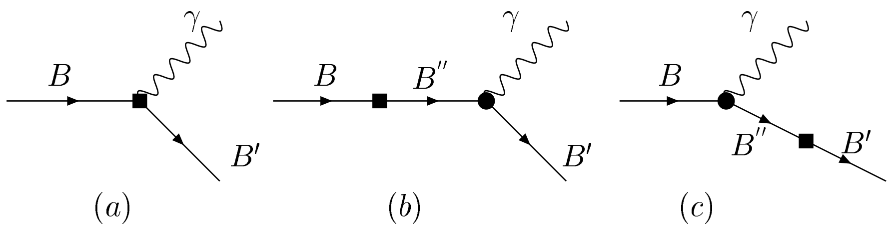

Figure 2. Diagrams for $B \to B'\gamma$. Here circles and squares represent strong and weak vertices respectively.

5 Magnetic Moments

The magnetic moments of the octet baryons were found to obey approximate $SU(3)$ symmetry a long time ago by Coleman and Glashow (CG) [30]. In the $SU(3)$-symmetric limit, the nine observable moments (including the transitional moment between Σ^0 and Λ) can be parameterized in terms of two parameters, and as a result obey approximate relationships. The two-parameter result of CG can in fact fit the observed magnetic moments up to about the 20% level. However, since at present the moments have been measured with an accuracy better than 1% [10], an improved theoretical understanding is clearly desirable.

Many attempts were made trying to improve the numerical predictions of CG by including $SU(3)$-breaking effects using ChPT [31, 32, 33, 34]. However, many of these efforts resulted in numerical fits *worse* than the leading-order $SU(3)$-invariant one by CG. For example, Caldi and Pagels [31] found that the leading $SU(3)$-breaking corrections, in their scheme for ChPT, appear in the nonanalytic forms of $\sqrt{m_s}$ and $m_s \ln m_s$. They showed that the $\sqrt{m_s}$ corrections are in fact at least as large as the $SU(3)$-invariant zeroth-order terms, which casts doubt on the applicability of ChPT. Caldi and Pagels suggested that this "failure" of ChPT might be attributed to the large mass of the kaon in the loops and the fact that the leading correction is of nonanalytic form. Such nonanalytic contributions were indeed pointed out earlier by Li and Pagels [37] and others [38]. However, the nonanalyticity appears only in the $SU(3)$-invariant chiral-symmetry-breaking mass, not in $SU(3)$-breaking parameters. More recently, similar large corrections to the baryon magnetic moments nonanalytic in m_s have been found by calculating them up to the one-loop level in ChPT [32, 26, 34]. By using only the $\sqrt{m_s}$ terms Jenkins *et al.* [26] could improve the accuracy of the Coleman-Glashow results from 20% to about 10%. However, this could only be achieved by using a *different* value of the meson decay constant in kaon loops than in pion loops, with the effect that the magnitude of the kaon loops is artificially reduced. In addition, Krause [32] showed that $m_s \ln m_s$ corrections are just as important, which disagrees with Refs. [26, 34]. Also, Krause further argued that the nonanalytic contributions are not at all a good approximation of the loop integrals.

5.1 Okubo relation

Shortly after CG, Okubo [39] derived a relation

$$6\mu_\Lambda + \mu_{\Sigma^-} - 4\sqrt{3}\mu_{\Lambda\Sigma^0} - 4\mu_n + \mu_{\Sigma^+} - 4\mu_{\Xi^0} = 0,$$

where $\mu_{\Lambda\Sigma^0}$ is the $\Sigma^0 \to \Lambda$ transition moment. This relation can be obtained if one assumes that $SU(3)$-breaking corrections to the moments are linear in the quark mass matrix σ, defined by $\sigma \equiv \mathrm{diag}(0, 0, m_s)$. The $SU(3)$-breaking terms introduce an additional 5 parameters. The resulting 7-parameter prediction can fit the current 9 high-precision observables to within 1.5% [35].

Within the context of ChPT, there is no unique way of treating the $SU(3)$ breaking yet. However, the nice fit of the Okubo relation can be taken as a hint. In Ref. [36], it was shown that, even within the scheme of ChPT used in Ref. [33], while the one-loop nonanalytic contributions of the form $\sqrt{m_s}$ satisfy the Okubo relation, the nonanalytic contributions of the form $m_s \ln m_s$ do not! In Ref. [35], a formulation of ChPT was suggested such that the leading $SU(3)$-breaking correction is indeed linear in m_s.

6 Summary

We have shown that there are still many issues that we do not understand associated with hyperon decays based only on fundamental principles. It is possible that some of the difficulties are such that, due to our inability to deal with the strong interaction, we cannot understand them without invoking some phenomenological models. It is also possible that some of the current data that make it so hard for us to understand may be changed by improved measurements in the future.

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