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Hyperon Physics—a Personal Overview

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A range of issues in the field of hyperon physics is presented, together with an assessment of where important challenges remain.

1. Introduction

The subject of hyperon physics is a vast one, as indicated by the fact that this workshop will run for three days, with presentations involving a range of different issues. Obviously it would be impossible for me to cover all of the interesting features in this introductory presentation. Instead, I will present a very personal picture of some of the issues in hyperon physics which *I* think need to be answered, and will trust the various speakers to fill in areas which I have omitted.

2. Hyperon Processes

I have divided my presentation into sections, which cover the various arenas which I think need attention:

2.1. Nonleptonic Hyperon Decay

The dominant decay mode of the $\frac{1}{2}^+$ hyperons is, of course, the pionic decay $B \to B'\pi$. On the theoretical side there remain two interesting and important issues which have been with us since the 1960's—the origin of the $\Delta I = 1/2$ rule and the S/P-wave problem[1]:

i) The former is the feature that $\Delta I = 3/2$ amplitudes are suppressed with respect to their $\Delta I = 1/2$ counterparts by factors of the order of twenty or so. This suppression exists in both hyperon as well as kaon nonleptonic decay and, despite a great deal of theoretical work, there is still no simple explanation for its existence. The lowest order weak nonleptonic $\Delta S = 1$ Hamiltonian possesses comparable $\Delta I = 1/2$ and $\Delta I = 3/2$ components and leading log gluonic effects can bring about a $\Delta I = 1/2$ enhancement of a factor of three to four or so[2]. The remaining factor of five seems to arise from the validity of what is called

the Pati-Woo theorem in the baryon sector[3] while for kaons it appears to be associated with detailed dynamical structure[4]. Interestingly the one piece of possible evidence for its violation comes from a hyperon reaction—-hypernuclear decay[5]. A hypernucleus is produced when a neutron in an atomic nucleus is replaced by a Λ. In this case the usual pionic decay mode is Pauli suppressed, and the hypernucleus primarily decays via the non-mesonic processes $\Lambda p \to np$ and $\Lambda n \to nn$. There does exist a rather preliminary indication here of a possibly significant $\Delta I = 1/2$ rule violation, but this has no Fermilab relevance and will have to be settled at other laboratories[6].

ii) The latter problem is not as well known but has been a longstanding difficulty to those of us theorists who try to calculate these things. Writing the general decay amplitude as

$$\text{Amp} = \bar{u}(p')(A + B\gamma_5)u(p) \qquad (1)$$

The standard approach to such decays goes back to current algebra days and expresses the S-wave (parity-violating) amplitude—A—as a contact term—the baryon-baryon matrix element of the axial-charge-weak Hamiltonian commutator. The corresponding P-wave (parity-conserving) amplitude—B—uses a simple pole model (*cf.* Figure 1) with the the weak baryon to baryon matrix element given by a fit to the S-wave sector. Parity violating BB' matrix elements are neglected in accord with the Lee-Swift theorem[7]. With this procedure one can obtain a good S-wave fit but finds P-wave amplitudes which are in very poor agreement with experiment. On the other hand,
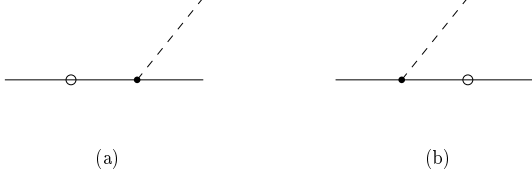
Figure 1. Pole diagrams used to calculated parity-conserving nonleptonic hyperon decay.

one can fit the P-waves, in which case the S-wave predictions are very bad[8]. Clearly the solution requires the input of additional physics, such as inclusion of $(70, 1^-)$ intermediate states as done in an SU(6) calculation by Le Youaunc et al.[9] or of intermediate $\frac{1}{2}^-$ and $\frac{1}{2}^+$ resonant states by Borasoy and myself in a chiral picture[10].

In either case, we do *not* require more and better data. The issues are already clear. What we need is more and better theory!

Where we *do* need data involves the possibility of testing the standard model prediction of CP violation, which predicts the presence of various asymmetries in the comparision of hyperon and antihyperon nonleptonic decays[11]. The basic idea is that one can write the decay amplitudes in the form

$$A = |A| \exp i(\delta_S + \phi_S), \quad B = |B| \exp i(\delta_P + i\phi_P) \tag{2}$$

where δ_S, δ_P are the strong S,P-wave phase shifts at the decay energy of the mode being considered and ϕ_S, ϕ_P are CP-violating phases which are expected to be of order 10^{-4} or so in standard model CP-violation. One can detect such phases by comparing hyperon and antihyperon decay parameters. Unfortunately nature is somewhat perverse here in that the larger the size of the expected effect, the more difficult the experiment. For example, the asymmetry in the overall decay rate, which is the easiest to measure, has the form[1]

$$
\begin{aligned}
C &= \frac{\Gamma - \bar{\Gamma}}{\Gamma + \bar{\Gamma}} \\
&\sim \left(-2(A_1 A_3 \sin(\delta_S^1 - \delta_S^3) \sin(\phi_S^1 - \phi_S^3)\right. \\
&+ \left. B_1^r B_3^r \sin(\delta_P^1 - \delta_P^3) \sin(\phi_P^1 - \phi_P^3)\right) \\
&\quad / \quad |A_1|^2 + |B_1^r|^2 \tag{3}
\end{aligned}
$$

[1]Here B^r indicates a reduced amplitude—$B^r = B(E' - M_{B'})/(E' + M_{B'})$.

where the subscripts, superscripts 1,3 indicate the $\Delta I = \frac{1}{2}, \frac{3}{2}$ component of the amplitude. We see then that there is indeed sensitivity to the CP-violating phases but that it is multiplicatively suppressed by *both* the the strong interaction phases ($\delta \sim 0.1$) as well as by the $\Delta I = \frac{3}{2}$ suppression $A_3/A_1 \sim B_3/B_1 \sim 1/20$. Thus we find $C \sim \phi/100 \sim 10^{-6}$ which is much too small to expect to measure in present generation experiments.

More sanguine, but still not optimal, is a comparison of the asymmetry parameters α, defined via

$$W(\theta) \sim 1 + \alpha \vec{P}_B \cdot \hat{p}_{B'} \tag{4}$$

In this case, one finds

$$
\begin{aligned}
A &= \frac{\alpha + \bar{\alpha}}{\alpha - \bar{\alpha}} = -\sin(\phi_S^1 - \phi_P^1) \sin(\delta_S^1 - \delta_P^1) \\
&\sim 0.1\phi \sim 4 \times 10^{-4} \tag{5}
\end{aligned}
$$

which is still extremely challenging.

Finally, the largest signal can be found in the combination

$$B = \frac{\beta + \bar{\beta}}{\beta - \bar{\beta}} = \cot(\delta_S^1 - \delta_P^1) \sin(\phi_S^1 - \phi_P^1) \sim \phi \tag{6}$$

Here, however, the parameter β is defined via the general expression for the final state baryon polarization

$$
\begin{aligned}
< \vec{P}_{B'} > &= \frac{1}{W(\theta)} \left((\alpha + \vec{P}_B \cdot \hat{p}_{B'}) \hat{p}_{B'} \right. \\
&+ \left. \beta \vec{P}_B \times \hat{p}_{B'} + \gamma(\hat{p}_{B'} \times (\vec{P}_B \times \hat{p}_{B'})) \right) \tag{7}
\end{aligned}
$$

and, although the size of the effect is largest— $B \sim 10^{-3}$—this measurement seems out of the question experimentally.

Despite the small size of these effects, the connection with standard model CP violation and the possibility of finding larger effects due to new physics demands a no-holds-barred effort to measure these parameters.

2.2. Nonleptonic Radiative Decay
Another longstanding thorn in the side of theorists attempting to understand weak decays of hyperons is the nonleptonic radiative mode $B \to B'\gamma$[12]. In this case one can write the most general decay amplitude as

$$
\begin{aligned}
\text{Amp} &= \frac{e}{M_B + M_{B'}} \epsilon^{*\mu} q^\nu \\
&\times \bar{u}(p') \left(-i\sigma_{\mu\nu} C - i\sigma_{\mu\nu}\gamma_5 D\right) u(p) \tag{8}
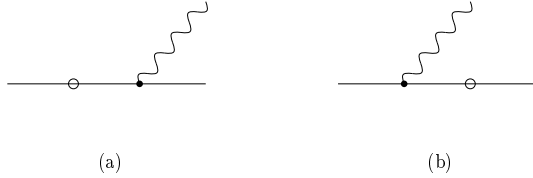\end{aligned}
$$

(a)　　　　　　　　　　(b)

Figure 2. Pole diagrams used to calculate radiative hyperon decay.

where C is the magnetic dipole (parity conserving) amplitude and D is its (parity conserving) electric dipole counterpart. There are two quantities of interest in the analysis of such decays—the decay rate and photon asymmetry, which go as

$$\Gamma \sim |C|^2 + |D|^2, \quad A_\gamma = \frac{2\mathrm{Re}C^*D}{|C|^2 + |D|^2} \qquad (9)$$

The difficulty here is associated with "Hara's Theorem" which requires that in the SU(3) limit the parity violating decay amplitude must vanish for decay between states of a common U-spin multiplet—*i.e.* $\Sigma^+ \to p\gamma$ and $\Xi \to \Sigma^-\gamma$[13]. (The proof here is very much analogous to the one which requires the vanishing of the axial tensor form factor in nuclear beta decay between members of a common isotopic spin multiplet[14].) Since one does not expect significant SU(3) breaking effects, we anticipate a relatively small photon asymmety parameter for such decays. However, in the case of $\Sigma^+ \to p\gamma$ the asymmetry is known to be large and negative[15]

$$A_\gamma(\Sigma^+ \to p\gamma) = -0.76 \pm 0.08 \qquad (10)$$

and for thirty years theorists have been struggling to explain this result. In leading order the amplitude is given by the simple pole diagrams, with the weak baryon-baryon matrix elements being those determined in the nonradiative decay analysis. The Lee-Swift theorem asserts that such matrix elements must be purely parity conserving in the SU(3) limit and this is the origin of Hara's theorem in such a model[7]. Although SU(3) breaking corrections have been calculated, none is large enough to explain the experimental result—Eq. 10[16]. As in the case of the S/P-wave puzzle, what is clearly required is the inclusion of additional physics and here too the inclusion of $(70, 1^-)$ states by Le Youaunc et al.[

9] or of $\frac{1}{2}^-$ and $\frac{1}{2}^+$ resonant states in a chiral framework by Borasoy and myself[17] appears to naturally predict a large negative asymmetry. However, in order to confirm the validity of these or any model what will be required is a set of measurements of *both* rates and asymmetries for such decays. In this regard, it should be noted that theoretically one expects all asymmetries to be negative in any realistic model[18]. It would be very difficult to accomodate a large positive asymmetry. Thus the present particle data group listing[15]

$$A_\gamma(\Xi^0 \to \Lambda\gamma) = +0.43 \pm 0.44 \qquad (11)$$

deserves to be carefully remeasured.

2.3. Hyperon Beta Decay

A mode that theory does well in predicting (in fact some would say *too* well) is that of hyperon beta decay—$B \to B'\ell\nu_\ell$, where ℓ is either an electron or a muon. Since this is a semileptonic weak interaction, the decays are described in general by matrix elements of the weak vector, axial-vector currents

$$
\begin{aligned}
< B'|V_\mu|B > &= \bar{u}(p')(f_1\gamma_\mu + \frac{-if_2}{M_B + M_{B'}}\sigma_{\mu\nu}q^\nu \\
&\quad + \frac{f_3}{M_B + M_{B'}}q_\mu)u(p) \\
< B'|A_\mu|B > &= \bar{u}(p')(g_1\gamma_\mu + \frac{-ig_2}{M_B + M_{B'}}\sigma_{\mu\nu}q^\nu \\
&\quad + \frac{g_3}{M_B + M_{B'}}q_\mu)\gamma_5 u(p)
\end{aligned}
$$

$$(12)$$

Here the dominant terms are the vector, axial couplings f_1, g_1 and the standard approach is simple Cabibbo theory, wherein one fits the g_1 in terms of SU(3) F,D coefficients and f_1 using CVC and simple F coupling. When this is done, one finds in general a very satisfactory fit— $\chi^2/d.o.f \sim 2.0$—which can be made even better by inclusion of simple quark model SU(3) breaking effects—$\chi^2/d.o.f. \sim 0.85$[19]. An output of such a fit is the value of the KM mixing parameter $V_{us} = 0.220 \pm 0.003$, which is in good agreement with the value $V_{us} = 0.2196 \pm 0.0023$ measured in K_{e3} decay. However, differing assumptions about SU(3) breaking will lead to slightly modified values.

The importance of such a measurement of V_{us} has to do with its use as an input to a test of the standard model via the unitarity prediction

$$|V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2 = 1 \qquad (13)$$

From an analysis of B-decay one obtains $|V_{ub}| \sim 0.003$, which when squared leads to a negligible contribution to the unitarity sum. So the dominant effect comes from V_{ud}, which is measured via $0^+ - 0^+$ superallowed nuclear beta decay—

$$V_{ud}^2 = \frac{2\pi^3 \ln 2 \, m_e^{-5}}{2 G_F^2 (1 + \Delta_R^V) \bar{\mathcal{F}} t} \qquad (14)$$

Here $\Delta_R^V = 2.40 \pm 0.08\%$ is the radiative correction and $\bar{\mathcal{F}} t = 3072.3 \pm 0.9$ sec. is the mean (modified) ft-value for such decays. Of course, there exist important issues in the analysis of such ft-values including the importance of isotopic spin breaking effects and of possible Z-dependence omitted from the radiative corrections, but if one takes the above-quoted number as being correct we obtain[20]

$$V_{ud} = 0.9740 \pm 0.0005$$
$$\text{and}$$
$$|V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2 = 0.9968 \pm 0.0014 \qquad (15)$$

which indicates a possible violation of unitarity. If correct, this would suggest the existence of non-standard-model physics, but clearly additional work, both theoretical and experimental, is needed before drawing this conclusion.

What is needed in the case of hyperon beta decay is good set of data including rates *and* asymmetries, both in order to produce a possibly improved value of V_{us} but also to study the interesting issue of SU(3) breaking effects, which *must* be present, but whose effects seem somehow to be hidden. A related focus of such studies should be the examination of higher order— recoil—form factors such as weak magnetism (f_2) and the axial tensor (g_2). In the latter case, Weinberg showed that in the standard quark model $G = C \exp(i\pi I_2)$-invariance requires $g_2 = 0$ in neutron beta decay $n \to p e^- \bar{\nu}_e$[21]. (This result usually is called the stricture arising from "no second class currents.") In the SU(3) limit one can use V-spin invariance to show that $g_2 = 0$ also obtains for $\Delta S = 1$ hyperon beta decay, but in the real world this condition will be violated. A simple quark model calculation suggests that $g_2/g_1 \sim -0.2$[22] but other calculations, such as a recent QCD sum rule estimate give a larger number—$g_2/g_1 \sim -0.5$. In any case good hyperon beta decay data—with rates and asymmetries—will be needed in order to extract the size of such effects.

2.4. Hyperon Polarizabilities

Since this subject is not familiar to many physicists, let me spend just few moments giving a bit of motivation. The idea goes back to simple classical physics. Parity and time reversal invariance, of course, forbid the existence of a permanent electric dipole moment for an elementary system. However, consider the application of a uniform electric field to a such a system. Then the positive charges will move in one direction and negative charges in the other—*i.e.* there will be a charge separation and an electric dipole moment will be induced. The size of the edm will be proportional to the applied field and the constant of proportionality between the applied field and the induced dipole moment is the electric polarizability α_E

$$\vec{p} = 4\pi \alpha_E \vec{E} \qquad (16)$$

The interaction of this dipole moment with the field leads to an interaction energy

$$U = -\frac{1}{2}\vec{p} \cdot \vec{E} = -\frac{1}{2} 4\pi \alpha_E \vec{E}^2, \qquad (17)$$

where the "extra" factor of $\frac{1}{2}$ compared to elementary physics result is due to the feature that the dipole moment is *induced*. Similarly in the presence of an applied magnetizing field $\vec{H}$ there will be generated an induced magnetic dipole moment

$$\vec{\mu} = 4\pi \beta_M \vec{H} \qquad (18)$$

with interaction energy

$$U = -\frac{1}{2}\vec{\mu} \cdot \vec{H} = -\frac{1}{2} 4\pi \beta_M \vec{H}^2. \qquad (19)$$

For wavelengths large compared to the size of the system, the effective Hamiltonian describing the interaction of a system of charge e and mass m with an electromagnetic field is, of course, given by

$$H^{(0)} = \frac{(\vec{p} - e\vec{A})^2}{2m} + e\phi, \qquad (20)$$

and the Compton scattering cross section has the simple Thomson form

$$\frac{d\sigma}{d\Omega} = \left(\frac{\alpha_{em}}{m}\right)^2 \left(\frac{\omega'}{\omega}\right)^2 [\frac{1}{2}(1 + \cos^2\theta)], \qquad (21)$$

where α_{em} is the fine structure constant and ω, ω' are the initial, final photon energies respectively. As the energy increases, however, so does the resolution and one must also take into account polarizability effects, whereby the effective Hamiltonian becomes

$$H_{\text{eff}} = H^{(0)} - \frac{1}{2} 4\pi(\alpha_E \vec{E}^2 + \beta_M \vec{H}^2). \qquad (22)$$

The Compton scattering cross section from such a system (taken, for simplicity, to be spinless) is then

$$
\begin{aligned}
\frac{d\sigma}{d\Omega} &= \left(\frac{\alpha_{em}}{m}\right)^2 \left(\frac{\omega'}{\omega}\right)^2 \left(\frac{1}{2}(1+\cos^2\theta)\right. \\
&- \frac{m\omega\omega'}{\alpha_{em}}[\frac{1}{2}(\alpha_E+\beta_M)(1+\cos\theta)^2 \\
&+ \left.\frac{1}{2}(\alpha_E-\beta_M)(1-\cos\theta)^2]\right) \qquad (23)
\end{aligned}
$$

It is clear from Eq. 23 that from careful measurement of the differential scattering cross section, extraction of these structure dependent polarizability terms is possible provided

  i) that the energy is large enough that these terms are significant compared to the leading Thomson piece and

  ii) that the energy is not so large that higher order corrections become important.

In this fashion the measurement of electric and magnetic polarizabilities for the proton has recently been accomplished at SAL and at MAMI using photons in the energy range 50 MeV $< \omega <$ 100 MeV, yielding[23] [2]

$$
\begin{aligned}
\alpha_E^p &= (12.1 \pm 0.8 \pm 0.5) \times 10^{-4} \text{ fm}^3 \\
\beta_M^p &= (2.1 \mp 0.8 \mp 0.5) \times 10^{-4} \text{ fm}^3. \quad (24)
\end{aligned}
$$

Note that in practice one generally exploits the strictures of causality and unitarity as manifested in the validity of the forward scattering dispersion relation, which yields the Baldin sum rule[27]

$$
\begin{aligned}
\alpha_E^{p,n} + \beta_M^{p,n} &= \frac{1}{2\pi^2}\int_0^\infty \frac{d\omega}{\omega^2}\sigma_{\text{tot}}^{p,n} \\
&= \begin{cases} (13.69 \pm 0.14)\times 10^{-4}\text{fm}^3 & \text{proton} \\ (14.40 \pm 0.66)\times 10^{-4}\text{fm}^3 & \text{neutron} \end{cases}
\end{aligned}
$$
$$(25)$$

as a rather precise constraint because of the small uncertainty associated with the photoabsorption cross section σ_{tot}^p.

As to the meaning of such results we can compare with the corresponding calculation of the electric polarizability of the hydrogen atom, which yields[28]

$$
\alpha_E^H = \frac{9}{2}a_0^2 \quad \text{vs.} \quad \alpha_E^p \sim 10^{-3} < r_p^2 >^{\frac{3}{2}} \qquad (26)
$$

---

[2]Results for the neutron extracted from $n - Pb$ scattering cross section measurements have been reported[24], but have been questioned[25]. Extraction via studies using a deuterium target may be possible in the future[26].

where a_0 is the Bohr radius. Thus the polarizability of the hydrogen atom is of order the atomic volume while that of the proton is only a thousandth of its volume, indicating that the proton is much more strongly bound.

The relevance to our workshop is that the polarizability of a hyperon can also be measured using Compton scattering, via the reaction $B+Z \to B + Z + \gamma$ extrapolated to the photon pole—i.e. the Primakoff effect. Of course, this is only feasible for charged hyperons—$\Sigma^\pm, \Xi^-$, and the size of such polarizabilities predicted theoretically are somewhat smaller than that of the proton[29]

$$
\alpha_E^{\Sigma^+} \sim 9.4 \times 10^{-4}\,\text{fm}^3, \qquad \alpha_E^{\Xi^-} \sim 2.1 \times 10^{-4}\,\text{fm}^3
$$
$$(27)$$

but their measurement would be of great interest.

2.5. Polarization and Hyperon Production

My final topic will be that of polarization in strong interaction production of hyperons, a field that began here at FNAL in 1976 with the discovery of Λ polarization in the reaction[30]

$$
p(300 \text{ GeV}) + Be \to \vec\Lambda + X \qquad (28)
$$

This process has been well studied in the intervening years[31] and we now know that in the fragmentation region the polarization is large and negative—$\vec{P}\cdot\hat{p}_{inc}\times\hat{p}_\Lambda < 0$—and that it satisfies scaling, i.e. is a function only of $x_F = \frac{p_\Lambda^\|}{p_p}, p_\Lambda^\perp$ and not of the center of mass energy . Various theoretical approaches have been applied in order to try to understand this phenomenon—e.g., Soffer and Törnqvist have developed a Reggized pion exchange picture[32], while DeGrand, Markkanen, and Miettinen have used a quark-parton approach wherein the origin of the polarization is related to the Thomas precession[33]—but none can be said to be definitive. One thing which seems to be clear is that there exists a close connection with the large negative polarizations seen in inclusive hyperon production and the large positive analyzing powers observed at FNAL in inclusive meson production with polarized protons[34]

$$
\vec{p} + p \to \pi^+ + X \qquad (29)
$$

Another input to the puzzle may be the availability in the lower energy region of new exclusive data from Saturne involving[35]

$$
\vec{p} + p \to p + \vec\Lambda + K^+ \qquad (30)
$$

which seems best described in terms of a kaon exchange mechanism. Clearly there is much more to do in this field.

3. Summary

I conclude by noting that, although the first hyperon was discovered more then half a century ago and much work has been done since, the study of hyperons remains an interesting and challenging field. As I have tried to indicate above, many questions exist as to their strong, weak, and electromagnetic interaction properties, and I suspect that these particles will remain choice targets for particle hunters well into the next century.

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