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I. The Nuclear Interactions of the Λ -particle

1. The s-wave Λ -N Interaction

The binding of Λ particles to the s-shell nuclei H^2 , H^3 , He^3 and He^4 is due dominantly to the s-wave Λ -N interactions V_S for the singlet ($S = 0$) and triplet ($S = 1$) configurations. From the spin determinations for ΛH^3 ($J = \frac{1}{2}$) and for ΛH^4 ($J = 0$), we know that the singlet interaction V_0 is more attractive than the triplet interaction V_1 . Our knowledge of their strengths depends on quantitative calculations on the structure of these light hypernuclei, as discussed below. However, the conclusion that the triplet interaction V_1 is attractive depends on rather qualitative considerations. The mean Λ -N attraction in ΛH^3 is given by $\frac{1}{4}(3V_0 + V_1)$, dominated by the singlet interaction; hence calculations on ΛH^3 alone provide a lower limit on the strength of V_0 . In ΛHe^5 , on the other hand, the mean Λ -N attraction $\frac{1}{4}(3V_1 + V_0)$ arises dominantly from the triplet interaction; owing to its small

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weighting, the singlet interaction cannot provide sufficient attraction to account for the ΛHe^5 binding and we must conclude that V_1 is also attractive.

In order to estimate the strength of the Λ -N interaction, we must make a definite assumption about its range. We shall assume that this interaction has the same intrinsic range (1.5 F) as a Yukawa potential of range parameter $(2m_\pi)^{-1}$. This corresponds to the assumption that the Λ -N potential arises dominantly from the exchange of pions; charge symmetry for the Λ -particle (which has isospin $I = 0$) forbids one pion exchange, but the two-pion exchange potential is allowed and is expected to be strong, in view of our current belief (based on the existence of a $p_{\frac{3}{2}} \pi\text{-}\Lambda$ resonance analogous to the π -N resonance, for example) that the $\pi\Sigma\Lambda$ coupling has strength comparable with that of the pion-nucleon coupling. This value for the intrinsic range is quite close to the values obtained for the Yukawa potentials which are equivalent (in the sense of having the same zero-energy scattering length a_s and effective range r_s) to the meson-theoretic Λ -N potentials calculated recently by de Swart and Iddings.¹ In fact, in view of the low binding energies of the light hypernuclei, it is reasonable to expect that these will not be greatly sensitive to the details of the potentials V_0 and V_1 , but that they will be determined primarily by the low-energy scattering parameters, (a_0, r_0) and (a_1, r_1) , of these potentials. As we shall see, the detailed calculations which have been made to date (but always with central potentials) do bear out this expectation.

$\Lambda^3\text{H}$ is the lightest hypernucleus, with Λ -d separation energy $B_\Lambda = 0.21 \pm 0.2$ MeV, and is known to be a singlet state, with isospin $I = 0$ and spin $J = \frac{1}{2}$. This $np\Lambda$ system is sufficiently simple to allow quite detailed and accurate variational calculations, even with hard-core potentials. For its space wavefunction, Downs, Smith and Truong² have recently used a 10-parameter trial function,

$$\psi = f(r_{\Lambda n}) f(r_{\Lambda p}) g(r_{np}), \quad (1.1)$$

where f, g both have the general form,

$$\left\{ (\exp(-\alpha(r-d)) - \exp(-\beta(r-d))) + x (\exp(-\mu(r-d)) - \exp(-\nu(r-d))) \right\}. \quad (1.2)$$

This calculation used a Λ -N potential of the exponential form

$$V_{\Lambda N}(r) = \bar{U} \exp\{-3.54(r-d)/(b-2d)\}, \quad r \geq d \quad (1.3)$$

outside the hard core radius d . The intrinsic range is given by $b = 1.5$ F. For $d = 0.4$ F, their result was $\bar{U} = 1168$ MeV, which corresponds to a mean well-depth parameter $\bar{s} = 0.825$. The scattering length which corresponds to this mean Λ -N potential is $\bar{a} = -2.0$ F. Calculations were also made for $d = 0.2$ F and 0.6 F, leading to scattering lengths $\bar{a} = -1.75$ F and 2.11 F, respectively. It is also of interest to compare these results with the value $\bar{a} = -1.55$ F which was obtained in an earlier calculation³ of comparable complexity, for Yukawa potentials of the same intrinsic range, without hard cores. These results illustrate the insensitivity of the low-energy parameters obtained (for fixed intrinsic

range) for the mean potential to the shape assumed for this potential. We may also compare this result of Downs et al. with that from ~~the~~^{their} earlier calculation² using a 4-parameter trial function of the form (1.1), with f, g of the form (1.2) with $x = 0$, which gave $\bar{U} = 1202$ MeV. The more elaborate trial function led only to 2.8% improvement in the value obtained for $\bar{U}$, so that we may assume their final value to be rather close (better than 1%) to the true value.

Muller⁴ has also made calculations on ΛH^3 for exponential hard-core potentials, using an 8-parameter trial function of the form (1.1). He assumed a potential of greater intrinsic range, of the form $\bar{U} \exp(-2.38(r-d))$ outside hard-core radius d , and obtained $\bar{U} = 228.2$ MeV. The corresponding calculation of Downs et al.,² for potential shape $\bar{U} \exp(-2.36(r-d))$, gave $\bar{U} = 226.1$ MeV, so that their results are in good agreement. The scattering length for this potential is $\bar{a} = -2.7$ F; we note that the scattering length obtained does increase^{with} the potential range assumed, as is reasonable.

For ΛHe^5 , it is a reasonable assumption to neglect the distortion of the He^4 core by the presence of the Λ -particle, in view of the tight binding of He^4 . Since the form of He^4 is known from electron-scattering experiments, the strength of the mean Λ -N potential in ΛHe^5 may be deduced directly, for an assumed Λ -N potential shape without hard core. This calculation was done quite early⁵ for Λ -N potentials of Gaussian form, and was combined with the results of calculations for ΛH^3 to give the

results shown in the first column of Table I. For later reference, we quote here the volume integrals of the singlet and triplet potentials obtained in this way: $\bar{V}_0 = 386 \text{ MeV F}^3$, $\bar{V}_1 = 170 \text{ MeV F}^3$. The errors quoted in the Table are those derived from the uncertainties in the B_Λ values and in the radius of He^4 .

With hard-core potential shapes, calculations have been made for ${}_\Lambda\text{He}^5$ and ${}_\Lambda\text{H}^4$ (with charge symmetry, ${}_\Lambda\text{He}^4$ is the mirror hypernucleus to ${}_\Lambda\text{H}^4$, with the same structure and the same B_Λ value, as is consistent with the data) by Dietrich et al.⁶ and by Gutsch,⁷ using a method proposed by Mang and Wild⁸ for light nuclei. This method takes into account accurately the two-body correlations in these light nuclei, rather in the spirit of the Brueckner method as developed for the discussion of nuclear matter, although the detailed justification of the method for these light nuclei depends on quite different physical factors (essentially on the large spacing of the one-particle states for the oscillator potential determined self-consistently) than does its justification (the Pauli principle and the high Fermi momentum) for the problem of nuclear matter. To date, these calculations have used square-well potentials of width ($b - 2d$) outside hard-core radius d ; the results available for $d = 0.2F$ and $d = 0.4 F$, with intrinsic range $b = 1.5 F$, are given in Table I. It is of interest to note the insensitivity of the conclusion to the hard-core radius assumed; for orientation, we remark that, for $d = 0$, the increase from $a_0 = 2.4 F$ to $a_0 = 4.2 F$ requires an increase by only 10% in the potential strength of V_0 . For $d = 0.4 F$, the scattering length a_0 corresponds to a well-depth

parameter $s = 0.8$ for the 1S_0 potential, so that there is no reason to expect a bound state for the Λ -N system.

Dietrich et al.⁶ and Gutsch⁷ have also made calculations for ΛH^3 , using the same method. For $d = 0.4$ F, Gutsch finds that the observed B_Λ requires a mean well-depth $\bar{V}_3 = 170$ MeV for the outer square-well potential; the value $\frac{1}{4}(3V_0 + V_1)$ given by the potentials V_0 and V_1 obtained from the discussion of ΛH^4 and ΛHe^5 is 168.0 MeV. The scattering length corresponding to this value $\bar{V}_3$ is $\bar{a} = 5.9$ F. This is appreciably larger than the value (2.0 F) found from the variational calculations mentioned above, and corresponds to a well-depth parameter about 12% too large; this is probably due to the fact that the Mang-Wild method is least accurate for small binding energies, as it does not give a correct account of the asymptotic region (which is of particular importance for small B_Λ). However, the internal comparisons between all of these calculations for the s-shell hypernuclei are rather satisfactory, in view of their approximate nature and the range of potential forms which have been used.

Finally, we consider whether there exists a $J = 1$ particle-stable state for the hypernuclei ΛH^{4*} and ΛHe^{4*} . This question is of great interest in view of its bearing on the interpretation of the ΛH^4 and ΛHe^4 hypernuclei observed following $K^- - He^4$ capture events, whether these are produced directly (which would imply that the $K\Lambda$ parity is odd) or whether they result from γ -decay of such directly-produced excited states, ^{with} direct production of the ground state hypernuclei being forbidden. The

situation is summarized in Table I. The early calculations^{3,5} for $d = 0$ did allow the possibility of a lightly-bound ($B_{\Lambda}^* \approx 1.3 \text{ MeV}$) excited state. The more recent calculations^{6,7} for hard-core potentials do not allow the existence of a $J = 1$ bound state. The attraction predicted for the $J = 1$ configuration actually decreases with increasing hard-core radius, the well-depth parameter of the Λ -He³ attraction being only 0.83 in this state for $d = 0.4 \text{ F}$. Further calculations along these lines for outer potential shapes more realistic than the square well, and which can be compared directly with the more accurate ^{variational} calculations on the Λ H³ system, would be rather desirable at this point.

The interpretation of these potentials V_0 and V_1 , together with the data available on the ratio $R = (\Sigma^0 n)/(\Lambda n)$ for $\Sigma^- p$ capture from rest, has been extensively discussed in terms of the pion-hyperon interactions $\Sigma\Sigma\pi$ and $\Sigma\Lambda\pi$ by de Swart and Iddings.¹ Their calculations are based on a two-channel approach, in which the ΛN and ΣN channels are treated on the same footing in view of the relatively small mass difference ($m_{\Sigma} - m_{\Lambda}$) $\approx 80 \text{ MeV}$. We confine our attention here to their essential conclusions for the case of even $\Sigma\Lambda$ parity. The 1S_0 potential depends primarily on the coupling $f_{\Sigma\Lambda}$ and its strength requires that $f_{\Sigma\Lambda}$ be comparable in magnitude with the pion-nucleon coupling f_{NN} . The 3S_1 potential strength and the ratio R both depend quite strongly on $f_{\Sigma\Sigma}$ and an adequate fit to all three data requires $f_{\Sigma\Sigma} \approx 0.1$, taking the hard-core radius to be $d = 0.5 \text{ F}$, as required for the corresponding NN potentials.

For small $f_{\Sigma\Sigma}$, the coupling between the ΛN and ΣN channels is not of major importance in the 1S_0 state, and the dominant interaction is the two-pion exchange potential in the ΛN channel, which depends only on the $\Sigma\Lambda\pi$ coupling. In the 3S_1 state, the one-pion exchange potential linking ΛN and ΣN channels is very strong and of the tensor form; there is a strong $f_{\Sigma\Sigma}$ dependence both because of the importance of the $\Sigma\Sigma\pi$ coupling in the ΣN channel and because $f_{\Sigma\Sigma}$ contributes to the fourth-order potential $V(\Lambda N, \Sigma N)$ which interferes with the strong one-pion exchange term in this potential. However, the low-energy ΛN scattering wavefunction does not have a large D-state, so that the tensor component of the effective ΛN potential is expected to be of relatively minor importance, in contrast with the situation for the NN system.

2. The s-wave Λ - Λ Interaction

Very recently Danysz et al.⁹ have reported the first example of a double Λ -hypernucleus, for which the most probable interpretation is $\Lambda\Lambda\text{Be}^{10}$, with separation energy $B_{\Lambda\Lambda} = 17.5 \pm 0.5$ MeV relative to $\Lambda + \Lambda + \text{Be}^8(\text{g.s.})$. Since $B_{\Lambda} = 6.5 \pm 0.15$ MeV ΛBe^9 , $B_{\Lambda\Lambda}$ exceeds $2 B_{\Lambda}(\text{Be}^9)$ by 4.5 ± 0.6 MeV, so that we conclude that the Λ - Λ interaction is attractive. Since both particles are in the ground s-orbit relative to the nuclear core, the Pauli principle requires that their spins be coupled in the singlet configuration, so that $B_{\Lambda\Lambda}$ is determined predominantly by the 1S_0 Λ - Λ potential.

An estimate of the Λ - Λ potential strength may be obtained using a simple product wavefunction,¹⁰

$$\psi = \varphi(|\underline{r}_1 - \underline{R}|) \varphi(|\underline{r}_2 - \underline{R}|) (u_1 v_2 - v_1 u_2) / \sqrt{2}, \quad (2.1)$$

where $\underline{R}$ denotes the c.m. position of the Be^8 core. If $E_\Lambda(\varphi)$ denotes the separation energy of ΛBe^9 (relative to $\Lambda + \text{Be}^8(\text{g.s.})$) calculated for the wavefunction φ , then

$$-B_{\Lambda\Lambda} \geq -2E_\Lambda(\varphi) + \langle V_{\Lambda\Lambda} \rangle. \quad (2.2)$$

We take the potential $V_{\Lambda\Lambda}$ to be of Gaussian form $-\bar{v}_{\Lambda\Lambda} (\lambda/\pi)^{3/2} \exp(-\lambda r^2)$, with $\lambda = 2.0604/b^2$ and intrinsic range $b = 1.5$ F corresponding to the process of two-pion exchange. If φ is taken to be the wavefunction for ΛBe^9 , then $-\langle V_{\Lambda\Lambda} \rangle$ can be identified with $(B_{\Lambda\Lambda} - 2B_\Lambda) = 4.5 \pm 0.6$ MeV, which leads to ^{a first} estimate 520 ± 70 MeV F³ for the volume integral $\bar{v}_{\Lambda\Lambda}$ of the 1S_0 Λ - Λ potential. However this estimate is considerably too large, since, owing to the additional binding energy, the wavefunction φ appropriate to ΛBe^{10} will actually fall off much more rapidly with r than φ for ΛBe^9 . To allow for this, we vary the form of φ , using the inequality (2.2) as a variational principle for $\bar{v}_{\Lambda\Lambda}$. The form

$$\varphi(r) = N \{ \exp(-ar^2) + y \exp(-cr^2) \} \quad (2.3)$$

provides a good fit to the ΛBe^9 wavefunction for $a = 0.25$ F⁻², $c = 0.07$ F⁻² and $y = 0.33$, the Λ - Be^8 potential being represented by a shell-model form with strength adjusted to fit the observed B_Λ value. The effect of interest is roughly represented by decreasing y ; the minimum occurs for $y = 0.15$, leading to the improved estimate,

$$\bar{v}_{\Lambda\Lambda} = 440 \pm 60 \text{ MeV F}^3. \quad (2.4)$$

Further improvements to this estimate are still necessary, with improved wavefunctions which take into account (a) the possibility of strong spatial correlations between the two Λ particles, and (b) the distortion of the nuclear core by the presence of two strongly bound Λ particles. The inclusion of these effects may be expected to lower our estimate for $\bar{v}_{\Lambda\Lambda}$ further, but certainly not by more than the error quoted in (2.4).

As pointed out by Danysz et al.⁹ there is possible an alternative interpretation for this event as $\Lambda\Lambda\text{Be}^{11}$ with separation energy $19.0 \pm 0.6 \text{ MeV}$ relative to $\Lambda + \Lambda + \text{Be}^9$. For ΛBe^{10} , $B_\Lambda = 7.9 \pm 0.4 \text{ MeV}$, but only the spin-average $\Lambda\text{-Be}^9$ interaction is effective in $\Lambda\Lambda\text{Be}^{11}$, since the Λ spins are in the singlet configuration. Using intermediate coupling wavefunctions for Be^9 , the spin-dependent term in B_Λ is $0.32 \langle \Delta \rangle$, where $\langle \Delta \rangle$ is an appropriate expectation value of the spin-dependent part of the $\Lambda\text{-N}$ interaction which has the value $1.8 \pm 0.4 \text{ MeV}$ for neighboring hypernuclei, as discussed in Sec. 3. Hence the quantity $(B_{\Lambda\Lambda} - 2\bar{B}_\Lambda)$ has the value $(19.0 - 2(7.9 - 0.6)) = 4.4 \pm 1.0 \text{ MeV}$, quite comparable with the value used above for $\Lambda\Lambda\text{Be}^{10}$, so that our estimate (2.4) of $\bar{v}_{\Lambda\Lambda}$ would not be appreciably modified if the alternative interpretation $\Lambda\Lambda\text{Be}^{11}$ were adopted.

As for the $\Lambda\text{-N}$ potential, the long-range part of the $\Lambda\text{-}\Lambda$ potential arises from two-pion exchange. For even $\Sigma\Lambda$ parity, the 1S_0 $\Lambda\text{-}\Lambda$ interaction arises dominantly from the

two-pion exchange potential in the $\Lambda\Lambda$ channel, since the potential coupling the $\Lambda\Lambda$ and $\Sigma\Sigma$ channels is relatively weak. The reason for this is essentially the same as for the Λ -N case, namely (a) for pseudoscalar mesons, the one-pion exchange potential is known to be quite weak, and (b) the two-pion contribution to $V(\Lambda\Lambda - \Sigma\Sigma)$ is necessarily proportional to $f_{\Sigma\Sigma}^2$, which appears to be a relatively small coupling parameter. The structure of the two-pion contribution to $V(\Lambda\Lambda, \Lambda\Lambda)$ is the same as that for $V(\Lambda N, \Lambda N)$, except for the additional factor $f_{\Sigma\Lambda}^2 / f_{NN}^2$. Hence, if the same hard-core radius is assumed for the Λ - Λ and Λ -N interactions, the comparison of the values $\bar{v}_{\Lambda\Lambda} = 440 \text{ MeV.F}^3$ indicates that $f_{\Sigma\Lambda} \approx f_{NN}$ quite closely.

3. Binding Energies for the p-shell Hypernuclei.

On the basis of a shell-model representation for the core nuclei, Lawson and Rotenberg¹¹ and Iwao¹² have proposed the following expression

$$B_{\Lambda} = C + N_p \langle \bar{V} \rangle + \alpha_N \langle \Delta \rangle \quad (3.1)$$

for the binding energies of hypernuclei with nuclear cores belonging to the 1p shell. In this expression, N_p denotes the number of p-shell nucleons, and $\langle \bar{V} \rangle$, $\langle \Delta \rangle$ denote the expectation values of $\frac{1}{4}(3V_1 + V_0)$ and $(V_0 - V_1)$, respectively, between the Λ -particle wavefunction and the wavefunction of a p-shell nucleon. The coefficient α_N depends on the spin of the hypernuclear state and on the details of the structure of the core nucleus; values of α_N have

been calculated by Dalitz and Soper¹³ using intermediate coupling wavefunctions for the core nuclei and including the admixture of excited nuclear states where these are of significance. Generally speaking, the values of α_N are close to those for L-S coupling at the beginning of the p-shell, for $A < 9$, and close to those for j-j coupling for $A > 9$. For j-j coupling, the value of α_N is⁷

$$\alpha_N = -\frac{1}{4} \left(1 - \frac{\ell(\ell+1) - \frac{3}{4}}{j(j+1)} \right) (J(J+1) - J_N(J_N+1) - \frac{3}{4}). \quad (3.2)$$

⁷We note here that the result established empirically for ${}^{\Lambda}\text{H}^3$, ${}^{\Lambda}\text{H}^4$ and ${}^{\Lambda}\text{Li}^8$, that the ground-state hypernuclear spin value is $|J_N - \frac{1}{2}|$ is not expected to be a general rule. For example, ${}^{\Lambda}\text{C}^{13}$ belongs to the $p_{\frac{1}{2}}$ shell, so that $j = \frac{1}{2}$, $\ell = 1$, $J_N = \frac{1}{2}$, and the coefficient α_N has the values $+\frac{1}{12}$ for $J = 1$, $-\frac{1}{4}$ for $J = 0$. Hence spin $J = 1$ is expected to hold for the ground state of ${}^{\Lambda}\text{C}^{14}$ and ${}^{\Lambda}\text{N}^{14}$, contrary to the above rule. With $J = 1$, the two-body decays ${}^{\Lambda}\text{C}^{14} \rightarrow \pi^- + \text{N}^{14} (J = 1+)$ and ${}^{\Lambda}\text{N}^{14} \rightarrow \pi^- + \text{O}^{14} (J = 0+)$ are both allowed through the dominant s-interaction. However, if $J = 0$ held for the $({}^{\Lambda}\text{C}^{14}, {}^{\Lambda}\text{N}^{14})$ doublet, the two-body decay ${}^{\Lambda}\text{N}^{14} \rightarrow \pi^- + \text{O}^{14}$ would be forbidden for the s-interaction, although still allowed through the p-interaction (about 10 times weaker). Since O^{14} has no known particle-stable excited states, this means that there can be no confusion from $(\pi^- + \text{O}^{14*})$ modes, and the two-body decay of ${}^{\Lambda}\text{N}^{14}$ would be quite rare; for $J = 0$, the two-body decay ${}^{\Lambda}\text{C}^{14} \rightarrow \pi^- + \text{N}^{14}$ is allowed through the s-interaction.

The derivation of expression (3.1) is based essentially on a variational estimate for B_{Λ} , in which the same trial function is used to describe the Λ -nucleus relative motion for all the p-shell hypernuclei. The wavefunctions directly calculated do not deviate

far from such a mean wavefunction, for as the B_Λ value increases and the tail of the wavefunction falls off more rapidly, so also does the radius of the Λ -nucleus potential increase, partly compensating this effect. The constant term C then results from the kinetic energy terms and the interaction of the Λ particle with the s-shell nucleons.

For hypernuclei whose core nuclei are spinless, expression (3.1) gives a linear relation between B_Λ and N_p . For the systems ${}^5_\Lambda\text{He}$ ($B_\Lambda = 3.1 \pm 0.05$ MeV), ${}^7_\Lambda\text{Be}$ ($B_\Lambda = 4.9 \pm 0.5$ MeV), ${}^9_\Lambda\text{Be}$ ($B_\Lambda = 6.5 \pm 0.15$ MeV) and ${}^{13}_\Lambda\text{C}$ ($B_\Lambda = 10.6 \pm 0.4$ MeV), this linear relation is well satisfied, with $C = 3.1 \pm 0.05$ MeV and $\langle \bar{V} \rangle = 0.90 \pm 0.04$ MeV. This value of $\langle \bar{V} \rangle$ is in good agreement with the value directly calculated from the s-wave Λ -N potentials (without hard core) discussed in Sec. 1.

The spin-dependent term $\langle \Delta \rangle$ can be estimated in four reliable cases, from ${}^7_\Lambda\text{Li}$ ($B_\Lambda = 5.52 \pm 0.45$ MeV), ${}^8_\Lambda\text{Li}$ and ${}^8_\Lambda\text{Be}$ ($B_\Lambda = 6.50 \pm 0.25$ MeV), ${}^9_\Lambda\text{Li}$ ($B_\Lambda = 8.0 \pm 0.3$ MeV), and ${}^{12}_\Lambda\text{B}$ ($B_\Lambda = 10.5 \pm 0.2$ MeV). The comparison with the expression (3.1) is given in Table III. In the ${}^7_\Lambda\text{Be}$ - ${}^7_\Lambda\text{Li}$ comparison, it is clear that $\Delta B_\Lambda = \langle \Delta \rangle$ cannot exceed $0.9(\pm 0.1)$ MeV, otherwise ${}^7_\Lambda\text{Be}$ would not be particle-stable. The value expected for $\langle \Delta \rangle$ on the basis of our knowledge of the Λ -N s-wave potentials is about 0.85 MeV, not incompatible with this value for $A = 7$, especially when the mean wavefunction assumed in (3.1) is replaced by the more diffuse wavefunction appropriate to the low B_Λ values for these systems. On the other hand, the large difference $\Delta B_\Lambda = 1.5$ MeV

between ΛBe^9 and ΛLi^9 requires a correspondingly large value for $\langle\Delta\rangle$, not really consistent with this estimate for $\langle\Delta\rangle$. Indeed, the values of $\langle\Delta\rangle$ given in Table II show a steady increase as N_p increases through the p-shell, roughly in proportion with N_p . The reason for this behavior of the B_Λ values is not yet known. Possibly it reflects the inadequacies of the approximations made in the derivation of the expression (3.1), or perhaps it reflects properties of the Λ -N forces (for example, the effect of a spin-orbit term) which have not yet been taken into account here.

At this point, accurate and reliable B_Λ values for further hypernuclear species in this p-shell region would be of great interest, to explore further the dependence of B_Λ on the spin and structure of the core nucleus. In order for such events to represent a useful addition to the present data, it is essential that they allow a unique and reliable identification of the species, either from analysis of its production process or from its decay process, or both.

II. The Decay Interactions of the Λ -particle

4. The Mesic Decay Processes

It is well known that, owing to the low momentum (100 MeV/c) of the proton resulting from Λ decay at rest, the rate of π^- decay of a Λ particle in nuclear matter is generally reduced as a result of the Pauli principle, since a large fraction

of the final states which can be reached by the final proton are already occupied. This effect is already quite large in ΛHe^5 , where the π^- decay rate is reduced by a factor of about 0.37; in ΛC^{13} , the calculated reduction factor is about 0.14.

It is perhaps less well-known that in the lightest hypernuclei, the effect of the Pauli principle can lead to an enhancement of the decay rate in suitable circumstances, namely if the spin configuration and the decay interaction are such that the final proton is necessarily emitted into a final state which satisfies the symmetry requirements of the Pauli principle. For example, in $\Lambda\text{H}^4 \rightarrow \pi^- + \text{He}^4$ decay with $J = 0$ for ΛH^4 , the s-interaction for Λ decay is spin-independent, so that the final proton is necessarily s-wave and in the spin configuration appropriate to He^4 , and the decay rate for this transition is enhanced by almost a factor 2 (the space wavefunctions do not overlap precisely) as a result of the antisymmetrization appropriate for a final state with two protons.

These same effects hold also for the π^0 decay modes, of course. Thus, with $J = 0$ for ΛHe^4 , the rate for π^0 decay through the s_0 -interaction is enhanced, whereas the rate for π^0 decay through the p_0 -interaction is strongly suppressed; the rate for π^- decay of ΛHe^4 is also strongly suppressed, by a factor about 0.35, since there are already two s-wave protons present in the initial state. Hence the π^0/π^- ratio in ΛHe^4 decay is strongly sensitive to the p_0/s_0 ratio in $\Lambda \rightarrow n + \pi^0$ decay. For free Λ decay, the π^0/π^- ratio is very close to 0.5; for ΛHe^4 decay,

Block et al.¹⁶ have observed the π^0/π^- ratio to be 2.0 ± 0.3 .

From this striking result, Block et al. have been able to deduce that the π^0 mode of Λ decay is dominantly through the s_0 -channel, with $p_0^2 / (p_0^2 + s_0^2) = 0.23 \pm 0.16$, in agreement with expectation (0.12 ± 0.03) from the $\Delta I = \frac{1}{2}$ rule.

5. The Non-mesic De-excitation of the Λ -particle.

For bound Λ particles, the presence of nucleons allows the weak interactions

$$\Lambda + p \longrightarrow n + p, \quad (5.1a)$$

$$\Lambda + n \longrightarrow n + n, \quad (5.1b)$$

to become effective, releasing the full energy difference of about 176 MeV between the Λ -particle and nucleon. This weak interaction can scarcely be investigated without appeal to studies of Λ -hyper-nuclear decay, and ^{rather} little is known of its detailed properties at this time.

These processes (5.1) represent the only strangeness-changing weak interaction which is readily accessible to observation and which involves four strongly-interacting fermions. [†] With

[†] strangeness-conserving
A weak interaction $n + p \longrightarrow n + p$ is expected to exist if the current-current hypothesis of Feynman and Gell-Mann (Phys. Rev. 109, 193 (1958)) ^{holds valid}. This weak interaction is expected to produce parity-violating effects of very small amplitude in nuclear forces and the properties of nuclear states, some of which have been quantitatively estimated by Blin-Stoyle and co-workers.

our present views on the nature of weak interactions, it is natural to expect the existence of a primary four-fermion interaction (5.1a). The validity of the $\Delta I = \frac{1}{2}$ rule for strangeness-changing weak interactions then requires also the existence of a primary interaction of the form (5.1b), although there is no direct evidence for a neutral weak interaction current (in fact, there is very strong evidence against neutral leptonic currents of strength comparable with the charged currents for leptons). In this view, it is natural to regard the $\Lambda \rightarrow N + \pi$ decay interactions as secondary in character since they can occur as a consequence of these four-fermion interactions; this connection is illustrated explicitly in Fig. 1, which depicts one sequence of processes by which the interaction (5.1a) can give rise to $\Lambda \rightarrow p + \pi^-$ decay. Since this four-fermion interaction involves four strongly-interacting particles, mesonic corrections may be expected to distort the form of the interaction quite strongly from the (V-A) form which the current-current theory would predict. This is illustrated in Fig. 2, which shows a series of processes contributing to the physically-observed interaction (5.1a). The terms (a) and (b) show the primary interactions (5.1a) derived from the charged and neutral currents, in turn, the vertices shown being distorted from their primary forms by the usual vertex corrections. The remaining graphs (c) - (f) show corrections which involve the exchange of various mesonic systems between the strangeness-changing and strangeness-conserving currents; the relation between these graphs and the primary interaction may be seen by analyzing each ($\Lambda \rightarrow N + \text{meson}$) vertex in the manner of Fig. 1. The terms (c) and (d) are

well-known, from their discussion by Karplus and Ruderman²⁴ who recognized the importance of their contribution if the Λ -spin were large and used its comparison with the experimental data to argue against the possibility of spin $\geq \frac{3}{2}$ for the Λ -particle. Now that $j_\Lambda = \frac{1}{2}$ is well-established and the $\Lambda \rightarrow N + \pi$ interaction is known to be dominantly s-wave, there is no reason to believe that these Karplus-Ruderman terms contribute dominantly to the observed process $\Lambda + N \rightarrow N + N$; these terms must be taken together with a large number of other mesonic correction terms, of which several are given explicitly in Fig. 2.

Rather little is known at present concerning the properties of this weak interaction

The simplest properties of interest are its charge-dependence and its spin-dependence. For bound Λ -particles, this interaction will dominantly occur for ΛN s-states, in view of the relatively low Λ -N relative momenta to be expected in hypernuclei. The final nucleon momentum for ΛN de-excitation at rest is about $q = 400$ MeV/c, so we shall simplify the discussion by using the non-relativistic form for the matrix-elements. For Λp capture from an s-state configuration, the general form of the matrix-element is

$$\begin{aligned} M(\Lambda p \rightarrow np) = & a_p P_0 + b_p P_2 + \frac{3}{\sqrt{8}} c_p (\underline{\sigma}_Y \cdot \underline{q} \underline{\sigma}_N \cdot \underline{q} - \frac{1}{3} \underline{\sigma}_Y \cdot \underline{\sigma}_N q^2) / M^2 \\ & + \sqrt{\frac{3}{8}} d_p (\underline{\sigma}_Y + \underline{\sigma}_N) \cdot \underline{q} / M + e_p (\underline{\sigma}_Y - \underline{\sigma}_N) \cdot \underline{q} P_0 / 2M \\ & + \sqrt{3} f_p (\underline{\sigma}_Y - \underline{\sigma}_N) \cdot \underline{q} P_2 / M, \end{aligned} \quad (5.4)$$

where $P_0 = (3 + \underline{\sigma}_Y \cdot \underline{\sigma}_N)/4$ and $P_1 = (1 + \underline{\sigma}_Y \cdot \underline{\sigma}_N)/4$ denote the singlet and triplet spin projection operators, $\underline{\sigma}_Y$ denotes the Λ or n spin and $\underline{\sigma}_N$ here denotes the proton spin. In this expression, the terms a_p and e_p denote the 1S_0 transition amplitudes, leading to the 1S_0 and 3P_0 final states, respectively, and the terms b_p, c_p, d_p and f_p denote the 3S_1 transition amplitudes, leading to the $^3S_1, ^3D_1, ^3P_1$ and 1P_1 final states. For Λn capture, the corresponding form is

$$M(\Lambda n \rightarrow nn) = a_n P_0 + \frac{\sqrt{3}}{\sqrt{8}} d_n (\underline{\sigma}_Y + \underline{\sigma}_N) \cdot \underline{q} / M + e_n (\underline{\sigma}_Y - \underline{\sigma}_N) \cdot \underline{q} P_0 / 2M, \quad (5.5)$$

where $\underline{\sigma}_Y$ here denotes the spin of the Λ particle and either of the final neutrons and $\underline{\sigma}_N$ denotes the spin of the other neutron. Terms of the form b, c and f are absent here, as they lead to final states which are forbidden for the n - n system. If the $\Delta I = \frac{1}{2}$ rule holds, then we have in addition the following equalities

$$a_n = \sqrt{2} a_p, \quad d_n = \sqrt{2} d_p, \quad e_n = \sqrt{2} e_p. \quad (5.6)$$

We note that the transitions a, b and c conserve parity, whereas the transitions d, e and f reverse the parity of the state. If time-reversal invariance holds, then the phases of each of these amplitudes may be determined from the known NN scattering phases, by the use of Watson's theorem. Obviously the determination of all these parameters, and the test of the equalities (5.6), will require polarization experiments, involving the measurement of the longitudinal polarization of the fast n and p emitted from unpolarized hypernuclei (or the angular distribution of the fast nucleons from non-mesic decay of polarized hypernuclei), and of the polarization correlation coefficients

for the fast (np) and (nn) pairs emitted in non-mesic decay of hypernuclei. These polarization coefficients could also be measured for initial ΛN states of definite spin by selecting the appropriate light hypernuclei (cf. below). Such experiments will be difficult to carry out and to interpret, and lie far in the future.

In principle, for the calculation of non-mesic decay rates for hypernuclei, matrix-elements of the forms (5.4) and (5.5) should be used for each of the nucleons and evaluated between the initial hypernuclear wavefunction and the final nuclear states. Here, we shall consider instead a simplified calculation for the non-mesic decay rates, which treats the Λ de-excitation by different nucleons as incoherent. This approximation is not strictly valid, because, owing to the identity of the neutrons and of the protons, the same final state can generally be reached through capture on any one of the nucleons; however, because of the large energy release, it is rather likely that these interference effects will generally be small and will tend to cancel out in the total de-excitation rates, summed over all final states.

Let us denote by R_{NS} the $\Lambda N \rightarrow NN$ transition rate, when the ΛN system is in an s-state of total spin S, for unit density of nucleons of type N at ^{the} Λ position. The non-mesic decay rate for a given hypernucleus is then given by the expression

$$\begin{aligned} R_{n.m.} &= \bar{R}(\Lambda Z^{A+1}) \rho_A \\ &= \bar{R}(\Lambda Z^{A+1}) \int \psi_\Lambda^2 \rho_N(\underline{r}) d^3r, \end{aligned} \quad (5.7)$$

where ρ_A is the mean density of nucleons at the Λ position, $\rho_N(\underline{r})$ being the nucleon density at position $\underline{r}$ and ψ_Λ the wavefunction for the Λ -nucleus relative motion, and $\bar{R}$ is the spin and charge average of the R_{NS} ^(which is) appropriate to the hypernucleus considered. In terms of the amplitudes (5.4) and (5.5), these quantities R_{NS} are given by

$$R_{po} = |a_p|^2 + |e_p|^2 (q/M)^2, \quad (5.8a)$$

$$R_{pi} = |b_p|^2 + |c_p|^2 (q/M)^4 + |d_p|^2 (q/M)^2 + |f_p|^2 (q/M)^2, \quad (5.8b)$$

$$R_{no} = (|a_n|^2 + |e_n|^2 (q/M)^2), \quad (5.8c)$$

$$R_{ni} = |d_n|^2 (q/M)^2. \quad (5.8d)$$

In the last expressions R_{NS} , we have taken into account the symmetry requirements of the final state and the identity of the two final neutrons. If the $\Delta I = \frac{1}{2}$ rule holds, the equalities (5.6) lead to the predictions

$$R_{po} = \frac{1}{2} R_{no}, \quad (5.9a)$$

$$R_{pi} \geq \frac{1}{2} R_{ni}. \quad (5.9b)$$

The quantity most readily measured is the ratio $Q =$ (non-mesic/ π^- -mesic). Here, we shall not include π^0 -mesonic decay events ²⁵ among the non-mesic decay rate, for the π^0 -mesonic events normally give an exceedingly small star with a visible energy release of at most several MeV and would very frequently be overlooked, whereas

the non-mesic decay events involve a large energy release. For Λ He hypernuclei in emulsion, estimates of this ratio have been reported as follows: $Q = 2.3 \pm 1.0$ by Fry,¹⁷ $1.5(\pm 0.4) \leq Q \leq 2.8$ by Silverstein,¹⁸ and $Q \geq 1.1 \pm 0.5$ by Schlein,¹⁹ whose non-mesic decay rate includes only two-prong events. For Λ He⁴, Block et al.¹⁶ have obtained a rather accurate value, $Q = 0.52 \pm 0.10$. About 80% of the π^- decays of Λ He in emulsion are due to Λ He⁵ decay; adopting the value $Q = 1.5$ for Λ He, the correction for admixture of Λ He⁴ non-mesic decay events leads to the estimate $Q = 1.8$ for Λ He⁵. On this basis, it is then possible to estimate Q for some of the heavier hypernuclei. This involves an estimate of $R_{n.m.}$ from (5.7), using the value $\rho_4 = 0.038 \text{ F}^{-3}$ for Λ He⁵ and corresponding calculated values ρ_A for the hypernuclei of interest, together with an estimate of the π^- decay rate $R(\pi^-)$ using the completeness-relation method.²⁰ For example, for Λ He⁵, the calculated value of $R(\pi^-)$ is $0.25 \tau_\Lambda$, where $1/\tau_\Lambda$ denotes the free Λ lifetime, so that our estimate of Q and the above value for ρ_4 lead to the estimate

$$\begin{aligned} \bar{R}(\Lambda\text{He}^5) &= \frac{1}{8}(3R_{p1} + R_{p0} + 3R_n + R_{n0}) \\ &= (1.8)(0.25)\tau_\Lambda / (0.038) = 12(\pm 3)\tau_\Lambda, \end{aligned} \quad (5.10)$$

where the error quoted represents only the statistical uncertainty on Q . With this value for $\bar{R}$, estimates have been made for Q for Λ Li⁷ (neglecting the spin-dependence of R_{NS}), Λ Be⁹ and Λ C¹³, as given in Table III. The rapid rise in Q with increasing A is due primarily to

the rapidly increasing suppression of the π^- decay process by the Pauli principle. However, the non-mesic decay does increase from 0.45 Λ for ΛHe^5 , through 1.5 Λ for ΛC^{13} , to a value of 2.0 Λ for a very heavy hypernucleus (say $A \approx 100$). Several independent estimates of $R(\pi^-)$ ^{we have made} for a hypernucleus of mass number $A \approx 100$, including pion re-absorption effects, lead to values about 0.015 Λ , and therefore to a large value $Q \approx 130$. We conclude that the decay lifetime of a heavy hypernucleus ($A \approx 100$) may be expected to be essentially independent of A , with the value $\approx 1.2 \times 10^{-10}$ sec.

The evidence on the charge dependence of the non-mesic process appears somewhat contradictory at first sight. For ΛHe^4 , Block et al.¹⁶ have classified non-mesic decay events with a recoil proton of momentum exceeding 250 MeV/c as due to Λp capture, leading to the estimate $(\Lambda p)/(\Lambda n) = 2.2 \pm 0.8$. Since ΛHe^4 has $J = 0$, this corresponds to the ratio

$$\frac{(3 R_{pl} + R_{po})/2}{R_{no}} = 2.2(\pm 0.8). \quad (5.11)$$

On the other hand, for spallation hyperfragments following high energy K^- interactions in emulsion, for which typically $A \approx 70$, Beniston²² and Renard²³ have compared the observed proton spectrum with those predicted by a Monte Carlo calculation for (Λp) and (Λn) capture processes, and have concluded that, in these heavy hypernuclei, the (Λn) capture process is the dominant non-mesic decay process. The ratio $(\Lambda n)/((\Lambda n) + (\Lambda p))$ obtained by

Beniston was $\gamma = 0.8 \pm 0.1$, and by Renard, $\gamma = 0.65 \pm 0.1$, the difference between these values being due, at least in part, to the differing assumptions made about the spatial distribution of the Λ particle. To continue, we shall adopt the value $\gamma = 0.7$, so that

$$\frac{3 R_{p1} + R_{po}}{3 R_{n1} + R_{no}} = \left(\frac{1}{\gamma} - 1 \right) \approx 0.4 \quad (5.12)$$

From this ratio and the value (5.11), we can conclude the value $3R_{n1} + R_{no} = 69(\pm 20) \Lambda$, to be used below. The striking difference between the $(\Lambda p)/(\Lambda n)$ ratios observed for ΛHe^4 and for heavy hypernuclei can be understood quite simply if the (Λn) de-excitation process has a strong spin-dependence, with

$$R_{no}/R_{n1} \approx 0.3, \quad (5.13)$$

the non-mesic (Λn) decay interaction being significantly weaker in the singlet configuration.

The remaining information to be discussed consists of the non-mesic rates for ΛHe^4 and ΛH^4 , for which the nucleon density ρ_3 is 0.019 F^{-3} . For ΛHe^4 , the calculated suppression²⁰ of $R(\pi^-)$ to 0.26Λ , and the *above value* $Q = 0.52 \pm 0.10$, lead to

$$3 R_{p1} + R_{po} + 2 R_{no} = 41 (\pm 10) \Lambda. \quad (5.14)$$

From (5.10), (5.13) and (5.14), we conclude that $R_{n1} = 21(\pm 7) \Lambda$, $R_{no} = 6(\pm 2) \Lambda$, and $(3 R_{p1} + R_{po}) = 28(\pm 8) \Lambda$. For ΛH^4 , the π^- decay rate $R(\pi^-)$ has been calculated²⁰ to be 0.74Λ , and the non-mesic rate is given by

$$0.31 (\pm 0.06) \geq (0.019) (R_{po} + 3R_{n1} + R_{no})/6 \geq 0.22 (\pm 0.06), \quad (5.15)$$

using the above value for $(3R_{n1} + R_{no})$ and the requirement $0 \leq R_{po} \leq 28(\pm 8) \Lambda$. On this basis, the ratio Q for ΛH^4 is expected to lie between $0.29(\pm 0.08)$ and $0.42(\pm 0.08)$. In the experiment of Block et al.,¹⁶ $120(\pm 11) \Lambda H^4 \rightarrow \pi^-$ events were seen. From the number of ΛHe^4 production events and from charge independence, it was concluded that the total number of ΛH^4 production events was 163 ± 10 . From the π^0/π^- ratio (2.5 ± 0.5) observed for ΛHe^4 decay, and from the prediction of the $\Delta I = \frac{1}{2}$ rule that²⁰

$$\frac{\pi^0}{\pi^-} (\Lambda He^4) \times \frac{\pi^0}{\pi^-} (\Lambda H^4) = \frac{1}{4}, \quad (5.16)$$

we may estimate the number of $\Lambda H^4 \rightarrow \pi^0$ events as 12 ± 1 . The remainder of the events, amounting to 31 ± 15 , must be attributed to non-mesic decay processes; this leads to the ratio $Q = 0.26 \pm 0.13$, which is not at all in disagreement with the value predicted above.

This small fraction of ΛH^4 decay events which proceed through the non-mesic mode may not be easy to identify in the helium bubble chamber experiment, especially because of the confusion possible with slow Σ^- hyperons following $K^- - He^4$ capture which come to rest and interact to give a one-prong star. If we assume (for lack of other knowledge²¹) that the Λp capture interaction is spin-

²¹With the $\Delta I = \frac{1}{2}$ rule for the $\Lambda N \rightarrow NN$ process, the value $R_{no} = 6(\pm 2) \Lambda$ and Eq.(5.9a) would require $R_{po} = \frac{3.0}{1.5}(\pm \frac{1.0}{0.5}) \Lambda$, a much smaller value than that assumed here. With this value the $(\Lambda p)/(\Lambda n)$ ratio for ΛH^4 would be $0.04(\pm 0.02)$, which would predict only $\frac{2}{one}$ non-mesic ΛH^4 with a fast proton in this experi-

ment. We may note here that these spin-dependences for the Λn and Λp processes are quite at variance with a simple (V-A) form for the $\Lambda N \rightarrow NN$ interaction; in the non-relativistic limit, the first term: (a_p) of (5.4) would be dominant, and the ratio R_{po}/R_{pl} would be predicted to be exceedingly large.

independent, i.e. $R_{po} = R_{pl}$, then R_{po} has the value $7 (\pm 2) \sqrt{\Lambda}$ and the $(\Lambda p)/(\Lambda n)$ ratio predicted for ΛH^4 has the value $2R_{po}/(3R_{nl} + R_{no}) = 14 (\pm 2)/69 (\pm 20) = 0.2 (\pm 0.06)$, so that only 6 non-mesic ΛH^4 decay events are expected to occur with a fast proton (say $\gtrsim 250$ MeV/c) in the experiment of Block et al.¹⁶ However, the majority of the (Λn) capture events are expected to give a visible prong. Some indication on this can be obtained from the $\pi^- + He^4$ interaction studies of Bortolani et al.,²³ which also proceeds through two-nucleon emission with a comparable energy release (≈ 140 MeV). Of 356 π^- interactions at rest in a helium chamber, only 23 gave no visible track; since the primary capture interaction $\pi^- + p + n \rightarrow n + n$ is believed to account for at least $3/4$ of these capture events (the remainder being due to the $\pi^- + p + p \rightarrow n + p$ interaction), we may estimate[†] that not more than 9% of the (Λn)

[†]We should emphasize here that the $\Lambda n \rightarrow nn$ interaction may have quite different spin dependence from the $\pi^- pn \rightarrow nn$ interaction, in which case this proportion might well turn out to be larger than this estimate. For example, H^3 production is found to occur in 22% of $\pi^- - He^4$ capture events; this appears to be a particular feature

of the (π^-pn) and (π^-pp) capture amplitudes. There is certainly no reason to expect the (Λn) and (Λp) capture amplitudes to have a similar spin and charge dependence, and it is quite unlikely that the mode $H^3 + n$ will prove to be as abundant in ΛH^4 non-mesic decay.

de-excitation events for ΛH^4 will give no visible track. Although the separation of the ΛH^4 non-mesic decay events is a difficult task, a value of the $(\Lambda p)/(\pi^-mesic)$ ratio for ΛH^4 , or even an upper limit on this would be very valuable in providing a limit on R_{po} , and therefore a direct indication of the spin dependence for the (Λp) de-excitation) process.

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TABLE I

Hard core radius d	d = 0.0 F (Gaussian Shape)	d = 0.2 F (Square well)	d = 0.4 F (Square well)
a_0 (F.)	$-2.4 \begin{smallmatrix} +1.2 \\ -0.6 \end{smallmatrix}$	$-4.2 \begin{smallmatrix} +3.0 \\ -1.3 \end{smallmatrix}$	-4.5
a_1 (F.)	-0.5 ± 0.1	-0.55 ± 0.05	-0.6
Notes	obtained from Λ^3 and ΛHe^5 (Refs. 3, 5)	obtained from Λ^4 and ΛHe^5 (Ref. 6)	obtained from Λ^4 and ΛHe^5 (Ref. 7)
Existence of J = 1 excited state Λ^4 ?	≈ 1.3 $B^*_{\Lambda} \begin{smallmatrix} \leftarrow 0.4 \\ \leftarrow 0.1 \end{smallmatrix} \text{ MeV}$ ($S^* = 1.25$)	Not bound ($s^* = 0.90$)	Not bound ($s^* = 0.83$)

TABLE II

Hypernuclei compared	ΔB_{Λ} (MeV)	Estimate from Eq. (3.1)*	Value of $\langle \Delta \rangle$
${}^{\Lambda}\text{Be}^7 - {}^{\Lambda}\text{Li}^7$	0.6 ± 0.5 (≤ 0.9)	$0.98 \langle \Delta \rangle$	$\langle \Delta \rangle = 0.6 \pm 0.5$ MeV
${}^{\Lambda}\text{Be}^8 - {}^{\Lambda}\text{Be}^9$	0.0 ± 0.3	$0.59 \langle \Delta \rangle - \langle \bar{V} \rangle$	$\langle \Delta \rangle = 1.5 \pm 0.5$ MeV
${}^{\Lambda}\text{Be}^9 - {}^{\Lambda}\text{Li}^9$	1.50 ± 0.3	$0.70 \langle \Delta \rangle$	$\langle \Delta \rangle = 2.1 \pm 0.4$ MeV
${}^{\Lambda}\text{B}^{12} - {}^{\Lambda}\text{C}^{13}$	-0.10 ± 0.4	$0.28 \langle \Delta \rangle - \langle \bar{V} \rangle$	$\langle \Delta \rangle = 2.7 \pm 1.3$ MeV

*The coefficients α_N used here are those computed from intermediate coupling nuclear wavefunctions constructed by J. Soper.

TABLE III $Q = (\text{non-mesic})/(\pi^- \text{-mesic})$

Hypernucleus	$\Lambda^4\text{H}$	$\Lambda^4\text{He}$	$\Lambda^5\text{He}$	$\Lambda^7\text{Li}$	$\Lambda^9\text{Be}$	$\Lambda^{13}\text{C}$	$\Lambda^{109}\text{Ag}$
$Q =$	<u>0.15</u>	<u>0.5</u>	<u>1.8</u>	2.7	6.2	11.3	130

The underlined entries represent the empirical values for Q ; the other entries are calculated values based on the entry for $\Lambda^5\text{He}$.

Fig. 1. The Decay $\Lambda \rightarrow p + \pi^-$ as a consequence of the Four-fermion Weak Interaction. (5.1a)

Fig. 2. Contributions to the Process $\Lambda + p \rightarrow n + p$.