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WHAT (IF ANYTHING) CAN FEW-BODY STRANGE SYSTEMS TEACH
US ABOUT QUARK-GLUON HADRONIC SUBSTRUCTURE?

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ABSTRACT

We discuss expectations, relevant to the proposed (π, K) program at PILAC, for the effects of hadronic quark-gluon substructure on the physics of few-body strangeness -1 systems, in the context of QCD-inspired models used previously to describe the hadron spectrum and short distance nucleon-nucleon scattering.

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INTRODUCTION

QCD-inspired quark models (eg., the non-relativistic quark, or cluster model, the MIT bag model, the cloudy bag model) incorporating what are believed to be the main relevant features of QCD, have proven capable of providing a good description not only of the spectrum of low-lying meson and baryon states, but also of the major features of their decay couplings¹⁻⁵. Numerous studies, moreover, of the low energy NN system, largely in the cluster model framework, have demonstrated that, at the level at which such models can be realistically constrained by the spectrum, those features of NN s- and p-wave scattering traditionally associated with vector meson exchange can be viewed equally well as arising from the underlying quark-quark interactions, vector meson exchange being completely suppressed⁶⁻⁸. This observation represents, to some at least, precisely what would be expected, given the known relative scales of the problem, vector meson exchange appearing rather problematic in view of the fact that vector meson ranges are considerably less than the rms proton charge radius. This statement, of course, rests on two underlying assumptions: 1) that mesons are degrees of freedom of the physical, non-perturbative, but not of the short-distance perturbative QCD vacuum, and 2) that the nucleon interior corresponds, at least approximately, to a configuration of quarks moving in the latter vacuum. This second assumption is by no means verified, though lattice simulations will, presumably, eventually shed light on the question. Even if it were to be verified (as appears likely, given that light quark color fields are known to suppress instanton configurations⁹), it is not clear that the “quark core” region of the nucleon, for example, need occupy a major fraction of the physical nucleon volume. Given these uncertainties, the quark-quark/vector meson exchange duality noted above is rather frustrating as it means that there exists no obvious candidates to distinguish between the two approaches. In particular, it means that, while one can demonstrate that, if the quark cores do occupy a sizeable fraction of the physical hadronic radius, there are, generally small, but non-negligible quark-induced effects on observables such as binding energies, form factors, response functions etc.¹⁰⁻¹² which would

not be present in the corresponding meson exchange picture, one cannot, in light of our inability to include ever higher mass meson exchanges, at ever higher order, involving ever higher mass intermediate baryonic excitations in that latter picture, ever demonstrate that such novel quark-induced effects are **required** by the data. This “theoretical hysteresis effect” plagues the field. The possibility of access to few-body strange systems, therefore, is a highly desirable one. Indeed, it has long been recognized that certain features of few-body strange systems, at least naively, appear particularly amenable to exposing possible quark effects. As examples let us mention 1) the question of whether the explicit u-d quark content of the Λ and Σ ’s could lead to Pauli blocking effects in hypernuclei not expected in the meson exchange picture, where hyperons are treated as distinct from nucleons¹³; 2) the possibility of strange quark delocalization in hypernuclei¹⁴; 3) the observation that, because of the structure of the quark level hyperfine interaction common to QCD-inspired models, exotic, dibaryon configurations are more likely in strange than in non-strange sectors¹⁵; 4) the fact that the high momentum transfer of the elementary non-mesonic $\Lambda N \rightarrow NN$ process means that explicit quark effects, if present at all in nuclei, might be expected to be especially evident in weak hypernuclear decay. We will examine such questions in more detail in what follows, our aim being to determine whether any interesting possibilities exist for significant quark effects accessible to a (π, K) physics program with 1.05 GeV/c pions.

HYPERON-NUCLEON INTERACTIONS

In this section we will assume, for argument’s sake, that the quark level description of the short-distance NN interaction represents the correct basic physics and ask what one would expect, in the same framework, for the low energy hyperon-nucleon interaction, in particular, whether there exist any features of these predictions which would distinguish them from those of meson exchange models. A few preliminary words, however, regarding the state of the NN scattering calculations, before proceeding.

The situation is, as stated above, that qq interactions are capable of generating the correct short range NN physics. The effects are induced primarily by the color hyperfine (hf) interaction of one gluon exchange (OGE), through the requirements of quark anti-symmetrization (AS'n). Confinement effects play a negligible role. In order to compare to experimental phase shifts, however, additional physics has to be added to the models to account for the long-range interaction and also the intermediate range NN attraction. The former is due to π exchange, which one necessarily expects to be present, on scale grounds, and is implemented by means of effective πqq couplings constrained to reproduce the known πNN strength. While there are indications that the NN attraction may be due, in the quark picture, to a small degree of quark delocalization associated with incomplete restoration of the physical vacuum in the nuclear medium¹⁶, it is not presently known how to model this effect in the context of scattering calculations, and the intermediate range attraction is implemented either by adding purely phenomenological contributions or by including the exchange of a fictitious scalar σ meson through the introduction of a σqq coupling. One may then obtain a good fit to the low energy s- and p-wave phase shifts (see Refs. 7 for the quality of the fits). The calculations, however, have been done generally in the approximation that only the symmetric spatial portion of the nucleon wavefunction is considered. In this approximation the nucleon quark core size cannot be unambiguously determined dynamically and this leads to quantitative (though not qualitative) uncertainties of order 20-30% in the calculated phase shifts. Further, while such an approximation is known to lead to an understanding of the spectrum which is quite compatible with that obtained from more sophisticated treatments (with one exception, discussed below), the inclusion of the mixed symmetry radially excited component of the nucleon wavefunction is known to produce non-trivial quantitative (though not qualitative) changes (again at the 20-30% level) in the resulting NN s-wave phase shifts¹⁷. This point bears particularly on the results for the OGE-induced NN spin-orbit force. Indeed, a OGE spin-orbit term of the strength employed to reproduce the vector meson exchange contribution to NN spin-orbit forces

would lead to unacceptably large spin-orbit splittings in the baryon spectrum, at least in the symmetric-spatial-configuration- only approximation. It is known that a more sophisticated treatment, which removes this approximation, produces acceptable baryon spin-orbit splittings¹⁸, but the corresponding NN scattering calculations have not been done using the resulting (more complicated) nucleon wavefunction structure. This is, no doubt, due to the fact that the required spin-orbit resonating group kernels are much more difficult to calculate for such a wavefunction. Thus, such a calculation, while highly desirable, does not appear likely in the foreseeable future. While the same (less than optimum) situation prevails for the hyperon-nucleon spin-orbit results discussed below, the treatment is, at least, compatible with that for NN.

Let us, then, turn to the case of hyperon-nucleon (YN) scattering. We quote results obtained from two recent resonating group method studies of YN phase shifts in the cluster model^{19,20}. These two studies differ slightly in their parametrizations of the effective QCD-inspired quark Hamiltonian, H_q ; see Refs. 19 and 20 (and references therein) for further details. The concomitant differences in the results quoted below are typical of those associated with the simplified treatment of the baryon structure, as discussed above for NN. If we consider only the short-distance, central YN forces resulting from H_q , two main features are evident. First, all s-wave YN residual interactions are repulsive and hard-core-like, and second, the induced $N\Lambda - N\Sigma$ transition is very weak, in both spin channels. The equivalent hard-core radii for the various YN channels are displayed in Table I. For comparison the equivalent hard-core radii for the NN channels $(J,I)=(1,0)$, $(0,1)$, for the parametrization of Ref. 19, are .44 and .50 fm, respectively (those for the parametrization of Ref. 20 are not quoted). Note that the $N\Sigma$ $(0, \frac{1}{2})$ and $(1, \frac{3}{2})$ channels are considerably more repulsive than the rest. This may be understood as due mainly to the Pauli principle. Indeed, the spins and isospins in the former case are both antiparallel while, in the latter, they are both parallel. Such a configuration would be Pauli forbidden for NN s-waves. While this is no longer the case for YN, one sees the residue of the effect in the existence

of small eigenvalues of the norm kernel ($1/4$ and $2/9$ respectively). The resulting “Pauli repulsion”, analogous to that in $\alpha\alpha$, then distinguishes these channels from the rest, for which the repulsion is, as in NN, dominated by the exchange color hf interaction. This prediction for the large difference in relative repulsive strengths of the two $N\Sigma$ channels relative to other YN channels is generic to the “OGE+quark AS’n” picture upon which the NN results are based and hence is of considerable interest. Unfortunately, the PILAC energy is not sufficient to access $N\Sigma$ scattering.

Table I Equivalent s-wave YN hard core radii from H_q (in fm)

	(J,I)	Ref. 19	Ref. 20
$N\Lambda$	$(0, \frac{1}{2})$	.44	.57
$N\Lambda$	$(1, \frac{1}{2})$	.37	.54
$N\Sigma$	$(0, \frac{1}{2})$	.72	.86
$N\Sigma$	$(1, \frac{1}{2})$	.30	.29
$N\Sigma$	$(0, \frac{3}{2})$	.40	.38
$N\Sigma$	$(1, \frac{3}{2})$	.77	.89

Concerning the results on the smallness of the short-range $N\Lambda - N\Sigma$ coupling, we see that the coupling is predicted to be dominated by the long-distance pionic mechanism. It is worthwhile noting that the Nijmegen one-boson exchange (OBE) YN potentials^{21,22}, including the above transition potential, have phenomenological hard cores of radius $\sim .5fm$. In Ref. 20 an SU(3) symmetric baryon-baryon σ coupling plus quark coupling to the octet of pseudoscalars (constrained to reproduce OBE model baryon-pseudoscalar couplings) have also been added in order to simulate intermediate and long range forces. Comparing the s-wave phase shifts of Ref. 22 for the Nijmegen D and F potentials to those of Ref. 20, one finds them to be compatible within the uncertainties of either framework.

Certain small systematic differences exist—for example the NA singlet channel on average lies slightly above the mean Nijmegen results whereas the triplet lies slightly below—but these should not be considered to be significant at this point since the results of Ref. 20 provide almost as good a fit to the existing (low quality) YN data as do the Nijmegen potentials (which are fit to these results) and with far fewer parameters. Less is known about the relation of YN spin-orbit forces in the quark and meson exchange pictures. This is due to the fact that the relevant calculations have yet to be performed in the quark framework. Morimatsu *et al.*²³ evaluate adiabatic YN spin-orbit potentials using only the Galilei-invariant piece of the qq spin-orbit interaction and find symmetric $NA(I = 1/2)$ and $N\Sigma(I = 3/2)$ potentials similar to those of the Nijmegen group but a $N\Sigma(I = 1/2)$ symmetric potential much smaller and opposite in sign. They also find strong antisymmetric spin-orbit potentials in the $NA(I = 1/2)$ and $N\Sigma(I = 1/2)$ systems (~ -30 and $+80 MeV$, respectively, at $r = 0$), whereas the Nijmegen results are small. Unfortunately, dropping the Galilei-non-invariant qq spin-orbit terms is not relativistically consistent, to the order of the calculation²⁴ and the calculation has not been redone to include them. In view of results for the Born approximation t -matrices mentioned below, however, it is probable that the qualitative features of the results may be unchanged. He, Wang and Wong²⁵ have also looked at YN spin-orbit forces, evaluating spin-orbit t -matrices in Born approximation and comparing to the corresponding Nijmegen t -matrices, having first constrained the quark level spin-orbit strength to reproduce the NN spin-orbit t -matrices of the Paris potential. They find that the symmetric YN spin-orbit t -matrices so generated are similar to those of the Nijmegen potentials, at least within variations already present between the D and F cases, though, for Σ^-p , they are $\sim 50\%$ smaller than the (smaller) F values (themselves $\sim 50\%$ smaller than the D values) and for ΛN they are peaked at somewhat higher momentum transfer. The antisymmetric spin-orbit t -matrices are less similar, being a factor of ~ 6 larger in the quark framework for Σ^-p and non-zero for Σ^0p where the Nijmegen results are small (and of different sign for D and F). If one makes a similar comparison, however, of the

results of Ref. 25 (in the quark model) with those of a different YN meson-exchange model (one generated, in analogy with the Bonn potential for NN, by considering both one-meson-exchange and two-meson-exchange-box diagrams (including Y^* intermediate states)²⁶), one finds that the quark model results are now rather similar to those of the meson-exchange model, i.e., that the results of Ref. 26 are closer to those of the quark model than they are to either the Nijmegen D or F results. This is of particular note for the antisymmetric spin-orbit potentials where one might otherwise have hoped to be able to distinguish the two approaches. Such a possibility appears unlikely given the present theoretical situation. It should be pointed out that there remains a potential systematic uncertainty in the quark model results of Ref. 25. This is because some of the NN spin-orbit interaction undoubtedly comes from intermediate range physics. The initial fitting procedure of Ref. 25 is, therefore, somewhat suspect, especially as the quark level spin-orbit strength required is considerably larger than that obtained from fits to the spectrum, even ignoring the spin-orbit problem there. The appropriate comparison, namely of the short-range quark-induced to vector-meson-induced contributions has not been made. The calculation of Ref. 25 does, however, display the effect of the Galilei-non-invariant terms dropped by Morimatsu *et al.*. One sees that the qualitative agreement with the symmetric spin-orbit meson exchange results is left unchanged, as are the differences between the quark and Nijmegen results and similarity between the quark and Ref. 26 results for the antisymmetric spin-orbit t-matrices. One may thus hope that the non-overlapping deficiencies of the two quark model calculations to some extent compensate. The conclusion about the absence of a clear signature distinguishing the quark and meson exchange approaches would then remain. The situation would, of course, be much more satisfactory if the intermediate range contributions were handled more plausibly in the quark model calculation. The approach of Ref. 20 provides a framework for a more satisfactory calculation than presently exists, in particular since, in comparing to the experimental YN observables, p-waves have been included, and, in fact, improve the agreement with experiment. The relevant comparisons

have not, however, been explicitly made. We stress again, though, that the fit to the YN data of Ref. 20 is quite reasonable. This is also true of both meson exchange model calculations, by construction.

We close this section with remarks on three aspects of hypernuclear physics from the quark perspective. First, contrary to naive hopes (expectations?), it is very unlikely that the strange quark of the Λ will be significantly delocalized in the hypernuclear environment. Indeed, we learn from Ref. 16 that, in the case of nucleons, at least, it is very important not to neglect the quark hf interaction, which is, after all, responsible for the fact that the octet baryons lie far below the decuplet in the spectrum. One finds that the induced hf repulsion strongly disfavors large quark delocalization, in the ordinary nuclear case. Since the $N\Lambda$ induced central repulsion is very similar to that for NN, the same conclusion will hold in the hypernuclear case. A second point of note is the observation of Hungerford and Beidenharn¹³ that, in ${}^5_\Lambda\text{He}$, the light quarks of the Λ are strongly Pauli blocked in the s-shell, thus providing a natural explanation of the “weak” binding of the system observed experimentally. Unfortunately, the notion that the binding in “weak” is based on an overly naive evaluation of long distance physics²⁷: if one includes the effects of the Σ degree of freedom, which is certainly present and coupled by known long-range physics, the binding is also naturally “weakened” (this is also true for the ${}_\Lambda A = 4$ system). The third point concerns the observed charge symmetry breaking (CSB) in the ${}_\Lambda A = 4$ system (experimentally $\Delta B_\Lambda = B({}^4_\Lambda\text{He}) - B({}^4_\Lambda\text{H})$ is $\sim 340\text{keV}$ for the $J^P = 0^+$ ground state and $\sim 230\text{keV}$ for the $J^P = 1^+$ excited state). These numbers are of interest because a microscopic meson exchange evaluation of the contributions due to particle mixing finds the ground state contributions to be dominated by $\Lambda - \Sigma$ mixing (which allows an isovector OBE component to what is nominally the $N\Lambda$ interaction). This contribution, however, changes sign in the 1^+ excited state, leading to a small negative net value for ΔB_Λ ²⁸. One should bear in mind that the results of Ref. 28 are obtained in an effective two-body folding model and that the excited state is assumed related to the ground state by a Λ spin flip. Since the folding approach is

known to underestimate ΔB_Λ by a factor of ~ 2 relative to the results of exact four-body calculations in the case of a separable approximation to the Nijmegen D potential²⁷, one should not place too much faith in the magnitudes of ΔB_Λ and ΔB_Λ^* of Ref. 28. The relative value, however, is of interest, as it does not appear bound to the shortcomings of the folding approximation. A similar comment holds for the absence of configurations other than the $(A = 3 \text{ core} + \Lambda)$ one in the folding model: while these are undoubtedly important in determining the hypernuclear structure, it appears very unlikely that their inclusion (with amplitudes $\sim .1$ relative to the leading one) could alter the conclusions of Ref. 28 concerning the dominant particle mixing CSB contributions. Exact four-body calculations for the excited state using CSB interactions constructed, as for the Nijmegen potentials, and in Ref. 28, in the same manner as those which enter the calculation of similar CSB contributions in the $A=3$ system, are, however, desirable, in order to check this supposition. CSB effects are of interest from the quark perspective because one can show, based on the constraints of chiral symmetry, that there are QED X QCD “interference” effects which play at least as significant a role in determining isomultiplet splittings in the spectrum as do those of ordinary 1-photon-exchange^{29,30}. These interactions lead to contributions to nuclear CSB which do not correspond to any of the sources considered in traditional nuclear calculations, and which apparently play an important role in the non-Coulombic CSB contribution to the $A=3$ binding energy difference^{30,31}. If one, however, considers the $(A = 3 \text{ core} + \Lambda)$ description of the $J^P = 0^+, 1^+$ ground and excited states, as above, one finds these novel contributions to be negligible in the ground state and to produce small negative contributions for the excited state³⁰. These results are not sensitive to the uncertainties present in such quark model calculations. These are not, however, the only quark level effects expected: there are also those due to the flavor dependence of the qq spin-dependent strong interactions and quark kinetic energy associated with the difference of up and down quark masses. Unfortunately, these contributions have not yet been estimated for the $\Lambda A = 4$ system, though this work is in progress. Experience in

other quark model calculations gives some hope that such effects could be of the same sign in both the $0^+, 1^+$ states and of the desired magnitude, but one must await the outcome of the calculations. While the whole question is an interesting one, and deserves further theoretical effort, there appears no reason to doubt existing experimental evidence. This means that, while improved hypernuclear data is always of interest, there is no strong role for PILAC to play in this matter.

EXOTIC HADRONS

One of the clearest signs for the necessity of explicitly including quark and gluon degrees of freedom in the description of few-body nuclear and hypernuclear systems would be the presence of exotic dibaryon states. According to standard QCD-inspired quark model arguments several such states might be expected to be present in the energy range accessible to PILAC. Three aspects of the basic physics of these models underlie this statement³². First, the quark hf interaction leads, for most baryon-baryon channels, on accounting for quark AS'n, to a residual repulsive interaction. This shields the system from access to short distances where exotic states would be expected, and makes channels with hf repulsion unfavorable to the existence of dibaryons. By studying the structure of the hf interaction for MIT-bag-like states in which all 6 quarks are in the bag ground state, one readily finds

$$\langle - \sum_{i < j} (\lambda \sigma)_i \cdot (\lambda \sigma)_j \rangle = -24 + 2C_3^c + \frac{4}{3}J(J+1) + 4C_3^F \quad (1)$$

where J is the dibaryon spin and $C_3^{c,F}$ are the quadratic color and flavor Casimirs. The expression (1) allows one to identify those few channels which are non-repulsive. The second point is that the one-body energies in the models are lowered in the 6q, as compared to the two-isolated-3q-clusters configuration, so that such a hf-non-repulsive channel then naturally tends to support a bound state. Finally, because the OGE-induced quark tensor force is weak, true dibaryons which can decay to lower-lying two baryon states only through the quark tensor interaction are likely to be narrow, so that the relevant threshold to

consider for the existence of narrow positive parity states is that consisting of two isolated baryons in relative s-waves. Bearing these points in mind one finds that, eg., the MIT bag model predicts three positive parity dibaryons multiplets having (spin, flavor) = (0,1), (2,8) and (3,10*). The first contains the H and is not accessible to PILAC unless the H is sufficiently bound that the channel $\Lambda N \rightarrow HK$ is open, in which case it will be seen as a branch for ordinary hypernuclear decay. The (3,10*) contains a non-strange state, the d^* , which the MIT bag would put at $\sim 2350 MeV$ ³³, the cluster model at $\sim 2200 MeV$ ³⁴, and the LAMP at $2150 \pm 40 MeV$ ³⁵, all of which lie well within the PILAC range for $\pi d \rightarrow \pi d^*$. The state may be narrow but has cross-sections of $< 1 \mu b$ ³⁵, so that its detection appears difficult, certainly without a spin-parity analysis. The strangeness -1 member of this multiplet, at least in the bag, is well out of reach. The (2,8) multiplet, in contrast, has no non-strange member. The $S = -1$ member is an isodoublet which would lie at $\sim 2225 MeV$ if one adjusted the location relative to $N\Sigma^*$ in the flavor symmetric version of the bag to the physical threshold³³, but is adversely affected by flavor symmetry breaking and, therefore, predicted to lie at $\sim 2295 MeV$ when this effect is taken into account³⁶. Flavor symmetry breaking would, therefore, take it outside the PILAC range. This leaves, as potential dibaryon candidates for PILAC the rotationally stabilized p-wave states D_s, D_t , having quark cluster structure $Q^4 \otimes Q^2$, the quartet cluster being flavor 3 and the doublet flavor 3*. These states have been discussed in detail by Aerts and Dover^{37,38}. They lie at ~ 2110 and $2150 MeV$, respectively, in the bag, and are expected to be narrow. One might look for the D_t in $d(\pi^+, K^+)D_t$ with $D_t \rightarrow N\Lambda, N\Sigma$, but this suffers from a large $N(\pi^+, K^+)\Lambda, \Sigma$ quasielastic background (see eg. Ref. 38 for the (K^-, π^-) version of the production mechanism). D_s production would require, in the absence of significant spin-flip amplitude, a spin singlet NN target, hence a three-body state like 3He . One would then want to use the (π^-, K^+) reaction to avoid the quasielastic $n(\pi^+, K^+)\Lambda$ background (D_s is predicted to lie below $N\Sigma$ threshold). There remains, however, the problem of $p(\pi^-, K^+)\Sigma^-$ followed by $\Sigma^- p \rightarrow n\Lambda$ which could contaminate the D_s signal. More details on the line reversed reactions may

be found in Ref. 38. An optimistic conclusion would, therefore, be that the d^* , D_s , and D_s^* may lie in the range accessible to PILAC, but none of the obvious detection experiments will be easy. If one allows for possible uncertainties in the bag model predictions the $S = -1$, $I = 1/2$ member of the $(2,8)$ multiplet might also be barely accessible. One very severe caveat, however, having to do with the reliability of the theoretical predictions, is in order before leaving this section. This caveat comes in two parts, one (potentially) much deadlier than the other. First, therefore, the less traumatic of the two. Despite standard claims, the dibaryon binding, in all cases of interest is attributable not to the hf interaction, for which there is considerable phenomenological evidence in the spectrum, but to the one-body (in which we include confinement) energy. This is because the discrete expectation of (1) is accompanied by a spatial matrix element which scales down in going from the $3q$ to $6q$ sector, at least for the bag, because of the increase of the bag radius. Thus, even for the H, the hyperfine binding relative to $\Lambda\Lambda$ is only $\sim 5MeV$, and the bulk of the binding is due to the shift in the one-body energies. This shift, moreover, comes from a very large cancellation between the decrease of the quark eigenenergies in the larger bag ($\sim 450MeV$) and the increase of the phenomenological $-Z/R$ term ($\sim 400MeV$). The hf interaction, being non-repulsive, merely facilitates access to this “delocalization” effect. A similar statement holds for the cluster model, in the one case (the d^* channel) for which the comparison is currently possible (cluster model calculations claiming only small binding or no binding in the H channel are **not** included here because they are incomplete: they neglect the short distance “hidden color” components accessible through the non-repulsive $N\Xi$ and $\Sigma\Sigma$ channels; adding such components in the d^* case, for which the $\Delta\Delta$ channel is non-repulsive, is known to increase the binding from $\sim 3MeV$ to $\sim 260MeV$ ³⁴). The existence or non-existence of the dibaryon states discussed above, therefore, tests primarily the multi-quark confinement physics and not, as usually claimed, the more reliably known hf physics. While this might give one pause, especially considering the phenomenological nature of the $-Z/R$ term in the bag model, any information on modelling the multiquark

confinement sector is, of course, of considerable interest. Unfortunately, a potentially much more serious theoretical uncertainty remains. We remind the reader that the vacuum energy density (B in the bag model) is related to the vacuum gluon condensate through the QCD trace anomaly³⁹

$$E = \frac{-b}{32\pi} \langle 0 | \alpha_s G_{\mu\nu}^a G^{a\mu\nu} | 0 \rangle \quad (2)$$

where $\beta(g) = -bg^2/16\pi^2 + \dots$. It turns out that the value of B so obtained (via the QCD sum rules extraction of the vacuum gluon condensate) is at least an order of magnitude greater than that associated with the phenomenological bag model fit $B_{MIT}^{1/4} = 145.5 \text{ MeV}$.⁴⁰ One is, therefore, lead to consider the stability of the dibaryon predictions with respect to changes in B_{MIT} . The astonishing thing is that these predictions are **extremely** sensitive to B : dibaryons such as the d^* , the H and the omegon (the $S = -3$ member of the $(2,8)$ multiplet) cease to exist for B values very little above the MIT fit. This is illustrated in Table II, where, for the d^* and omegon, the quoted binding is relative to the relevant s-wave threshold states $\Delta\Delta$ and $N\Omega$.

Table II Dibaryon binding (in MeV) relative to relevant thresholds

Dibaryon	$B^{1/4} = 145.5 \text{ MeV}$	$B^{1/4} = 160 \text{ MeV}$	$B^{1/4} = 170 \text{ MeV}$
H	82	-26	-43
d^*	119	-112	-131
omegon	79	-23	-36

The source of this sensitivity is, of course, the high degree of cancellation in the one-body contributions to binding discussed above and the fact that the value of Z obtained by first fitting to the baryon spectrum scales differently with B than does the ratio of 6q to 3q bag radii. Since the dibaryons of interest all vanish long before the sum rules value of B

is reached (by which point, in any case, the bag is in severe trouble owing to the collapse of the meson spectrum) it is not clear what degree of confidence, if any, to place in the predictions. (The $D_{s,1}$ states will also be moved up as B is increased since they are built by rotational excitation on the s-wave band-head state for which the above effect will also occur.) One may, of course, still call for the relevant experiments, since a positive outcome would serve to dispel the above caveats, but one must bear in mind that a negative outcome may, because of these problems, effectively teach us nothing.

NON-MESONIC HYPERNUCLEAR DECAY AND THE $\Delta I = 1/2$ RULE

It has long been recognized that, owing to its large ($\sim 420 MeV/c$) momentum transfer, the non-leptonic weak decay process $\Lambda N \rightarrow NN$ probes short distances and might, therefore, expose effects associated with the explicit quark-gluon substructure of baryons, if such exist. This observation is especially intriguing in view of the fact that the meson exchange picture does not appear to provide a good description of the existing experimental data. Unfortunately, only one quark model calculation (in the hybrid quark hadron model) exists, the results of which appear open to question. We discuss each of these points, briefly, below, and, in addition, comment on the status of the $\Delta I = 1/2$ rule in light of the apparent recent understanding of its origin in $K \rightarrow \pi\pi$.

First, the experimental situation. The results quoted are based on an analysis of data for ${}^4_\Lambda He$, ${}^4_\Lambda H$ and ${}^6_\Lambda He$ in which it is assumed that Λ conversion on different nucleons is incoherent and that final state interactions (eg. nucleon charge exchange) are negligible (so that the experimental data reflect the spin-isospin dependence of the basic conversion process). This analysis is originally due to Dalitz and collaborators⁴¹ and is reviewed by Dover and Walker⁴² (see also Ref. 43 for the experimental references and an up-to-date version of the analysis). One finds that 1) 3S_1 dominates 1S_0 conversion, 2) the final state is dominantly $I_f = 1$, 3) the neutron- to proton-stimulated decay ratio $\Gamma_n/\Gamma_p = \Gamma(n\Lambda \rightarrow nn)/\Gamma(p\Lambda \rightarrow np)$

is $\sim 1 - 2$ (pure $I_f = 1$ and the $\Delta I = 1/2$ rule would imply $\Gamma_n/\Gamma_p = 2$), 4) the 1S_0 neutron conversion (evidenced by $n\Lambda \rightarrow nn$ in $^4_\Lambda\text{He}$, $\gamma_{np}^4 = \Gamma_n(^4_\Lambda\text{He})/\Gamma_p(^4_\Lambda\text{He}) \simeq 1/2$), though smaller than the 3S_1 conversion, is non-zero, and 5) $^3S_1 \rightarrow ^3P_1$ conversion $> ^1S_0 \rightarrow ^1S_0, ^3P_0 > \dots$, so that parity violating dominate parity conserving transitions.

For the predictions of the meson exchange picture we quote the results of Dubach *et al.*⁴⁴, based on $\pi, \rho, \eta, \omega, K$ and K^* exchange, the required weak vertices being related via PCAC and $SU(6)_w$ to those for $K \rightarrow \pi$ and $\Lambda \rightarrow p\pi^-, \Sigma^+ \rightarrow p\pi^0$. They find 1) $\Gamma(^3S_1 \rightarrow \text{all})/\Gamma(^1S_0 \rightarrow \text{all}) \sim 60$, 2) $\Gamma(I_f = 0)/\Gamma(I_f = 1) \simeq 1.6$, 3) $\Gamma_n/\Gamma_p \simeq .42$, 4) **very small** 1S_0 conversion ($\gamma_{np}^4 \simeq 0$) and 5) $\Gamma(^3S_1 \rightarrow ^3D_1) \simeq \Gamma(^3S_1 \rightarrow ^3P_1) > \Gamma(^3S_1 \rightarrow ^3S_1) > \dots$ and $\Gamma_{p.c.}/\Gamma_{p.v.} \simeq 1.1$. The agreement with experiment is clearly not good. We note, however, that retaining only the long-range one-pion-exchange mechanism leads to even worse results (strong dominance by the parity conserving $^3S_1 \rightarrow ^3D_1$ transition and hence $\Gamma_n/\Gamma_p \sim 0$) so that the inclusion of heavier meson exchanges has improved the situation. This improvement comes largely from K exchange, which reduces the $^3S_1 \rightarrow ^3D_1$ and enhances the $^3S_1 \rightarrow ^3P_1$ rate.

The only quark model calculation of which I am aware does little to improve the situation⁴⁵. It is based on the hybrid quark hadron model in which the one-pion-exchange-induced transition provides the only contribution for separations $r \geq r_0 \sim 1fm$, while at shorter distances, where a 6q bag structure is assumed, the weak transition is taken to be mediated by perturbative W exchange between quarks. QCD box graph corrections are included here, formally at least, but not those associated with penguin graphs; indeed, in the preferred fit of Ref. 45, the coefficients of the non-penguin operators are adjusted heuristically to reproduce the full phenomenological $\Delta I = 1/2$ rule. We will comment more on this procedure below. The result of the calculation is that the rate is determined slightly more by the exterior than the interior region. While the overall non-mesonic rate turns out to be reasonable, this means that the results will suffer from many of the same deficiencies as the "pion-only" meson exchange results mentioned above, i.e. $\Gamma_{p.c.}/\Gamma_{p.v.} > 1$ (though this result is not quoted in Ref. 45), Γ_n/Γ_p small (using simple quark counting rules, which,

however, ignore possible spin dependences, for the interior region, the combined result, not contained in the published papers, but quoted in Ref. 43, is .37), $\Gamma(I_f = 0)/\Gamma(I_f = 1) > 1$ (again this number is not quoted). There are, however, a number of points to be made about this calculation. Perhaps the most worrisome is the observation that the rate associated with the interior region is due entirely to the $\Delta I = 3/2$ combination, $c_1 + c_2$, of coefficients in the effective weak Hamiltonian. There is (to me at least) no obvious reason why the $\Delta I = 1/2$ amplitude should vanish and, as it is intrinsically much larger than the $\Delta I = 3/2$ amplitude, such a contribution could radically alter the quoted results. This point needs clarification. A second point (ignoring, for the moment, the potential problem above) concerns the method used for enforcing the $\Delta I = 1/2$ rule. It is known that the penguin graphs produce non-trivial contributions to the $\Delta I = 1/2$ rule, yet these are not included in the calculation of Ref. 45. Rather, their role, and that of any possible long distance enhancements, is incorporated into the two structures arising from the bare W-exchange and W-gluon-box graphs. Even ignoring the long-distance enhancements, this is a dangerous procedure since the penguins contain left-right pieces, unlike the box- and W-exchange graphs, which are pure left-left, and hence will have different spin properties. This, for example, is why the penguins are potentially significant in the $K \rightarrow \pi$ transition, the left-left operators being helicity suppressed, while the left-right ones are not. Finally, we point out that it has now been argued (rather convincingly, to my mind) that a factor of ~ 4 in the relative enhancement of the $\Delta I = 1/2$ to $3/2$ amplitude in $K \rightarrow \pi\pi$ is due to long distance, final state interaction effects⁴⁶⁻⁴⁸. This means that enforcing the observed $\Delta I = 1/2$ rule via the coefficients of the short distance operators is an incorrect procedure, even if one includes there the effect of the penguins. Ref. 45 also quotes results based on the unmodified Gilman-Wise⁴⁹ short distance enhancements. In this case the 6q short distance contributions dominate the long distance pion-exchange ones. (This, incidentally, shows that the appearance of the factor $c_1 + c_2$ is not just a typographical error since, as the scale for the QCD corrections is increased, it is the $\Delta I = 3/2$ strength that increases, not

the $\Delta I = 1/2$ strength.) This would presumably, if the calculation is correct, ameliorate the “pion-exchange-only” problems, but leads to a non-mesonic rate which appears too large. These results are all, of course, still subject to the queries raised above. In conclusion, the calculation of Ref. 45 needs to be re-checked. A cluster model calculation would also be of interest since the antisymmetrized configuration of two baryons has, in general, rather different spin-isospin structure from that of a ground state 6q bag and hence might produce rather different spin-isospin selectivity. The possibility of the exchange enhancement of only certain of the quark level weak operators (for example, any operator having a symmetric spin-color structure like that responsible for generating the induced residual hf repulsion in NN) is also an intriguing one. Unfortunately, there is little more to be said, at present, regarding predictions based on quark level physics.

We close this section with a comment on the status of the $\Delta I = 1/2$ rule in $\Lambda N \rightarrow NN$. Here the quark model is able to make a clear prediction. As mentioned above, there is good reason to believe that the short distance $\Delta I = 1/2$ to $\Delta I = 3/2$ ratio is $\sim 4 - 5$ and not the value ~ 20 observed empirically, eg. in $K \rightarrow \pi\pi$. If this is the case, one might hope to see non-trivial $\Delta I = 3/2$ components in $\Lambda N \rightarrow NN$. One finds, however, that the leading dimension $\Delta I = 3/2$ operator appearing in the effective weak Hamiltonian is symmetric in its color labels. (See Ref. 50 for a discussion of various aspects of the quark model and non-leptonic weak interactions, including this one.) Its expectation, therefore, automatically vanishes if the two active quarks come from a single baryon. Moreover, in a two baryon state, there is twice the amplitude to find a pair from different baryons in an antisymmetric, as opposed to symmetric, color state, even before accounting for the decreased weight of such terms resulting from the overlap of the different baryon wavefunctions in the nuclear environment. Thus one sees that the quark picture predicts the $\Delta I = 1/2$ rule to remain valid for $\Lambda N \rightarrow NN$. Unfortunately, so does the meson exchange model of Ref. 44, so that, in this respect, there is no distinction between the two pictures. Comparisons of predictions for other observables must await the results of future quark model calculations. An improvement

in the experimental situation would, of course, provide considerable encouragement for theorists. PILAC would certainly be useful in this regard. Moreover, since the hypernuclear experimental program would be pursued in any case, one need call for no new experiments in order to obtain useful constraints. Whether these constraints will tell us anything about the explicit quark substructure of baryons, however, is unknown at this point.

CONCLUSIONS

There are, of course, many interesting facets of (π, K) physics accessible to a $1.05\text{GeV}/c$ pion beam. Most of them, however, appear unlikely to shed much light on the question of the role of explicit quark-gluon degrees of freedom in few-hadron systems. The most promising area, in this regard, is non-leptonic weak hyperon decay, though considerable theoretical effect is still required, and one is not necessarily guaranteed that the results will fulfill this promise. Fortunately, the experimental program is of considerable interest in its own right and will, therefore, provide useful constraints for theorists, regardless. Less promising are dibaryon searches. While, in QCD-inspired models, there are dibaryons predicted to be accessible to the PILAC energy range, the theoretical uncertainties associated with the models are very large, and the experiments required non-trivial. Finally, though any data on low energy YN scattering will be gratefully received by one and all, such data, from PILAC, will apparently be unable to tell us anything about the existence of possible non-negligible quark-gluon effects.

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