NARROW DIBARYONS OF STRANGENESS $S = -1$?

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In a simple model consisting of wave function rearrangement followed by a final state interaction, we estimate the widths for the decay of $S = -1$ six-quark dibaryons into hyperon+nucleon (YN). For the lowest-lying dibaryons $D_{s,t}$, in the bag model of quark cluster structure $Q^4 \otimes Q^2$ (1P_1 or $^3P_{0,1,2}$), we predict very small widths ($\Gamma \lesssim 10$ MeV for $m_D \lesssim 2.2$ GeV).

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The subject of dibaryon resonances has been of prime interest for the last few years^{1),2)}. A central part of the problem consists in exhibiting the influence of confined six-quark bag states³⁾⁻⁵⁾ on the observable baryon-baryon scattering amplitudes. Much effort has been devoted to the nucleon-nucleon (NN) problem, since many precise experimental data are available. Here, one has difficulty in isolating unambiguously the effect of the "primitives" (the six-quark states) in the presence of very strong coupled channel effects ($NN \leftrightarrow N\Delta$) of conventional nature, which can also give rise to apparent "resonance" structures⁶⁾. For the NN system, the "primitives" of the bag model [poles of the P matrix⁷⁾] lie well above the $NN\pi$ threshold. The poles of the S matrix are considerably shifted in real part with respect to those of the P matrix, and they generally acquire a sizeable imaginary part of order $\Gamma \approx 100$ MeV due to the very strong channel coupling. Some possible exceptions to this are discussed by Mulders⁴⁾.

Narrow dibaryon structures are intrinsically more likely in strangeness $S \neq 0$ systems, for several reasons. The colour-magnetic forces of QCD operate with maximum attraction for the case of an SU(3) flavour singlet configuration; this requires the quark combination uuddss of $S = -2$. The lowest-lying $S = -2$ six-quark state, called the H, could even be stable⁸⁾ with respect to strong decay. Methods for producing the H in the (K^-, K^+) reaction have been discussed⁹⁾. For $S = -1$ or 0 a singlet $\{1\}$ of flavour cannot occur. However, a six-quark $S = -1$ system permits a lower dimensional flavour representation than for $S = 0$; hence the $S = -1$ configurations lie closer than the $S = 0$ states to the relevant strong decay thresholds (ΔN at 2055 MeV and ΣN at 2130 MeV). For the lowest-lying $S = -1$ states considered here (see below), the decay channel $YN\pi$ is closed and there is no strongly coupled (virtual) Y^*N channel analogous to ΔN for $S = 0$. Thus we may anticipate $S = -1$ six-quark states D for which the S and P matrix poles lie close together, and the strong decay width $\Gamma_{D \rightarrow YN}$ is relatively small. In this letter, we obtain estimates of Γ for the two $S = -1$ dibaryons which have "stable" quark clusters, and comment on the prospects for finding them experimentally in the (K^-, π) reaction.

We consider $S = -1$ six-quark "primitives" with cluster structure $D = Q^4 \otimes Q^2$, where the two- and four-quark clusters have a relative orbital angular momentum L . We focus our attention on the two lowest-lying configurations with $L = 1$ obtained in Ref. 3), namely

$$D_s = \left\{ \left[\underline{3} \left(\frac{1}{3}, \frac{1}{2} \right) 0 \right] \otimes \left[\underline{3}^* \left(\frac{2}{3}, 0 \right) 0 \right] \right\}_{^1P_1, I = \frac{1}{2}} \quad (1)$$

$$D_t = \left\{ \left[\underline{3} \left(\frac{1}{3}, \frac{1}{2} \right) 1 \right] \otimes \left[\underline{3}^* \left(\frac{2}{3}, 0 \right) 0 \right] \right\}_{^3P_J, I = \frac{1}{2}}$$

where the notation $\underline{f}(y, i)_s$ refers to the flavour $\underline{f}$, hypercharge y , isospin i and spin s of each cluster. The labels s and t on D specify the spin singlet or triplet nature of the dibaryon. The quantum numbers of the YN channel, into which D decays, are also indicated. Note that the strange quark resides in the Q^4 cluster. In Ref. 3), the masses of D_s and D_t are estimated to be about 2110 and 2150 MeV, respectively. For D_s , only the ΛN decay channel would then be energetically allowed; for both D_s and D_t , the $\Lambda N \pi$ channel is closed. Many other (Q^4 - Q^2) "primitives" are given in Ref. 3), these (generally $L > 1$) all lie above the $\Lambda N \pi$ threshold, and are expected to be broader. For example, the lowest $I = 3/2$, Q^4 - Q^2 configuration is predicted³⁾ at 2340, and contains resonance admixtures and mixtures (Y^*N and $Y\Delta$) in its wave function, leading directly to the $YN \pi$ decay channel. Strange dibaryon states of other cluster structures also exist³⁾; the lowest of these (of Q^6 structure), around $m_D = 2170$ MeV, couples to the 3S_1 , $I = \frac{1}{2}$ channel of YN . However, this configuration, as well as others of type $Q^3 \otimes Q^3$ and $Q^5 \otimes Q$, are unstable with respect to "quark tunnelling", as we discuss later.

The quarks in each cluster ($\underline{r}_1$ through $\underline{r}_4$ in Q^4 , $\underline{r}_5$ and $\underline{r}_6$ in Q^2) are taken to be in relative S states, with $L = 1$ relative motion between the clusters. We assume a dibaryon space wave function of oscillator form

$$\psi_{\underline{k}}(\underline{r}_1, \dots, \underline{r}_6) = N e^{i \underline{k} \cdot \underline{R}} \frac{r}{R_P} Y_{lm}(\hat{r}) * e^{-\frac{1}{2}(\rho^2 + \lambda^2 + \sigma^2 + \xi^2)/R_D^2} e^{-r^2/2R_P^2} \quad (2)$$

where N is a normalization factor, $\underline{k}$ is the total momentum of D , and the relative co-ordinates are taken to be

$$\begin{aligned}
 \tilde{R} &= (\tilde{r}_1 + \tilde{r}_2 + \tilde{r}_3 + \tilde{r}_4 + \tilde{r}_5 + \tilde{r}_6) / 6 \\
 \tilde{L} &= (\tilde{r}_1 - \tilde{r}_2) / \sqrt{2} \\
 \tilde{\lambda} &= (\tilde{r}_1 + \tilde{r}_2 - 2\tilde{r}_3) / \sqrt{6} \\
 \tilde{\sigma} &= (\tilde{r}_1 + \tilde{r}_2 + \tilde{r}_3 - 3\tilde{r}_4) / \sqrt{12} \\
 \tilde{\xi} &= (\tilde{r}_5 - \tilde{r}_6) / \sqrt{2} \\
 \tilde{\tau} &= (\tilde{r}_1 + \tilde{r}_2 + \tilde{r}_3 + \tilde{r}_4 - 2(\tilde{r}_5 + \tilde{r}_6)) / \sqrt{12}
 \end{aligned} \tag{3}$$

A similar oscillator wave function, with radius parameter R_B , is used for a Q^3 cluster. In momentum space, the overlap of D with a state of two Q^3 baryons [$\{56\}$ -plet of flavour-spin $SU(6)$] Y and N of momenta $\underline{k}_Y$ and $\underline{k}_N$ in a relative P state is given by:

$$\begin{aligned}
 \langle \underline{k}_Y, \underline{k}_N | \psi_{\tilde{N}} \rangle &= (2E_D \cdot 2E_Y \cdot 2E_N)^{1/2} \times \\
 &\times (2\pi)^3 \delta(\underline{k} - \underline{k}_Y - \underline{k}_N) C(YN, D) \Gamma(\underline{k}_Y - \underline{k}_N)
 \end{aligned} \tag{4}$$

with

$$\begin{aligned}
 \Gamma(\underline{k}) &= k N_P Y_{lm}(\hat{k}) e^{-\frac{1}{12} \bar{R}^2 k^2} \\
 \bar{R}^2 &= (R_B^2 R_P^2 + R_B^2 R_D^2 + 2 R_D^2 R_P^2) / (R_D^2 + R_P^2 + 2 R_B^2)
 \end{aligned}$$

where $C(YN, D)$ is the space-spin-colour-flavour recoupling coefficient. We have chosen the bag radii $R_B = 0.83$ fm for the hyperon and nucleon to ensure the correct rms radius, [see Ref.9)] and taken $R_P = 1.02$ fm and $R_D = 0.83$ fm in Eq. (2) so that the Q^4 and Q^2 clusters are non-overlapping, i.e.,

$$\langle r^2 \rangle_D^{1/2} \gg \langle r^2 \rangle_{Q^2}^{1/2} + \langle r^2 \rangle_{Q^4}^{1/2}$$

The value $\bar{R} = 0.92$ fm then appears in $\Gamma(k)$.

The colour-flavour spin wave functions for D_s and D_t are totally antisymmetric for each cluster. Standard recoupling techniques can be applied to write a $Q^4 \otimes Q^2$ structure in a basis of $Q^3 \otimes Q^3$ states. We find

$$\begin{aligned} |D_s\rangle = \sum_{i,j} c(ij, D_s) |ij\rangle = \frac{1}{16} \{ \sqrt{3} (p\Lambda - \Lambda p) - \sqrt{2} (u\Sigma^+ - \bar{u}\Sigma^+) \\ + (p\Sigma^0 - \Sigma^0 p) \} + \dots \end{aligned} \quad (5)$$

$$\begin{aligned} |D_t\rangle = \sum_{i,j} c(ij, D_t) |ij\rangle = \frac{1}{16\sqrt{3}} \{ (p\Lambda + \Lambda p) + \sqrt{6} (u\Sigma^+ + \bar{u}\Sigma^+) \\ - \sqrt{3} (p\Sigma^0 + \Sigma^0 p) \} + \dots \end{aligned}$$

where we have isolated the channels containing baryons in the $\{56\}$ -plet of $SU(6)$, i.e., the ones that we assume are energetically allowed for D-decay. The dots in Eq. (5) indicate all other couplings, involving baryons in the $\{70\}$ -plet and also colour $\{8\} \otimes \{8\}$ configurations of $Q^3 \otimes Q^3$. The wave functions (5) include a colour recoupling factor of $1/\sqrt{3}$, spin and space recoupling factors of $1/\sqrt{2}$, as well as flavour recouplings. Note that the YN channels in Eq. (5) correspond to only a relatively small part of the "dissociation probability" of $D_{s,t}$ into three-quark clusters. For large radial separations r between the Q^3 clusters, colour confinement and energy conservation (assuming m_D lies below the Y^*N and YN^* thresholds) require that the virtual components [the dots in Eq. (5)] vanish. The decomposition (5) which is appropriate for small r would be multiplied by a normalization factor $N_s = 8/\sqrt{3}$ for D_s and $N_t = 8\sqrt{3}/5$ for D_t . In principle, a calculation of the widths $\Gamma_{D \rightarrow YN}$ involves an integration of the D wave function (times other factors) over r . In the absence of a dynamical calculation of the D, Y and N wave functions beyond the simple geometrical overlap of Eq. (4), we assume that such integrals are dominated by the small r region where Eq. (5) prevails. One would include the factors $N_{s,t}$ in estimating upper limits on

the widths. This still gives rather small widths as we see later. One general point to note is that the ΛN channel, which has the largest phase space, represents only a small piece of the $D_{s,t}$ wave function. This is a dominant factor in obtaining relatively long-lived dibaryons.

The widths $\Gamma_{D \rightarrow YN}$ are calculated according to the mechanism of Fig. 1; here the vertex Γ describes the wave function overlap Eq. (4) between D and the channels $Y'N$, and the transition matrix T represents a final state interaction which puts the YN system on-shell. In principle, the virtual intermediate state $Y'N$ could contain all the components of the D wave function. We have suppressed intermediate configurations other than ΛN or ΣN , arguing that T is small for off-shell transitions of the type $Y^*N \rightarrow YN$ or $8_c \otimes 8_c \rightarrow YN$. In the absence of strong coherences among these off-shell amplitudes, the effect of wave function components other than YN on the width is expected to be bounded by the factors $N_{s,t}$ defined above.

Two different models have been adopted: a) on-shell scattering amplitudes, constructed from the one-boson exchange (OBE) models of the Nijmegen group¹⁰⁾ are used for T ; these reproduce the available ΛN and ΣN scattering data, and b) Born approximation, in which T is replaced by a transition potential due to meson exchange. The width Γ (in the rest frame of the D), corresponding to Fig. 1 can be cast in the form

$$\Gamma_{D \rightarrow YN} = \frac{1}{m_D} \int d^3 p_Y |\tilde{T}|^2 \delta(m_D - E_Y - E_N) \quad (6)$$

where m_D is the mass of the D and

$$\begin{aligned} \tilde{T} = \lambda \sum_{Y'N'} C(Y'N'; D) \int \frac{d^3 p}{(2\pi)^3} \frac{P_{Y'N'}(p)}{(2E_{Y'} 2E_N)^{1/2}} e^{-\frac{1}{3} p^2 R^2} * \\ * T_{Y'N' \rightarrow YN}(\underline{p}, -\underline{p}; \underline{p}_Y, -\underline{p}_Y) \end{aligned} \quad (7)$$

In Eq. (7), λ is a factor independent of m_D and $C(Y'N'; D)$ is the recoupling coefficient from Eq. (5). We now expand $T_{Y'N' \rightarrow YN}$ in terms of partial wave amplitudes $T^{LL'SJ} = \langle L'SJ | T | LSJ \rangle$; only the amplitudes corresponding to the initial $\{LSJ\}$ of the dibaryon contribute. In the calculations employing a

meson exchange potential V in place of T , we have instead used the helicity representation, in order to have an independent check on the normalizations, etc. All spin dependent parts of the interaction are included.

After performing the momentum integrals, and inserting the bag radii given above, we find for a dibaryon with quantum numbers $\{L=1, S, J\}$

$$\Gamma_{D \rightarrow \Upsilon N} \approx (270 \text{ MeV}) \sum_{L'=1,3} \left| \sum_{\Upsilon N'} C(\Upsilon N', D) T_{\Upsilon N' \rightarrow \Upsilon N}^{1L'SJ} \right|^2 \quad (8)$$

where we have used typical values $m_D \approx 2150 \text{ MeV}$, $E_{Y_1} \approx E_N \approx m_D/2$. An upper limit for the total width Γ_D results if one invokes the unitarity limit for the T matrix elements

$$\Gamma_D = \sum_{\Upsilon N} \Gamma_{D \rightarrow \Upsilon N} \leq (270 \text{ MeV}) \sum_{\Upsilon N} |C(\Upsilon N, D)|^2 \quad (9)$$

where the ΥN sum includes the open channels only.

The result Eq. (9) is instructive since it displays Γ_D as a typical strong decay width of order 200 MeV multiplied by the probability that the D wave function couples to the energetically available channels. For the strange dibaryons of Eq. (5), $\sum_{\Upsilon N} |C(\Upsilon N, D)|^2$ is $3/64$ for D_s and $5/192$ for D_t (assuming we are above the ΣN threshold) and thus Γ_D is expected to be 10 MeV or less.

A sample of our numerical results for Γ_D as a function of m_D is shown in Fig. 2. The widths of D_s and D_t remain small ($\Gamma_D < 10 \text{ MeV}$) for the mass range of interest ($m_D < 2.2 \text{ GeV}$). Note that Γ_D is considerably smaller than the unitarity limit of Eq. (9), if on-shell T matrices are used, due to the fact that the ΥN amplitudes in the P wave are rather weak. If one uses $T \approx V$, that is the Born approximation, without a hard core, the calculated widths are larger, but still remain small. Note that this approximation can in some cases violate the unitarity limit. The effect of the hard core is to cancel off some of the attraction coming from the long and medium range Yukawa exchanges, and thus reduce the effective P wave scattering volumes. We have also done some calculations using Hulthén rather than oscillator wave functions; the widths then increase by about a factor of two.

Our results are sensitive to the model chosen for V . For instance, if $m_D = 2.14$ GeV we find widths $\Gamma(^3P_0) = \{5.2, 1.9\}$, $\Gamma(^3P_1) = \{0.1, 1.3\}$, $\Gamma(^3P_2) = \{0.4, 0.2\}$ and $\Gamma(^1P_1) = \{2.0, 0.4\}$ MeV in Born approximation for models D and F of Ref. 10). However, the widths remain small in all cases we have considered. The preference for $D_{s,t}$ to decay into the ΣN rather than the ΛN channel is clear from the rapid growth of Γ seen in Fig. 2 at the ΣN thresholds. We also found the effect of tensor forces to be important.

Note that the dibaryon D_t^J corresponds to three states with $J = 0, 1, 2$. These will be split in mass by small parts of the QCD residual interaction which we have neglected here. Since we find very small widths for each J , this splitting of $J = 0, 1, 2$ states is likely to exceed Γ .

It is useful to compare our results for $S = -1$ dibaryons with the corresponding expectations for the $S = 0$ case. The key notion³⁾ is that of "cluster stability" against emission of a quark, and stability of the six-quark system against dissociation into two colour singlet three-quark clusters. The colour-magnetic interaction energy of a cluster³⁾ of N quarks is proportional to

$$\Delta = -\frac{1}{4}N(10-N) + \frac{1}{3}S(S+1) + F_f^2 + \frac{1}{2}F_c^2 \quad (10)$$

where S is the cluster spin and $F_{f,c}^2$ are the quadratic Casimir operator eigenvalues for $SU(3)$ flavour and colour. The most stable clusters have the largest negative values of Δ ; for $(Q^n)_c$ configurations (c =colour), these are the ones with the flavour representations of lowest dimension:

$$\begin{aligned}
 n=2, \quad c=3^*, \quad \underline{f} = \underline{3}^*, \quad S=0, \quad \text{with } \Delta = -2 \\
 n=3, \quad c=1, \quad \underline{f} = \underline{8}, \quad S=1/2, \quad \text{with } \Delta = -2 \\
 n=4, \quad c=3, \quad \underline{f} = \underline{3}, \quad S=0, \quad \text{with } \Delta = -4 \\
 n=4, \quad c=3, \quad \underline{f} = \underline{3}, \quad S=1, \quad \text{with } \Delta = -10/3
 \end{aligned}
 \tag{11}$$

From Eq. (11) we see that the Q^4 clusters in $D_{s,t}$ [see Eq. (1)] are stable with respect to $Q^4 \rightarrow Q^3+Q$ dissociation. D_s and D_t are then both also stable with respect to quark tunnelling³⁾, i.e., the process $(Q^4)_3+(Q^2)_3^* \rightarrow (Q^3)_1+(Q^3)_1$ corresponds to an increase in Δ . Thus for these configurations, it is sensible to write a wave function of the form Eq. (2), with two spatially localized clusters. For the non-strange sector, in contrast, the lowest-lying $L=1$ $(Q^4)_3-(Q^2)_3^*$ state has the structure $[6^*(4/3,0)1 \otimes 3^*(2/3,0)0]$, which is unstable with respect to dissociation into $(Q^3)_1+(Q^3)_1$. In addition, this object corresponds to $I=0$, $S=1$, $L=1$, which is a Pauli forbidden channel for NN and also not allowed for ΔN , because of isospin conservation. It would be expected to have a sizeable decay width into $NN\pi$. The calculation of this width is outside the scope of our model based on Fig. 1. Other relatively low-lying non-strange six-quark states, such as $(Q^6)_1$ in the $10^*(2,0)$, and $27(2,1)0$ configurations³⁾, are also unstable with respect to $(Q^3)_1+(Q^3)_1$ dissociation, and couple strongly to the $^{13}S_1$ and $^{31}S_0$ NN channels. We thus expect no narrow $Y=2$ states analogous to the $Y=1$ states of Eq. (1). Also, no stable $(Q^6)_1$ states with $Y=1$ are expected. Note that for $Y=0$, ($S=-2$), in addition to the H, of structure $(Q^6)_1[1(0,0)0]$, there are four other configurations built of stable clusters $[(Q^4)_3 \otimes (Q^2)_3^*]$ which are energetically disfavoured from dissociation into $(Q^3)_1 \otimes (Q^3)_1$. The two lowest-lying ones ($m_D \approx 2.30$ GeV) are

$$H' = \left\{ \left[\underline{3} \left(-\frac{2}{3}, 0 \right) 0 \right] \otimes \left[\underline{3}^* \left(\frac{2}{3}, 0 \right) 0 \right] \right\}_{1P_1, I=0} \quad (12)$$

$$H'' = \left\{ \left[\underline{3} \left(\frac{1}{3}, \frac{1}{2} \right) 0 \right] \otimes \left[\underline{3}^* \left(-\frac{1}{3}, \frac{1}{2} \right) 0 \right] \right\}_{1P_1, I=1,0}$$

In our model these states would also be narrow, particularly the $I = 1$ state, which can only decay into NE . As for the H , these could be searched for in the ${}^3\text{He}(K^-, K^+)$ reaction⁹⁾.

What are the prospects for finding the predicted narrow $S = -1$ dibaryons? Already in some older data¹¹⁾ on the $d(K^-, \pi^-)$ reaction, there is some hint of a "shoulder" in the Λp missing mass spectrum. This appears as an energy shift and broadening at larger θ of the sharp peak seen at small θ very near the ΣN threshold [the cusp effect¹²⁾]. Such a phenomenon is expected if D_t is formed, since this requires orbital angular momentum transfer $\Delta L = 1$; the cross-section for $D_{s,t}$ will peak at non-zero values of θ , where the cusp effect, corresponding to $\Delta L = 0$, has dropped off somewhat. The angular dependence of the cross-section for the reaction $d(K^-, \pi^-) D_t$ has been studied theoretically in Ref. 13); peak cross-sections of 1-2 $\mu\text{b/sr}$ are anticipated around $\cos\theta \approx 0.8-0.9$, a level of cross-section which could probably be separated¹⁴⁾ from the tail of the ΣN cusp effect. To form the spin singlet dibaryon D_s , the reaction ${}^3\text{He}(K^-, \pi^+) D_s$ is appropriate. Since the diproton in ${}^3\text{He}$ is automatically in the 1S_0 state, the formation of $D_s({}^1P_1)$ proceeds via $\Delta L = 1$, without spin flip. Cross-sections of a few $\mu\text{b/sr}$ (at peak) have also been estimated¹³⁾ for this reaction.

In summary, we have shown that the two lowest-lying $S = -1$ dibaryon states of quark cluster structure $[Q^4 \otimes Q^2]_{L=1}$ are expected to be very narrow. It is important to search for these objects experimentally, since they represent a significant prediction of the multiquark bag model.

R E F E R E N C E S

- 1) A. Yokosawa, Phys. Reports 64 (1980) 47;
N. Hoshizaki, Progr.Theor.Phys. 60 (1978) 1796 and 61 (1978) 129.
- 2) See the review talks by D. Bugg and K.K. Seth, in Proceedings of 10th International Conference on Few-Body Problems in Physics, Karlsruhe, RFA, August 1983 (ed. B. Zeitnitz).
- 3) P.J. Mulders, A.T. Aerts and J.J. de Swart, Phys.Rev. D21 (1980) 2653.
- 4) P.J. Mulders, Phys.Rev. D25 (1982) 1269 and D26 (1982) 3039.
- 5) I.L. Grach, I.M. Narodetskii and M.Zh. Shmatikov, Z.Phys. C21 (1984) 229.
- 6) W.M. Kloet, J.A. Tjon and R.R. Silbar, Phys.Letters 99B (1981) 80;
J.A. Tjon and E. van Faessen, Phys.Letters 120B (1983) 39;
E.L. Lomon, Phys.Rev. D26 (1982) 576.
- 7) R.L. Jaffe and F.E. Low, Phys.Rev. D19 (1979) 2105.
- 8) R.L. Jaffe, Phys.Rev.Letters 38 (1977) 195 and 617(E).
- 9) A.T.M. Aerts and C.B. Dover, Phys.Rev. D28 (1983) 450 and D29 (1984) 433.
- 10) M.M. Nagels, T.A. Rijken and J.J. de Swart, Phys.Rev. D15 (1977) 2547 (model D) and Phys.Rev. D20 (1979) 1633 (model F). We used Model F for the calculations presented. We would like to thank J.J. de Swart and T.A. Rijken for supplying the YN T matrix elements in numerical form.
- 11) D. Cline, R. Laumann and J. Mapp, Phys.Rev.Letters 20 (1968) 1452;
D. Eastwood et al., Phys.Rev. D3 (1971) 2603;
O. Braun et al., Nuclear Phys. B124 (1977) 45.
- 12) R.H. Dalitz and A. Deloff, Czech.J.of Phys. B32 (1982) 1021.
- 13) A.T.M. Aerts and C.B. Dover, in preparation.
- 14) H. Piekarz et al., Brookhaven AGS proposal 773, September, 1982; we would like to thank Dr. Piekarz for useful discussions on various aspects of this proposal.

FIGURE CAPTIONS

Figure 1: Mechanism for strange dibaryon decay into hyperon-nucleon channels ($D \rightarrow YN$). The vertex (wave function) Γ refers to the dissociation $Q^4 \otimes Q^2 \rightarrow Q^3 \otimes Q^3$ of D into two three-quark clusters and T is a final state interaction, taken either to be either a one-boson-exchange potential or an on-shell scattering amplitude.

Figure 2: Total width Γ_D of the strange dibaryons D_s (1P_1) and D_t (3P_J) as a function of their mass m_D . The curves labelled 1P_1 , $^3P_{0,1,2}$ refer to a width calculation which employs on-shell T matrices for the YN final state interactions, taken from Model F of Ref. 10). The curve labelled 3P_1 (OBE) corresponds to Γ_D calculated with the full one-boson exchange potential of Model F, but without hard core cut-offs. The effects of all channel couplings $Y'N \rightarrow YN$, with $Y', Y = \{\Lambda; \Sigma\}$, are included, as well as tensor potentials. The width Γ_D is summed over the final channels Λp , $\Sigma^0 p$ and $\Sigma^+ n$ which are energetically allowed; in the OBE calculation, all intermediate YN states are included, while in the T matrix model, only open channels occur as intermediate configurations. The widths Γ_D are calculated with the small r wave function of Eq. (5); if the large r wave function is used instead, we multiply the curves by $64/3$ or $192/5$ for D_s or D_t , respectively, to give a (pessimistic) upper limit for Γ_D .

Fig. 1

Fig. 2