S=-1 DIBARYON PRODUCTION WITH KAON BEAMS

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Production cross-sections for strangeness $S=-1$ six-quark dibaryon states $D_{s,t}$ in the $(K^-, \pi^\pm)$ reactions on deuterium and ${}^3\text{He}$ targets have been estimated. For the lowest-lying configurations (mass range 2100-2200 MeV in the MIT Bag model) of quark structure Q^4-Q^2 , namely 1P_1 and 3P_J , peak differential cross-sections of order $1 \mu\text{b/sr}$ are predicted. The angular distributions are peaked in the region of $20^\circ \leq \theta \leq 30^\circ$, characteristic of angular momentum transfer $\Delta L=1$. We evaluate the prospects for detecting a dibaryon signal in the presence of background processes such as $d(K^-, \pi^-)YN$ or ${}^3\text{He}(K^-, \pi^+)\Sigma^-pn$.

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1 INTRODUCTION

The phenomenology of bag models^[1,2] has been generally successful in reproducing the observed spectra of mesons (quark-antiquark $Q\bar{Q}$ states) and baryons (three quark Q^3 composites). Once the parameters of the model have been adjusted to reproduce the masses of known particles, one can predict the spectrum of all multi-quark states^[3-8] of type $Q^n\bar{Q}^m$.

The dibaryon spectrum ($n=6, m=0$) has received a great deal of attention, both theoretically^[9,10] and experimentally^[11-18]. The search for non-strange dibaryon resonances has been most intense, due to the copious amount of precise data on the nucleon-nucleon amplitudes over a wide range of energy. Strangeness $S=-1$ and $S=-2$ dibaryons have received less attention, although there have been several experiments^[14-17] which indicate structure in the Λp missing mass spectrum for the reaction $d(K^-, \pi^-)\Lambda p$. For instance, in the experiment of Braun et al.^[17], a narrow, symmetric peak was seen at $M_{\Lambda p} = 2129 \pm 0.4$ MeV, with a width of $\Gamma = 5.9 \pm 1.6$ MeV. This peak is most strongly excited at minimum values of the four-momentum transfer $|t|$ to the recoiling Λp system. For larger negative t values, a broad asymmetric shoulder is observed^[17], extending to a mass of 2180 MeV. The interpretation of these structures has been the object of several detailed theoretical studies^[19-22]. A marked cusp behaviour in the Λp mass spectrum is expected near the ΣN threshold, since intermediate $\pi \Sigma N$ channels play a dominant role in the $d(K^-, \pi^-)\Lambda p$ reaction. Toker et al.^[20] have shown how a shoulder in the Λp mass spectrum may also arise for K^- capture at rest because of the interference between direct Λ production and second and higher order processes involving an intermediate Σ .

In the present paper, we entertain the notion that some of the narrow structure in the Λp mass spectrum may be due to six-quark dibaryon resonances rather than multi-channel effects in the conventional picture. In Chap. 2, we recall the main features of the $S=-1$ dibaryon spectrum expected in the bag model. It is argued that narrow structures in the 3S_1 channel ($S=-1$) are not expected, and hence the narrowness of a cusp effect has nothing to do with a dibaryon resonance. However, the shoulder seen for larger $|t|$ could possibly be a dibaryon signature, since narrow^[23] six-quark states communicating with the 3P_J hyperon-nucleon channels may exist. Such long-lived strangeness $S=-1$ states would have a $Q^4 \otimes Q^2$ cluster structure, with a relative P-wave centrifugal barrier between the two

clusters. They could be produced in the reactions $d(K^-, \pi^-)D_t$, for a spin-triplet dibaryon D_t , or ${}^3\text{He}(K^-, \pi^+)nD_s$ for a spin-singlet dibaryon D_s . The production cross-section should have an angular shape characteristic of a $\Delta L=1$ transition (peak at finite angle) rather than $\Delta L=0$ as for the cusp phenomenon (forward peaked). This feature is the only hope for distinguishing the dibaryon signal from backgrounds ("noise").

The purpose of the present paper is to estimate cross-sections for $S=-1$ dibaryon production, in order to obtain some idea of the expected "signal to noise" ratio. Since all $S=-1$ dibaryon candidates lie above the ΛN (and perhaps also ΣN) thresholds, it is important to estimate the amount of "noise", i.e. (K^-, π) processes, such as $d(K^-, \pi^-)\Lambda p$, in which the target nucleus breaks up without dibaryon formation. We do these estimates with the same model for the $\bar{K}N \rightarrow \pi Y$ amplitudes used to obtain the dibaryon cross-sections.

The outline of the article is as follows. We first review briefly, in Chap. 2, the properties of the $S=-1$ six-quark spectrum predicted by bag models. Although there is a large number of expected structures, most of these correspond to highly unstable objects. By looking at the colour magnetic interactions (in the one gluon exchange approximation) within a quark cluster, one can determine which clusters are stable with respect to spontaneous quark emission ($Q^n \rightarrow Q + Q^{n-1}$) and which dibaryon configurations $Q^n \otimes Q^m$ ($n+m=6$) are stable against dissociation into two colour singlet Q^3 clusters. When we apply these two stability criteria, only two types of $S=-1$ dibaryons survive the test. These have $Q^4 \otimes Q^2$ cluster structure, in a spin singlet (D_s) or spin triplet (D_t) state.

We focus on the production mechanisms for the dibaryon D_s and D_t in Chap. 3. A two step mechanism is discussed, where the $\bar{K}N \rightarrow \pi Y$ reaction first produces a hyperon Y , which then fuses with another nucleon to form the dibaryon ($YN \rightarrow D_{s,t}$). The two reactions which we consider in detail are

$$\begin{aligned} d(K^-, \pi^-)D_t \\ {}^3\text{He}(K^-, \pi^+)nD_s \end{aligned} \tag{1.1}$$

One could also contemplate the use of heavier nuclear targets, but the theoretical analysis is complicated by the presence of nuclear excitations.

In Chaps. 4 and 5, we display our numerical results for the reactions (1.1). We plot angular distributions ($d\sigma/d\Omega$) and the excitation functions ($d^2\sigma/d\Omega dM$ for fixed θ), as a function of K^- lab momentum k_L and dibaryon mass m_D . Attention is paid to the problem of choosing the range of k_L and θ so as to maximize the ratio R of dibaryon to background cross-sections. A promising region $800 \leq k_L \leq 900$ MeV/c and $\theta \approx 20-30^\circ$ is found where R could attain values as large as 0.2. It is only in this rather restricted angular range (matched to $\Delta L=1$) that one could hope to resolve a dibaryon signal. Even here our estimated values for R are not large, so good energy resolution is important in an experiment. We also consider the effects of a cut-off on the momentum of the "spectator" proton in the reaction $d(K^-, \pi^-)\Lambda p$. By such a cut-off, one can diminish the size of the background, but the dibaryon signal is also altered (here the proton comes from $D \rightarrow \Lambda p$). A choice $k_{\min} = 100 - 150$ MeV/c improves the "signal to noise" ratio R somewhat at finite angle. Note that in an experiment which measures only the outgoing π^- , one cannot impose this cut-off (one needs to observe a $\pi^- + p$ coincidence).

We present our conclusions in Chap. 6. If such long-lived objects exist, it does seem feasible to search for them in the reactions (1.1). The dibaryon cross-sections are only of the order of a few $\mu\text{b/sr}$ at most, however, and the experiments will be difficult.

It is a fundamental question whether such multi-quark configurations really exist as particles, or whether they are some manifestation of the short range behaviour of the baryon-baryon wave function. If they exist, can the bag model predict their masses correctly^[24]? Topological soliton models (see Balachandran et al. in Ref.(4)) also imply some dibaryon states. Good experiments are required to decide whether such theoretical models have any predictive power.

2 THE $S=-1$ DIBARYON SPECTRUM

Theoretical predictions for the spectra of six-quark states have been obtained by a number of authors. Among the most extensive calculations are those of the Nijmegen group^[5,7,9,25,26], which are based on the MIT bag model. These authors have considered all colour-singlet arrangements of six-quarks into two clusters of m and n quarks ($m+n=6$) with $SU(3)$ colour representations c and c' , i.e., dibaryons D of structure

$$D = Q_c^m \otimes Q_c^n, \quad (2.1)$$

In addition to colour, each cluster carries labels $f(y,i)s$ referring to its flavour representation f , hypercharge y , isospin i and intrinsic spin s . The mass formula for dibaryons is assumed^[7] to have the form

$$m_D = m_l + m_{\text{mag}} \quad (2.2)$$

where m_{mag} is the contribution due to the colour magnetic interactions, treated in the one gluon exchange approximation and m_l has the form of a linear Regge trajectory,

$$m_l^2 = m_0^2 + l/\alpha \quad (2.3)$$

where l is the relative orbital angular momentum between the two clusters. The intercept m_0^2 is obtained from the spherical cavity approximation^[1] to the bag model in the form

$$m_0 = \frac{4\pi B R^3}{3} - \frac{Z_0}{R} + \sum_i \frac{N_i \alpha_i(R)}{R} \quad (2.4)$$

The constants B and Z_0 are taken from a fit to ordinary baryon and meson masses^[1,7], N_i counts the number of quarks of flavour $i \in (u,d,s)$, and α_i measures^[1] the energy of a quark in a spherical bag of radius $R=r_0 N^{1/3}$. The slope $1/\alpha$ of the Regge trajectories depends on the colour structure of the quark clusters via

$$1/\alpha = (1.1 \text{ GeV}^2)^{3/4} f_c^2)^{1/2} \quad (2.5)$$

where f_c^2 is the eigenvalue^[7] of the quadratic Casimir operator for colour representation c (for $c=3$, $1/\alpha$ is the usual slope of meson and baryon trajectories). The colour magnetic energy is assumed^[7] to have the form

$$m_{\text{mag}} = m_1 \Delta_1 + m_2 \Delta_2 \quad (2.6)$$

where

$$m_i = aN_i^{-1/3} - bN_i^S/N_i \quad (2.7)$$

$$\begin{aligned} \Delta_i &= -1/4 N_i(10-N_i) + 1/3 s_i(s_i+1) + (f_f^i)^2 + 1/2 (f_c^i)^2 \\ &= -\sum_{i>j} (F^C \sigma)_i \cdot (F^C \sigma)_j \end{aligned}$$

Here $a=107$ MeV, $b=28$ MeV, N_i^S is the number of strange quarks in cluster i , and f_f^i is the (root of the) eigenvalue of the quadratic Casimir operator for flavour SU(3). Interactions between quarks in different clusters are small^[6] and neglected here.

Using the formulae (2.2)-(2.7), one can compute the masses of all six-quark dibaryon states as done for example in ref.(7). This gives rise to a huge number of configurations of total hypercharge $Y=0,1,2$ (strangeness $S=-2,-1,0$, respectively). These correspond to the "primitives" or P-matrix poles, in the notation of Jaffe and Low^[27]. In most cases, these "primitives" do not correspond to S-matrix poles with a small imaginary part, and hence are not interpretable as long-lived multi-quark excitations. For instance, in a recent study of the coupled $^1D_2(pp)-^5S_2(N\Delta)$ channels in the region above the Δ -isobar production threshold, Grach et al.^[28] have shown how a $J^\pi = 2^+$ six-quark "primitive" gives rise to an S-matrix pole displaced by some 200 MeV, with a width of 100-200 MeV. The observed anomaly in the 1D_2 pp amplitude then occurs about 200 MeV below the energy of the "primitive", and corresponds to a very broad structure.

In the present paper, we focus on the question of whether any of the multitude of bag-model "primitives" are likely to correspond to long-lived excitations with a small decay width into the energetically available baryon-baryon channels. This requires a study of the stability of individual quark clusters against spontaneous dissociation by quark emission, i.e. $Q^n \rightarrow Q^{n-1} + Q$, as well as the stability of composites $Q^n \otimes Q^m$ against dissociation into $(Q^3)_1 + (Q^3)_1$ (two colour-singlet baryons). Our discussion is based on the colour magnetic energy m_{mag} of Eqs. (2.6) and (2.7). For cluster dissociation (virtual) of the type

$$Q^n \rightarrow Q^{n-p} + Q^p \quad (2.8)$$

or dibaryon decay (real)

$$D = Q^n \otimes Q^m \rightarrow (Q^3)_1 + (Q^3)_1 = B_1 + B_2 \quad (2.9)$$

one can define an energy difference

$$\Delta m_{\text{mag}} = \begin{cases} m_{\text{mag}}(Q^n) - m_{\text{mag}}(Q^{n-P}) - m_{\text{mag}}(Q^P) \\ m_{\text{mag}}(D) - m_{\text{mag}}(B_1) - m_{\text{mag}}(B_2) \end{cases} \quad (2.10)$$

If $\Delta m_{\text{mag}} < 0$, the cluster Q^n or the dibaryon D is "stable" against quark emission or quark rearrangement, respectively. We now note^[7] that m_i does not depend very strongly on N_i and N_i^S , so for a qualitative discussion we can replace m_i by an average value $\bar{m}$ and thus obtain

$$\Delta m_{\text{mag}} = \bar{m} \delta \Delta \quad (2.11)$$

where $\delta \Delta$ is the difference of colour-spin expectation values given by Eq.(2.7), and tabulated in Table 3 of ref.(7). A quark configuration D is then "stable" against dissociation into some channel $A + B$ if $\delta \Delta = \Delta_D - \Delta_A - \Delta_B$ is negative.

As an example, we show the $S=-1$ and $S=-2$ dibaryon spectrum from ref.(7) in Fig. 1. Only configurations lying below the $\Lambda N \pi$ ($S=-1$) and $\Lambda \Lambda \pi$ ($S=-2$) thresholds are shown. These correspond to S and P wave states in the baryon-baryon decay channels. States lying above these pion emission thresholds are likely to be very broad, and we will not consider them further here. In Table 1, we display in more detail the flavour-spin structure of these states, as well as the differences $\delta \Delta$ for $D \rightarrow (Q^3)_1 + (Q^3)_1$ and

$$\delta M = \begin{cases} m_D - m_\Lambda - m_N & (S=-1) \\ m_D - 2m_\Lambda & (S=-2) \end{cases} \quad (2.12)$$

δM provides a measure of the energy available for the strong decay modes $D \rightarrow B_1 B_2$. For $(Q^6)_1$ configurations, only the $Y=0$, 0^+ state is stable against dissociation ($\delta \Delta < 0$). This object is the H particle discussed by Jaffe^[4]. It is the only candidate for a dibaryon stable against strong decays ($\delta M < 0$). Mechanisms for its production have been discussed in detail^[29]. An object stable against strong decay with the quantum numbers of the H also emerges in a chiral model with $SU(3)$ flavour symmetry (see Balachandran et al. in ref.(4)). This model

treats a baryon as a topological soliton and makes no explicit reference to quarks and their colour magnetic interactions. In the context of the quark bag model, SU(3) flavour representations of low dimensionality are favoured energetically. For a flavour singlet (the H, with two strange quarks), the colour magnetic interaction is attractive and maximal.

From Fig. 1 and Table 1, we note that the $(Q^5)_1$ structures which couple to the 3S_1 baryon-baryon (B-B) channels are expected to be very broad ($\delta\Delta > 0$ and δM rather sizeable). Thus, except for the H, all candidates for narrow strange dibaryons are of $Q^4 \otimes Q^2$ structure, coupling to P waves of the B-B system. In the bag model, states of cluster structure $(Q^5)_3 * \otimes (Q)_3$, $(Q^4)_6 * \otimes (Q^2)_6$ and $(Q^3)_8 \otimes (Q^3)_8$ all lie at least 100 MeV above the $BB\pi$ threshold and have $\delta\Delta > 0$, so they are intrinsically very unstable.

The clusters $(Q^2)_3^*$ and $(Q^4)_3$ which build the P-wave low-lying strange dibaryons are examined in more detail in Table 2. We see that they both correspond to flavour representations of low dimensionality (3^* or 3). This gives these clusters a considerable amount of colour magnetic attraction, which stabilizes them against quark dissociation processes, as shown in Table 2. Except for the transition $(Q^4)_3 \rightarrow (Q^3)_8 [1, 1/2]^+ Q$, discussed later, all other dissociation processes of $(Q^2)_3^*$ or $(Q^4)_3$ have $\delta\Delta < 0$. Most importantly, $(Q^4)_3$ is stable against emitting a quark plus a colour singlet baryon.

We intend to focus on the $S=-1$ dibaryons in this paper, but a few remarks on the $S=-2$ states of $(Q^4)_3 \otimes (Q^2)_3^*$ character are in order. These objects occur in both spin singlet or triplet forms which we denote by H'_s and H'_t , respectively. The triplet states of isospin $I=1$ should be excited in the reaction

$$d(K^-, K^+) H'_t(I=1) \quad (2.13)$$

In an experiment at CERN, the reaction (2.13) was studied^[30] at 1.4 GeV/c and $\theta_{K^+} \approx 0^\circ$. No evidence was found for a narrow ($\Gamma \lesssim 8$ MeV) structure at the level of 5-40 nb/sr, depending on its mass. In Fig. 1, we note that a state of type $H'_t(I=1)$ is predicted in the vicinity of $m_D \approx 2340$ MeV, above both ΞN and $\Lambda \Sigma$ decay thresholds. At 0° , it will thus be embedded in a large background coming from the reaction $d(K^-, K^+) \Xi^- n$, and hence would be very difficult to see.

The production of H'_t requires an orbital angular momentum transfer $\Delta L=1$ (it is a P-wave baryon-baryon state), and hence the cross-section peaks at finite angle, where the background has also dropped off considerably. The "signal / noise" ratio is thus dramatically improved at some optimal non-zero angle for the K^+ . We discuss these kinematical features in detail later in connection with the reaction $d(K^-, \pi^-)D_t$. Note that the spin singlet $S=-2$ dibaryons H'_s could be produced in the process

$${}^3\text{He} (K^-, K^+) H'_s n \quad (2.14)$$

The (K^-, K^+) reaction has been discussed in detail for H production in ref.(29). For the H, which is an 1S_0 BB configuration, the (K^-, K^+) cross-section peaks at forward angles, while for H'_s , the maximum will be found at non-zero angle as for H'_t production.

Let us now return to the $S=-1$ spectrum. As per Fig. 1, the only candidates for narrow states are the configurations

$$\begin{aligned} D_s &= \{[3(1/3, 1/2)0] \otimes [3^*(2/3, 0)0]\} {}_1P_1, \quad I=1/2. \\ D_t &= \{[3(1/3, 1/2)1] \otimes [3^*(2/3, 0)0]\} {}_3P_J, \quad I=1/2. \end{aligned} \quad (2.15)$$

of $(Q^4)_3 \otimes (Q^2)_3^*$ structure. In D_s and D_t , the strange quark resides in the $(Q^4)_3$ cluster. Note that both D_s and D_t are $I=1/2$ states. The lowest $S=-1$, $I=3/2$ states are spin triplets predicted^[7] around 2330 MeV $[(Q^6)_1$ with $f(\gamma, i)s=10(1, 3/2)1]$ and 2340 MeV $[(Q^4)_3 \otimes (Q^2)_3^*$ of $3(1/3, 1/2)0 \otimes 6(2/3, 1)1$ structure]. These lie well above the $\Lambda N\pi$ threshold, and are also unstable to dissociation into $(Q^3)_1 + (Q^3)_1$, for which $\delta\Delta = +14/3, +2/3$, respectively. They will certainly be very broad. There is no experimental indication of any spin triplet $I=3/2, S=-1$ dibaryons. These have been searched for^[31] in the $d(\pi^-, K^+)D_t(I=3/2)$ and $d(K^-, \pi^+)D_t(I=3/2)$ reactions at 1.4 GeV/c.

The possible quark dissociation modes of D_s and D_t are given in Table 3, along with the values of $\delta\Delta$. These are processes in which a quark tunnels from the Q^4 to the Q^2 cluster or vice versa. Note that both D_s and D_t are stable ($\delta\Delta < 0$) against quark tunneling into $(Q^3)_1 + (Q^3)_1$, i.e., two colour singlet baryons. This of course does not imply that D_s and D_t are absolutely stable against strong decay into YN channels (since $D_{s,t}$ certainly lie above the ΛN and perhaps also the

ΣN thresholds), but simply that decays $D_{s,t} \rightarrow YN$ are not super-allowed, and could give rise to narrow widths. The decay widths of $D_{s,t}$ have been estimated in ref.(23), in a model involving wave function rearrangement (virtual) of $D_{s,t}$ into YN channels, followed by a final state interaction which puts the YN system on-shell. The widths were found to be rather small (of the order of 1 MeV), indicating that D_s and D_t are indeed good candidates to be long-lived dibaryons.

Comparing Tables 2 and 3, we see that the instability of the $(Q^4)_3[3,1]$ cluster for fissioning into $(Q^3)_8[1,1/2] + Q$ is of little consequence. It is unfavourable for any of the non-strange quarks to recombine with the $(Q^2)_3^*$ cluster to form a colour octet, flavour octet baryon.

In the present article, we concentrate on the strange dibaryons D_s and D_t , but one might ask if stable structures can also exist for the non-strange sector. It is clear in general that the low-lying strange dibaryons have stronger colour magnetic attraction than their non-strange counterparts. This is because lower dimensional flavour representations are permitted for systems containing strange quarks, and hence the colour magnetic interaction Δ of Eqs.(2.6) and (2.7) becomes more attractive (smaller values of $(f_F)^2$) than for $S=0$ systems. For non-strange dibaryons the lowest-lying configurations^[7] are listed in Table 4. The $(Q^6)_1$ structures lie well above the NN and $NN\pi$ thresholds. They are also very unstable in terms of colour magnetic interactions ($\delta\Delta$ is positive and large), and hence they cannot correspond to narrow structures. The first two $(Q^4)_3 \otimes (Q^2)_3^*$ configurations in Table 4 are less unstable ($\delta\Delta$ is smaller) in the quark tunneling sense, but lie about 95 or 185 MeV above the $NN\pi$ threshold. Note that these objects cannot couple to the NN channel (1P_1 and 3P_1 are forbidden by the Pauli principle) and also cannot couple to $N\Delta$ except through interactions which violate isospin or intrinsic spin conservation. They have substantial phase space for decay into the $(NN\pi)_{I=0}$ continuum, however.

Recently, Tatischeff et al.^[13] have studied the reaction $p(^3\text{He}, d)D$ at a ^3He kinetic energy of 2.7 GeV, and presented evidence for a narrow $I=1$ structure with $m_D=2.24$ GeV and a width $\Gamma \approx 16$ MeV. A study^[32] of the process $d(d, d)D$, which picks out $I=0$ dibaryons D , did not show any indication of narrow peaks.

Since the structure $D(2240)$ of Tatischeff et al.^[13] lies 70 MeV above the $N\Delta$ threshold, it is difficult to imagine how it can be narrow

if the decays $D \rightarrow NN, N\Delta$ are allowed. Hence, in the context of the bag model, the only reasonable assignment would seem to be

$$D(2240) = (Q^4)_3 [15(^4/3, 1)0] \otimes (Q^2)_3 [3^*(^2/3, 0)0] \quad (2.16)$$

with quantum numbers 3P_1 in LS coupling. The predicted mass is 2200 MeV. This is one of the "extraneous" configurations^[33], i.e. a six-quark system whose quantum numbers are forbidden for NN and $N\Delta$ channels. However, it is not clear that such an extraneous state would be narrow, since it can decay into the $NN\pi$ continuum (non- $N\Delta$ part). If the 3P_1 state is seen in the $p(^3\text{He}, d)D$ reaction, one might expect that the lower-lying 3P_J "extraneous" state would be observed in the process $d(d, d)D$. Although the identification (2.16) is an intriguing and exciting possibility, one must certainly await experimental confirmation in the $D(2240)$. Note that the last entry in Table 4 has $l=1$ and reproduces almost exactly the observed^[13] mass of $D(2240)$. However, this object corresponds to the 3P_J channel, which is Pauli-allowed for NN and also allowed for ΔN ; hence one expects a large width.

In two recent experiments, namely studies of the reactions $p(d, p)D$ at 3.3 GeV/c^[34] and $^3\text{He}(p, d)D$ at 0.925 GeV/c^[35], evidence for a narrow ($\Gamma \approx 20$ MeV) non-strange dibaryon at $m_D \approx 2120$ MeV has been presented. On the basis of the (p, d) measurements^[35], the $D(2120)$ must have $l=1$. This object lies some 50 MeV below the ΔN threshold, which could possibly account for its narrow width. In the six-quark bag spectrum^[7], no $l=1$ states are predicted below 2200 MeV ("extraneous" or not), so there appears no natural identification of the $D(2120)$. In the context of non-relativistic potential models, bound states of $N\Delta$ or $\Delta\Delta$ character have been discussed^[36,37]. Arenhövel^[36] has shown how the decay width of the (off-shell) Δ in a specific ΔN bound state can be reduced by a factor of three to five by binding effects and the Pauli principle. For a $J=2, l=2$ ΔN state, he obtains^[36] a binding energy of the order of 30 MeV, and a width $\Gamma \approx 30$ MeV, but no narrow deeply bound ΔN state is found for $l=1$. Sato and Saito^[37], on the other hand, obtain many $\Delta\Delta$ bound states in a one boson exchange model ($l=0, 1, 2, 3$), but have not estimated the decay widths.

As an alternative to the six-quark bag calculations emphasized here, one could also consider the possibility of strangeness $S=-1$ or $S=-2$ bound states in potential models. A recent calculation^[38] indicates the possibility of some weakly bound states (analogous to the deuteron)

lying near the YN thresholds. In the context of potential models, it may also be interesting to consider whether bound states of type $Y\Delta$ or Y^*N might in some cases be quasi-stable.

It is worth mentioning another approach to the dibaryon spectrum, namely that of dual topological unitarization (DTU). Balázs and Nicolescu^[39] have shown in DTU that the Regge trajectories of dibaryons and six-quark mesons ($QQQ\bar{Q}\bar{Q}\bar{Q}$) are the same. The two lowest-lying states are at 2180 MeV ($J=0$, intrinsic spin $S=0$) and 2310 MeV ($I=1$, $S=1$). The isospin $I=0,1,2$, and 3 dibaryons are degenerate in mass in this approach. This feature is quite distinct from the bag model, where this isospin degeneracy is badly broken. For instance, in ref(7) the lowest $I=2,3$ dibaryons lie several hundred MeV above the $I=0,1$ states of Table 4. The DTU dibaryons have the unusual feature that they only couple to BB channels through higher order graphs in the topological expansion. Rather they prefer to decay into a baryon and a particle of $Q^4\bar{Q}$ structure. Since these $Q^4\bar{Q}$ objects are inhibited from direct decay into meson+nucleon by a topological selection rule, it would be very difficult to form these DTU dibaryons in a reaction such as $p(^3\text{He},d)D$. However, once formed, they are expected to have narrow decay widths into BB channels, unlike six-quark bag states, which are not subject to a topological selection rule.

3 PRODUCTION MECHANISMS OF THE STRANGE DIBARYONS D_s AND D_t .

The mechanisms we envisage for the production of D_s and D_t are shown in Fig. 2. A hyperon Y is first produced via the strangeness exchange reaction $\bar{K}N \rightarrow \pi Y$ and then the hyperon fuses with a second nucleon of the target to form D_s or D_t . The "background" or "quasi-elastic" processes which lead to Λp continuum states are displayed in Fig. 3 for the reaction $d(K^-, \pi^-)\Lambda p$. Similar diagrams were evaluated^[29] for H production in the $^3\text{He}(K^-, K^+)Hn$ reaction.

The processes of Fig. 2 are favourable for strange dibaryon formation, since the necessary strange quark is already present in the initial state. However, the line reversed reactions (π^+, K^+) or (π^-, K^+) could also serve as well. Although the momentum transfer q for associated production is substantial, even at $\theta_{K^+}=0^\circ$, this is not unfavourable for $D_{s,t}$ formation, which requires orbital angular momentum transfer $\Delta L=1$ in any case. The (π, K) reaction was suggested^[40] as a general means of producing Λ and Σ hypernuclei in

higher spin excited states. The particular case $^{12}\text{C} (\pi^+, K^+) ^{12}_{\Lambda}\text{C}$ was later studied experimentally^[41]. The (π, K) channel would of course benefit relative to $(\bar{K}, \pi)$ from the greater available intensity of pion beams, which compensates for the fact that the two-body $\pi N \rightarrow KY$ cross-sections are smaller than those for $\bar{K}N \rightarrow \pi Y$. The form factors $F(q)$ for the $(\bar{K}, \pi)$ and (π, K) reactions^[40], in the absence of distortion effects, are the same. However, $\theta=0^\circ$ corresponds to different values for q for the two reactions. Thus (π, K) cross-sections for $D_{s,t}$ formation will be forward peaked, while $(\bar{K}, \pi)$ cross-sections peak at finite angle where q is optimally matched for a $\Delta L=1$ transition.

To produce dibaryons of different spin, different nuclear targets are required, as shown in Fig. 1. For the spin triplet dibaryon D_t , a deuteron target (also spin triplet) is needed, since the $(\bar{K}, \pi)$ or (π, K) reactions produce no spin-flip at forward angles (and very little at finite angles). For D_s formation, on the other hand, a spin singlet nucleon-nucleon pair is needed in the nuclear target. A di-proton is thus the optimum target, since by the Pauli principle it is automatically in the spin singlet state (S-wave). For D_s formation, the lightest available nuclear target is then ^3He (Fig. 2(b)). The $d(K^-, \pi^-)D_t$ reaction has the advantage of a two-body final state, so that a narrow D_t shows up as a sharp peak in the YN missing mass spectrum, superimposed on a continuum background arising from the graphs in Fig. 3. For D_s , a three-(or more) body final state is always required, so one would in principle have to measure π^+n coincidences to obtain a peak with the width of D_s in a missing mass plot. However, as our numerical results (Chap. 5) show, even if only the π^+ is detected, the (distributed) $D_s n$ missing mass still displays a peak reasonably well separated from the Σ^-pn continuum, at least if one accepts the theoretical estimates of the mass D_s .

The techniques for evaluating the cross-sections for the processes of Figs. 2 and 3 have already been discussed in some detail previously^[29], so we provide only a skeleton summary here. The ingredients of the calculations are

- i) amplitudes T for the reactions $\bar{K}N \rightarrow \pi Y$,
- ii) a model for the d or ^3He ground state wave functions, and
- iii) a description of the fusion vertex Γ_D for $YN \rightarrow D$.

For the $\bar{K}N \rightarrow \pi Y$ amplitudes, we use the model of Gopal et al.^[42], which provides a good fit to the available on-shell two-body data. We replace the off-shell T -matrices in Figs. 2 and 3 by their on-shell limits. The overall calculation is done in the K^-d (or $K^-^3\text{He}$) lab

system. The struck nucleon (assumed on-shell) has a certain momentum distribution (integrated over) in the lab frame. Given the nucleon lab momentum, we perform a Lorentz boost to the K^-N c.m. system and compute the energy $\sqrt{s}$ of the collision. The momentum transfer $\sqrt{|t|}$ is computed from the pion lab angle and momentum. In this way, the full "Fermi averaging" of the $\bar{K}N \rightarrow \pi Y$ amplitudes is included.

The $d(K^-, \pi^-)D_t$ amplitude is proportional to a coherent sum of amplitudes for different charge states ($K^-n \rightarrow \pi^- \Lambda, \pi^- \Sigma^0$ or $K^-p \rightarrow \pi^- \Sigma^+$) weighted by recoupling coefficients for the fusion processes $\Lambda p, \Sigma^0 p$ or $\Sigma^+ n \rightarrow D_t$. The cross-section thus depends sensitively on the relative phase of the $\bar{K}N \rightarrow \pi \Lambda$ and $\pi \Sigma$ amplitudes, a quantity which is not fixed directly by the experimental data^[19]. The phase adopted by Gopal et al.^[42] is consistent with that expected from SU(6) symmetry. This gives destructive interference between the $\bar{K}N \rightarrow \pi \Lambda$ and $\pi \Sigma$ channels (the $\pi^- \Sigma^0 p$ and $\pi^- \Sigma^+ n$ channels interfere constructively). We have also performed calculations in which this relative phase is changed arbitrarily. Such phase rotations were also considered by Dalitz and Deloff^[19]. If the phase is rotated by 180° , the $\pi \Lambda, \pi \Sigma$ interference is maximally constructive, leading to the largest $d(K^-, \pi^-)D_t$ cross-section. These calculations are labelled AG ("anti-Gopal") as opposed to G ("Gopal"). For D_s production on ^3He , this phase does not enter, since only the $\pi^+ \Sigma^-$ charge combination contributes.

For the $d \rightarrow np$ vertex Γ_d (or equivalently the deuteron wave function) we have adopted the standard form

$$|d, \underline{p}_d\rangle = \Gamma_d = \frac{1}{(2\pi)^3} \sum_{s_i, t_i} \int d\underline{k}_1 \int d\underline{k}_2 \frac{(2E_d)^{1/2} \delta^{(3)}(\underline{p}_d - \underline{k}_1 - \underline{k}_2) \phi_d(\underline{k}_1 - \underline{k}_2)}{(4E_n(\underline{k}_1)E_p(\underline{k}_2))^{1/2}} \quad (3.1)$$

$$* \sqrt{1/2} C(s_1 s_2 1) C(t_1 t_2 0) |\underline{k}_1 s_1 t_1\rangle |\underline{k}_2 s_2 t_2\rangle$$

where the C's couple the n and p spins and isospins s_i and t_i to $S=1, I=0$, and the normalization is

$$\langle d, \underline{p}'_d | d, \underline{p}_d \rangle = (2\pi)^3 (2E_d) \delta^{(3)}(\underline{p}_d - \underline{p}'_d) \quad (3.2)$$

In Eq.(3.1), E_n, E_p and E_d are total energies and ϕ_d is the relative wave function, which we take to be of Hulthén form

$$\phi_d(k) = C_0 [(\beta^2 + k^2)^{-1} - (\gamma^2 + k^2)^{-1}] \quad (3.3)$$

with $\beta^2=2.096 \times 10^{-3} \text{ GeV}^2$ and $\gamma^2=0.10269 \text{ GeV}^2$.

For the wave function of ^3He (the vertex Γ_{He} in Fig. 2(b)), we have adopted an oscillator form

$$\psi_{\text{He}}^{\underline{k}}(\underline{r}_1, \underline{r}_2, \underline{r}_3) = \sqrt{2} E e^{i \underline{k} \cdot (\underline{r}_1 + \underline{r}_2 + \underline{r}_3)} \psi_{\text{He}}(\underline{\rho}, \underline{\lambda}) \quad (3.4)$$

with a relative wave function

$$\psi_{\text{He}}(\underline{\rho}, \underline{\lambda}) = 3^{-3/4} (\pi b^2)^{-3/2} e^{-(\underline{\rho}^2 + \underline{\lambda}^2)/2b^2} \quad (3.5)$$

depending on relative co-ordinates $\underline{\rho} = (\underline{r}_1 - \underline{r}_2)/\sqrt{2}$ and $\underline{\lambda} = (\underline{r}_1 + \underline{r}_2 - 2\underline{r}_3)/\sqrt{6}$. The oscillator parameter $b=1.7 \text{ fm}$ is chosen to reproduce the r.m.s. charge radius of ^3He . In ref.(29), the effect of short range correlations in ψ_{He} on the H production cross-section was explored. Correlation effects were not found to be large, and are neglected here. Our numerical calculations are performed in momentum space, so we utilize the Fourier transform $\psi_{\text{He}}(\underline{k}_{\underline{\rho}}, \underline{k}_{\underline{\lambda}})$ of Eq.(3.5).

To obtain a model of the $\text{YN} \rightarrow \text{D}$ fusion vertex Γ_{D} , we need a description of the dibaryon wave function $\psi_{\underline{k}}(\underline{r}_1, \dots, \underline{r}_6)$. Here $\underline{r}_1$ through $\underline{r}_4$ label the quarks in the Q^4 cluster, while $\underline{r}_5$ and $\underline{r}_6$ refer to those in the Q^2 cluster. The quarks in each cluster are in relative S-states while the two clusters are in a relative P-state ($L=1$). We assume an oscillator form

$$\begin{aligned} \psi_{\underline{k}}(\underline{r}_1, \dots, \underline{r}_6) = & N e^{i \underline{k} \cdot \underline{R}} (r/R_p) Y_{1m}(\hat{\underline{r}}) \\ & * e^{-1/2(\underline{\rho}^2 + \underline{\lambda}^2 + \underline{\sigma}^2 + \underline{\xi}^2)/R^2} e^{-\underline{r}^2/2R_p^2} \end{aligned} \quad (3.6)$$

where the co-ordinates $\{\underline{R}, \underline{\rho}, \underline{\lambda}, \underline{\sigma}, \underline{\xi}, \underline{r}\}$ are defined by

$$\begin{aligned} \underline{r} &= (\underline{r}_1 + \underline{r}_2 + \underline{r}_3 + \underline{r}_4 - 2(\underline{r}_5 + \underline{r}_6))/\sqrt{12} \\ \underline{R} &= (\underline{r}_1 + \underline{r}_2 + \underline{r}_3 + \underline{r}_4 + \underline{r}_5 + \underline{r}_6)/6 \\ \underline{\rho} &= (\underline{r}_1 - \underline{r}_2)/\sqrt{2} \\ \underline{\lambda} &= (\underline{r}_1 + \underline{r}_2 - 2\underline{r}_3)/\sqrt{6} \\ \underline{\sigma} &= (\underline{r}_1 + \underline{r}_2 + \underline{r}_3 - 3\underline{r}_4)/\sqrt{12} \\ \underline{\xi} &= (\underline{r}_5 - \underline{r}_6)/\sqrt{2} \end{aligned} \quad (3.7)$$

We now consider the rearrangement process $Q^4 \otimes Q^2 \rightarrow Q^3 + Q^3$. For the Q^3 clusters, we use an oscillator wave function as in Eqs. (3.4) and (3.5), with radius parameter R_B . The overlap of the dibaryon D with a configuration of two Q^3 baryons Y and N of momenta $\underline{k}_Y$ and $\underline{k}_N$ is given by

$$\begin{aligned} \langle \underline{k}_Y, \underline{k}_N | \psi_K \rangle = & (2E_D 2E_Y 2E_N)^{1/2} (2\pi)^3 \delta^{(3)}(\underline{K} - \underline{k}_Y - \underline{k}_N) \\ & * C(YN, D) \Gamma_D(\underline{k}_Y - \underline{k}_N) \end{aligned} \quad (3.8)$$

with a vertex function Γ_D of the form

$$\begin{aligned} \Gamma_D(k) = & N_{\Gamma} k Y_{1m}(\hat{k}) e^{-1/12 R'^2 k^2} \\ R'^2 = & (R_B^2 R_p^2 + R_B^2 R_D^2 + 2R_D^2 R_p^2) / (R_D^2 + R_p^2 + 2R_B^2) \end{aligned} \quad (3.9)$$

The bag radius $R_B = 0.83$ fm has been chosen^[29] to obtain the correct r.m.s. radii for the nucleon and the hyperon (the latter assumed to have the same size as the nucleon). Further, we take $R_D = 0.83$ fm and pick $R_p = 1.02$ fm to be somewhat larger than the other radii in our dibaryon to ensure that the Q^4 and Q^2 clusters are non-overlapping. This gives $R' = 0.92$ fm as the radius which appears in Eq.(3.9).

In Eq.(3.8), $C(YN, D)$ is a recoupling coefficient for the space, spin, colour and flavour parts of the dibaryon wave function, when expressed in a basis of $Q^3 \otimes Q^3$ states. For D_s and D_t , we find^[23]

$$\begin{aligned} |D_s\rangle = & \sum_{i,j} C(ij, D_s) |ij\rangle = \frac{1}{16} \{ \sqrt{3}(p\Lambda - \Lambda p) - \sqrt{2}(n\Sigma^+ - \Sigma^+ n) + (p\Sigma^0 - \Sigma^0 p) \} + \dots \\ |D_t\rangle = & \sum_{i,j} C(ij, D_t) |ij\rangle = \frac{1}{16\sqrt{3}} \{ (p\Lambda + \Lambda p) + \sqrt{6}(n\Sigma^+ + \Sigma^+ n) - \sqrt{3}(p\Sigma^0 + \Sigma^0 p) \} + \dots \end{aligned} \quad (3.10)$$

where we have explicitly isolated the channels where both Q^3 clusters belong to the {56}-plet of $SU(6)$. These correspond to the energetically available channels for the decay of $D_{s,t}$. The dots in Eq.(3.10) refer to configurations involving the {70}-plet of $SU(6)$, colour octet Q^3 clusters, etc.

The decomposition (3.10) is appropriate for the short distance part of the $D_{s,t}$ wave functions, i.e. it contains a number of virtual components which cannot be present asymptotically (large r) because of colour confinement and/or energy conservation. In the absence of a dynamical understanding of how the wave function evolves from the small to large r regions, we simply adopt Eq.(3.10) for computing the $D_{s,t}$ formation cross-sections as per Figs. 2(a) and 2(b). Note that the fraction of the $D_{s,t}$ wave functions (3.10) involving the physically observable ΛN and ΣN decay channels is rather small, so our dibaryon cross-section estimates are likely to be rather conservative. If, on the other hand, we chose to write a normalized (large r) wave function for $D_{s,t}$ involving only YN channels, then the decomposition (3.10) (omitting now the dots) is multiplied by factors $N_s=8/\sqrt{3}$ for D_s and $N_t=8\sqrt{3}/\sqrt{5}$ for D_t . Our cross-section estimates (Chaps. 4 and 5) would then be multiplied by $N_s^2 \approx 21.3$ and $N_t^2 \approx 38.4$; this would certainly represent an upper limit for $D_{s,t}$ production cross-sections.

We have now assembled the ingredients for a calculation of the amplitudes of Figs. 2(a) and 2(b). The procedure involves standard techniques already discussed in ref.(29) for the case of H production, so we omit details. For the reaction $d(K^-, \pi^-)D_t$, for instance, we have a T-matrix element of the form

$$\langle \pi^- D | T | K^- d \rangle = \alpha \int d^3 \underline{\eta} \, \eta Y_{1m}(\hat{\eta}) e^{-\eta^2} T_{K^- n \rightarrow \pi^- \Lambda}(\underline{k}_L, \underline{\xi}_-, \underline{k}_{\pi L}, \underline{\xi}_+) \quad (3.11)$$

$$\frac{1}{(E_n(\underline{\xi}_-) E_\Lambda(\underline{\xi}_+))^{1/2}}$$

where $\underline{\eta} = R'(\underline{k}_L - \underline{k}_{\pi L} + 2\underline{k}_n)/2\sqrt{3}$ and $\underline{\xi}_\pm = \sqrt{3}\underline{\eta}/R' \pm \underline{k}_D/2$. Here, $\underline{k}_L$, $\underline{k}_{\pi L}$ and $\underline{k}_D$ are the kaon, pion and dibaryon momenta, respectively, in the K^-d lab system. We integrate over the momentum $\underline{k}_n$ of the nucleon which participates in the $K^-N \rightarrow \pi^-Y$ reaction, or, equivalently over the dimensionless variable $\underline{\eta}$. The constant α includes geometrical factors arising from wave function and vertex (Γ_D) normalizations as well as the recoupling weight $C(YN, D)$. The factor $\eta Y_{1m}(\hat{\eta}) e^{-\eta^2}$ comes from the p-wave fusion vertex (3.9). We include the full momentum dependence of the two-body amplitude $T_{K^-N \rightarrow \pi^-Y}$, although our results are not significantly altered if we factor $T_{K^-N \rightarrow \pi^-Y}$ out of the integral (3.11), and evaluate it at some average nucleon momentum. In Eq.(3.11), we have given the case $K^-n \rightarrow \pi^- \Lambda$ as an example. The complete calculation involves the coherent sum

$$C(\Lambda p, D_t) T_{K^- n \rightarrow \pi^- \Lambda} + C(\Sigma^0 p, D_t) T_{K^- n \rightarrow \pi^- \Sigma^0} + \\ + C(\Sigma^+ n, D_t) T_{K^- p \rightarrow \pi^- \Sigma^+} \quad (3.12)$$

The differential lab cross-section for the $d(K^-, \pi^-)D_t$ reaction is then

$$\frac{d\sigma}{d\Omega_\pi} = \frac{1}{2\lambda^{1/2}(s, m_K^2, m_d^2)} \frac{1}{(2\pi)^2} \int_0^{k_{\pi L}^{\max}} k_{\pi L}^2 dk_{\pi L} \delta(E_K + m_d - E_\pi - E_D) \\ * |\langle \pi^- D | T | K^- d \rangle|^2 \quad (3.13)$$

where $\lambda(a, b, c) = a^2 + b^2 + c^2 - 2(ab + bc + ca)$ is the usual triangle function^[43].

Similar expressions may be written for the "quasi-elastic" (QE) processes of Fig. 3. The total QE cross-section for $K^- d \rightarrow \pi^- \Lambda p$, for instance, is of the form

$$\sigma_{K^- d \rightarrow \pi^- \Lambda p}^{\text{QE}} = \frac{1}{(2\pi)^6} \int \frac{d^3 k_{\pi L}}{2E_\pi} \int \frac{d^3 k_p}{2E_p} \int \frac{d^3 k_\Lambda}{2E_\Lambda} \delta^{(4)}(K - k_{\pi L} - k_p - k_\Lambda) \\ * |M^{(3)}|^2 / 2\lambda^{1/2}(s, m_K^2, m_d^2) \quad (3.14)$$

where $K \equiv \{k_K, E_K + m_d\}$ and the matrix element $M^{(3)}$ for the three-body final state is proportional to the product of the amplitude $T_{K^- n \rightarrow \pi^- \Lambda}$ and the deuteron wave function. Differential cross-sections $d\sigma^{\text{QE}}/d\Omega_\pi$ and "missing mass" spectra $d^2\sigma^{\text{QE}}/d\Omega_\pi dM_{\Lambda p}$ may be readily obtained^[29] from Eq.(3.14) by differentiations. For the reactions ${}^3\text{He}(K^-, \pi^+)D_s n$ and ${}^3\text{He}(K^-, \pi^+)\Sigma^- pn$, formulae completely analogous to Eqs. (3.11)-(3.14) may be written. We refer the reader to ref.(29) for further details, and proceed here with a discussion of the numerical results.

4 PRODUCTION CROSS-SECTIONS FOR D_t IN THE $D(K^-, \pi^-)D_t$ REACTION.

We now present our estimates of the cross-sections for the formation of the 3P_J dibaryon D_t of quark structure $Q^4 \otimes Q^2$ (see Eq.(2.15)) in the reaction $d(K^-, \pi^-)D_t$. These predictions are compared with the

cross-sections for quasi-elastic (QE) processes $d(K^-, \pi^-)YN$, which provide a "background" for the dibaryon "signal". We discuss to what extent the signal/background ratio R in a missing mass plot can be improved by imposing a cut-off on the momentum of a decay proton. The choice of optimum kinematics (kaon energy, pion angle) for dibaryon formation is also discussed. We note that the dibaryon D_t actually consists of three states with $J=0,1,2$. We expect^[6,7] that their mass splitting is small (of the order of a few MeV), so that they will appear as a single peak in a missing mass plot. The theoretical cross-sections presented here are summed over J .

The predicted differential cross-sections $d\sigma/d\Omega$ for the $d(K^-, \pi^-)D_t$ reaction are shown in Fig. 4 for a range of incident K^- lab momentum k_L . The dashed curves are obtained when the $\bar{K}N \rightarrow \pi\Lambda$ amplitudes of Gopal et al.^[42] are used. Here, the $\pi\Lambda$ and $\pi\Sigma$ channels interfere destructively in D_t formation. The solid curves refer to a calculation where the relative phase of the $\bar{K}N \rightarrow \pi\Lambda$ and $\bar{K}N \rightarrow \pi\Sigma$ amplitudes is rotated by 180° , producing the maximum constructive interference. Some change of relative phase was required by Dalitz and Deloff^[19] to understand semi-quantitatively the cusp effect in the $d(K^-, \pi^-)\Lambda p$ reaction, i.e. the appearance of a sharp peak at the Σ^+n threshold. The difference between the dashed and the solid lines in Fig. 4 indicates the sizeable uncertainty introduced by the choice of the $\pi\Lambda, \pi\Sigma$ relative phase. This phase is not determined by $\bar{K}N \rightarrow \pi Y$ data at one energy, and is fixed only if one imposes some theoretical prejudices as to the energy dependence of the amplitudes.

Fig. 4 shows the typical peaking at finite angle, due to the fact that the transition from the 3S_1 deuteron to the 3P_J dibaryon D_t requires orbital angular momentum transfer $\Delta L=1$, and hence finite momentum transfer q . This peaking is pronounced when the $\pi\Lambda$ and $\pi\Sigma$ amplitudes are in phase, and somewhat washed out if they interfere destructively. Note that the cusp effect, which involves $\Lambda N \rightarrow \Sigma N$ coupling in the 3S_1 state, is forward peaked, since $\Delta L=0$ is optimum for a $^3S_1 \rightarrow ^3S_1$ transition. The characteristically different angular shapes for the cusp effect and the dibaryon D_t afford a natural way of disentangling them.

In Fig. 5, we plot the peak differential cross-section for D_t formation as a function of k_L . The lab angle corresponding to the peak decreases as k_L increases. The dashed and solid curves correspond to the choices of $\pi\Lambda$ - $\pi\Sigma$ relative phases as in Fig. 4. For curve AG ("anti-Gopal" phases) $k_L \approx 0.95 \text{ GeV}/c$ seems an optimum choice, since the

peak angle ($\theta_L \approx 24^\circ$) is large enough so that the forward peaked cusp effect has largely disappeared. The dashed curves illustrate the point that the $\bar{K}N \rightarrow \pi \Sigma$ channel, rather than the $\pi \Lambda$, provides the largest contribution to the D_t cross-section. This is due to the fact that the recoupling coefficients $C(\Sigma N, D_t)$ are significantly larger than $C(\Lambda p, D_t)$, as per Eq.(3.10). Note that the opposite sign of $C(\Sigma^+ n, D_t)$ and $C(\Sigma^0 p, D_t)$ is compensated by a relative minus sign between $T_{K^- n \rightarrow \pi^- \Sigma^0}$ and $T_{K^- p \rightarrow \pi^- \Sigma^+}$, so the two ΣN channels always add coherently. Thus the peak D_t cross-section is not attained at a momentum $k_L \approx 0.87$ GeV/c, where the $K^- n \rightarrow \pi^- \Lambda$ cross-section is maximum, but at a somewhat higher value.

In Fig. 6, we display angular distributions for the $d(K^-, \pi^-)D_t$ and $d(K^-, \pi^-)YN$ reactions at $k_L = 0.95$ GeV/c. Even though the $d(K^-, \pi^-)\Lambda p$ cross-section is strongly forward peaked, its value at $\cos\theta_L \approx 0.9$, where D_t production is maximum, is still very large. Summing over the YN quasi-elastic channels, we obtain

$$\frac{(d\sigma/d\Omega)_{D_t}}{(d\sigma/d\Omega)_{QE}} \approx 2-4 \times 10^{-3} \quad (4.1)$$

in the range $0.95 \geq \cos\theta_L \geq 0.8$. Eq.(4.1) gives essentially the "fusion probability" for D_t formation via YN final state interactions. A very similar number was found in ref.(29) for doubly-strange H dibaryon production. Clearly the D_t signal cannot be seen by simply measuring the angle of the π^- in the final state. The π^- momentum must also be measured with good resolution and used to construct the YN missing mass spectrum.

The dependence of the D_t differential cross-section on the assumed dibaryon mass m_D is shown in Fig. 7. The predicted value from ref.(7) is $m_D \approx 2150$ MeV/c². The dependence of $d\sigma/d\Omega$ on m_D is seen to be rather weak. As m_D increases, the peak moves to smaller angles, but the magnitude of the cross-section remains essentially unchanged.

In Fig. 8, the cross-sections $d^2\sigma/d\Omega dm_{YN}^2$ for D_t formation and QE processes are shown, as a function of the hyperon-nucleon "missing mass" M_{YN} . Favourable kinematical conditions ($k_L = 0.95$ GeV/c, $\cos\theta_L = 0.85$) have been chosen to optimize D_t production. The signal for D_t has been spread out over an interval of 10 MeV, reflecting the currently attainable limit of experimental energy resolution (quoted as 5

MeV in ref.(18)) as well as the intrinsic width (estimated to be of the order of 1 MeV in ref.(23)) and the energy splitting between the $J=0,1,2$ components of D_t . We can now define two types of "signal/noise" ratios R , namely

$$R_1 = \frac{(d^2\sigma/d\Omega dM^2_{YN})_{D_t}}{(d^2\sigma/d\Omega dM^2_{YN})_{\Lambda p + \Sigma^0 p + \Sigma^+ n}} t \quad (4.2)$$

$$R_2 = \frac{(d^2\sigma/d\Omega dM^2_{YN})_{D_t}}{(d^2\sigma/d\Omega dM^2_{YN})_{\Lambda p}} t.$$

In principle, one should include the influence of the cusp at the $\Sigma^+ n$ threshold in the denominators of Eq.(4.2). Taking the parametrization of Braun et al.^[17] for the angular dependence of the cusp effect, we find that it does not significantly influence $R_{1,2}$ except near $\theta_L=0$. We include here only the first order QE process of Fig. 3(a) and neglect 3(b). For three choices of $\cos\theta_L$ and $k_L=0.95$ GeV/c², we list the values of $R_{1,2}$ in Table 5. Constructive ΛN - ΣN interference (AG phases) has been assumed, and a width $\Delta M_{YN}=10$ MeV for the D_t signal. If we had instead assumed $\Delta M_{YN}=5$ MeV, corresponding to the optimum energy resolution claimed in ref.(18), the entries in Table 5 are multiplied by a factor of four. From Table 5, we see that R_1 and R_2 remain less than unity if only the pion angle and momentum (and hence M_{YN}) are measured, so detection of the dibaryon signal might appear marginal. In Fig. 8 and Table 5, there is no cut-off imposed on the nucleon momentum in the final state.

If one detects recoil protons as well as the π^- in the final state, the situation can be improved considerably. As is well known^[15,17,18], if one cuts out events corresponding to final proton momenta $k_p \leq k_{min}$, much of the background due to QE processes for low M_{YN} can be removed. In Fig. 9, we demonstrate the effect of imposing a cut-off $k_{min}=150$ MeV/c on the momentum of the nucleon in the $d(K^-, \pi^-) YN$ reactions. For both Λp and ΣN channels the cut-off has the effect of removing the main part of the QE peaks. Note that a fraction of the protons which arise from Λ or Σ decays is of low momentum, so they are also removed by such a cut-off; this effect has not been included in Fig. 9.

Since a cut-off $k_p \leq k_{\min}$ is a kinematical constraint, it also modifies the dibaryon signal, by removing a certain fraction of the events due to the decay $D_t \rightarrow \Lambda p, \Sigma^0 p$. For the $D_t \rightarrow \Lambda p$ decay mode, the proton momentum is about 300 MeV/c (kinetic energy 50 MeV). At forward angles the dibaryon system has a momentum of 200 MeV/c and so the signal is not altered much by the cut-off (See Fig. 9, we have assumed $m_D = 2.14$ GeV/c²). For $D_t \rightarrow \Sigma^0 p$, on the other hand, the proton momentum is around 100 MeV/c, so these decays are much more sensitive to the cut-off. Fig. 9 also shows that the effect of the cut-off on the Λp and ΣN signals depends strongly on the pion angle (as well as M_{YN}). At $\cos\theta_L = 0.75$ the dibaryon with mass $M_D = 2.14$ GeV/c² will have a momentum of about 600 MeV/c. Now the ΣN signal is not altered much by the cut-off, whereas the Λp signal is cut very strongly. The effect of imposing various cut-off momenta on the dibaryon signal is shown in Table 6. Since the preferred decay mode of D_t is into the ΣN rather than the Λp channel (see Eq.(3.10)!), the effects of the cut-off are especially marked for $k_{\min} = 150$ MeV/c. A cut-off $k_{\min} \approx 100 - 150$ MeV/c seems indicated, since otherwise one removes too much of the dominant $D_t \rightarrow \Sigma N$ decays. In Table 7, we show how the "signal/noise" ratio R_1 changes as $k_{\min}$ is varied. By comparing Tables 5 and 7, we see that at forward angles there is no particular advantage in imposing a nucleon momentum cut-off. At angles $\theta_L \approx 35^\circ$, this will enhance the ΣN part of the dibaryon signal relative to the Λp part, which is favourable if the dibaryon D_t decays predominantly into ΣN channels. In the proposal of ref.(18), in contrast, it was assumed that the main decay mode is $D_t \rightarrow \Lambda p$, and hence the effects of a momentum cut-off are not so beneficial.

In summary, our calculations for the reactions $d(K^-, \pi^-)D_t$ and $d(K^-, \pi^-)YN$ indicate that the search for a dibaryon D_t will be experimentally difficult, although perhaps not unfeasible. The predicted "signal/noise" ratios R are rather small, even when the kinematical variables (k_L, θ_L) are optimally chosen. Good energy resolution is extremely important (assuming the intrinsic width of D_t is very small, as predicted in ref.(23). If D_t has an intrinsic width greater than about 10 MeV, it will be difficult indeed to disentangle it from the quasi-elastic background, particularly if its dominant decay mode is ΣN rather than Λp .

5 CROSS SECTIONS FOR D_s PRODUCTION IN THE ${}^3\text{He}(K^-, \pi^+)nD_s$ REACTION.

The dibaryon D_s of Eq.(2.15) couples to the spin-singlet 1P_1 YN channel. Since the spin-flip parts of the $\bar{K}N \rightarrow \pi Y$ amplitudes are generally weak even at finite angles, D_s cannot be produced with any measurable cross-section via the $d(K^-, \pi)D$ reaction. As discussed in Chap. 3, a ${}^3\text{He}$ target is appropriate, since the di-proton of ${}^3\text{He}$ is automatically in the desired spin-singlet state. In this section, we present some estimates of the cross-sections for the ${}^3\text{He}(K^-, \pi^+)nD_s$ reaction.

We have chosen the (K^-, π^+) rather than the (K^-, π^-) reaction to eliminate first order QE processes involving a Λ in the final state. That is, the first order QE background arises from the reaction ${}^3\text{He}(K^-, \pi^+)np\Sigma^-$ induced by the two-body reaction $K^-p \rightarrow \pi^+\Sigma^-$. If one accepts the predicted^[7] mass $m_{D_s} \approx 2.11 \text{ GeV}/c^2$, the D_s lies below the ΣN thresholds at $2.3 \text{ GeV}/c^2$. It could well be very narrow^[23], since only the ΛN channel is available for its decay. The D_s should then show up clearly in a missing mass plot, if the main QE background is due to Σ formation. Note that the reaction ${}^3\text{He}(K^-, \pi^+)nn\Lambda$ will also occur in general, and this gives events in the same region of missing mass as the D_s . However, this reaction can only occur through a second order process $K^-p \rightarrow \pi^+\Sigma^-$ followed by $\Sigma^-p \rightarrow \Lambda n$, and we argue that this will only be a small part of the QE background. We have neglected such processes in our numerical estimates. It is worth noting that the same idea for eliminating first order Λ background does not work for the $d(K^-, \pi^+)D$ reaction, since in this case only $I=3/2$ dibaryons D can be produced, and not the $I=1/2$ state D_t desired. For the ${}^3\text{He}(K^-, \pi^+)nD_s$ reaction, in contrast, the presence of the spectator neutron enables us to produce both $I=1/2$ and $3/2$ dibaryons D_s .

The predicted differential cross-sections for the ${}^3\text{He}(K^-, \pi^+)nD_s$ reaction are plotted in Fig. 10 for several values of k_L . Here there is no interference between different $\bar{K}N \rightarrow \pi Y$ channels as in the case of $d(K^-, \pi^-)D_t$, so the cross-sections on ${}^3\text{He}$ reflect more directly the $K^-p \rightarrow \pi^+\Sigma^-$ cross-section, which has a peak at about $700 \text{ MeV}/c$ (at 0°). Since D_s formation also requires $\Delta L=1$, as for the D_t dibaryon, the cross-sections peak at a finite angle which decreases as k_L increases. The QE cross section $d\sigma/d\Omega$ is forward peaked and much larger than the D_s signal. As in Eq.(4.1), we find a small overall "fusion probability" for D_s formation :

$$\frac{(d\sigma/d\Omega)_{D_s}}{(d\sigma/d\Omega)_{QE}} \approx 4-6 \times 10^{-3} \quad (5.1)$$

for $k_L \approx 0.8-0.9 \text{ GeV}/c$ and θ_L chosen such that $(d\sigma/d\Omega)_{D_s}$ is maximum.

The missing mass spectrum $d^2\sigma/d\Omega dM_{nX}^2$ is displayed in Fig. 11. Here we measure the angle and momentum of the π^+ only, so M_{nX} is the total invariant mass of the nD_s or $np\Sigma^-$ systems. The peak due to D_s production is well separated from the (much larger) first order QE background. However, to estimate the true "signal/noise" ratio R , we would have to compute the second order contributions to the $nn\Lambda$ background. Even a "leakage" of a few percent of the $np\Sigma^-$ QE into the $nn\Lambda$ channel would reduce R to less than unity.

In Fig. 12, we show a missing mass spectrum appropriate to an experiment in which the π^+ and neutron are measured in coincidence. The missing mass M_X then no longer includes the neutron energy. In Fig. 12, the cross-section $d^2\sigma/d\Omega dM_X^2$ has been integrated over the neutron angles. Here, D_s appears as a sharp peak which is separated from the Σ^-p quasi-elastic background. Fig. 12 also illustrates how the QE part diminishes with angle. Again, the computation of R requires a detailed calculation of how many QE events leak into the region of the D_s due to $\Sigma^-p \rightarrow \Lambda n$ conversion in the final state. Thus it is difficult to evaluate the relative chances of success of the $d(K^-, \pi^-)D_t$ and ${}^3\text{He}(K^-, \pi^+)nD_s$ experiments. The ${}^3\text{He}$ case may be more favourable if D_s indeed lies below the ΣN threshold. However, a π^+-n coincidence measurement is required if one is to isolate D_s as the "missing mass". This will of course suffer from very low counting rates, given the presently available kaon beams of low intensity. The ${}^3\text{He}(K^-, \pi^+)nD_s$ coincidence experiment would benefit enormously from the intense kaon beams of a future "kaon factory".

6 OUTLOOK.

The bag model predicts an abundant spectrum of six-quark dibaryon states of cluster structure $Q^m \otimes Q^n$. Most of these "primitives" are associated with poles of the S-matrix having very large imaginary parts, and thus they do not correspond to long-lived "particles". By invoking simple stability arguments for quark clusters, we are able to reduce this plethora of six-quark states to only a very few which are likely to have small widths for decay into baryon-baryon channels. For

strangeness $S=-1$, we find only two candidates D_s and D_t , of Q^4Q^2 cluster structure, which couple to the 1P_1 and $^3P_{0,1,2}$ hyperon-nucleon channels, respectively.

The goal of this paper is to estimate production cross-sections for the strange D_t and D_s dibaryons. We argue that the $(\bar{K}, \pi)$ and (π, K) reactions on deuterium and ^3He targets are especially suited to the fabrication of D_t and D_s , respectively, and cross-section estimates for the $d(K^-, \pi^-)D_t$ and $^3\text{He}(K^-, \pi^+)nD_s$ reactions have been obtained. Due to the small probability (few $\times 10^{-3}$) of forming a strange dibaryon from the fusion of a hyperon and a nucleon, production cross-sections are small; we find peak values of order $2 \mu\text{b/sr}$ for D_t and $0.5 \mu\text{b/sr}$ for D_s . Since $D_{s,t}$ are P-wave dibaryons, orbital angular momentum transfer $\Delta L=1$ is required for their formation. Thus the cross-sections peak at a characteristic finite angle of order $\theta_L \approx 20-25^\circ$ for incident momenta in the optimum range $k_L \approx 0.8-1.0 \text{ GeV/c}$. This feature makes it conceivable that one could separate the dibaryon signal from the quasi-elastic (QE) background from processes $d(K^-, \pi^-)YN$ or $^3\text{He}(K^-, \pi^+)np\Sigma^-$, which is almost three orders of magnitude larger in the differential cross-section $d\sigma/d\Omega$. This separation is effected by means of a "missing mass" plot $[d^2\sigma/d\Omega dM_{YN}^2]$. If, as we suspect, D_s and D_t have narrow widths, the dibaryon events are concentrated in a narrow region of the hyperon-nucleon invariant mass M_{YN} . By an optimum choice of pion angle θ_L (as above), one passes to a region where the forward peaked QE processes have decreased considerably, and the "signal/noise" ratio R is enhanced. Nevertheless, predicted values of R remain small (order of 0.1 for a width $\Delta M_{YN} \approx 10 \text{ MeV}$ of the dibaryon signal), and the experiments require very good energy resolution to have any chance of success. Also, if $D_{s,t}$ have intrinsic widths larger than 10 MeV or so, it would seem unlikely that they can be detected above the QE background.

Our cross-section estimates are perhaps useful in providing some guidance for experimental strange dibaryon searches. It should be noted, however, that we have made several simplifying assumptions regarding the reaction mechanism and the dibaryon wave functions. For instance, we have used harmonic oscillator wave functions to estimate the strength of the fusion vertex Γ_D for $YN \rightarrow D$. Our model thus contains only geometrical overlap factors (related to the r.m.s. size of Y , N and D) and colour-flavour-spin recoupling coefficients. We have neglected the dynamics of how the quark bags deform during the fusion process (we calculate the overlap of spherical bags). Ours is essentially a non-relativistic treatment, i.e. Γ_D depends only on three momenta,

and does not contain any off-shell factors depending on the energy variables. There is considerable additional model dependence which arises because of the uncertainty in the relative phase of $\bar{K}N \rightarrow \pi\Lambda$ and $\bar{K}N \rightarrow \pi\Sigma$ amplitudes, which enter coherently in the $d(K^-, \pi^-)D_t$ reaction. If future experiments provide more specific hints of the existence of $S=-1$ dibaryons, it would certainly be worthwhile to refine the theoretical calculations.

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Quark structure	Y	f(y,i)s	$2I+1, 2S+1 L_J$	δM	$\delta \Delta$
$(Q^6)_1$	1	$8(1, 1/2)_1$	$^{23}S_1$	115	$+5/3$
$(Q^4)_3 \otimes (Q^2)_3^*$	1	$3(1/3, 1/2)_0 \otimes 3^*(2/3, 0)_0$	$^{21}P_1$	55	-2
$(Q^4)_3 \otimes (Q^2)_3^*$	1	$3(1/3, 1/2)_1 \otimes 3^*(2/3, 0)_0$	$^{23}P_J$	95	$-4/3$
$(Q^6)_1$	0	$1(0, 0)_0$	$^{11}S_0$	-30	-2
$(Q^6)_1$	0	$8(0, 0)_1$ or $8(0, 1)_1$	$^{13}S_1$ or $^{33}S_1$	120	$+5/3$
$(Q^4)_3 \otimes (Q^2)_3^*$	0	$3(-2/3, 0)_0 \otimes 3^*(2/3, 0)_0$	$^{11}P_1$	65	-2
$(Q^4)_3 \otimes (Q^2)_3^*$	0	$3(1/3, 1/2)_0 \otimes 3^*(-1/3, 1/2)_0$	$^{11}P_1, ^{31}P_1$	65	-2
$(Q^4)_3 \otimes (Q^2)_3^*$	0	$3(-2/3, 0)_1 \otimes 3^*(2/3, 0)_0$	$^{13}P_J$	100	$-4/3$
$(Q^4)_3 \otimes (Q^2)_3^*$	0	$3(1/3, 1/2)_1 \otimes 3^*(-1/3, 1/2)_0$	$^{13}P_J, ^{33}P_J$	105	$-4/3$

Table 1: Strange Six-Quark Dibaryons D predicted^[7] to lie below the $\Lambda N\pi$ or $\Lambda \Lambda\pi$ Thresholds. $\delta\Delta$ refers to the decay mode $D \rightarrow (Q^3)_1 + (Q^3)_1$ into two colour singlet baryons. δM is in MeV.

Cluster	[f,s]	Dissociation Channel	$\delta \Delta$
$(Q^2)_3^*$	$[3^*, 0]$	$Q + Q$	-2
$(Q^4)_3$	$[3, 0]$	$(Q^3)_1[8, 1/2] + Q$	-2
		$(Q^3)_8[1, 1/2] + Q$	$-1/2$
		$(Q^2)_3^*[3^*, 0] + (Q^2)_3^*[3^*, 0]$	0
$(Q^4)_3$	$[3, 1]$	$(Q^3)_1[8, 1/2] + Q$	$-4/3$
		$(Q^3)_8[1, 1/2] + Q$	$+1/6$

Table 2: Dissociation Modes of Quark Clusters $(Q^2)_3^*$ and $(Q^4)_3$ used to build the Dibaryons D_s and D_t .

Dissociation Channel	$\delta\Delta(D_s)$	$\delta\Delta(D_t)$
$(Q^3)_1[8, 1/2] + (Q^3)_1[8, 1/2]$	-2	$-4/3$
$(Q^3)_8[1, 1/2] + (Q^3)_8[8, 1/2]$	-2	$-4/3$
$(Q^5)_3^*[3^*, 1/2] + Q$	-2	$-4/3$

Table 3: Quark Tunneling Modes of the $(Q^4)_3 \otimes (Q^2)_3^*$ Dibaryons D_s and D_t .

Quark Structure	$f(y, i)s$	M	$\delta\Delta$	BB Channel
$(Q^6)_1$	$10^*(2, 0)1$	2165	$10/3$	1^3S_1
$(Q^6)_1$	$27(2, 1)0$	2245	6	3^1S_0
$(Q^4)_3 \otimes (Q^2)_3^*$	$6^*(4/3, 0)1 \otimes 3^*(2/3, 0)0$	2110	$2/3$	1^3P_J (forbidden)
$(Q^4)_3 \otimes (Q^2)_3^*$	$15(4/3, 1)0 \otimes 3^*(2/3, 0)0$	2200	2	3^1P_1 (forbidden)
$(Q^4)_3 \otimes (Q^2)_3^*$	$15(4/3, 1)1 \otimes 3^*(2/3, 0)0$	2245	$8/3$	3^3P_J

Table 4: Low-lying non-strange dibaryons in the Bag Model^[7].

$\cos\theta_L$	R_1	R_2
0.95	0.04	0.16
0.85	0.12	0.17
0.75	0.07	0.08

Table 5: Cross-Section Ratios R_1 and R_2 of Eq. (4.2) for $k_L = 0.95$ GeV/c (no cut-off on recoiling nucleon momentum).

$\cos\theta_L$	$d^2\sigma/d\Omega dM^2_{YN}(k_{\min}=0)$	$k_{\min}=50\text{MeV}/c$	$k_{\min}=100\text{MeV}/c$	$k_{\min}=150\text{MeV}/c$
1.00	18.0	10.4	2.7	1.1
0.95	46.7	43.0	16.8	8.4
0.90	53.3	53.3	43.2	13.9
0.85	49.1	49.1	40.3	21.6
0.80	40.2	36.6	29.7	26.5
0.75	31.5	25.8	21.4	19.8

Table 6: Cross-Section $d^2\sigma/d\Omega/dM^2$ (in $\mu\text{b}/\text{sr}/\text{GeV}^2$) for dibaryon formation in $d(K^-, \pi^-)D_t$ at 0.95 GeV/c, as a function of angle and nucleon cut-off momentum $k_{\min}$ (in MeV/c). We take $m_D=2.14$ GeV/c² and a mass resolution of $\delta m_{YN}=10$ MeV.

$\cos\theta_L$	$R_1(k_{\min}=50\text{MeV}/c)$	$R_1(k_{\min}=100\text{MeV}/c)$	$R_1(k_{\min}=150\text{MeV}/c)$
0.95	.04	.04	.02
0.85	.12	.17	.22
0.75	.11	.24	.39

Table 7: Cross-Section Ratios R_1 of Eq.(4.2) for $k_L=0.95$ GeV/c and different values for the nucleon cut-off momentum $k_{\min}$

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FIGURE CAPTIONS

Figure 1: Predicted spectrum of strange six quark dibaryon states taken from Mulders et al.^[7]. The mass scale on the left refers to the $S=-1$ configurations, and that on the right to the $S=-2$ spectrum. All the predicted levels below the $\Lambda N\pi$ ($S=-1$) or $\Lambda\Lambda\pi$ ($S=-2$) thresholds are shown as solid lines, labelled by the quark cluster structure and the quantum numbers of the baryon-baryon channels to which they can decay. The various baryon-baryon thresholds are shown as dashed lines. Here, we focus on the production cross-sections for the lowest-lying $S=-1$ dibaryons D_s and D_t , which are expected to be long-lived^[23].

Figure 2: Production mechanisms for the $S=-1$ dibaryons D_t and D_s in the reactions (K^-, π^+) on deuterium and Helium-3 targets, respectively. The vertices Γ_d , Γ_{He} and Γ_D represent wave function overlaps and the amplitudes T are constructed from the model of Gopal et al.^[42], which fits the two-body $\bar{K}N \rightarrow \pi Y$ data.

Figure 3: Background processes of first and second order for the (K^-, π^-) reaction on deuterium, producing $\pi^- \Lambda p$ final states. These processes can produce a Λp missing mass in the same kinematical region as the reaction $d(K^-, \pi^-)D_t$ producing a dibaryon D_t . Reaction (b) can produce peaking in the Λp missing mass^[20], but with a width larger than that expected^[23] for D_t .

Figure 4: Lab differential cross-sections $d\sigma/d\Omega$ for the reaction $d(K^-, \pi^-)D_t$ for various incident K^- lab momenta (as labelled), as a function of the cosine of the π^- laboratory angle ($\cos\theta_L$). The mass of the dibaryon D_t was taken to be $2.14 \text{ GeV}/c^2$. The dashed curves correspond to the use of the $\bar{K}N \rightarrow \pi\Lambda$, $\pi\Sigma$ amplitudes of Gopal et al.^[42], while for the solid curves the relative phase of the Σ and Λ production amplitudes is changed by 180° , as discussed in the text.

Figure 5: Peak lab differential cross-sections for the process $d(K^-, \pi^-)D_t$ as a function of lab momentum k_L . The dashed curve labelled G refers to a calculation employing the $\bar{K}N \rightarrow \pi\Lambda, \pi\Sigma$ amplitudes of Gopal et al.^[42], while for the solid

curve (AG for "anti-Gopal") the relative Λ - Σ phase is arbitrarily rotated by 180° . The separate contributions of Λ and Σ intermediate states (see Fig. 2(a)) are shown as dashed lines. For the Gopal phases, the Λ - Σ interference is seen to be destructive, while the curve AG represents the maximum amount of constructive Λ - Σ interference. For curve AG, the lab angles corresponding to the peak differential cross-sections are indicated for two values of k_L .

Figure 6: Lab angular distributions for the $d(K^-, \pi^-)$ reaction at 0.95 GeV/c. The solid curves refer to the process $d(K^-, \pi^-)D_t$, for Gopal (G) or anti-Gopal (AG) phases as in Fig.5. The dashed curves correspond to the "background" processes $d(K^-, \pi^-)YN$, with no cut-off on the momentum of the nucleon in the final state. These curves have been scaled down by a factor of 10^3 .

Figure 7: Dependence of the lab cross-section for the $d(K^-, \pi^-)D_t$ reaction at 0.95 GeV/c on the mass m_D of the dibaryon. Anti-Gopal phases are used for each of the curves.

Figure 8: Missing mass spectra for the reactions $d(K^-, \pi^-)YN$ and $d(K^-, \pi^-)D_t$ at 0.95 GeV/c, for $\cos\theta_L = 0.85$. The double differential lab crosssection $d^2\sigma/d\Omega dM^2_{YN}$ is shown as a function of the hyperon-nucleon invariant mass M_{YN} . The background cross-sections for the Λp , $\Sigma^0 p$ and $\Sigma^+ n$ channels are shown separately, as well as the signal corresponding to dibaryon (D_t) formation. For the latter, we have assumed anti-Gopal (AG) phases for the $\bar{K}N \rightarrow \pi Y$ amplitudes, a mass $m_D = 2.14$ GeV/c², and spread the D_t cross-section over a mass interval of 10 MeV, to simulate the effects of experimental resolution. No cut-off on the nucleon momentum in the final state has been imposed.

Figure 9: Missing mass spectra for the reactions $d(K^-, \pi^-)YN$ at 0.95 GeV/c, for three different choices of lab angle (as labelled). The solid curves refer to the cross-section for the Λp , $\Sigma^+ n$ and $\Sigma^0 p$ channels, after imposing the requirement that the nucleon momentum exceed 150 MeV/c. The dashed curves represent calculations where no momentum cut-off is applied.

Figure 10: The lab differential cross-sections for D_s formation in the reaction ${}^3\text{He}(K^-, \pi^+)nD_s$ are shown as solid curves labelled by the K^- lab momentum. For 0.8 and 0.9 GeV/c, the background cross-sections for the process ${}^3\text{He}(K^-, \pi^+)n\Sigma^-p$ are shown as dashed lines; these cross-sections have been multiplied by 10^{-3} .

Figure 11: Missing mass spectra for the reactions ${}^3\text{He}(K^-, \pi^+)nX$ for incident K^- momentum 0.8 GeV/c and angle $\theta_L \approx 26^\circ$. For the curve labelled $X=D_s$, we have chosen a mass $m_D = 2.1 \text{ GeV}/c^2$. Since $m_D < m_\Sigma + m_N$, the D_s peak is separated from the background curve with $X=\Sigma^-p$. M_{nX} corresponds to the invariant mass of the $n + X$ system.

Figure 12: Missing mass plots for the process ${}^3\text{He}(K^-, \pi^+)nX$ at 0.8 GeV/c. The solid curves for $X = \Sigma^-p$ are labelled by the cosine of the pion lab angle, while the shaded rectangles correspond to $X = D_s$ with $m_D = 2.06$ and $2.10 \text{ GeV}/c^2$. The rectangles represent the cross-sections for $\cos\theta_L = 0.95$, while the arrows indicate the cross-sections for $\cos\theta_L = 0.85$ and 0.75 . In contrast to Fig. 11, M_X does not include the neutron mass.

Figure 2

Figure 3

Figure 4

Figure 5

Figure 6

Figure 7

Figure 8

Figure 9

Figure 10

Figure 11

Figure 12