S=-1 DIBARYON FORMATION IN THE Σ^- - D ATOMA.T.M. Aerts⁺⁾

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A B S T R A C T

An estimate of the formation rate of the strangeness S=-1 dibaryons D_s and D_t via " Σ^- -capture" in a Σ^- -d atom is presented. Reasonable branching ratios are expected for formation from the atomic P orbitals. The Σ^- -d atom experiment is found to be sensitive to the formation of D_s dibaryons in a mass range of only 10 MeV starting at the Σ^- -p mass, while for the D_t dibaryons a range of 30 MeV is expected.

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1 INTRODUCTION

The mass spectrum of the six-quark (Q^6) system has been studied extensively in search for stable particles^[1] in a variety of models, such as the MIT bag model^[2-6] and the Skyrmin model^[7,8]. An important result from such studies^[4] is that ground state dibaryons containing strange quarks can be more stable than dibaryons made up from only nonstrange quarks. This is a consequence^[3] of the assumption that the hyperfine splitting of the hadronic mass spectra is dominated by colour magnetic forces, which are strongly attractive in systems with small spin, colour and flavour quantum numbers^[4]. The attraction is maximal for the dibaryon H (assuming only u,d and s quarks^[9]) which has spin-parity $J^\pi=0^+$ and is a flavour singlet (isospin $I=0$, strangeness $S=-2$). These quantum numbers clearly cannot be realized by a six-quark system containing only nonstrange quarks. The H is unique in that it is predicted to be absolutely stable having a mass approximately 60 MeV below the lowest strong decay threshold^[4-6]. (For a discussion of more recent mass estimates for the H see Refs.10 and 11.)

From the phenomenological analysis of Ref.12 it appears that there may also be relatively stable dibaryons with strangeness $S=-1$. These states, however, are not dibaryons in their ground state. In order to achieve strong enough colour magnetic attraction one has to split the Q^6 system up in smaller (albeit coloured) clusters, such as in a Q^4 - Q^2 configuration. One way to do this is by introducing an angular momentum barrier between the clusters^[13]. The colour-magnetic attraction gained this way may in some cases even compensate the energy it costs to orbitally excite the system. Examples of such a system are two states, a spin-singlet dibaryon D_s and a spin-triplet dibaryon D_t , both with quark-structure Q^4 - Q^2 and orbital angular momentum $L=1$. These states have recently been the subject of both experimental^[14,15] and theoretical^[16,17] attention. They are predicted^[12] to lie 40 MeV apart just below (D_s at 2.11 GeV/c²) and above (D_t at 2.15 GeV/c²) the ΣN threshold. There have been several^[18] experiments which indicate structure in the Λp missing mass spectrum in this mass range for the reaction $d(K^-, \pi^-)\Lambda p$ at finite pion angles (i.e. distinct from the ΣN threshold cusp^[19-21] which shows up at forward pion angles). A theoretical analysis^[17] of the decay widths of D_s and D_t confirms that these states can indeed be sufficiently stable ($\Gamma < 3$ MeV) to cause this finite angle structure. (Various other explanations for this phenomenon have been discussed in Refs.22 and 23.)

Assuming then that there are indeed relatively stable dibaryons with strangeness $S=-1$, one should ask, how to produce them? One way is to use the (K,π) strangeness exchange reaction on a suitable target and optimize the kinematics of the reaction for the dibaryon production process^[16]. It then appears that the reaction :

$$K^- + d \rightarrow \pi^- + D \quad (1.1)$$

is well suited for the production of the spin-triplet dibaryon. This follows from the observations that the deuteron is also a spin-triplet state and that the non-spin flip part of the elementary reaction $KN \rightarrow \pi Y$ driving the dibaryon production dominates the spin flip part by an order of magnitude^[24]. Therefore, it was suggested in Ref.16 that the reaction

$$K^- + {}^3\text{He} \rightarrow \pi^+ + D_s + n \quad (1.2)$$

where the target nucleus contains a 1S_0 proton pair, might be preferable to one on deuterium for D_s production. In order to isolate the D_s as the "missing mass", in which variable it would show up as a sharp peak^[17], one needs to perform a π^+-n coincidence experiment with very good energy resolution. This, however, is not very easy, considering the low intensity of the presently available kaon beams. Also, if one were to detect only the neutron, a "signal to noise" ratio of about only 1/200 is expected under experimental circumstances which are optimized for the D_s formation.

In this paper we propose an alternative to the processes of Eqs.(1.1) and (1.2). We shall study a two-step process for producing low-lying $S=-1$ dibaryons. In the first step a Σ^- is produced by stopping kaons in LH_2 , through the reaction:

$$K^- + p \rightarrow \Sigma^- + \pi^+ \quad (1.3)$$

which has a large branching ratio ($\approx 44\%$). One detects the (222 MeV/c) π^+ , which serves as a tag for the Σ^- production. The Σ^- then is slowed down and stopped in a LD_2 target where a Σ^- -d atom is formed. The dibaryon, finally, is created via " Σ^- capture" :

$$(\Sigma^- - d)_{\text{Atom}} \rightarrow D + n \quad (1.4)$$

Note that a similar approach for the formation of the H dibaryon using the Ξ^- -d atom, has been studied theoretically^[25,26] and that the experiment has been approved^[27].

The two-step process of Eqs.(1.3) and (1.4) can be used for the formation of dibaryons as far as their mass lies below the Σ^-p threshold. This is expected to be the case for the D_s [12], but could also be true for one or more of the triplet dibaryon states (total spin $J=0,1$ or 2) in case of sufficiently strong spin-orbit splittings.

We note that both stages (1.3) and (1.4) involve only simple two-body kinematics. This situation has to be contrasted with that of the process of Eq.(1.2), where a π^+-n coincidence measurement of sufficient resolution has to be performed to construct the dibaryon signal. The exotic atom approach should thus be more attractive experimentally.

A disadvantage of the two-step process is the loss of Σ^- 's to weak decay during the time when they are slowed down and bound in the atom. Repeating the analysis of Ref.26 to estimate the loss of Ξ^- 's due to weak decay for the case of the Σ^- , we estimate a time $t_{\text{cap}} \approx 4 \cdot 10^{-10}$ sec. for a Σ^- with an initial momentum of 173 MeV/c to be captured in an atomic orbital in a liquid hydrogen target at 20° K and one atm pressure. This is a few times the Σ^- lifetime ($\tau_{\Sigma^-} = 1.48 \cdot 10^{-10}$ sec. [28]) and in an LH_2 moderator only one in about 15 Σ^- 's would reach the strong interaction stage (other times involved, such as the atomic deexcitation time t_{Aug} , are short compared to t_{cap}). However, this number can be increased easily by a factor of two or more by letting the Σ^- 's, once produced in LH_2 , slow down in the liquid deuterium target. The number of Σ^- 's available for study of the Σ^-d atom would then be quite acceptable (see also Ref. 29). (Note, that by producing the Σ^- 's directly in the LD_2 target one may have more Σ^- 's available for strong interactions, but one would not have the simple two-body kinematics and monoenergetic pion tag.)

It, therefore, depends on the rate of dibaryon formation as compared to the background processes:

whether the two-step process is preferable to the one-step process. This branching ratio R when expressed as the ratio:

$$R = \frac{\Gamma_{\Sigma^-d \rightarrow Dn}}{\Gamma_{\Sigma^-d \rightarrow Dn} + \Gamma_{\Sigma^-d \rightarrow \Lambda nn} + \Gamma_{\Sigma^-d \rightarrow \Sigma^0 nn}} \tag{1.6}$$

is a function of the atomic orbital angular momentum ℓ , but it is practically independent of the principal quantum number n . When R is larger than 0.1 for formation from atomic S orbitals, the two-step experiment is quite feasible, and could even be more straightforward than that of Ref.27. For $0.1 > R > 0.01$, the exotic atom experiment still has a better "signal to noise" ratio for the formation of the singlet dibaryon and thus is preferred to the simplest version for the process of Eq.(1.2). The experiment will not be very easy, though. For $R < 0.01$, we regard the experiment as hard under the present experimental circumstances. As will be discussed at some length in Section 3, we expect a branching ratio $R \ll 0.01$ for the formation of the D_s dibaryon from S orbitals. However, we find $R \approx 0.01$ for formation from a P orbital. Although strong interaction processes take dominantly place from the S orbitals, it is the better branching ratio that makes forming the $L=1$ dibaryons from the P orbitals the most attractive experimental approach. Such an experiment we would consider competitive with the coincidence experiment of Eq.(1.2).

The experimental situation for forming triplet dibaryons appears to be easier than it does for the singlet one. These states could even be formed with reasonable branching ratio's from the S orbitals, if they have a sufficiently low mass.

Given the prediction from the masses of D_s and D_t , and the expected suppression^[24] of the spin flip amplitudes in the reaction of Eq.(1.1), the process of Eq.(1.4) could serve as an independent check for establishing the dibaryon character of some of the signals found in Ref.18.

The paper is organized as follows. In Section 2 we give the models for dibaryon formation and the calculation of the background processes of Eq.(1.5). In Section 3 we give our results for the rates and in Section 4 we present some concluding remarks.

2 D_S AND D_T FORMATION AND BACKGROUND CALCULATION

The formation and background rates are calculated in the impulse approximation. The atom is taken to dissociate into a Σ^- and a deuteron. The deuteron in turn dissociates into a (spectator) neutron, which is detected later on, and a proton. In the process of Fig. 1a, displaying the mechanism for dibaryon formation, the Σ^- and the proton interact to form a hyperon-nucleon pair, which fuses to form the

dibaryon. In Fig. 1b, showing the mechanism for the "quasi-elastic" or "background" process, the Σ^- and the proton interact to form a hyperon and a neutron, thus leading to a continuum of Λnn and $\Sigma^0 nn$ states. All particles in Fig. 1, except those occurring in the final state, are in principle off shell.

Many of the details of the calculations we present here have been discussed at some length before^[16,17,26] and where appropriate we shall refer to the previous work. The main ingredients of the calculations are:

- (i) a model for the dissociation vertex Γ_d of the deuteron into two nucleons
- (ii) a model for the fusion vertex Γ_D of the hyperon and the nucleon into the dibaryon.
- (iii) a model for the transition amplitudes T .

Relativistic kinematics is used throughout, but the models used in the calculation are essentially nonrelativistic. The atomic wave functions $\psi_{n\ell}$ are expanded in the variable $(r/a)^{[26]}$, where $a=37$ fm is the Bohr radius and r the Σ^- -d separation, and only the lowest order term is kept, i.e.,

$$\psi_{n\ell m}(\vec{r}) = c_{n\ell} Y_{\ell m}(\hat{r}) r^\ell \quad (2.1)$$

We imagine the coupling of the baryon number $B=2$ systems (d, D) to the two baryons in Figs. 1a and 1b to be a virtual dissociation process. We furthermore assume that the two baryons can be described by plane waves. The vertex functions Γ then are given by the expansion of the wave functions of the $B=2$ systems into plane waves.

For the $d \rightarrow np$ decomposition we have adopted the form^[16]:

$$|d, \vec{p}_d, s_d\rangle = C_S(s_1, s_2, 1) C_I(t_1, t_2, 0) * \int \frac{d^3 q}{(2\pi)^3} \frac{\sqrt{2m_d}}{2E_N(q)} \phi_d(\vec{q} - \vec{p}_d/2) |\vec{q}, s_1, t_1\rangle |\vec{p}_d - \vec{q}, s_2, t_2\rangle \quad (2.2)$$

where the C_S couples the nucleon spins s_i to $S=1$ and the C_I the isospins t_i to $I=0$. The normalization is:

$$\langle d, \vec{p}_{d'}, s_{d'} | d, \vec{p}_d, s_d \rangle = (2\pi)^3 (2E_d) \delta^{(3)}(\vec{p}_{d'} - \vec{p}_d) \delta_{s_{d'} s_d} \quad (2.3)$$

In Eqs.(2.2) and (2.3) E stands for the energy $E=(m^2+p^2)^{1/2}$ and ϕ_d is the relative wave function, which we take to be the Hulthén form:

$$\phi_d(\vec{k}) = C_0[(\beta^2 + \vec{k}^2)^{-1} - (\gamma^2 + \vec{k}^2)^{-1}] \quad (2.4)$$

with $\beta^2 = 2.096 \cdot 10^{-3} \text{ GeV}^2$ and $\gamma = 7\beta$.

To describe the $YN \rightarrow D$ fusion vertex we use the model of Ref.16. The dibaryon consists of a Q^4 and a Q^2 cluster separated by an angular momentum barrier. The transition of the two baryon state to the dibaryon state is considered a quark-rearrangement process:

$$Q^3 + Q^3 \rightarrow Q^4 \otimes Q^2 \quad (2.5)$$

which leads to :

$$\begin{aligned} \langle \vec{k}_Y, s_Y, \vec{k}_N, s_N | D, J_D, s_D, M_D \rangle = \\ = (2E_D 2E_Y 2E_N)^{1/2} (2\pi)^3 \delta^{(3)}(\vec{k}_D - \vec{k}_Y - \vec{k}_N) C(YN, D) \gamma_D(\vec{k}_Y - \vec{k}_N) \end{aligned} \quad (2.6)$$

with an overlap function γ_D of the form:

$$\gamma_D(\vec{k}) = N_D k Y_{1m}(\hat{k}) e^{-(1/12)R^2 k^2} \quad (2.7)$$

Here $R=0.92 \text{ fm}$ represents an effective radius. In Eq.(2.6), $C(YN, D)$ is a (re)coupling coefficient for the space, spin, colour and flavour parts of the dibaryon wave function in a basis of $Q^3 \otimes Q^3$ states. For D_s and D_t we find^[17] (using de Swart's conventions^[30])

$$\begin{aligned} |D_s\rangle = \sum_{i,j} C(ij, D_s) |ij\rangle = \\ (1/16\sqrt{3}) \{ 3(n\Lambda - \Lambda n) + \sqrt{6}(p\Sigma^- - \Sigma^-p) - \sqrt{3}(n\Sigma^0 - \Sigma^0 n) \} + \dots \end{aligned} \quad (2.8)$$

and

$$\begin{aligned} |D_t\rangle = \sum_{i,j} C(ij, D_t) |ij\rangle = \\ -(1/16\sqrt{3}) \{ -(n\Lambda + \Lambda n) + \sqrt{6}(p\Sigma^+ - \Sigma^+p) - \sqrt{3}(n\Sigma^0 + \Sigma^0 n) \} + \dots \end{aligned} \quad (2.9)$$

where we have listed the energetically available channels for the decay of $D_{s,t}$ only. The dots in Eqs.(2.8) and (2.9) refer to configurations involving orbitally excited baryons, colour octet Q^3 clusters, etc.

Finally, the operators T , which drive the dibaryon formation (Fig.1a) are approximated by static transition potentials. These hyperon-nucleon potentials are taken from Refs.20 (Model D) and 21 (Model F). Although these two models describe the same YN data with comparable accuracy, they differ considerably in details. The potentials of model D and F are local potentials, due to one-boson exchanges. In model D, all members of the pseudoscalar and vector meson nonets and the scalar meson ϵ , taken to be a unitary singlet, are exchanged. Model F includes a full nonet of scalar mesons. The YN coupling constants in model D are related via SU(3) symmetry to those of the NN system. In model F, SU(3) relations are assumed for the axial-vector coupling of the pseudoscalar mesons, the electric and magnetic coupling constants of the vector mesons, the direct couplings of the scalar mesons, and the hard core radii, which cut off the potentials at short distances. In model D, the hard core radii are adjusted per partial wave (actually only four different ones are needed). These differences are reflected in the description of the individual YN partial waves. For instance, in model F a resonance in the 3S_1 - 3D_1 wave is predicted close to the ΣN threshold, whereas in model D the 3S_1 - 3D_1 wave is also quite strong at this threshold, but not resonant. (The S-waves are the main ingredient for our background calculation.) Also, the potentials in the P-waves, which we need for the calculation of the dibaryon formation rate, are quite different and this fact introduces a considerable model dependence in our results. Detailed comparisons will be made in the next section.

The transition matrix element T in Fig.1b is modeled differently. Since this matrix element receives its most important contributions from the YN 1S_0 and 3S_1 - 3D_1 waves, which are strong and in one of the cases considered here even resonant near the ΣN threshold, it is not adequate to replace here T by V . Therefore, we use a K-matrix parametrization of T with the K-matrix given in terms of an effective range expansion at the Σ^-p threshold^[31] This procedure should yield a good description of the Σ^-p to Σ^0n and Λn transitions, especially just below the Σ^-p threshold. Since the effective range expansion only has been given for model F, we have just this model for the reaction mechanism of Fig.1b.

We are now in a position to write down expressions for the widths. The width for the process of Eq.(1.4) (Fig. 1a, two-body final state) is given by:

$$\Gamma(2) = \frac{E_D(p_n)}{2E_Y(p_n/2)2E_N(p_n/2)m_\Sigma} \frac{m_D^2 c_n^2}{24} \frac{m_Y^2 m_N^2}{(m_Y + m_N)^2} * \\ * \sum |G_\ell(L)|^2 p_n / v_s \quad (2.10)$$

where D, Y and N refer to the dibaryon, hyperon and nucleon, respectively and p_n is the spectator neutron momentum. The average hyperon mass (occurring in the potential V_{NY}) is $m_Y = 1.174 \text{ GeV}/c^2$. The matrix element $G_\ell(L)$ is defined as

$$G_\ell(L) = \sum \frac{\{(2\ell+1)2\ell!\}^{1/2}}{\{2\lambda!2(\ell-\lambda)!\}^{1/2}} g^{\lambda'}(p_n) f^{\ell-\lambda}(p_n) (2\lambda'+1)^{1/2} * \\ * W(1, \lambda', \ell, \ell-\lambda; \lambda L) C_0^{\ell-\lambda} \begin{matrix} \lambda' & L \\ 0 & 0 \end{matrix} \quad (2.11)$$

with $\lambda' = \lambda \pm 1$. W is the Racah recoupling coefficient^[32]. Here

$$g^{\lambda'}(p_n) = N_g(\lambda) \int dr r^{3+\lambda} j_{\lambda'}(p_n r/2) V_{NY}(r) e^{-3r^2/4R^2} \quad (2.12)$$

involves the overlap between the dibaryon wave function and the potential V_{NY} . And

$$f^\lambda(p_n) = N_f(\lambda) \int dr r^{\lambda+1} j_\lambda(p_n r) \phi_d(r) \quad (2.13)$$

contains the deuteron wave function. $N_g(\lambda)$ and $N_f(\lambda)$ are constants containing the wave function normalizations and depending on λ . The integrations in Eqs.(2.12) and (2.13) extend from 0 to infinity, but in Eq.(2.12) the integrand is only nonzero for values of $r > x_c$, the hard core radius of the potential V_{NY} .

In order to arrive at the simple result of Eq.(2.10) we have assumed that the total spin of the atomic system is a good quantum number^[33]. We then assumed a statistical population of the atomic $S=1/2$ and $S=3/2$ states and summed over all final state polarizations and averaged over the initial state polarizations. Furthermore, we have neglected the rather weak Σ^- -momentum dependence of the baryon energies in the intermediate state to simplify the calculations.

The corresponding expression for the background process (Fig. 1b, three-body final state) is

$$\Gamma(3) = \frac{C_{n\ell}^2 N_f^2 \int ds_d s_d \lambda(s, s_d, m_n^2)^{1/2}}{8m_\Sigma M_{n\ell}^2 2\pi^2 \lambda(s_d, m_\Sigma^2, m_p^2)^{1/2}} |f^\ell(p_n)|^2 \quad * \quad (2.14)$$

$$* \left[\frac{1}{4} |T(^1S_0)|^2 + \frac{3}{4} \{ |T(^3S_1 \rightarrow ^3S_1)|^2 + |T(^3S_1 \rightarrow ^3D_1)|^2 \} \right]$$

Here $\lambda(a,b,c)=a^2+b^2+c^2-2(ab+bc+ca)$. $M_{n\ell}$ is the mass of the Σ^-d atom. The transition matrix elements T are evaluated at the YN invariant mass $\sqrt{s_d}$. Eq.(2.14) also has been obtained by averaging over the initial state polarizations and summing over the final state polarizations. The neutron wave functions have been taken to be plane waves and a statistical factor of $1/2$ has been included in the phase space.

Eq.(2.14) clearly reflects that we have only considered incoming S-waves for the T matrix elements. The orbital angular momentum of the atom has thus to be absorbed by the spectator neutron, which gives rise to the ℓ 'th "deuteron form factor" $|f^\ell(p_n)|^2$.

3 RESULTS

Our results for the absolute rates for the background processes are given in Table 1. Since the dibaryon formation process never is found to dominate the background process but at best is competing with it, Table 1, combined with our results for the relative rate R displayed in Figs. 2, is sufficient to calculate all rates of interest.

Table 1 shows that the decays of the atomic S and P states are dominated by strong interaction processes, and that the decay of the D states is mainly radiative. (Remember that the dependence on the atomic principal quantum number n is contained in the constant $c_{n\ell}$ which is common to all rates.) We further see that the strong decay rates for the D states compare unfavourably also with the weak decay rate of the Σ^- 's. We will, therefore, only include the D state results for R in our discussion when they are of theoretical interest.

By far the most strong interaction processes will take place in the higher S orbitals, which are continuously populated through the Stark mixing due to the collisions of the Σ^-d atom with the deuterium molecules^[34,26]. The population of the P orbitals can occur both through Stark mixing and through radiative transitions from the D orbitals. The X-ray accompanying the latter (E_1) transitions is interesting in two ways. It gives information about the strong

interaction width and shift of the nP level. This information is the usual motivation for detecting these X-rays. We expect the width and shifts to be rather small ($\Gamma \lesssim 0.1$ eV, see Table 1). Secondly, the X-ray can be used as a trigger to signal the population of the P-levels^[35], from which then the strong decay is observed which has a branching ratio of more than 99 %. This can be done by looking for the spectator neutron.

Note the the branching ratio r :

$$r = N(\Sigma^0) / [N(\Sigma^0) + N(\Lambda)] \quad (3.1)$$

where $N(Y^0)$ is the number of final state hyperons Y^0 , is found to be $r=0.053$, in agreement with the experimental result $r_b=0.037 \pm 0.022$ ^[36].

Our results for the branching ratio R of Eq.(1.5) are displayed in Figs. 2a-2d for the $D_s(^1P_1)$ and the $D_t(^3P_0, ^3P_1 \text{ and } ^3P_2)$ states. We have added to each graph a horizontal line to indicate the 1 % level for R . When R falls below this level, the two-target experiment no longer is preferred to the one-target experiment of Eq.(1.2) on the basis of the "signal to noise" ratio alone, since this ratio would be comparable for the two experiments in their simplest form.

Figs. 2a-2d have several features in common. First of all, the ratio R is largest (within a given model for V) for the $L=1$ atomic states. This is not surprising, since in this state the orbital angular momentum in the atomic initial state and the final state containing the dibaryon are best matched, whereas the matching of the orbital angular momenta is less favourable for the background process. The peak value of R for the S wave is seen to be at least a factor of 20 smaller than that for the P wave, whereas the peak value of R for the D waves is actually quite close to the P wave value. This phenomenon is associated with the large strength of the background processes in the YN S waves, which comes out the best in the atomic S orbital case.

Secondly, R decreases sharply when the dibaryon mass $m_D \rightarrow M_{n\bar{p}}$, which is a kinematical effect. R also decreases when m_D decreases. In this case p_n , the spectator neutron momentum, becomes large and R reflects the behaviour of the overlap between the deuteron wave function and the atomic wave function as a function of p_n .

All curves are seen to peak around $m_D \approx 2.13 \text{ GeV}/c^2$. This implies that an exotic atom experiment will be most sensitive to the formation of dibaryons with masses in the range $M_{n\ell} > m_D > M_{n\ell} - 30 \text{ MeV}$. The peak region for dibaryon formation falls in the spectator momentum range $30 \text{ MeV}/c < p_n < 70 \text{ MeV}/c$ (see Table 2). Fig. 3 shows that this range of neutron momenta also covers the region where the background neutron spectrum peaks. Fortunately, the $L=1$ dibaryons are expected to be narrow^[17] ($\Gamma < 3 \text{ MeV}$), and good momentum resolution is expected to allow one to resolve the dibaryon signal from the background.

Finally, we see that the predictions for R from models F and D differ generally by a factor of about 10, with the model F result the larger one in the states with $J=1$ or 2 and that of model D being larger for the $J=0$ state. Such large differences are not found between the individual channels in models D^[20] and F^[21]. They are a consequence of our model for dibaryon formation, in which the amplitudes for formation in the channels Λn , $\Sigma^0 n$ and $\Sigma^- p$ (the latter two of which are closed for most dibaryon masses we consider here) are added coherently according to Eqs. (2.8) and (2.9).

There are, as discussed in Ref.37, several reasons for preferring model D over model F. For instance, when one computes the effective Ξ^- -nucleus potential, model F leads to a repulsive potential, whereas model D gives an attractive one, with a well depth of the right strength to explain the existence of several, relatively long-lived Ξ^- -nucleus states. In view of this, one should attach more weight to our model D results.

We will now discuss the individual waves. The singlet dibaryon $\sim ({}^1P_1)$ which is not expected to be easily formed in the one-step processes of Eqs.(1.1) and (1.2), also has a small branching ratio for formation from an atomic orbit. The model D, $L=1$ prediction for the peak value of the relative rate R is about 3 %, while the model F result is roughly 1 %. This result is due to the destructive interference of the ΛN and ΣN contributions to the formation amplitude. Omitting the ΣN channels, the model D prediction goes up to : $R=0.06$ (0.30, 0.24) for formation from an atomic S (P, D) state. A similar result holds for model F (0.12, 0.45, and 0.38 for the S, P, and D states respectively). This implies that, were we to use a model for T based on the physical YN phase shifts (i.e., only contributions from the ΛN channel below the ΣN threshold), we would find that the singlet dibaryon would be produced copiously, even from the atomic S wave states. Using our present model^[10], which allows both the ΛN and ΣN channels to couple

to the dibaryon, the 1P_1 state will hardly be seen in a Σ^- -d atom experiment, if only the neutron is detected and no special arrangements are made. This point is illustrated for the atomic 1S state in Fig. 3, where we plot the neutron spectrum coming from the background channels Λnn and $\Sigma^0 nn$, and the signal due to a monoenergetic neutron accompanying a spin-singlet dibaryon with a mass $m_D = 2.10 \text{ GeV}/c^2$ (i.e., 40 MeV below the "triplet-signature" at $2.14 \text{ GeV}/c^2$.) The neutron momentum has been spread out by an amount $\Delta k/k = \pm 5 \%$ to simulate a typical momentum resolution. Even if we use the larger model F prediction, the neutron peak has to be multiplied by a factor 100 to make it stick out well from the background.

The situation of Fig.3 can be greatly improved, however, by studying dibaryon formation from atomic P orbitals. In order to obtain approximately the neutron spectrum for the 2P state from that of the 1S state, one has to divide the units in Fig. 3 by another factor $3 \cdot 10^4$. The factor multiplying the dibaryon signal now becomes 25. When we decrease the neutron momentum to $\sim 50 \text{ MeV}/c$, this factor decreases to 1! The selection of P orbitals can be done by requiring X-rays from the atomic D to P transitions to precede the neutron detection. Observing X-rays is common practice in exotic atom experiments and for the $3D \rightarrow 2P$ (and even the $2P \rightarrow 1S$) transitions in the Σ^- -d atom reasonable yields are expected^[29]. By increasing the resolution of the spectator neutron momentum (the dibaryons are expected to be narrow^[17]) and selecting the atomic doublet states as initial states, one can further raise the signal to background ratio to an acceptable level of (in the case of model D) about a few per cent around the peak. The mass range in which the 1P_1 dibaryon can be detected at present probably will stay limited to roughly the first 10 MeV below the Σ^-p threshold.

It will be considerably easier to detect the spin-triplet dibaryons, if their mass lies not more than 30 MeV below the mass of the Σ^- -d atom. In model F the 3P_1 and 3P_2 states, and in model D the 3P_0 state could even be detected from the atomic S orbitals. In all cases, selection of the $L=1$ orbitals as initial state will ensure reasonable branching ratios of the order of 10 % or more, and, with good momentum resolution for the spectator neutron, the dibaryon signal should stand out well from the background.

Omitting the ΣN channels from the dibaryon formation amplitude reduces the branching ratios in model D for the spin-triplet dibaryons by factors of 2 to 4, but does not have the same dramatic effect it has

on the 1P_1 dibaryon branching ratio. Similar results are found for model F.

4 CONCLUSIONS AND COMMENTS

In the two previous sections, we presented our model for dibaryon formation from the Σ^- -d atom, and calculated its predictions for the formation rates for the D_s and D_t dibaryons. The spin-singlet level is predicted to have a mass of $2.11 \text{ GeV}/c^2$ and the spin-triplet level is expected to lie around $2.15 \text{ GeV}/c^2$ ^[12]. Accounting for the experimental structure at $2.14 \text{ GeV}/c^2$, these masses could change to 2.14 and $2.10 \text{ GeV}/c^2$, respectively. Due to residual spin-dependent interactions, the triplet level presumably will be split up into three states, with $J=0, 1$ and 2 , respectively, some of which may lie below the Σ^-p threshold. We found that the Σ^- -d atom experiment at present would be sensitive to the formation of dibaryons in a mass range of about 30 MeV below the atom's mass for the spin-triplet dibaryons and of perhaps 10 MeV for the spin-singlet dibaryon, provided formation from the atomic P orbitals is studied. In view of the predicted mass of the D_s , this mass region may be too small. It is found that good momentum resolution for the spectator neutron is quite important in improving the signal to background ratio.

Our model neglects off-shell effects in the formation amplitude. If such effects lead to the suppression of the ΣN contribution with respect to the ΛN contribution to the dibaryon formation amplitude, the prospects for detecting the spin-singlet dibaryon improve, while those for the spin-triplet dibaryons become slightly worse. Also, we have assumed that^[17] the $L=1$ $Q^4 \otimes Q^2$ dibaryon can be considered as a superposition of two Q^3 clusters. This would probably be correct when we could consider a zero-range approximation for the interaction matrix element T . In reality, the model for T is a finite range potential, and, therefore, the weight of the ΛN and ΣN components in the dibaryon wave function will presumably be larger than given in Eqs.(2.8) and (2.9). (The overall weight of the coloured and other energetically not allowed components will decrease when one takes the overlap of the dibaryon with two Q^3 clusters at relative separation r .) This effect will increase R for both the singlet and triplet dibaryons. Our results thus represent a conservative estimate of R , in particular for the case of the singlet dibaryon.

From the point of view of the signal to background ratio, the Σ^- -d atom experiment is in a large number of situations preferred to the version of experiment of Eq.(1.2) where only the neutron is detected. Comparing its chance of success to that for the coincidence measurement version of the one-step experiment is much harder. The experiment of Eq.(1.2) has as an advantage that it is sensitive over a much larger mass range than the exotic atom experiment. On the other hand, the Σ^- -d atom experiment, even with the X-ray tag, appears to be simpler. Since the two- and the one-step experiments involve different reaction dynamics (YN vs. (K, π)), they will yield complementary information. Both experiments are, therefore, well worth considering.

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Tabel 1. Decay widths Γ (in eV) of the $n=1,2,3$ levels of the Σ^- -d atom to background modes. For the radiative decays only the E_1 transition has been considered. $\Gamma_{\text{weak}}=4.4 \cdot 10^{-6}$ eV.

Level (n ℓ)	Λnn	$\Sigma^0 nn$	$\gamma \Sigma^- d$
1S	2900	155	-
2S	360	19	-
3S	27	6	$5.9 \cdot 10^{-6}$
2P	$87 \cdot 10^{-3}$	$4.1 \cdot 10^{-3}$	$.59 \cdot 10^{-3}$
3P	$31 \cdot 10^{-3}$	$1.4 \cdot 10^{-3}$	$.16 \cdot 10^{-3}$
3D	$2 \cdot 10^{-6}$	$.09 \cdot 10^{-6}$	$60. \cdot 10^{-6}$

Tabel 2. Relation between the spectator neutron momentum p_n (in MeV/c) and the mass of the recoiling YN or dibaryon system (in GeV/c²). The thresholds are at $m=2.1357$ GeV/c² (Σ^-p), 2.1324 GeV/c² ($\Sigma^0 n$) and 2.0552 GeV/c² (Λn). The maximal neutron momentum for a background event is $p_n=321$ MeV/c.

p_n	m_{NY}	p_n	m_{NY}
30	2.1327	150	2.1161
50	2.1314	200	2.1028
70	2.1296	250	2.0857
100	2.1257	300	2.0649

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FIGURE CAPTIONS

Figure 1: In (a) we display the mechanism for dibaryon formation in the process $\Sigma^- + d \rightarrow D + n$. The hyperon-nucleon intermediate states which couple to the dibaryons are indicated. In (b) we show the mechanism for the 'background' processes $\Sigma^- d \rightarrow \Lambda nn, \Sigma^0 nn$. Note that $\Sigma^- pn$ (and for small enough dibaryon mass also $\Sigma^0 nn$) appears as intermediate state in (a), although it would energetically not be allowed as a final state in (b). Γ indicates a fusion or dissociation vertex, T an interaction.

Figure 2: The branching ratio R for formation of the D_s (a) and D_t ($J=0$ (b), $J=1$ (c) and $J=2$ (d)) dibaryons, as a function of the dibaryon mass. Results for models D (Ref.20, dashed curve) and F (Ref.21, solid curve) are shown. The horizontal line indicates the 1 % level below which the $\Sigma^- d$ atom experiment becomes unattractive.

Figure 3: The neutron momentum spectrum $d\Gamma/dp$ as a function of p , arising from the reactions $\Sigma^- d \rightarrow \Lambda nn, \Sigma^0 nn, Dn$ via the $1S$ state of the $\Sigma^- d$ atom. The signal from the monoenergetic spectator neutron from the reaction $\Sigma^- d \rightarrow Dn$ (calculated in model F) has been spread out by an amount $\Delta p/p = \pm 5\%$ to simulate a typical experimental resolution. It has also been multiplied by a factor of 100 to make it visible.

(a)

(b)

- Figure 1 -

- Figure 2a -

- Figure 2b -

- Figure 2c -

- Figure 2d -

- Figure 3 -