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1. INTRODUCTION

The standard model of strong interactions, the QCD, requires that the low-lying baryons and mesons are composites of three quarks (q^3) or a quark and an antiquark $(q\bar{q})$, respectively. However, few observables can actually be calculated at low energies in the framework of QCD. In particular the theory of the quark confinement mechanism is inexact. Lacking a more complete theory, phenomenological quark models such as the MIT bag model[1] provide an explanation of the hadron mass spectrum. With a limited number of free parameters the low-lying masses are calculated using a simple model for the interactions of confined quark triplets and quark-antiquark pairs.

In addition to fitting the known mass states, the bag models predict the existence of other more massive states of structure $q^m\bar{q}^n$, with $m+n\geq4$. As example, we mention the $q^2\bar{q}^2$ baryonium configurations, the Z^* resonances $(q^4\bar{q})$ of strangeness $S=+1$, and the dibaryons (q^6). There is as yet no conclusive experimental evidence for long-lived objects of these structures. Some candidates, in turn, focus the attention, in particular hunting the nonstrange dibaryons, on which a great deal of effort has been dedicated, since these are the most accessible experimentally.

The problem in interpreting such structures as six quark states is that each lies well above the threshold for decay into two baryons, making the interpretation of the results more difficult. For instance, the lightest nonstrange six-quark particle in the MIT model has the same spin and isospin quantum numbers of the deuteron but is 300 MeV more massive. In this case the dibaryon would be instable against deformations and the mass calculation may simply reflect the repulsive core of the short range n-p interaction. Moreover, it has proven difficult in many cases to disentangle quark aspects from ordinary hadronic mechanism, when the effects due to the latter are strong. It is, for example, the case of the coupling of NN to $N\Delta$ channel, which does not involve explicit consideration of quark degrees of freedom.

The situation turns out significantly different if we consider a six-quark system containing one strange quark. On the basis of rather general arguments, it can be shown[2] how the addition of strange quarks to a system containing up and down quarks in certain cases is energetically favorable, leading to the possibility of stable or quasi-stable dibaryons. The lightest $S=-1$ dibaryon resonance is predicted to be only 100 MeV above the ΛN mass. This increased binding is due to the form of the color-magnetic interaction. Although the strength of this term is model dependent, certain features of the level ordering appear to be more general: the maximum amount of color-magnetic attraction is generated by flavor representations of minimum dimensionality. A tendency for larger binding with higher symmetry is in fact a property of QCD.

Additional symmetry is gained - as it has been firstly pointed out by Jaffe[3] - by going to the strangeness $S=-2$ system. This gives enough color-magnetic attraction to produce a six-quark object which is lighter than the mass of two Λ hyperons. Thus this object - dubbed the *H* particle - an $SU(3)_f$ singlet configuration of structure *uu dd ss*, with all the six quarks in the *1s* ground state, with quantum numbers $J^{\pi}=0^{+}$ and isospin $I=0$, - a configuration which takes maximum advantage of the color-magnetic binding forces of quantum chromodynamics - is the best candidate for being stable against strong decays and is therefore unique among the possible dibaryons. The *H* could be formed by the fusion of two three-quark bags, corresponding to the $\Lambda\Lambda$, ΞN or $\Sigma\Sigma$ channels, without the need of any quark to be promoted to higher orbitals in the composite system.

The original *H* particle mass value calculated by Jaffe is $m_H \approx 2150$ MeV, some 80 MeV below the $\Lambda\Lambda$ threshold. Later, the Nijmegen group explored[4] the (q^6) spectrum for all allowed values of hypercharge, and again found the *H* as the only candidate for a stable dibaryon. Other calculations which include, always within the context of the MIT bag model, center-of-mass[5] or pionic cloud[6] corrections, all predict a particle which is near or below the mass of the $\Lambda\Lambda$ system.

The increased binding energy due to higher symmetry is seen in other models as well. The mass of the *H* has been estimated in the context of topological soliton (Skyrmion) models[7÷9]. Soliton theory, while not capable of reproducing the known hadron mass spectrum as accurately as the bag models, is thought by some to be more directly connected with basic QCD principles It also predicts a bound *H* particle with a mass lying in the range $m_H \approx 1.5 \div 2.2$ GeV

Quark matter calculations[10,11] indicate that the stable state of condensed matter may be "strange matter", in which the number of strange quarks approximately equals the number of up and down quarks. Such matter may be found in the core of neutron stars: in this context the H may be simply the lightest of a family of $-S \approx B$ particles.

More recently, some preliminary evidence against the stability of the H with respect to decay into two lambdas has been given using the technique of lattice QCD[12]. However, this result is not yet firm needing more work on the lattice calculation.

If the H exists and is indeed stable, it will be a unique object in multiquark spectroscopy (n>3). In ordinary potential models involving boson exchange, *deeply bound* states in the $\Lambda\Lambda$, ΞN and $\Sigma\Sigma$ channels, which can share the quantum numbers of the H, *cannot* be generated[13]. Thus the H - although it has become traditional to refer to it as a "dibaryon" - ought to be thought of as a genuine six-quark object, with all the quarks in one bag, rather than a weakly bound "quasi-molecule" such as the deuteron, where the two rather well separated three-quark bags retain much of their identity. The H is somewhat analogous to the α particle in nuclear physics but containing *three* different types of constituents. The discovery of this unique prediction of the quark model would provide the first definitive example of an n-quark state with $n=6$ as well as a dramatic confirmation of the model itself. Its implications would go beyond: from the confinement theory to the possible astrophysical consequences.

2. THE H IN THE QUARK MODEL

2.1. The MIT bag model

In the MIT bag model[1], the color gauge theory of quarks and gluons (QCD) has been adapted to conventional meson and baryon spectroscopies. The strongly interacting particles are considered as the finite region of space to which the interaction fields are confined. Thus the hadron is an extended object (bag) which contains hadronic fields consisting of fractionally charged colored quarks which interact with massless colored gluons. As both the quark and gluon fields are confined to the bag, only color singlet hadronic states are observable. The bag is taken to be a sphere of radius R and the quarks are placed in the $1s$ ground state. Inside the bag the quarks can move freely, except for a weak one-gluon-exchange interaction between the color charges, proportional to $g\lambda^{\xi}_{\alpha}$, where g is the gluon coupling constant and λ^{ξ}_{α} ($\alpha=1...8$) are the eight generators of SU(3,C) in the irrep $\underline{3}$, and between the color magnetic moments, proportional to $g\lambda^{\xi}_{\alpha}\sigma_k$, where σ_k $(k=1...3)$ are the three generators of SU(2,J) in the $J=1/2$ irrep.

The MIT model was acclaimed a success when it was shown that by adjusting few parameters of fundamental nature (bag pressure B, quark-gluon coupling constant α_c, zero-point energy Z_o and mass of the strange quark m_s) the masses of known baryons and mesons could be well reproduced. Without additional parameters, the model may be extended to any multiquark color singlet system $q^n\bar{q}^m$, with $n+m = N>3$, by providing a method to calculate the masses of various N-quark states. The states have an SU(2,J) classification for the space-spin part, an SU(3,C) classification for the color part, and an SU(3,F) classification for the flavor part (assuming only three flavors). Because of generalized Fermi statistics the N-quark states must be totally antisymmetric. Therefore up to 18 colored quarks can be put in the states of the bag.

2.2. The color-magnetic interaction between quarks

Let's consider a fixed number of quarks and antiquarks in a bag, all in the ground state and altogether forming a color singlet. The ordering of states is dictated by the color-magnetic interaction between quarks. Because of it, 0^- mesons are lighter than 1^- mesons, and $(1/2)^+$ baryons are lighter than $(3/2)^+$ baryons.

This consequence of QCD for the forces between quarks in clusters compares with a similar behaviour in QED for the electromagnetic forces between the electrons and nuclei of atoms. Such similarities are indeed observed, most noticeably in the hyperfine splittings between certain levels of the atom or quark clusters. In hydrogen there is a magnetic interaction between the electron and proton which contributes a positive or negative amount to the total energy depending upon whether the electron and proton are spinning parallel or antiparallel (the frequency of the photons emitted in the hyperfine transition from the singlet $F=0$ to the triplet $F=1$ is known with the extraordinary precision of one part over 10^{12}). In quark clusters there is an analogous "chromomagnetic" interaction between pairs of quarks which adds to the energy of an aligned spin pair, and depletes that of an antiparallel spin pair: so the $^3S_1$ combinations (ρ, K^*, ϕ) are some hundreds of MeV heavier than their $^1S_0$ counterparts (π, K, η) and the Δ^{++} splits 300 MeV over the proton mass.

The mass operator of an N-quark system can be written in the bag model as:

$$M = E_B + E_q + E_c + E_m$$

where

E_B is the energy associated with the bag (bag pressure, zero-point energy),

E_q is the rest energy and kinetic energy of the quarks,

E_c is the energy due to the color-electric interaction,

E_m is the energy due to the color-magnetic interaction.

The color-magnetic interaction term at the lowest-order gluon-exchange has the form:

$$E_m = -(\alpha_c/R) \sum_{ij} M_{ij}(R)(\vec{\lambda}^c \cdot \vec{\sigma})_i (\vec{\lambda}^c \cdot \vec{\sigma})_j$$

where

$\alpha_c = g^2/4\pi$ is the QCD coupling constant,

$M_{ij}(R) = M(m_i R, m_j R)$ is a function of the product of R and the quark masses: it measures the strenght of the interaction,

$\vec{\lambda}_i$ and $\vec{\sigma}_i$ are the color and spin operator of quark i.

If we ignore for the moment differences in the quark masses between nonstrange and strange quarks, we see that E_m is proportional to the term

$$-\Sigma\, (\vec{\lambda}^c \cdot \vec{\sigma})_i (\vec{\lambda}^c \cdot \vec{\sigma})_j.$$

Here

$$\vec{\lambda}_i \vec{\lambda}_j = \Sigma_\alpha (\lambda_\alpha)_i\, (\lambda_\alpha)_j$$
$$\vec{\sigma}_i \vec{\sigma}_j = \Sigma_k (\sigma_k)_i\, (\sigma_k)_j$$

The sum over all the quarks gives:

$$E_m \propto\ 1/4\, N(N-10) + 1/3\, J(J+1) + (1/2) f_C^2 + f_F^2$$

where f^2 denotes the eingenvalue of the quadratic Casimir operator for SU(3).

Clearly a hadron which is a color, flavor and spin singlet $(f_C^2 = f_F^2 = J = 0)$, such as the H, has the largest color-magnetic attraction.

To see in more detail which parts of the H wave function contribute most to the color-magnetic attraction one can decompose $|H\rangle$ into $q^4 \otimes q^2$ components[14] and evaluate the eingenvalue of E_m.

Assuming

$$\mu = m_s R \qquad m_u = m_d = 0$$
$$\langle H | E_m | H \rangle = -(\alpha_c/4R)\, [5M(0,0) + 22M(0,\mu) - 3M(\mu,\mu)]$$

where

$M(0,0)$ describes the nonstrange quark interaction,

$M(0,\mu)$ describes the strange quark-nonstrange quark interaction,

$M(\mu,\mu)$ describes the strange quark interaction.

The information we get is that the exchange term is the larger one. This follows from the fact that the color-magnetic interaction is attractive when both color and spin wave functions have the same symmetry, and is repulsive (but half as strong) otherwise. The Pauli principle requires the former configuration to pair up with the antisymmetric flavor representation, which is why the effect of the strange quark-nonstrange quark interaction is so large.

2.3. The two-baryon content of the H

If the H is stable and relatively deeply bound (say 25 MeV), it cannot be confused with a deuteronlike object bound weakly by ordinary meson-exchange forces. To exhibit the two-baryon content of the H, one can write its wave function in terms of $q^3 \otimes q^3$ components[14,15]:

$$|H\rangle = \sqrt{4/5}\,|8_c \otimes 8_c\rangle + \sqrt{1/10}\,|\Xi N\rangle_{I=0} - \sqrt{1/40}\,|\Lambda\Lambda\rangle_{I=0} + \sqrt{3/40}\,|\Sigma\Sigma\rangle_{I=0}.$$

Note that 80% of the time, the H falls apart *virtually* into two color-octet three-quark systems and only 20% into two-particle channels.

Using the $\Lambda\Lambda$, $\Sigma\Sigma$ and ΞN meson-exchange potentials of the Nijmegen group[16] the authors of Ref. 13 have found that *no* two-body bound states with $S=-2$ are expected. A different approach[17] using a quark-pairing mechanism to arrive at an one-boson-exchange model yields a lightly bound (1÷2 MeV) $\Lambda\Lambda$ $(^1S_0)$ state. This model, however, does not yield as high quality a fit to NN and YN data as that of Ref. 16. In any reasonable one-boson-exchange model, there is *no mechanism* for producing a deeply bound (tens of MeV) state with $S=-2$. This can occur *only* for six quarks in one bag.

2.4. The spin content of the H

The presence of color in QCD creates a fundamental difference from the nuclear physics analog of the H, the α particle. In the α, like particles (p,n) are paired into spin singlets, and similarly it happens in the electromagnetic analog, the Cooper pair of the BCS theory of superconductivity, due to statistics. In the H, the *three* different pairs of constituents (u,d,s) can couple into *triplet* states, due to color combined with the Pauli principle.

The spin content of the $H(J^\pi=0^+)$ in the various virtual baryon-baryon channels is as follows:

$$
\begin{aligned}
|\Lambda\Lambda\rangle &= |\uparrow u_1 \downarrow u_2\ \downarrow d_1 \uparrow d_2\ \uparrow s_1 \downarrow s_2\rangle \\
|\Xi^- p\rangle &= |\downarrow u_1 \downarrow u_2\ \downarrow d_1 \uparrow d_2\ \uparrow s_1 \uparrow s_2\rangle \\
|\Xi^o n\rangle &= |\downarrow u_1 \uparrow u_2\ \downarrow d_1 \downarrow d_2\ \uparrow s_1 \uparrow s_2\rangle \\
|\Sigma^o \Sigma^o\rangle &= |\uparrow u_1 \downarrow u_2\ \uparrow d_1 \downarrow d_2\ \downarrow s_1 \uparrow s_2\rangle \\
|\Sigma^+ \Sigma^-\rangle &= |\uparrow u_1 \uparrow u_2\ \downarrow d_1 \downarrow d_2\ \downarrow s_1 \uparrow s_2\rangle
\end{aligned}
$$

3. MODE DECAY AND LIFETIME OF THE H

For masses exceeding $2m_\Lambda$ *(2231 MeV)* the H is unbound. In particular, for $m_H > 2m_\Sigma$ *(2386 MeV)* the H is above all two-baryon thresholds to which it couples strongly. It represents, therefore, a repulsive interaction in all two-baryon channels to which it couples.

In the case $2231 < m_H < 2386$ *MeV* the H is a $\Sigma\Sigma$ state decaying strongly into $\Lambda\Lambda$. The H would appear as a bump in invariant mass plots.

For masses below $2m_N$ *(1879 MeV)* the *H* is strictly stable (up to processes which violate baryon number conservation).

The MIT bag model predicts the mass of the *H* to lie within the interval $(m_\Lambda + m_N, 2m_\Lambda)$, i.e.

$$2055 < m_H < 2231\ MeV.$$

Specifically, the mass prediction for the *S*-wave flavor-singlet with $J^\pi = 0^+$ gives

$$m_H = 2150\ MeV.$$

For this region the *H* would decay predominantly by nonleptonic weak decay to two baryons in $\Delta S = 1$ transitions. The more conventional nonleptonic decay involving a pion

$$H \to \Lambda N \pi$$

is forbidden if $m_H < 2195$ *MeV*, and is probably unimportant over the mass range of interest.

In the calculation of Jaffe[3], by assuming the nonleptonic weak Hamiltonian to be octet-dominated and taking into account the S-wave phase space, the three two-baryon decay channels (see Fig. 1) are related by SU(3) according to the branching ratios

$$\begin{aligned} H &\to \Sigma^- p \quad 50\% \\ &\to \Sigma^o n \quad 30\% \\ &\to \Lambda n \quad 20\% \end{aligned}$$

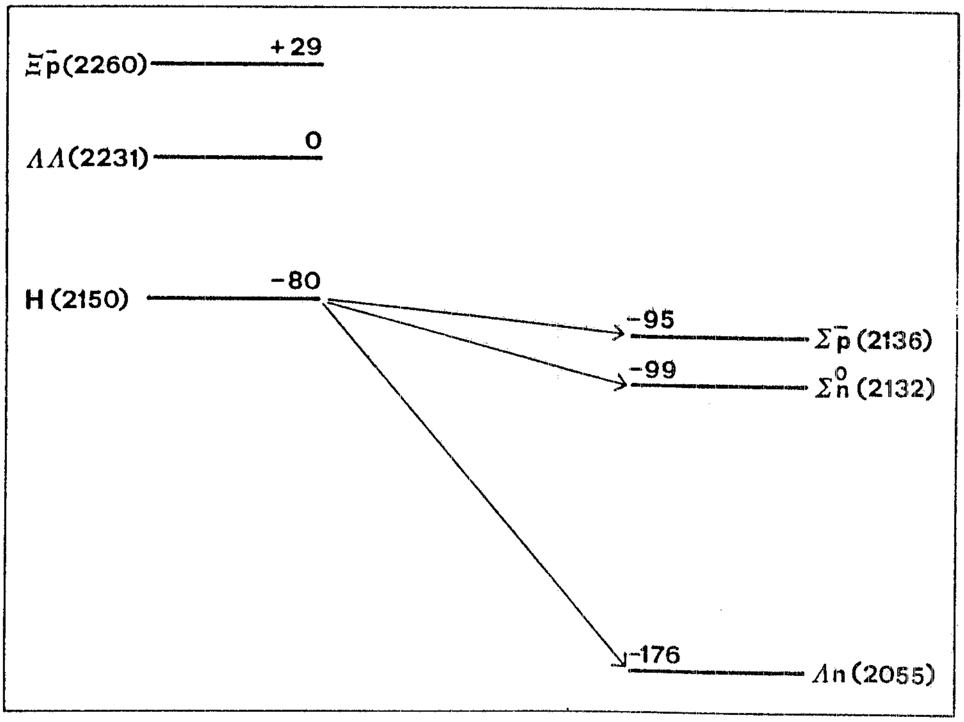


FIGURE 1
Two-baryon weak decay channels of the H particle

As the H mass decreases down to $2055\ MeV$, in particular under the threshold $m_{\Sigma}+m_N$ $(2135\ MeV)$, the Λn mode comes to dominate.

If the H were very tightly bound, i.e. if its mass were below $m_{\Lambda}+m_N$ $(2055\ MeV)$, all $\Delta S=1$ decay channels would be kinematically forbidden and the H would need to decay via $\Delta S=2$ to an nn final state. This possibility seems unlikely, but it has been raised recently as a possible way to explain the unusual events associated with radiation from Cygnus X-3.

According to the naive estimate[3], the H may be expected to live somewhat longer than the Σ^- or Λ (which are connected to members of the same flavor multiplet by H^{NL}_{weak}):

$$\tau_H \approx 10^{-10}\,s.$$

At present, the only possible H-particle candidate in literature, the supposed $p+d \rightarrow K^+K^opH$ event found by Shahbazian and Kechechyan at Dubna[18], should decay weakly to Σ^-+p with a lifetime

$$\tau_H = 0.668 \cdot 10^{-10}\,s.$$

A very recent study of weak decays of the H calculates[19] decay rates and lifetimes both for the expected $\Delta S=1$ channels and for the possible $\Delta S=2$ mode and finds a substantially different scenario. The calculation proceeds in three stages:

(i) The H wave function is constructed in terms of quark creation operators following what is dubbed the "Jaffe's theorem". The latter simply follows from the fact that SU(3) flavor representations in the spectroscopy of all possible color singlet six-quark configurations of S-wave quarks each appear just *once* in the original Jaffe's formulation[3]. In particular, while the singlet corresponds to the H particle, the 27-plet contains the possible baryon final states to which the H can couple via the weak Hamiltonian. This can be expressed by the statement that "any six-quark configuration having the correct quantum numbers must be a valid wave function for the associated state" (Jaffe's theorem).

(ii) The relevant weak decay amplitudes are computed in the framework of the MIT bag model.

(iii) The P-matrix formalism of Jaffe and Low[20] is used in order to relate the ensuing quark configuration to dibaryon scattering states, and thus obtain the associated decay rate.

It is noteworthy to comment the point (iii). Each of the H decay modes mentioned thus far involves a weak nonleptonic transition to a pair of baryons in the continuum. However, in quark models, such as the bag model, one *cannot* treat a continuum state directly, since the six-quark state representing the dibaryon is permanently confined as a result of the boundary conditions. The *P*-matrix represents a rigorous way to connect this artificially confined six-quark state with the real strongly interacting two-baryon final state. Jaffe and Low have shown[20] that in searching for states of the quark bag it is more correct to study the poles of the *P*-matrix (the "primitives") rather than the *S*-matrix. The coupling of the primitives to the allowable baryon-baryon channels must be performed before identifying m_H with the physically observable mass of the *H* (i.e. an *S*-matrix pole). The *S*-matrix pole corresponding to the *H* generally lies *below* the *P*-matrix pole for a wide class of parametrization of the *P*-matrix[14].

The results of Ref. 19 are reported in Figs. 2 and 3, and can be summarized as follows:

a) For $\Delta S=1$ transitions the *H* lifetime is (depending on its mass) in the 10^{-8} *s* range, which is quite *longer* then the naive estimate and has obvious implication in planning experiments.

b) For $\Delta S=2$ transition the *H* has a lifetime of the order of days, which poses difficulties for the *H* explanation of the Cygnus X-3 events (a

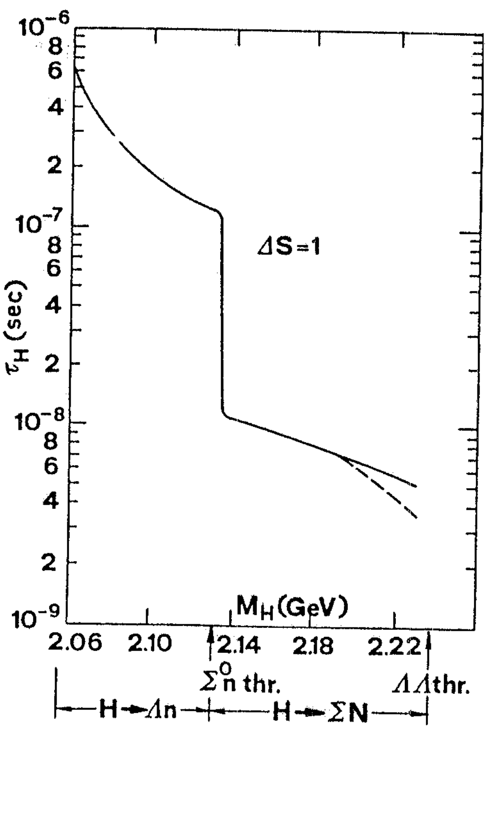


FIGURE 2

Predicted S-wave $\Delta S=1$ $H\to B_1B_2$ lifetime as a function of H mass. The dashed line displays the effect of $H\to\Lambda N\pi$

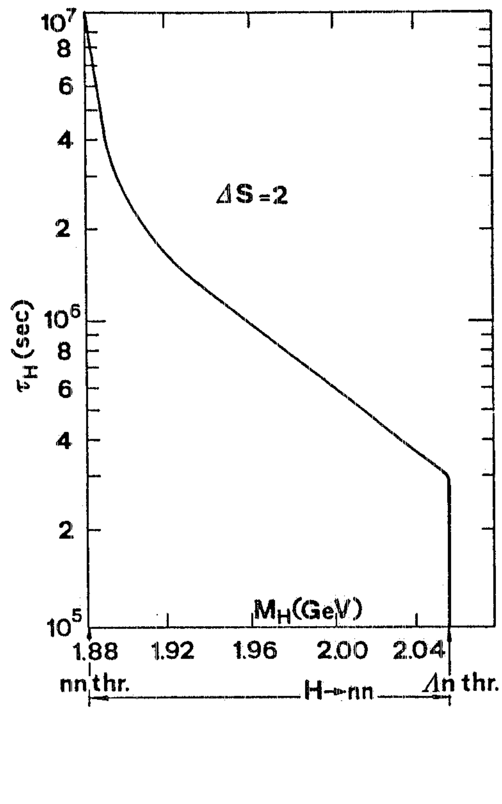


FIGURE 3

Predicted S-wave $\Delta S=2$ $H\to nn$ as a function of H mass lifetime of years should be required).

c) In $\Delta S=1$ transitions the H can decay into ΣN as well as ΛN. What is found is that the $\Delta I=3/2$ mode dominates the $\Delta I=1/2$ mode, due to the Pauli principle which forces the six-quark final state to be in an SU(3) 27-plet. If observed, the $H\rightarrow\Sigma N$ decays would constitute the first breakdown of the famous $\Delta I=1/2$ rule in baryonic nonleptonic transitions. Needless to say, this phenomenon alone would attract substantial attention.

4. HUNTING THE H

Production of the H requires strangeness $S=-2$ exchange. Part or all of this may be broken up in the choice of beam. The more natural way is to utilize an $S=-1$ incident beam (K^- or $\overline{K^0}$) , and to transfer two unit of strangeness to a nuclear system (exit particle is K^+ or K^0 with $S=+1$), which then recoils with $S=-2$. By bringing in one unit of strangeness, one obviates the need for double associated production. Strangeness *exchange* reactions (both single and double) have considerably larger cross sections than associated-production processes - as, for instance, those involving p or $\bar{p}$ - and are thus intrinsically *favorable* for H production.

Among the methods for producing the H , we mention the following reactions which actually have been experimentally investigated or have been object of a proposal or have been suggested:

a) $K^-+{}^3He\rightarrow K^++n+H$

b) $K^-+d\rightarrow K^{0*}+H$

$\Xi^-p)_{atom}\rightarrow\gamma+H$

c) $\Xi^-d)_{atom}\rightarrow n+H$

$\Xi^-{}^4He)_{atom}\rightarrow t+H$

d) $p+p\rightarrow K^++K^++H$

e) $p+d\rightarrow K^++K^0+p+H$

f) $p+A\rightarrow X+H$

g) $\bar{p}+A\rightarrow X+H$

To these "direct evidence" searches one must add a possible "indirect evidence" following the speculative suggestion that the radiation from Cygnus X-3 has a hadronic component which produces high energy muons in the Earth's atmosphere and mantle and which consists of H dibaryons.

Despite the variety of suggestions, experimental data on the subject are extremely limited. The main reason is that a kaon beam optimized for this goal in energy and intensity does not exist at this time.

Only *three* experimental results which deal explicitly with the existence of the H are known at present. Two out of them simply fix upper limits for the production cross section or for the frequency, the third one is a bubble chamber event which constitutes the only *H*-particle candidate up to now found. The three experiments concern the reactions *d)*, *e)* and *g)* of the above list.

The reaction *d)* was investigated by Carroll *et al.*[21] at the Brookhaven Alternating Gradient Synchrotron some months after the Jaffe's prediction for a stable dibaryon[3]. One expects the dominant process to be $p \to K^+ + \Lambda$ dissociation (twice) followed by $\Lambda\Lambda \to H$ recombination. In the K^+K^+X missing - mass spectrum (incident beam momentum 5÷6 GeV/c) no narrow structure was observed. The upper limit of *40÷50 nb* for the *H* cross section in the mass range $2.1 < m_H < 2.23$ *GeV* is too crude to rule out its existence, in particular if one accepts the estimate of *1 nb* for reaction *d)* due to Badalyan and Simonov[22]. These authors give

$$\sigma_H/\sigma_{\Lambda\Lambda} \approx 4 \cdot 10^{-3},$$

where $\sigma_{\Lambda\Lambda}$ is the cross section for $p+p \to K^+K^+\Lambda\Lambda$.

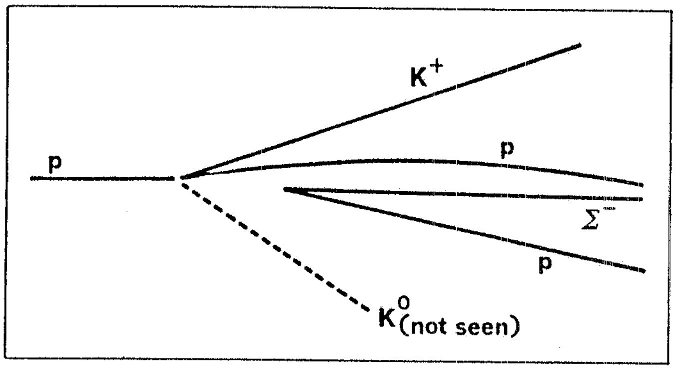


FIGURE 4
A diagram of the event reported in Ref. 18. The proton is assumed to interact with a quasi-deuteron in a carbon nucleus, producing the H particle which then decay Σ^-p

The production of the *H* particle through this process is indeed inhibited by the large momentum transfer involved which suppresses the $\Lambda\Lambda$ recombination probability leading to the *H* (see also § 6).

In a recent experiment Condo *et al.*[23] searched for *H* particles produced from multinuclear capture of slow antiprotons in complex nuclei according to the reaction *g)*. If an incident low-energy antiproton interacts with a cluster of nucleons some unusual reaction can occur, as those with doubly strange final states. Because of the destruction of a unit of baryon number of the capturing cluster, these reactions possess in fact a larger *Q*-value than would otherwise be the case. For instance, the *Q*-value of the double associated strangeness producing reaction $(\bar{p},3N) \to \Lambda\Lambda KK$ is ≈ *500 MeV*, not substantially less than the value for ordinary associated strangeness producing annihilation reaction $\bar{p}p \to KK\pi\pi$. In order to search for any possible double strangeness annihilations the group examined a sample of ≈ *80.000* antiproton annihilations *($\bar{p}$ momentum < 400 MeV/c)* occurring in four thin elemental *(C,Ti, Ta, Pb)* plates placed in the BNL 30 inch hydrogen bubble chamber.

At the 90% confidence level, the probability of the emission of an *H* with subsequently decay to $\Sigma^- p$ turned out to be less than $9 \cdot 10^{-5}$, assuming an *H* lifetime comparable to that of the Λ. Although the negative result, it should be noted that only upper limits were obtained for other double strangeness producing channels as well. For example, an upper limit of $4 \cdot 10^{-4}$ was found for $\Lambda\Lambda$ production.

The only case of an event reported as a possible *H* particle can be found in a Dubna preprint (1984) and in a presentation at the recent Panic Conference (Kyoto 1987) by Shahbazian *et al.*[18]. These authors give the results of a systematic search of the *H* undertaken in negative pion-, neutron-, ^{12}C nucleus-, and proton-propane interactions at 4.0, 7.0, 50.4 and 10 GeV/c, respectively, using 55 cm and 2 m propane bubble chambers. They report an event, reproduced schematically in Fig. 4, in which a 10 GeV/c proton travels 3/4 of the way across a 2 m propane bubble chamber and produces a two-prong star followed by a neutral "Vee". The two prongs of the star are two positive particles, one of which is heavily ionizing and therefore identified as a proton. The "Vee" might be the decay of an *H.* Failing to find a less exotic explanation, the authors hypothesize that the event may be explained by assuming that the incoming proton interacts with a quasi-deuteron producing an *H* through the reaction *e).* The mass of the event treated as an $H \rightarrow \Sigma^- p$ weak decay appeared to be *(2173.94 ± 1.32)MeV* , its lifetime being $0.668 \cdot 10^{-10}$ *s.* Both are in agreement with MIT model calculation of Jaffe[3].

Besides the usual problems with interpreting *single* event emulsion of bubble chamber data, this explanation requires a number of assumptions such as the existence of the unseen K^0 and the fact that the decay of the Σ^- is *not* seen.

Although the authors calculate a 10% probability that the Σ^- would survive over the observed track length, the lack of positive particle identification certainly reduces the significance of the event. The discovery of additional events with more complete kinematics would be of great interest.

5. DOUBLE ΛΛ HYPERNUCLEI AND THE H

The observation of ΛΛ hypernuclei is usually taken as escluding the possibility of a bound *H.* Moreover, the identification of a ΛΛ hypernucleus in emulsion requires that both Λ's decay pionically. Since in heavier Λ hypernuclei *(A>5)* pionic decay is strongly suppressed, this implies that $\Lambda\Lambda$ hypernuclei may indeed be produced quite abundantly but cannot be identified in emulsion. Since the mere existence of a Λ hypernucleus is relevant to the question of a bound *H*, the problem to establish some signature of a $\Lambda\Lambda$ hypernucleus which decay non pionically in emulsion should be of great interest.

There are *two* event in the emulsion data which have been interpreted as the production and subsequent decays of $\Lambda\Lambda$ hypernuclei. The first such event was reported by Danysz *et al.*[24] and has been identified as the production of ${}^{10}_{\Lambda\Lambda}Be$. A second event was later reported by Prowse[25] to show the formation of ${}^{6}_{\Lambda\Lambda}He$.

The ${}^{6}_{\Lambda\Lambda}He$ event recorded by Prowse[25] corresponds to the following chain of processes:

$$\Xi^- + {}^{12}C \rightarrow {}^{7}Li + {}^{6}_{\Lambda\Lambda}He$$

$${}^{6}_{\Lambda\Lambda}He \rightarrow \pi^- + p + {}^{5}_{\Lambda\Lambda}He$$

$${}^{5}_{\Lambda\Lambda}He \rightarrow \pi^- + p + \alpha$$

Since ${}^{6}_{\Lambda\Lambda}He$ is observed to decay sequentially through two successive weak decays $\Lambda \rightarrow \pi^- + p$, this kind of event has been used to argue against the existence of the *H*. In fact, if the *H* exists as a bound state with respect to the $\Lambda\Lambda$ threshold, the strong decay

$${}^{6}_{\Lambda\Lambda}He \rightarrow \alpha + H$$

should be much more rapid than the weak decay chain. This argument has been quantified by Kerbikov[26] who evaluated the width Γ_s for the strong decay:

$$\Gamma_s \propto (2\mu Q)^{1/2} exp(-2\mu Q)$$

where μ is the reduced mass of the $H+\alpha$ system and $Q=2m_\Lambda - m_H - B_{\Lambda\Lambda}$ is the energy released in the decay. $B_{\Lambda\Lambda}$ is the separation energy of two Λ's from the core nucleus: $B_{\Lambda\Lambda} \approx 10$ *MeV* for $A \geq 6$. For $m_H = 2150$ *MeV* $Q \approx 70$ *MeV*. According to Kerbikov[26] the strong decay time $\tau_s = \Gamma_s^{-1}$ is generally much smaller of the time $\tau_w (\approx 10^{-10} s)$ characteristic of weak decay. However, for an *H* which is *weakly bound* by $10 \div 20$ *MeV* the factor $Q^{1/2}$ renders τ_s comparable with τ_w and the observed weak decay is not inconsistent with the presence of an alternate strong mode governed by the *H*. The same conclusion can be reached in the opposite case of a *tightly bound H* ($m_H < 1900$ *MeV*). Now the width Γ_s is suppressed by the exponential factor containing a large *Q* and one finds $\tau_s > \tau_w$ and again the observed $\Lambda\Lambda$ hyperon decay does *not* appear to rule out the existence of the *H* particle.

6. ON THE H PRODUCTION IN $(\bar{K}, K)$ REACTIONS

Various kinematical and dynamical arguments indicate that the $(\bar{K}, K)$ double-strangeness-exchange reactions on nuclear targets affords the best chance for observing the *H*. In a two-step mechanism, a Ξ^- is firstly produced and then a $\Xi^- p$ fusion yields an *H*. If the fusion takes place within the same nuclear volume, the overall processes are described by the reactions a) and b) of § 4.

An alternative approach - often referred to as the "two-target" approach - consists in using the $K^- p \to K^+ \Xi^-$ reaction in a hydrogen target to tag the production of a Ξ^-, which is then slowed down electromagnetically and captured in an atomic state in a hydrogen, deuterium or ^4He target. The *H* is subsequently produced by annihilation of the atoms via the fusion processes described by the reactions c) of § 4.

6.1. (K^-, K^+) reaction on a bound diproton

The reaction

$$K^- + (pp)\,^1S_0 \to K^+ + (\Xi^- p)$$

followed by

$$(\Xi^- p) \to H$$

has been investigated in detail in Ref. 14. The system $(\Xi^- p)$ is already in an isospin singlet $(I=0)$ as required for the *H*. It must be prepared also in a spin-singlet $(S=0)$ state. This will be the case if the process is studied by revealing the K^+ in the forward direction $(\theta_{K^+}=0^0)$. Then spin-flip amplitude will be absent and if the two protons are in $(L=0)$ they will be automatically in the singlet state, due to the Pauli principle.

It can be shown that *H* production via *bound* diproton ($L=0$ as well as $S=0$) is favored over production in *pp* collisions by a factor of *16* from the spin-flavor-color factors alone.

A bound diproton appears therefore as the "optimal" target for making an *H*. The lightest bound nucleus containing a diproton is ^3He, so the prototype reaction is

$$^3He(K^-, K^+)nH$$

This reaction has also the advantage of having a simple three-body final state: measuring K^+ and neutron in coincidence, in a "missing-mass" spectrum the only "missing-mass" is the *H* itself. It should appear as a sharp peak, the width of which is the experimental resolution (the *H* natural width is essentially zero).

The vertex $\Gamma(\vec{k}_1, \vec{k}_2)$ for *H* formation in the ΞN collision can be obtained by an overlap of quark wave functions. It is proportional to the quantity

$$\Gamma(\vec{k}_1 - \vec{k}_2) = \Gamma_0 (8\pi R^3/3)^{3/4} exp[-R^2/12(\vec{k}_1 - \vec{k}_2)]$$

where Γ_0 is the color-spin-flavor recoupling coefficient and $\vec{k}_1$, $\vec{k}_2$ are the momenta of the two particles.

The salient feature of Γ is its strong dependence on the relative momentum $\vec{k}_R=\vec{k}_p-\vec{k}_\Xi$. Small values of $\vec{k}_R$ are preferred for H production in a baryon fusion process. In the case of a nuclear target the Fermi motion of the nucleus ought to provide a region of phase space where $\vec{k}_R$ is small. The ^{3}He wave function, on the other hand, provides only relatively small momenta (of order of *250 MeV/c* or less) for the proton to match the typically large (of order of *400 MeV/c* at the "optimal" K^- momentum of *1.8 GeV/c*) value of $\vec{k}_\Xi$. The possibility of achieving small $\vec{k}_R$ is enhanced if the overall momentum transfer $q=k_L-k'(k_L$ is the $S=-1$ incident-particle momentum and k' is the $S=+1$ outgoing-particle momentum) is made as small as possible. Which again means that forward K^+ production is the optimal case.

The choice of the momentum beam is straightforward considering where peaks the two-body reaction cross section $K^-p\rightarrow K^+\Xi^-$. From the available measurements, the maximum Ξ^- flux occurs for $k_L\approx 1.8$ *GeV/c*.

The lowest-order mechanism for H formation on the pp pair in ^3He is shown in Fig. 5. Several higher order processes could also contribute to H production on a diproton pair. These processes lead to the H via $\Lambda\Lambda$ or $\Sigma\Sigma$ fusion and are expected to be much less important (about three orders of magnitude) than the leading graph of Fig. 5.

The predicted differential cross section for $^3He(K^-,K^+)nH$ is shown in Fig. 6. It is noticeably the fact that it increases as the H is more bound.

The reaction is thus appropriate over a wide range of m_H values. If one

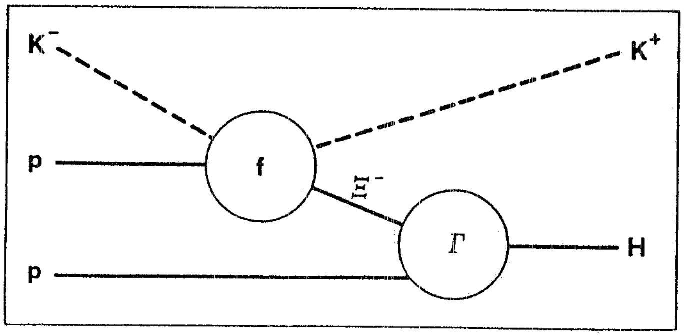


FIGURE 5
Lowest-order process for H production on a diproton pair via the (K⁻,K⁺) reaction.

compares the momentum dependence of $d\sigma/d\Omega_{K^+}$ for H production with the total 0^0 cross section for $\Xi^- p$ continuum pair production $[\approx 2(d\sigma/d\Omega)(K^- p \to K^+ \Xi^-)]$ one gets a rough estimate for the "fusion probability":

$P(\Xi^- p \to H) \approx 4 \cdot 10^{-3}$ at $m_H \approx 2200$ MeV

This small value reflects the fact that the small relative momentum condition is satisfied only in a very restricted region of the phase space supplied by the ^3He wave function.

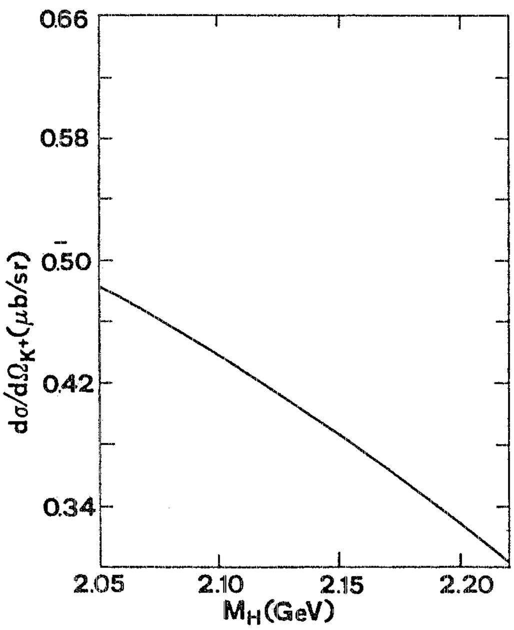


FIGURE 6

Predicted differential cross section $d\sigma/d\Omega_{K^+}$ at $\theta_{K^+} = 0°$ for H production in the ^{3}He $(K^-,K^+)nH$ reaction as a function of the mass m_H at laboratory momentum k_L= 1.8 GeV/c

The virtue of measuring the K^+ and neutron in coincidence in the $^3He(K^-K^+)nH$ reaction is that the only "missing-mass" is the H itself. In the predicted "missing-mass" plot of Fig. 7, if the H is bound it should be clearly separable from the continuum background. On the contrary, if it lies above threshold, it would be very difficult to disentagle from the much larger background cross section.

Two major sources of background must deserve a careful investigation:

(i) K^+ quasi-elastic production.

It concerns K^+ in the final state associated with typical processes such as Ξ^- production *not* followed by fusion or $(\Xi^- p)$ scattering with production of two Λ's (see Fig. 8).

In Fig. 9 the predicted K^+ spectrum corresponding to both H and $\Xi^- p$ continuum production is plotted. Although the continuum cross section is roughly *100* times larger at peak than that for H production, the figure suggests that the two processes can be distinguished if one can measure the K^+ momentum to 2% or better $(\Delta k'/k' \approx 30/1400)$.

(ii) π^+ in the final state misidentified as K^+

For an incident momentum $k_L \approx 1.8$ GeV/c the K^+ spectrum is peaked at $k' \approx 1.42$ GeV/c (for m_H=2200 MeV) which corresponds to β=0.9445. A π^+ with

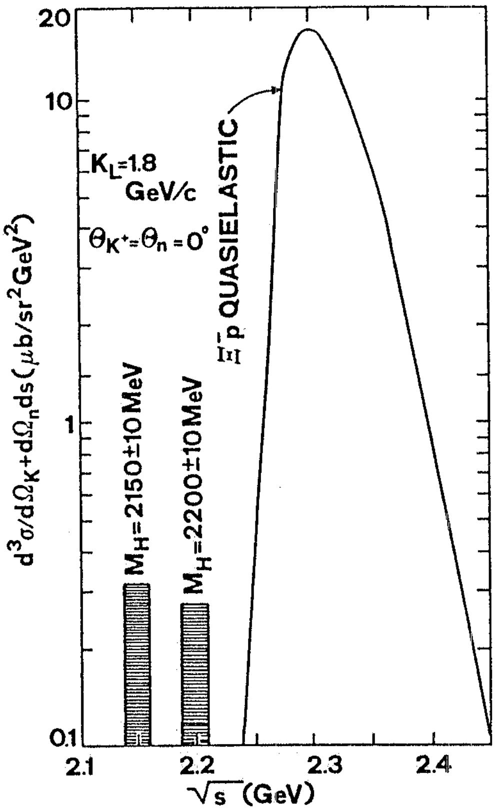


FIGURE 7

The predicted "missing-mass" spectrum for the ^3He $(K^-,K^+)nH$ or ^3He $(K^-,K^+)n\ \Xi^-p$ reactions. The invariant mass of the Ξ^-p system is $\sqrt{s}$.

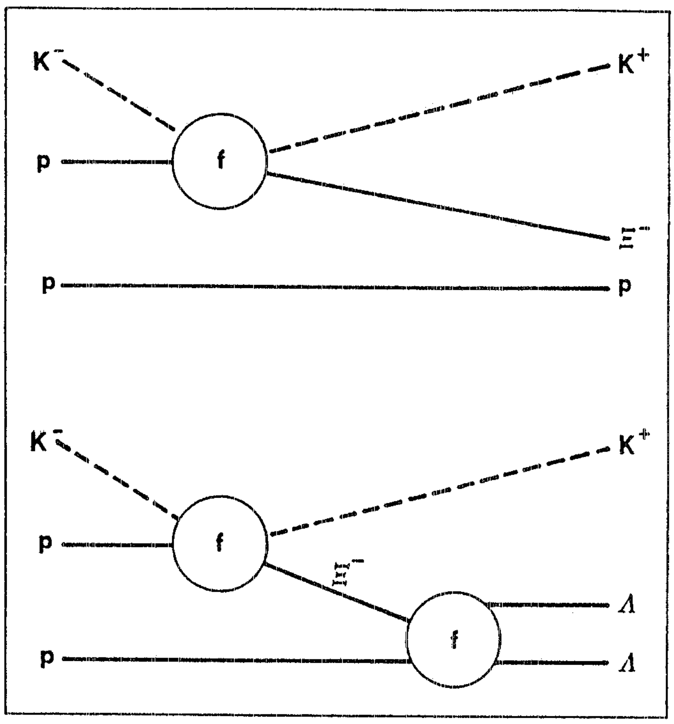


FIGURE 8

Processes involving the (K^-,K^+) reaction on a diproton pair which do not involve H formation. These contribute to the "quasi-elastic" background.

the same momentum (but with β=*0.9952*) may occasionally be misidentified as a K^+ providing an other background for *H* production. Thus, reactions of the type

$$K^- + {}^3He \rightarrow \pi^+ + X$$

have to be considered. The calculated π^+ combined spectrum is shown in Fig. 9. The sum displays a minimum around $k_L - k' \approx 0.38$ *MeV/c* corresponding to the K^+ peak for *H* production for m_H=*2.2 GeV*. The π^+ background spectrum is still roughly *50* times as large as the *H* signal in this domain.

The conclusion is that a study of *H* production by using the $^3He(K^-,K^+)nH$

reaction is feasible if careful consideration is given to the suppression of backgrounds due to misidentified $\pi^{+\prime}s$. Discrimination of $\pi^{+\prime}s$ from K^+'s at the level *at least* of *1:100* is needed in order to achieve an observable *H* signal.

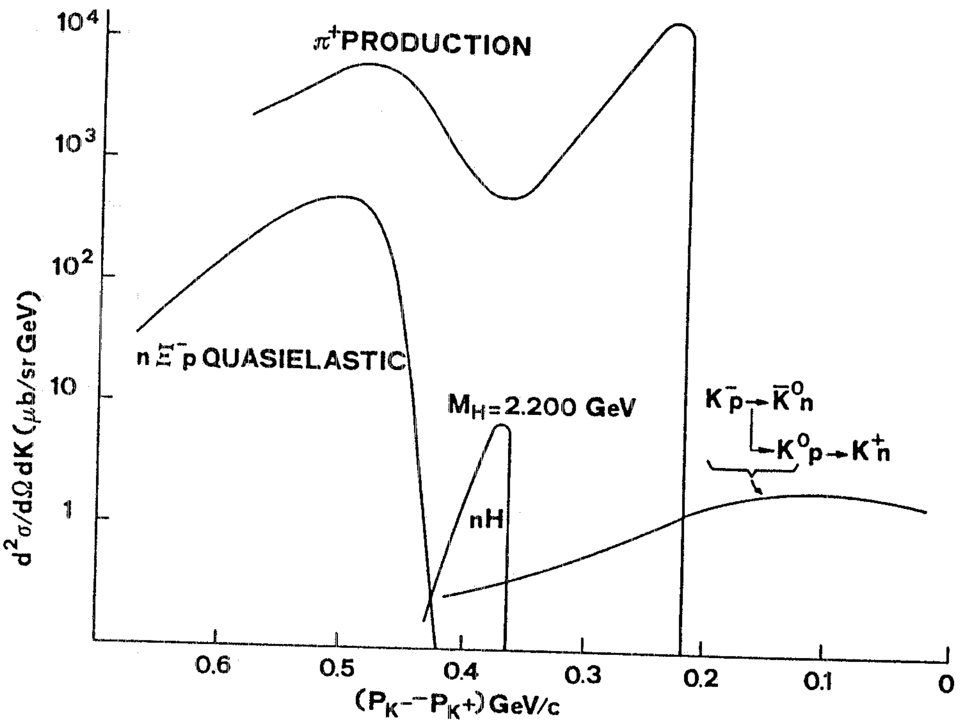


FIGURE 9
Predicted H signal and main background sources

6.2 *(K^-,K^{0*})* reaction on deuteron

The possibility to produce the *H* via the reaction

$$K^- + d \to K^{0*} + H$$

has been recently suggested by Fitch[27] and is investigated in detail by C. Dover in Ref. 2.

Note that - unlike the case of ^{3}He target - *H* production on deuterium (or any *np* pairs with $I=0$ and $S=1$) through the "prototype" reactions

$$K^- + d \to K^+(K^0) + (\Xi N)$$

is *unfavorable* since they have the wrong spin-isospin-flip properties. In fact, if Ξ^- production comes from the proton it gives the $(\Xi^- n)$ pair the *wrong isospin* to fuse in an *H*:

$$K^- + d \to K^+ + (\Xi^- n)_{I=1},$$

while a Ξ^- produced from a neutron (detecting the K^0 at 0^0) gives the $(\Xi^- p)$

pair the *wrong spin* to become an *H*:

$$K^- + d \rightarrow K^0 + (\Xi^- p)_{S=1}$$

On the contrary, to detect a $K^{0*}(892)$ instead of a K allows for a spin-flip transition $(\Delta S=1)$ with no isospin-flip $(\Delta I=0)$, just right to transform the deuteron into the H.

The data on K^{0*} production are rather limited, but using the results of Scheuer *et al.*[28] at 3 GeV/c, and with some approximations, Dover[2] extracts a value for the cross section:

$$d\sigma/d\Omega_{K^{0*}} (K^- d \rightarrow K^{0*} H) \approx (1/10)\, d\sigma/d\Omega_{K^+} (K^- {}^3He \rightarrow K^+ n H).$$

Despite of the uncertainties in the estimate and although the cross section is about an order of magnitude smaller than that on ^3He, a (K^-,K^{0*}) experiment seems to present two advantages: i) there is no need to detect the K^+ and n in coincidence, which would be an important experimental advantage; the charged decay product from $K^{0*} \rightarrow K^+ \pi^-$ are detected in a forward cone, and ii) K^- beams in the $3 \div 5$ *GeV/c* region are available at the Brookhaven AGS.

6.3 *H* formation from Ξ^- atoms

The "two-target" scheme for *H* production, originally suggested by Barnes *et al.*[29], is designed to produce the *H* through a *low momentum transfer* reaction in order to obtain the maximum production cross section. Branching ratios for *H* formation from Ξ^- atoms have been predicted in Ref. 30.

After production in a first hydrogen target, there are several time scales relevant to the slowing down of a Ξ^- particle in matter which have to be compared with the proper lifetime $(\tau_\Xi = 1.641 \cdot 10^{-10} s)$. Following the classic discussion of capture processes leading to exotic atoms due to Leon and Bethe[31] and Wightman[32], one can distinguish:

i) slowing down in a moderator *(Pb,W)* by ionization of the atoms from the initial production momentum of about *500 MeV/c* down to the characteristic velocities β_e of electrons in atomic orbits $(\beta_e \approx \alpha)$ in a time

$$t_{mod} \approx 1.3 \cdot 10^{-10} s$$

ii) capture in a Bohr orbit around a nucleus of the second target and deexcitation down to a state with principal quantum number $n \approx \sqrt{\mu}$, where μ is the reduced mass of the Ξ^- atom (in units of m_e) in a time

$$t_{chem} \approx 3.2 \cdot 10^{-11} s$$

(iii) deexcitation by Auger processes from $n \approx \sqrt{\mu}$ to $n=5$ in a time

$$t_{Aug} \approx 4 \cdot 10^{-11} s$$

(iv) deexcitation by radiative processes from $n \leq 4$ states in

$$t_{rad} \approx 4 \cdot 10^{-11} s$$

If liquid targets are used, the Stark effect produces a continuous population of high-n states, from which annihilation takes place. The low-n circular orbits for which radiative processes dominate are scarcely produced and hence t_{rad} plays a little role. Thus the relevant cascade time is that for Auger processes. Only a few transitions are needed between states of relatively high n so that

$$t_{Aug}(total) \approx t_{weak}$$

It is clear that a considerable number of Ξ^- atoms will be lost to weak decay during the cascade process, as well as during moderation.

We consider in particular a second target of deuterium. The mechanism for H formation in the process

$$\Xi^- d \rightarrow Hn$$

consists in the coupling followed by fusion of baryon-baryon virtual intermediate states, as shown in Fig. 10 a). The background processes are the energetically allowed annihilations shown in Fig. 10 b) (*Q-values = 26.18 MeV* for $\Lambda\Lambda n$ and *2.91 MeV* for $\Xi^0 nn$, respectively). As with other hadronic probes such as the π^-, K^-, $\bar{p}$ and Σ^-, annihilation for the Ξ^- will occur dominantly from the S-state, with some contribution from the P-states (although not as much as for a very strongly absorbed probe such as the $\bar{p}$).

The branching ratio R defined by

$$R = \Gamma(\Xi^- d \rightarrow Hn)/\Gamma_{tot}$$

has been calculated[30] describing the amplitudes $\Xi N \rightarrow \Xi N$, $\Lambda\Lambda$, $\Sigma\Sigma$ which enter coherently in Fig. 10 in terms of OBE real transition potentials, since there are no experimental data on these processes at low energies. It is displayed in Fig. 11 for the S-state of $\Xi^- d$ atoms. Note that R is rather large ($\approx$0.9) when m_H lies near $2m_\Lambda$, and decreases rapidly as the H becomes more deeply bound. This behaviour is evident in the approximate formula

$$R \approx (Q+E_b/Q)^{1/2} |\phi_d(k_0)|^2$$

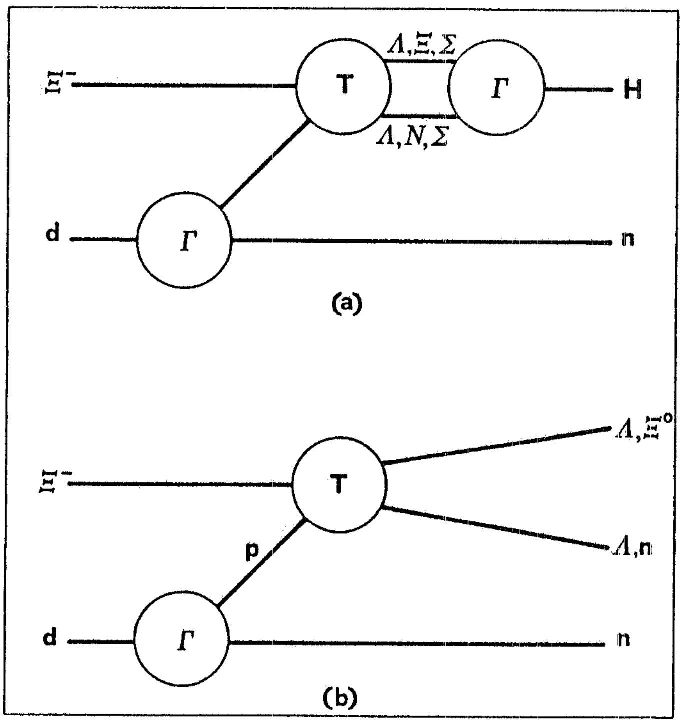
where $Q = m_\Xi + m_p - 2m_\Lambda$ and $E_b = 2m_\Lambda - m_H$ is the binding energy of the H; $\phi_d(k_0) \propto 1/k_0$ is the deuteron wave function evaluated at the momentum k_0 required for energy-momentum conservation in the process $\Xi^- d \rightarrow Hn$

$$k_0 = [2m_n m_H(Q+E_b)/m_n + m_H)]^{1/2}$$

As m_H decreases, k_0 becomes larger (k_0=*194*, *280* and *378 MeV/c* respectively, for m_H=*2231*, *2200*, *2150 MeV*).

The B_H dependence of R seen in Fig. 11 is now easily understood. As m_H decreases, the factor $[(Q+E_b)/Q]^{1/2}$ increases, since more phase space becomes available. However this is more than compensated by the rapid *decrease* in $|\phi_d(k_0)|^2$. Thus R is determined essentially by the probability that the neutron in the deuteron has the momentum k_0 required by energy-momentum conservation.

The neutron momentum spectrum predicted from the annihilations $(\Xi^- d)_{1s} \to \Lambda\Lambda n$, $\Xi^0 nn$, Hn is shown in Fig. 12. The backgrounds neutrons from $\Lambda\Lambda n$ and $\Xi^0 nn$ are concentrated at very low momenta of the order of *50 MeV/c*. The maximum momentum for background neutrons is $k_{n,max} \approx 194$ *MeV/c*. If $m_H < 2m_\Lambda$ the neutrons from the $\Xi^- d \to Hn$ reaction are *kinematically distinct* from background neutrons since they have necessarily $k_n < k_{n,max}$. This clear signal of relatively high-momentum neutrons in a region of small background looks very promising from an experimental point of view.



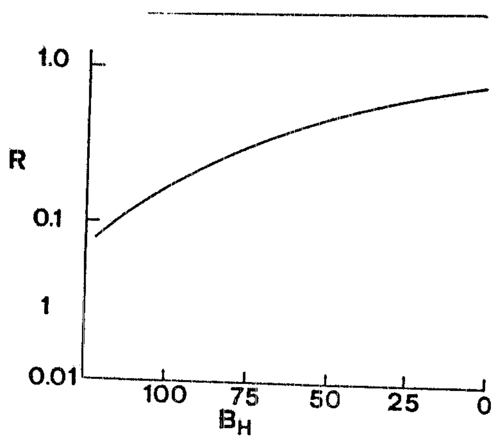


FIGURE 10

In a), we show the mechanism for H formation in the process $\Xi^- d \to Hn$. The baryon-baryon virtual intermediate states which couple to the H are indicate. In b), we display the "background" processes $\Xi^- d \to \Lambda\Lambda n$, $\Xi^\circ nn$

FIGURE 11

Predicted ratio of the $(\Xi^- d)_{atom} \to Hn$ process as a function of binding energy

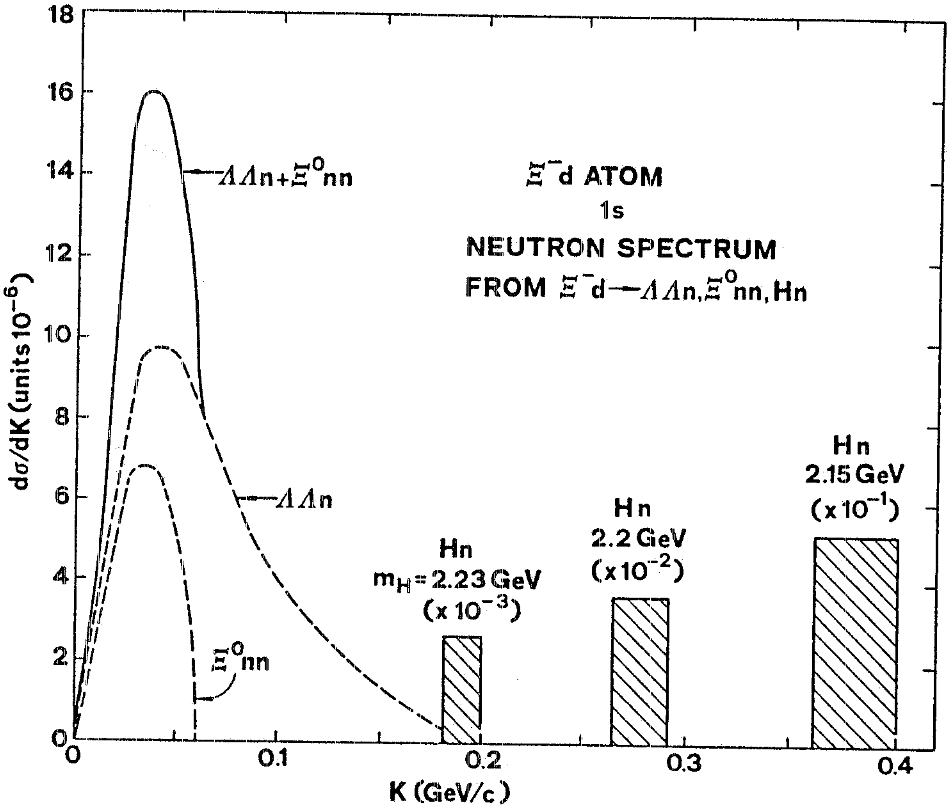


FIGURE 12

The predicted neutron momentum spectrum for the reactions $\Xi^- d \rightarrow \Lambda\Lambda n$, $\Xi^\circ nn$ and Hn

In conclusion, looking for a comparison between single-volume and two-target methods to produce the *H*, one can say that for an *H* close to the ΛΛ threshold, Ξ^-d atom formation appears to be more favorable [*R* increases as m_H increases, while $d\sigma/d\Omega_{K^+}$ decreases]. For a tightly bound *H* (m_H<2100 *MeV*) production on a ^3He target appears as the optimal method.

7. FUTURE EXPERIMENTAL PROGRAMMES ON THE H PARTICLE

7.1. Kaon research

Both Λ and Σ hypernuclei have been studied using negative kaon beams available at BNL, CERN and KEK. However, current hypernuclear programmes with kaons exist only at BNL and KEK[33]. The BNL LESB-I beam line is limited to kaon beams with less than 1 GeV/c. The KEK K3 beam line deliver kaons in the momentum range 0.5÷1 GeV/c. The K2 beam line at KEK is currently the only existing kaon beam line capable of delivering a separated kaon beam with a momentum in the range 1-2 GeV/c, therefore useful for exploring doubly-strange systems as the *H* particle. The measured beam intensity at *2.1 GeV/c* is $1\cdot 10^5 K^-$ per 10^{12} primary protons of *12 GeV* energy.

A new beam line has been proposed[34] and recently financed for the BNL AGS which will deliver kaon beams in the 1-2 GeV/c momentum range and will allow, for the first time, detailed studies of $S=-2$ hypernuclear systems. In particular, the beam line will provide a minimum of 10^6 negative *1.8 GeV/c* kaons per AGS beam spill. The kaon beam purity should be about *1:1* unwanted particles to negative kaons.

There is currently a conditionally approved experiment at BNL AGS (contingent upon the construction of the new beam line) designed to search for the H using the $\Xi^- d$ atom formation method. A schematic diagram of the proposed experiment[35] (AGS Experiment 813) is shown in Fig. 13.

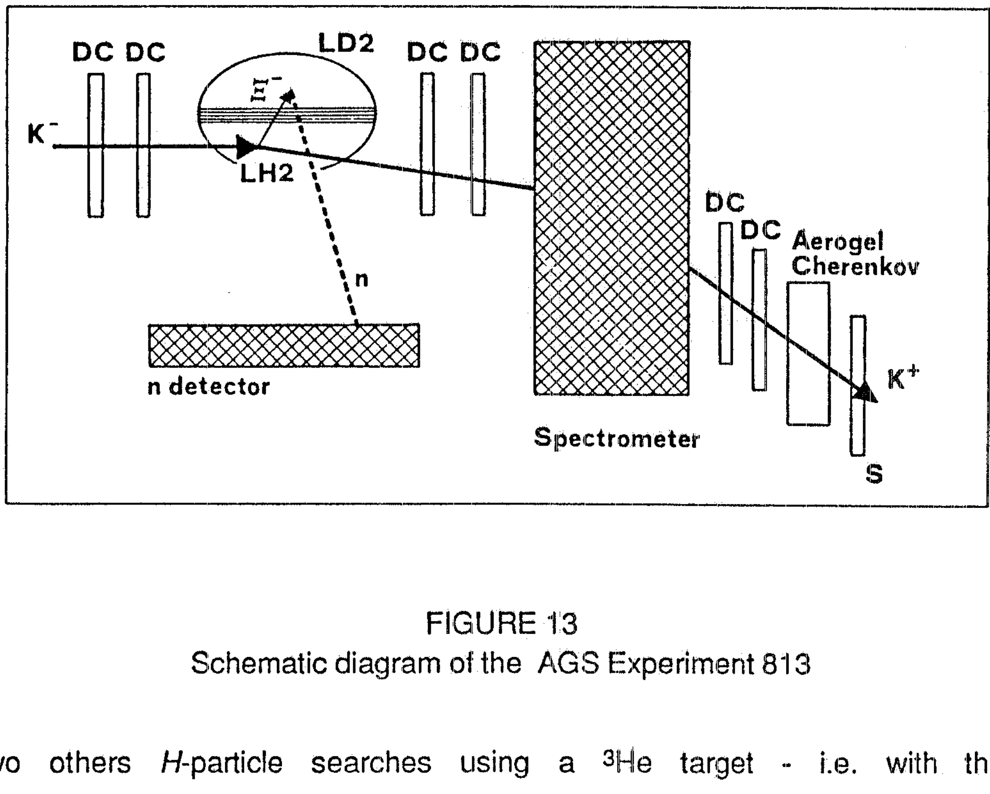


FIGURE 13
Schematic diagram of the AGS Experiment 813

Two others *H*-particle searches using a ^3He target - i.e. with the final $(\Xi^- N)$ fusion in the same nuclear volume - have ben proposed at the BNL AGS, one by the proponents of the quoted Ξ^- d atom experiment[36] (AGS Proposal 836) and one by the collaboration working on $S=-1$ dibaryons[37] (AGS Proposal 830).

It can be useful to get an approximate idea of the effective rates involved in both kind of experiments.

A) ^{3}He production

The number of produced H is given by:

$$N_H = N_{K^-} (d^2\sigma/d\Omega_K d\Omega_n)\Delta\Omega_K \Delta\Omega_n \rho (N/A)s \varepsilon_n$$

By assuming

$$N_{K^-} = 4 \cdot 10^5/s$$

$$d^2\sigma/d\Omega_K d\Omega_n \approx 24 nb/sr^2$$

$$\Delta\Omega_K = 20 msr \qquad \Delta\Omega_n = 1\pi sr$$

ε_n neutron detection efficiency $\approx 20\%$ for $E_n = 10 \div 100 MeV$
Liquid ^{3}He target: *10 cm ~ 0.9 gr/cm²*

one gets a K^+ neutron coincidence rate of $\sim 2 \cdot 10^{-5}$ events/second= $7.2 \cdot 10^{-2}$ ev/hour. To obtain *100* events would thus require ~ *60 days* of running. Compare this figure with the rate obtainable at the European Hadron Facility[38] (EHF)

$$N_{K^-} (EHF) \approx 10^8/s \quad \text{at 2 GeV/c.}$$

Therefore *100* events are obtainable in ~*1.4 hours.*

B) $(\Xi^- d)_{atom}$ formation

1) Ξ^- beam production

$$N_{\Xi^-} = N_{K^-} (d\sigma/d\Omega_{K^+}) \Delta\Omega_K \rho(N/A)s$$
$$N_{K^-} = 4 \cdot 10^5/s$$

By assuming

$(d\sigma/d\Omega_{K^+}) \approx 52 \mu b/sr$ $\quad \Delta\Omega_{K^+} = 50 msr$

LH_2 *target: 1.2 cm ~0.07 gr/cm²* one obtains $N_{\Xi^-} \approx 150/hour$

2) $(\Xi^- d)_{atom}$ formation and annihilation

$$N_H = N_{\Xi^-} S R \Delta\Omega_n \varepsilon_n$$

where S is the "survival coefficient" after the beam losses due to Ξ^- decay, absorption and out-scattering in degrader, range straggling. For a 0.5 cm tungsten degrader and 0.7 cm LD_2 stopping target Barnes[29] gives a value $S \approx 13\%$.

Assuming $R \approx 0.1$ $\quad \Delta\Omega_n = 1\pi sr$ $\quad \varepsilon_n = 20\%$ one gets a K^+ neutron coincidence rate of ~ 0.4 events/hour. To obtain *100 events* would thus require ~ *10 days* of running. Compare this figure with the rate obtainable at EHF[38]: for $N_{K^-}(EHF) \approx 10^8/s$ to obtain *100 events* requires ~*15 min* of running.

7.2. Proton research

Bjorken[39] has recently suggested to search for the H particle in high-energy neutral beams at Fermilab. Such beams *(200<E<900 GeV)* should contain appreciable numbers of H if $m_H < 2m_\Lambda$.

The idea of the experiment is shown in Fig. 14. The primary high-energy proton beam by interaction with a target produces H particles. Both charged particles from the reactions and the primary beam are swept away by a magnet. The H particles undergo diffraction dissociation into $\Lambda\Lambda$ pairs in an active dissociator. The dilambdas are then analyzed in a magnetic spectrometer. The very crude estimate of rate indicates that about 10^3 H particles could be detected with $2 \cdot 10^{13}$ protons. The main uncertainty of estimated rate is due to insufficient knowledge of the relative production rate of the two-baryon system over the single baryon at these high energies.

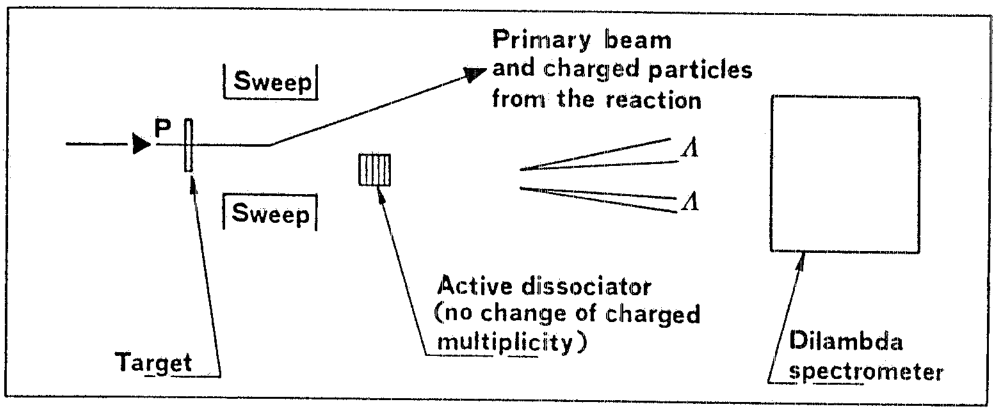


FIGURE 14
Schematic diagram of the suggested experiment searching for H particles in high-energy neutral beams at Fermilab.

8. THE CASE OF CYGNUS X-3

According to a speculative suggestion[40] H particles are emitted by Cygnus X-3 and propagate towards Earth thus giving a possible explanation about the TeV underground muons observed from Cygnus X-3.

Cygnus is a binary system consisting of a compact heavy star and a main sequence star companion. The binary period is 4.8 hours. It is very near the edge of our galaxy, situated in the galactic plane at a distance of $3.7 \cdot 10^4$ parsec from Earth. The compact star is most likely a pulsar accelerating particles up to 10^5 TeV in the 10^{12} Gauss magnetic field near its polar caps. Twice per binary cycle this beam interacts with the atmosphere of the companion producing π° $\pi^\pm$ whose decay γ and ν are beamed towards Earth.

The first hint that perhaps Cygnus X-3 might be opening a new chapter in particle physics came from the Kiel experiment[41] studying the muon content of the γ-showers from Cygnus X-3. The main source of muon in a γ shower is $\gamma N \to \pi X$ photoproduction, followed by $\pi \to \mu$ decay. As photoproduction is a rare process, the muon content of a γ-initiated shower is typically a few percent of that in a hadron-initiated shower. They measured roughly 70% instead. The speculations that the very high energy radiation emitted by Cygnus X-3 is not carried by photons were further encouraged by the observation of the underground muons arriving in direction and in time with the Cygnus X-3 accelerator from the Soudan[42] and NUSEX[43] experiments. The results are listed in Table I from Ref. 44.

TABLE I - Summary of muon energies and fluxes from Cygnus X-3 (units of flux are particles cm^{-2} sec^{-1}).

Detectors	*Soudan 1*	*NUSEX*
Depth of the detector	~2 km.w.e.	~4 km.w.e.
Muon threshold energy	0.65 TeV	3.4 TeV
Primary energy	~6.5 TeV	~34 TeV
CR background rate	$8\cdot 10^{5}$/year	10^{4}/year
Observed flux	$7\cdot 10^{-11}$	$\sim 10^{-11}$
Observed # events	90/year	13/year
Expected μ flux if:		
Primary = hadron	$3\cdot 10^{-12}$	10^{-13}
Primary = photon	$3\cdot 10^{-14}$	$4\cdot 10^{-15}$
$\nu\rightarrow\mu$ flux from Cygnus	$\sim 2\cdot 10^{-15}$	$\sim 2\cdot 10^{-15}$

The assumption that the primary is a hadron or a photon falls short of accomodating the data by over one and three order of magnitude, respectively. Exclusion of neutrinos is presumably the weakest of the negations. It is difficult but not impossible to imagine a "beam dump" capable of producing a sufficiently intense neutrino flux, and lasting for more than a few years. However in Table I is reported the result of a calculation[44] based on the model of a binary system according to which also neutrinos fall short of about 10^4.

Summarizing, it has been already excluded that the parents of the detected muons are: (i) a particle with regular hadronic interactions (proton, neutron, gluino, etc.); (ii) a photon; (iii) a neutrino.

At the same time the properties of the unknown *X*-particles implied by observations are:

(i) *X* is neutral (charged particles forget direction in intergalactic fields);

(ii) $M(X) \leq$ a few GeV;

(iii) $\tau(X) \geq 10^8 s$ enough to cover the 10^4 parsec distance;

(iv) $10\mu b \leq \sigma(XN) \leq 1\ mb$

From these boundaries the following NO-GO theorem must follows: "A particle with the properties (i)÷(iv) should have been discovered by accelerators".

An alternative type of explanation is that X is actually not a particle but a composite system: the parents of the observed muons are now fragments of a larger system. One possibility is that they are neutral bits of quark matter[10,11]. A neutron star may have a large core of strange quark matter[45] but how to extract it and direct it to the Earth is no simple engineering. Quark stars are generally behaved to be covered by a layer of conventional neutron star stuff and a skin of "Coulomb-crystal" matter. This hard shield would not be easily penetrated by a high-energy particle rain from a companion, a step in a possible mechanism for excavating quark matter from the quark-star. Maybe violent astrothermal quark tectonics could, as in Fig. 15, produce jets of strange quark matter from a kind of volcanos or geisers. Quark globs of the right mass and energy can penetrate the atmosphere. They are strongly bound and can interact several times with air nuclei with very little exchange of energy. Finally, the high content of strange quarks would lead to enhanced Λ or K production and thus to a relatively high yield of muons.

A specific version of a stable ensemble of quarks which could be relevant in the context of underground signals from Cygnus X-3 was suggested[40]. This is the bound dihyperon state called the H particle.

Although attractive, some objections can be raised to this hypotesis. There is a puzzling feature of the NUSEX data which is not explained by the proposal. The NUSEX detector sees the muon signal from a region many degrees on a side around Cygnus X-3. There is no mechanism for dispersing H particles over such a large angular range. Moreover, we have already shown in Fig. 3 that, according to a recent calculation[19], also a very tightly bound $H(m_H<2055\ MeV)$ has a lifetime of the order of days rather than years. The latter might be accomodated in the $\Delta S=2$ decay only if m_H is so close to nn threshold to render the particle almost stable. All this poses serious difficulties for the H explanation of the Cignus X-3 events[40].

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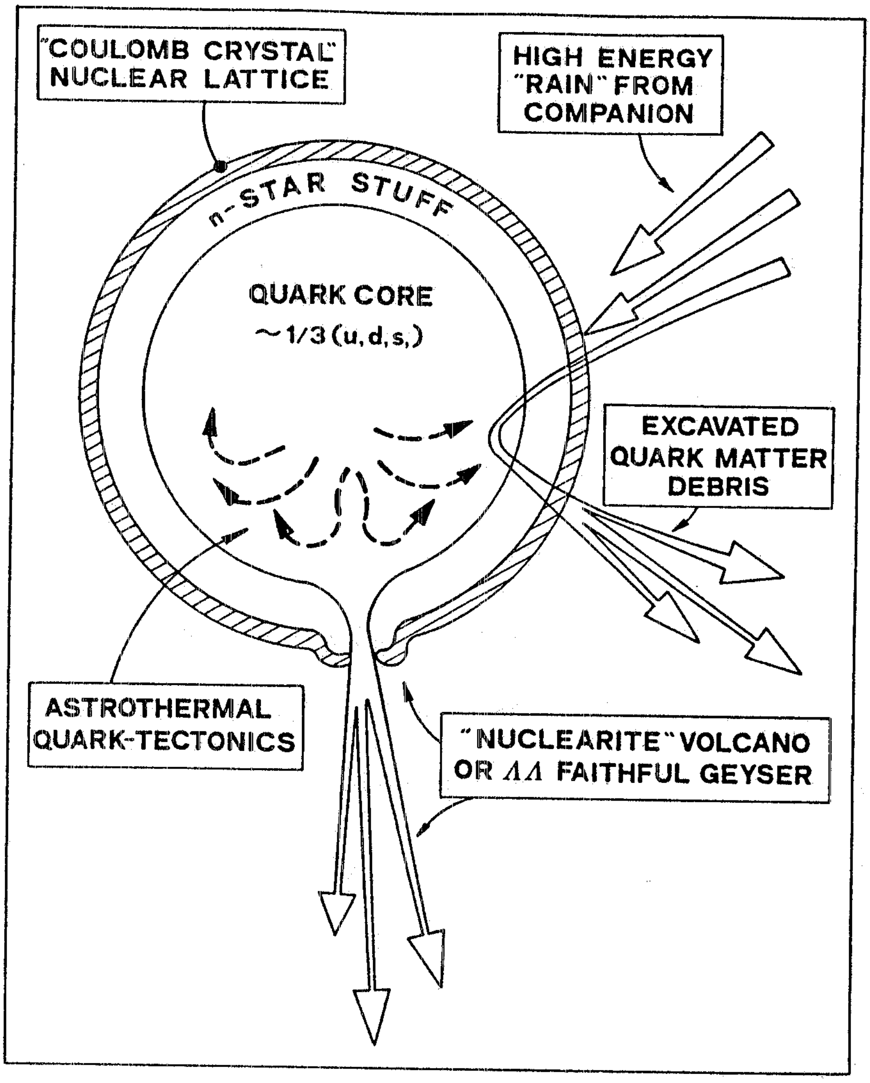

FIGURE 15
Models for a strange-matter reservoir and beams from Cygnus X-3. (From Ref. 46)

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