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Production of the H dibaryon in relativistic heavy-ion collisions

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We discuss mechanisms for the production of exotic hadrons in heavy-ion collisions at high energy, focusing on the doubly strange H dibaryon. In a hadron gas phase, we estimate the H formation rate, per central Si+Au collision, to be of order 0.01 for temperatures $T=120-140$ MeV, based on a conventional coalescence model. If the quark-gluon plasma phase is excited, we expect an enhancement relative to the rate for a hadron gas. We explore various schemes for detecting the H in the high multiplicity background characteristic of heavy-ion collisions. The proposed detection techniques include the measurement of weak decays such as $H \rightarrow \Sigma^- p$, the search for nuclear fragments with anomalous charge/mass ratio, and diffractive dissociation processes such as $Hp \rightarrow \Xi^- pp$.

I. INTRODUCTION

There has been much theoretical attention devoted to the possible existence of stable multiquark systems with baryon number $B=1,2$. For $B=1$, certain pentaquark states ($\bar{c}suud$) are predicted[1] to be stable with respect to strong decay, while for $B=2$, strangeness $S=-2,-3$ there are also candidates.[2,3] The system which has received the most scrutiny is the H particle,[2] with quark structure $uuddss$ ($J^\pi=0^+$, $I=0$). In the context of the Massachusetts Institute of Technology (MIT) bag model,[4] as well as the Skyrmion picture,[5] the H mass m_H lies below the $\Lambda\Lambda$ threshold, so it is stable with respect to strong decays. In the most recent estimate of m_H in lattice gauge theory,[6] it was also found that $m_H < 2m_\Lambda$.

There have been numerous suggestions for experiments to search for the H dibaryon.[7] Two of these, involving the double strangeness exchange (K^-,K^+) reaction,[8] will be carried out at the Brookhaven Alternating Gradient Synchrotron (AGS). In the present article, we suggest that the H can also be produced with measurable rate in relativistic heavy-ion collisions. We find that such collisions at AGS energies would be a suitable place to search for the H, because of the recently reported[9,10] high stopping power in nuclear collisions at 14.5 GeV/nucleon. Another suitable place would be the target fragmentation region corresponding to nuclear collisions at CERN energies 60 and 200 GeV/nucleon.

Unlike other methods, where one detects another particle produced in conjunction with the H (e.g., the monoenergetic neutron in $\Xi^- d \rightarrow Hn$),[8] in heavy-ion collisions it is necessary to detect the H directly, because of the high multiplicity of particles produced in a typical heavy-ion collision event. We address the experimental problems associated with detecting the H in each region of m_H. An encouraging feature of heavy-ion collisions is the rather substantial fraction of events at AGS energies which involve creation of more than one unit of strangeness.[10] The E802 collaboration[10] reports a K^+/π^+ ratio of $\sim 0.25 \pm 0.05$ in the ^{28}Si+^{197}Au collision at 14.5 GeV/nucleon. Preliminary data on the π^+/p ratio indicate a value ~ 0.4. Since the data indicate nearly complete stopping of the projectile in target nuclei heavier than Ag in a typical central collision, one might expect that nearly all target nucleons are finally involved in the reaction process. The total number of active protons is then approximately given by $Z(\mathrm{Si})+Z(\mathrm{Au})=14+79=93$, yielding about 37 π^+'s. If we now assume that the number of K^+'s is approximately the same as the number of K^0's, we find ~ 15 $\bar{s}$ quarks produced which have to be balanced by $\bar{K}$'s and hyperons. Since the number of $\bar{K}$'s turns out to be small as deduced from the measured[10] ratio $K^-/\pi^- \sim 0.05$, we expect roughly ten hyperons with $N_\Lambda \sim 2-4$ in a typical central collision. A different collision scenario is depicted in Fig. 1. Here, the incoming projectile interacts with the nucleons it cuts out of the target. The number of participating nucleons is then $\simeq 108$ for the central Si+Au collision. Since this is approximately one half of the total number of nucleons, we assume that the π^+/p ratio for the "hot" fireball would be ~ 1. Assuming further that the number of participating protons is $\sim \frac{93}{2}$, we end up with the same estimate for hyperon production as before. Although the internal dynamics of these two scenarios would be different, a large amount of negative strangeness is produced in both cases. We predict that a measurable fraction (of order 10^{-2}) of such events involve H formation through baryon-baryon coalescence.

Our goal in this paper is to address H formation in heavy-ion collisions in two different regimes: (i) from a hadronic fireball in thermodynamic equilibrium, and (ii) from a quark-gluon plasma (QGP) phase. If the "strangeness enhancement" signature[11] of the transient QGP phase is indeed preserved through the complicated dynamical evolution of the hadronization process, we expect the production of multistrange objects such as the H to be similarly augmented. For the H, enhanced production in the QGP phase can be driven by the large degeneracy factor associated with color octet three-quark clusters, 8_c, which in the deconfined phase can fuse to form a H.

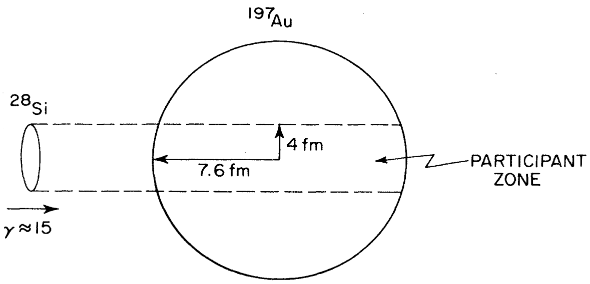

FIG. 1. Sketch of the collision geometry assumed for a typical central collision between Si and Au at ~ 14.5 GeV/c beam momentum.

The outline of this paper is as follows. In Sec. II, we briefly summarize the theoretical predictions regarding the H. In Sec. III, we present our estimates for the H production rate from a hadronic fireball. A system in thermodynamic equilibrium leads to a spectrum of produced hyperons which is peaked at low momentum. This is favorable for H formation via the baryon-baryon fusion mechanism. We calculate the overlap of the momentum space wave function of the H with baryon-baryon plane wave states whose momenta are distributed according to a Maxwell-Boltzmann form. This yields the rate for H formation via baryon-baryon fusion.

Section IV is devoted to a discussion of H production from the QGP phase. We argue that the H rate can be enhanced in the plasma phase due to $8_c \otimes 8_c$ cluster fusion at short distances. We also explore the role of $s\bar{s}$ separation and multibody forces in the baryon-rich regime.

A variety of considerations relevant to experimental searches for the H are given in Sec. V. As the mass m_H drops below successive thresholds ($\Lambda N\pi, \Sigma N, \Lambda N$), the weak decay lifetime τ_H increases, and different experimental techniques are required to detect the H. These are developed in detail for each range of m_H values.

Our perspective and conclusions are given in Sec. VI. We predict an experimentally accessible rate for H production in heavy-ion collisions at Brookhaven AGS energies. Heavy-ion experiments appear to be competitive to K^- induced reactions[8] for H production. The (K^-,K^+) experiments offer the simplicity of two-body kinematics, so direct H detection is not necessary, but they suffer from the low intensity of available K^- beams. Heavy-ion beams have higher luminosity, but one must search for H decay in the debris of complicated multibody final states (a typical charged particle multiplicity for a central heavy-ion collision at AGS energies is about 100–300). This search, although difficult, does seem to be feasible.

II. THE H DIBARYON

The H particle, a dibaryon with the quantum numbers of two Λ particles in a singlet s state, is the most promising candidate for a deeply bound six-quark state.[2] In the spectrum of six-quark bag states, the $S=-2$ sector plays a special role. Indeed, only a six-quark system containing two up (u), two down (d), and two strange (s) quarks can exist in an SU(3)-flavor singlet with spin zero, a configuration which takes maximum advantage of the attraction due to the color-magnetic interactions of quantum chromodynamics (QCD).

The color-magnetic interaction term in the MIT bag model is given by[2,4]

$$H_m = -\frac{\alpha_s}{R}\sum_{i<j} M(m_i R, m_j R)(\lambda\sigma)_i(\lambda\sigma)_j \ , \tag{1}$$

with α_s the QCD coupling constant, R the bag radius, m_i the mass of quark flavor $i=(u,d,s)$, and $\{\lambda_i,\sigma_i\}$ are the color and spin operators of quark i. Ignoring SU(3)-flavor breaking and assuming $m_u=m_s$, one finds a mass term

$$\langle H_m\rangle \sim \tfrac{1}{4}N(N-10)+\tfrac{1}{3}S(S+1)+\tfrac{1}{2}C_c^2+C_F^2 \tag{2}$$

for systems composed of N quarks. The C_i^2 denote the eigenvalues of the quadratic Casimir operators of the corresponding SU(3)$_i$. Obviously, a hadron which is a color, flavor and spin singlet ($C_c^2=C_F^2=S=0$) has the largest color-magnetic attraction.

Within the bag model, it was first noted by Jaffe[2] that the ($uuddss$) state with spin parity $J^\pi=0^+$, isospin $I=0$, which he called the H, would be stable against strong decays, with a mass approximately 80 MeV below the $\Lambda\Lambda$ threshold ($2m_\Lambda=2231$ MeV/c^2).

The color-spin-flavor wave function of the quarks in the H is often written as a linear combination of $8_c\otimes 8_c$ and $1_c\otimes 1_c$ structures, with weights $(\frac{4}{5})^{1/2}$ and $(\frac{1}{5})^{1/2}$ given by SU(3)$_c$ Clebsch-Gordan coefficients.[12] However, it has been pointed out by Donoghue *et al.*[13] and Vento and Gonzalez[14] that the color singlet (1_c) three-quark clusters already form a complete basis of states for the expansion of the H-wave function. Donoghue *et al.*[13] discuss in detail how $8_c\otimes 8_c$ can be reexpressed in terms of the hadron-hadron basis ($1_c\otimes 1_c$). For our evaluation of H production from a hadron gas, we omit the $8_c\otimes 8_c$ "hidden color" component, and expand the H-wave function in the baryon-baryon basis as

$$\Psi_H = \frac{1}{\sqrt{8}}(\Lambda\Lambda+\Sigma^0\Sigma^0+\Sigma^+\Sigma^-+\Sigma^-\Sigma^+ +\Xi^-p+p\Xi^-+\Xi^0 n+n\Xi^0) \tag{3}$$

displaying explicitly the SU(3) flavor singlet character of the H.

A variety of predictions for the mass spectra of strange dibaryons are available. Most early calculations are based on some version of the quark bag model[2,15–20] or the Skyrme-type soliton model.[21–26] A review of the work can be found in Ref. 7. More recently, the H dibaryon mass has been calculated in lattice QCD,[6,27] as well as in the quark cluster/resonating group method.[28,29] Weakly bound deuteronlike $S=-2$ dibaryons can also exist in meson exchange models.[30] In the most complete lattice QCD (Ref. 6) and quark cluster[29] calculations, the H dibaryon lies below the $\Lambda\Lambda$ threshold, and is thus stable against strong decay. However, the predicted binding energy $B_H \equiv 2m_\Lambda - m_H$ is strongly model dependent: The converged lattice QCD result[6] gives a deeply bound $H(m_H \approx 2m_n)$, while the quark/gluon plus meson exchange model of Straub *et al.*[29] yields a weakly bound system ($B_H \approx 20$ MeV).

In the first estimates of the H mass in the context of the Skyrme model,[21–24] all SU(3) coordinates were treated as collective modes, so that the pion and kaon fields were treated on the same footing. This extension of the SU(2) Skyrmion to SU(3) fails to reproduce the $S=-1$, $B=1$ spectrum, unless one adopts an unacceptably small value of the pion decay constant F_π. An improvement was made by Callan and Klebanov,[25] who incorporate strong SU(3) breaking from the beginning by treating the K and π fields on a separate footing: The Lagrangian is expanded to second order in the K field (vibrational approximation), while the π field retains a classical limit. The $B=1$, $S=-1$ configurations then correspond to kaon bound states in the background field provided by the SU(2) Skyrmion.[25] The $B=2$ spectrum in this picture has been calculated by Kunz and Mulders.[26] It is worthwhile noting that the treatment of the K field as a vibration leads to a lowest-lying state $\tilde{H}$ with $J^\pi=0^+$, $I=1$, SU(3) flavor $\{f\}=\{10^*\}$ and mass $m_{\tilde{H}} \approx 2350$ MeV/c^2. The first excited state $\tilde{H}'$ has $J^\pi=0^+$, $I=0$ (which corresponds to the quantum numbers of the H in the Q^6 model), but $\{f\}=\{27\}$ rather than $\{f\}=1$.

Numerous methods have been suggested (for discussion and references see Refs. 7, 31, and 32) for producing the H. The H has been searched for in the reactions

$$pp \to K^+ + K^+ + H \ , \tag{4a}$$

$$\bar{p} + A \to X + H \ . \tag{4b}$$

A cross section limit of $\sigma_H < 50$ nb was obtained by Carroll *et al.*[33] for reaction (4a) in the mass range $2.1 \le m_H \le 2.23$ GeV/c^2. However, if one accepts the estimate $\sigma_H \approx 1$ nb for reaction (4a) due to Badalyan and Simonov,[34] this experiment[33] is not sensitive enough to rule out the H. Similarly, Condo *et al.*[35] quote an upper limit of 9×10^{-5} for the branching ratio of (4b) at rest. Given that $\Lambda\Lambda$ production was not observed in this experiment,[35] it is not surprising that the H, which would have a much smaller rate, is also not seen.

An event which could be interpreted in terms of the formation and decay of an H has been reported by Shahbazian *et al.*,[36] based on their study of high-energy proton-propane interactions. The event is consistent with $H \to \Sigma^- p$ decay, but unfortunately the $\Sigma^- \to \pi^- n$ decay was not seen, so the interpretation is not very convincing.

For the future, there are two approved experiments[8] for the Brookhaven Alternating Gradient Synchrotron (AGS) and another at the KEK laboratory in Japan[37] which utilize the double strangeness exchange (K^-,K^+) reaction to produce an $S=-2$ system. In one AGS experiment[8] the $K^- + p \to K^+ + \Xi^-$ reaction is used to tag the formation of a Ξ^-, which is then slowed down electromagnetically and captured in deuterium. One then looks for a monoenergetic neutron from the process $\Xi^- + d \to n + H$. In the other, the reaction $K^- + {}^3\mathrm{He} \to K^+ + H + n$ is studied.[8] Status reports on these two (K^-,K^+) experiments have been given by Franklin[8] and Barnes.[32] It has also been proposed[38] that the H may be found in $\bar{p}$-nuclear reactions. The case of $\bar{p} + {}^3\mathrm{He} \to (KK\pi)^+ + H$ has been discussed in detail by Guaraldo.[31]

The observation of two $\Lambda\Lambda$ hypernuclear events[39,40] is often taken as an argument against the existence of the H. However, Kerbikov[19] has argued that the sequential $\Lambda \to p\pi^-$ decays observed for these events does not rule out a weakly bound H. A similar conclusion follows if the H is deeply bound ($m_H \le 1900$ MeV/c^2), since then the strong decay ${}^6_{\Lambda\Lambda}\mathrm{He} \to H + \alpha$, for instance, is greatly suppressed by wave function effects.[7] We regard the existence of the H as an open question, and urge further experimental searches utilizing relativistic heavy-ion collisions.

III. PRODUCTION OF THE H DIBARYON FROM A BARYON-RICH HADRON GAS

We now estimate the production probability of the H dibaryon in a baryon-rich gas of hadrons. The observation of large transverse energy production at AGS energies ($E_{\mathrm{lab}} \sim 10-15$) GeV/nucleon suggests a large degree of stopping. Accordingly, we assume that complicated multibody interactions lead to the formation of a hot dense hadron fireball, which could consist either of deconfined quarks and gluons or matter composed of individual hadrons, depending on the energy or baryon density reached during the initial stage of the collision process. Here we will assume that the fireball consists of hadrons only.

In our approach, we assume that, for example, the collision process for a ${}^{28}\mathrm{Si} + {}^{197}\mathrm{Au}$ collision proceeds roughly as indicated in the sketch in Fig. 1. Depending on the impact parameter, a certain fraction of projectile and target nuclei will interact and form a fireball. We assume now that in a typical central collision the 28 projectile nucleons will initially interact with the number of nucleons the projectile cuts out from the target. If the projectile is stopped by the nuclear matter in that tube, we can estimate the number of baryons contained in the fireball. From simple geometrical considerations, one finds the number of participating nucleons to be

$$A_P + 1.5(A_P^2 A_T)^{1/3} \simeq 108 \ .$$

Assuming further that the hadrons in the fireball reach chemical equilibrium, we can calculate the mean number of hyperons for a given baryon density as a function of temperature. This is shown in Fig. 2 for two different baryon densities $\rho_B/\rho_0 = \frac{3}{4}$ and $2.5(\rho_0 = 0.15\ \mathrm{fm}^{-3})$. We see that for any reasonable temperature larger than $T \gtrsim 100-110$ MeV the number of produced Λ's and Σ's is larger than two. One should note that this number is very close to what we have extrapolated from the AGS data. We regard this as a check that our assumption of chemical equilibrium is reasonable. We note that the chemically equilibrated hadron gas would even predict a quite appreciable number of Ξ's.

In the case of a Au+Au collision (see Fig. 3), we have ~ 394 participating nucleons in a central collision. As-

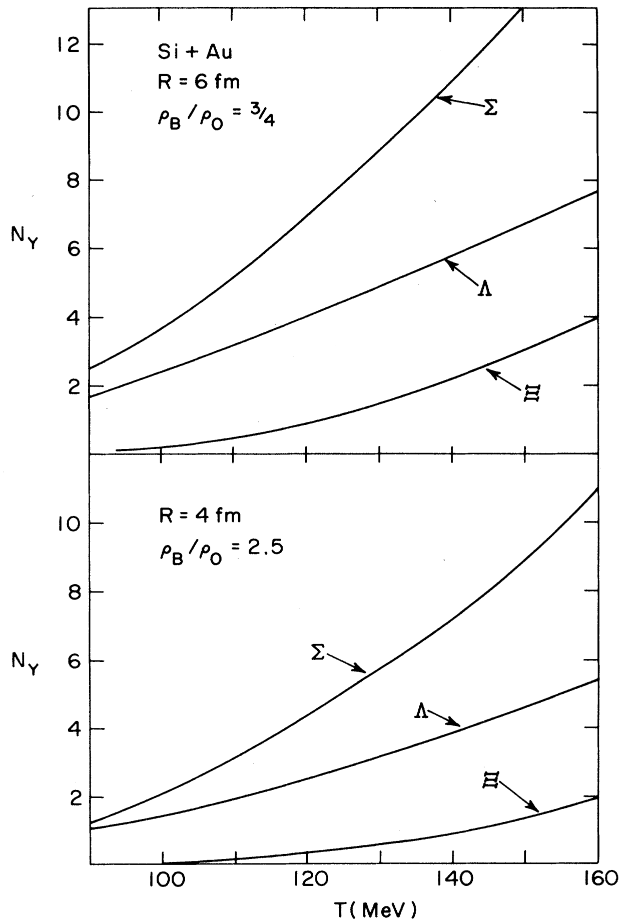

FIG. 2. Abundances of hyperons as calculated in the chemically equilibrated hadron gas as a function of temperature T. In the top half of the plot, we have fixed the baryon density to be $\rho_B \simeq \frac{3}{4}\rho_0$ ($\rho_0 = 0.15$ fm^{-3}= nuclear matter ground-state density), which corresponds to a fireball radius of 6 fm when the number of participating nucleons is $\simeq 108$, appropriate to a Si+Au collision. In the bottom half, $\rho_B \simeq 2.5\rho_0$ is chosen.

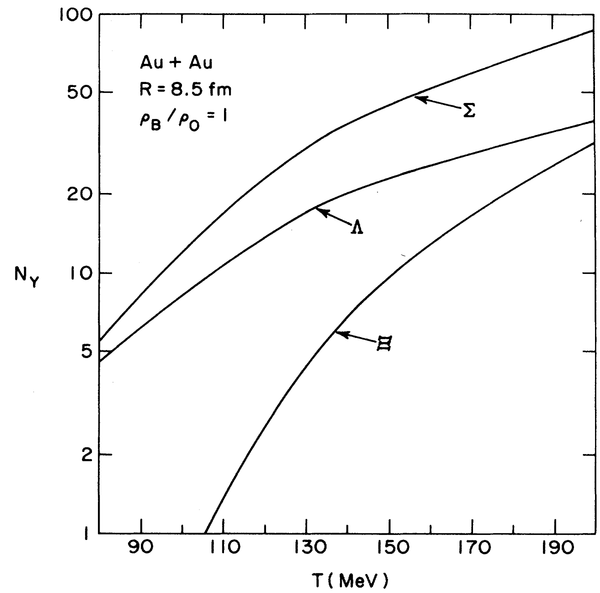

FIG. 3. The same as Fig. 2, but now for a Au+Au collision with $\rho_B = \rho_0$, resulting in a fireball radius of $R \simeq 8.5$ fm.

suming a baryon density $\rho_B \simeq \rho_0$, we estimate that roughly 10–20 Λ's 15–40 Σ's (summed over all charged states), and even 1–5 Ξ's will be produced at reasonable temperatures.

Thus, we can expect a relatively large number of events where at least two hyperons are produced, opening up the possibility of H production via baryon-baryon fusion. This process is described in terms of the momentum space overlap of the baryon wave functions with that of the H. Microscopically, the formation of the H results from the fusion of two three-quark bags into a six-quark bag. Following Aerts and Dover,[41] the amplitude for this fusion process is approximated by,

$$A_{B_1B_2} = \frac{1}{(2\pi)^{18}} \int \frac{d^3k_1}{2E_1} \cdots \frac{d^3k_6}{2E_6} \Psi_{H,\mathbf{k}_H}(\mathbf{k}_1 \cdots \mathbf{k}_6) \otimes \Psi_{B_1,\mathbf{k}_{B_1}}(\mathbf{k}_1 \cdots \mathbf{k}_3) \Psi_{B_2,\mathbf{k}_{B_2}}(\mathbf{k}_4 \cdots \mathbf{k}_6) \tag{5}$$

with the momenta labeled 1–6 now referring to the quark momenta. For the wave functions Ψ an oscillator approximation for N quarks confined to a bag of radius R was used:

$$\Psi_{N,\mathbf{k}}(\mathbf{k}_1 \cdots \mathbf{k}_N) = \delta^{(3)}\left[\mathbf{k} - \sum_i^N \mathbf{k}_i\right] N^{3/4} \left[\frac{R^2}{\pi}\right]^{3(N-1)/4} \exp\left[-\frac{R^2}{2}\sum_i^{N-1} \bar{k}_i^2\right], \tag{6}$$

where $\{\bar{k}_i\}$ is an appropriate set of $N-1$ relative momenta and $\mathbf{k}$ is the corresponding total momentum. If the H is not far below the $\Lambda\Lambda$ threshold, the violation of energy conservation implicit in our approach is not a serious problem. Hence, one finds[41]

$$A(\mathbf{k}_{B_1}, \mathbf{k}_{B_2}) = (2\pi)^3 (2E_H 2E_{B_1} 2E_{B_2})^{1/2} \delta^{(3)}(\mathbf{k}_H - \mathbf{k}_{B_1} - \mathbf{k}_{B_2}) \Gamma(\mathbf{k}_{B_1} - \mathbf{k}_{B_2}), \tag{7}$$

where

$$\Gamma(\mathbf{k}_{B_1} - \mathbf{k}_{B_2}) = \Gamma_0 \left[\frac{2R_{B_1}R_H}{R_H^2 + R_{B_1}^2}\right]^3 \left[\frac{2R_{B_2}R_H}{R_H^2 + R_{B_2}^2}\right]^3 \left[\frac{8\pi R_H^2}{3}\right]^{3/4} \exp\left[\frac{-R_H^2}{12}(\mathbf{k}_{B_1} - \mathbf{k}_{B_2})^2\right]. \tag{8}$$

The factor Γ_0 is due to color-flavor-spin recoupling.

We have neglected changes in the wave functions due to medium effects as caused by the surrounding hot and dense hadron matter, as well as dynamical distortions of the wave functions as the bags approach each other. These approximation seem adequate for our first-order estimate.

In Eq. (8) R_H, R_{B_1}, and R_{B_2} represent the oscillator parameters for the H, and the two fusing baryons. We follow Ref. 41 and choose all parameters $R_{B_1} \simeq R_{B_2} \simeq R_H = 0.8$ fm. Within these approximations we calculate the total number of H's in the rest system of the fireball as

$$N_H = \sum_{i,j} \int \frac{d^3k_i}{2E_i} \frac{d^3k_j}{2E_j} \frac{d^3k_H}{2E_H} f_i(\mathbf{k}_i) f_j(\mathbf{k}_j) |A(\mathbf{k}_{B_i}, \mathbf{k}_{B_j})|^2 , \tag{9}$$

where the sum over $\{i,j\}$ includes $\Lambda\Lambda$, $\Sigma\Sigma$, and ΞN pairs.

The $f(\mathbf{k})$ are the phase space distribution functions of the fusing baryons. In squaring the amplitude we have to take care of the δ function. We use the prescription

$$[\delta^{(3)}(\mathbf{k}_H - \mathbf{k}_{B_1} - \mathbf{k}_{B_2})]^2 = \frac{V}{(2\pi)^3} \delta^{(3)}(\mathbf{k}_H - \mathbf{k}_{B_1} - \mathbf{k}_{B_2}) , \tag{10}$$

where V is taken to be the volume of the fireball.

Consistent with our assumption that the projectile and target nuclei form a complex intermediate state (fireball) in kinetic and chemical equilibrium, we approximate $f(\mathbf{k})$ by a thermal distribution. For our estimate it is sufficient to utilize the nonrelativistic Boltzmann approximation, which in the rest system of the fireball is given by

$$f_i(\mathbf{k}_i) = \frac{1}{(2\pi)^3} \exp\left[-\frac{\mathbf{k}_i^2}{2m_i T} - \frac{m_i}{T} \right] \exp\left[\frac{\mu_B}{T} + \frac{\mu_S}{T} S_i \right] \tag{11}$$

with μ_B, μ_S the chemical potentials introduced to control the baryon number and strangeness content of the fireball. S_i corresponds to the strangeness quantum number of the baryon under consideration. We then obtain

$$N_H = 4 \left[\frac{2\pi}{3} \right]^{1/2} \left[\frac{R_H}{R} \right]^3 \times \sum_{ij} N_i N_j \Gamma_{ij}^2 \left[1 + \frac{R_H^2 m_i T}{3} + \frac{R_H^2 m_j T}{3} \right]^{-3/2} , \tag{12}$$

where

$$N_i \equiv \frac{\left[\frac{4\pi}{3} R^3 \right]}{(2\pi)^{3/2}} (m_i T)^{3/2} \exp\left[-\frac{m_i}{T} + \frac{\mu_B}{T} + \frac{\mu_S}{T} S_i \right] \tag{13}$$

is the total number of baryon species i in the nonrelativistic Boltzmann approximation and R is the radius of the fireball. Formula (12) is not difficult to understand. The total number of H's grows proportionally to the product of the number densities of the baryons from which it is produced. The factor $(R_H/R)^3$ reflects the fact that the baryons must be found in the volume occupied by the H. For given radius R_H and large volume, the density of baryons decreases for fixed total number. Hence, also the probability to form a H decreases. A more subtle limit to understand is the case where the H is very large, i.e., $R_H \to \infty$. In this case the formation probability depends on $(m_i T)^{-1}$, which can be interpreted as a penalty factor one has to pay for embedding high momentum modes into a spatially extended and massive particle. This suppresses H formation by particles with high momenta or high temperatures.

In Fig. 4 we have shown the number of H's produced for some typical parameters (baryon density, temperature) we expect to be relevant for a hadron gas produced in a nuclear collision. Here we have assumed that each central collision leads to formation of a hadronic fireball. For illustration we have chosen Si+Au and Au+Au collisions. We have also shown the separate $\Lambda\Lambda$, $\Sigma\Sigma$, and ΞN contributions to H formation. Although the ΞN channel is the largest part of the H-wave function, the fusion rate is dominated at low temperatures by the Λ's and Σ's. For $T > 130$ MeV, the ΞN channel becomes significant, and for very large T it contributes $\frac{1}{2}$ of the H formation rate.

One should note from Eq. (12) that the number of H's is proportional to the factor

$$\exp\left[-\frac{1}{T}(m_i + m_j) \right]$$

coming from the Boltzmann factor of the two fusing baryons. If we had taken energy conservation into account properly, we would expect to find a factor $\sim e^{-m_H/T}$.

The effect of our approximation is clearly shown in Fig. 5, where we have plotted the H abundance as calcu-

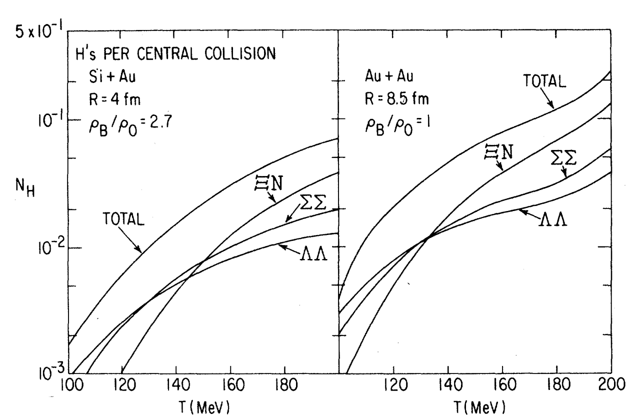

FIG. 4. The total number of H's per central collision and the individual baryon-baryon fusion contributions are shown as a function of temperature for Si+Au collisions (left-hand side) and Au+Au collisions (right-hand side).

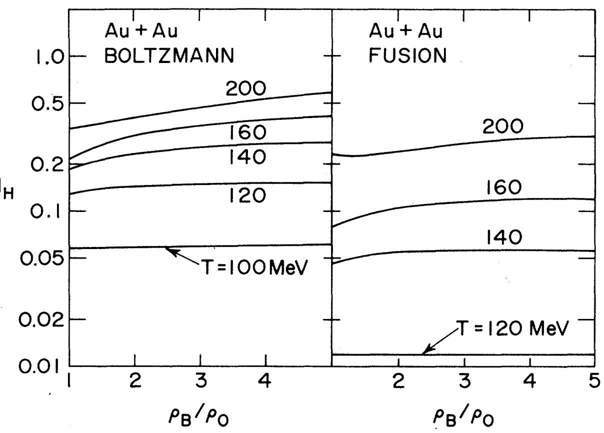

FIG. 5. The abundance of H's as calculated in the (a) Boltzmann particle approximation and (b) in the fusion model as a function of baryon density ρ_B, for various values of the temperature T. As an example, a Au+Au central collision was chosen.

lated from our fusion process, compared with that obtained if one treats the H as a Boltzmann particle with a mass of 2.1 GeV. The error introduced in our approach is most pronounced for low temperatures because T is of the order of $m_i+m_j-m_H$. Our fusion estimate can be viewed as a lower limit for H production, especially in the low temperature domain. The Boltzmann particle limit for the H might be relevant if the H is considered to be a hadron and not a loosely bound state as, for example, the deuteron. However, it is not clear that the collisions of the H with other hadrons in the fireball are sufficient for the H to reach its chemical equilibrium value (i.e., Boltzmann particle case). Keeping all these uncertainties in mind, we expect our fusion model estimate to be adequate as an order of magnitude estimate if a pure hadron gas is formed. One might speculate that the Boltzmann particle limit for the H might be relevant for the case when a quark-gluon plasma is formed during the initial stage of the collision process. In this situation it is possible that the H would be produced directly from the plasma with its chemical equilibrium abundance. In this case, the H would stay in equilibrium with the rest of the hadrons until freeze out. At reasonable freeze-out temperatures in the range 100–120 MeV, the Boltzmann H is approximately a factor of ~ 10 more abundant than the "fusion" H. Although this cannot prove plasma formation, it constitutes a strong hint that an initial state existed where a number of quarks occupied a small space-time region.

Assuming that essentially each central collision leads to the formation of a hadronic fireball, we only have to multiply the number of H's produced for given fireball parameters by the cross section σ_c for having a central collision. For example, for a Si beam incident on a heavy target, we estimate $\sigma_c \approx \pi R^2$, with $R = (\frac{5}{3})^{1/2}\langle r^2\rangle_{\mathrm{Si}}^{1/2}$, or $\sigma_c \approx 500$ mb. For $T \approx 130-150$ MeV, appropriate to a fireball formed at AGS energies, we obtain $\sigma_H \approx N_H \sigma_c \approx 5-10$ mb. This level of H production should be measurable at the AGS. Note that σ_H will drop rapidly as the energy decreases: At 5 GeV/nucleon, Sano *et al.*[42] have estimated $\sigma_H \approx 2.6$ μb for Ne+Ne collisions. Our model cannot be sensibly extrapolated to such low energies, since the assumption of chemical equilibrium for the hyperons breaks down.

IV. HADRONIZATION AND EXOTICS

Quantum chromodynamics lattice calculations[43] predict that, under conditions of high temperature and baryon density, hadronic matter undergoes a transition to a deconfined chirally symmetric phase composed of quarks, antiquarks, and gluons. According to current expectations, the required conditions can be attained in ultrarelativistic collisions of sufficiently heavy nuclei. First experiments searching for a quark-gluon plasma have been already performed at CERN and the Brookhaven AGS.[44] A plasma produced in such a way, however, is expected to be a highly unstable state which undergoes a complex dynamical evolution. In the past several years, a considerable effort has been devoted to obtaining the space-time picture of quark matter formation and decay.

Unfortunately, not much is known about the hadronization of the plasma on the microscopic level, although some model calculations based on phenomenological descriptions of color confinement have been attempted to describe hadronization from the outer surface[45,46] (meson radiation) as well as meson formation during hadronization in bulk based on a simple combinatorial picture.[11,47]

Within these models, the crucial assumption is made that the plasma hadronizes directly into the low mass hadronic states. Although such an assumption seems reasonable for the radiation of mesons from the plasma surface, it is no longer obvious for that portion of quark-gluon matter which hadronizes in bulk at a late stage when the plasma has reached the phase boundary with hadron matter. At this stage, one can no longer expect that color screening is operative in such a way that it forbids short-range color correlations.

Guided by the success of phenomenological models in hadron spectroscopy based on the short-range color magnetic interaction arising from quark-gluon exchange, one might expect that close to the phase transition region the interaction mediated by gluons favors certain quark configurations due to their color-magnetic attraction. The relevance of diquark clustering in a quark-gluon plasma has been discussed recently,[48] and it was found that in cold matter a buildup of diquark clusters might be favored at certain densities. Gleeson *et al.*[49] have discussed the special role of three- and four-point gluon vertices and argued that these could even dominate the interaction between certain color configurations at high density. It was found that the three-quark interaction annihilates overall color-singlet and decouplet states, and therefore does not contribute to the energy of baryonic states[49] but it does not vanish in the case of the color octet three-quark state.

Even more interesting from our point of view is the idea by Barrois[50] that in dense quark matter an unusual variant of superconductivity might show up. Instead of pairing of fermions (quarks), the quarks may clump into units of $6N$ quarks in the case of SU(N) flavor quark matter. Adding the requirement that, in the absence of any symmetry breaking, the condensing unit must be a singlet under every symmetry group, this clumping would favor the formation of multiquark states like the H in the case of SU(3) flavor symmetric quark matter.

From these considerations, one might expect that in a quark-gluon plasma close to the hadronization point, a considerable amount of clustering takes place. The bulk of all possible singlet configurations will not represent observable resonant states. These unstable states will be so broad (and plentiful) that one cannot expect to isolate any specific one of them: They will build up a rather structureless continuum of final particles. In this sense they might serve as "doorway" states for hadronization from the plasma. The H, if it is stable, stands out with respect to the continuum, and could be formed in the plasma through the fusion of color octet three-quark clusters. This could lead to an enhanced production rate.

Another mechanism which would be relevant for the formation probability of the H dibaryon from quark matter is the recent speculation[51] that in baryon-rich quark-gluon plasma an enrichment in the s-quark content of the plasma can result. This mechanism was found to be further supported by cooling and net strangeness enrichment due to prefreezeout evaporation of pions and K^+,K^0, which carry away entropy and antistrangeness from the system. Thus, it was found that metastable droplets (which decay by weak interactions) of strange matter ("strangelets") can be formed during the phase transition. Even if large clumps of strange quark matter are not stable with respect to strong interactions, a stable H dibaryon would be one of the most favored products of "strangelet" decay. In this case, H production would not require one to bring together two hyperons which are separated in phase space, and again one expects some enhancement of H dibaryon production relative to that from a gas of individual hadrons.

V. EXPERIMENTAL DETECTION OF THE H

The experimental challenge is to detect the H in the high multiplicity background of particles produced in central collisions of relativistic heavy ions. If we are looking for the weak decay of the H, we must take cognizance of the lifetime τ_H and the various branching ratios, which depend sensitively on the mass m_H. For each range of m_H, different techniques are appropriate. In this chapter, we outline the experimental requirements.

The thresholds for the possible decay modes of the H are given in Table I, together with the strangeness change ΔS for each channel. For $m_H > 2m_\Lambda$, the H decays strongly. In this domain, one could look for broad enhancements in $\Lambda\Lambda$, ΞN, or $\Sigma\Sigma$ invariant mass plots. This would be a difficult endeavor in a high multiplicity environment. Here we focus on the case $m_H < 2m_\Lambda$, where the H decays weakly. Note that for $m_H < 2m_n$, the H would be absolutely stable (neglecting processes which violate baryon number conservation). However, in this case, deuterium would decay via the $\Delta S = 2$ weak process $d \to He^+\bar{\nu}_e$ or, phase space permitting, the strong process $d \to H\pi^+$. We do not consider this case further, and restrict our attention to the range $2m_n < m_H < 2m_\Lambda$.

TABLE I. Possible decay modes of the H.

Channel	Threshold mass (MeV/c^2)	ΔS
$\Lambda\Lambda$	2231	0
$\Lambda N\pi$	2192	1
$NN\pi\pi$	2152	2
ΣN	2134	1
Λn	2055	1
$NN\pi$	2016	2
nn	1879	2

Several estimates of weak decay lifetimes for the H exist in the literature. The most systematic approach is due to Donoghue *et al.*,[13] who have predicted τ_H in the entire range $2m_n < m_H < 2m_\Lambda$. A plot of their results is shown in Fig. 6. They may be summarized as follows.

(i) For $\Delta S = 1$ with $m_H > m_\Sigma + m_N, \tau_H$ is in the range $10^{-8}-10^{-9}$ sec, at least an order of magnitude longer than the free space Λ lifetime $\tau_\Lambda \approx 2.5\times10^{-10}$ sec. Even near the $\Lambda\Lambda$ threshold, where the mode $H \to \Lambda N\pi$ comes into play, $\tau_H > 10^{-9}$ sec is predicted by Donoghue *et al*,[13] in contrast to the naive estimate $\tau_H \approx \tau_\Lambda$.

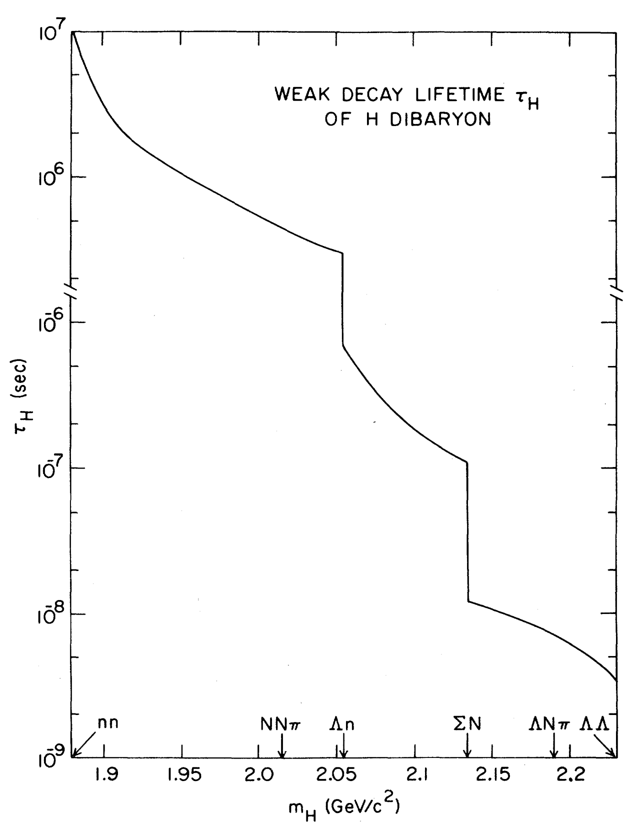

FIG. 6. The weak decay lifetime τ_H of the H dibaryon as a function of the H mass, m_H (adapted from Ref. 13).

(ii) For $\Delta S=1$ parity conserving decays, the $\Delta I=\frac{3}{2}$ transition is found to dominate the $\Delta I=\frac{1}{2}$ one,[13] and thus the ΣN final state is preferred over Λn. This is due to the Pauli principle, which forces the *s*-wave six-quark final state to be the $\{27\}$-plet of SU(3). If confirmed experimentally, the decay $H\rightarrow\Sigma N$ would represent the first example in which the well-known $\Delta I=\frac{1}{2}$ rule for the nonleptonic decays is violated. As a consequence, the value of τ_H becomes markedly longer (10^{-6} to 10^{-7} sec) below the ΣN threshold.

(iii) For $\Delta S=2$, the H lifetime is of the order of several days (τ_H starts from 3×10^5 sec just below the Λn threshold and exceeds 10^6 sec for $m_H<1.96$ GeV/c^2).

We now discuss several classes of experiments for detecting the H. These are (i) detection of the decay mode $H\rightarrow\Sigma^-p$; (ii) diffractive dissociation of the H through collision with a nuclear target far from the production point; (iii) search for charged nuclear fragments containing the H in a bound state. Note that any trigger which can select events with multiple strange particle production (e.g., multiple K^+ mesons) will enhance H production.

In the mass range $2134<m_H<2231$ GeV/c^2, the decay channel $H\rightarrow\Sigma^-p$ is present. Even close to the $\Lambda\Lambda$ threshold, the YN decay modes are predicted by Donoghue *et al.*[13] to be larger than $\Lambda N\pi$. They find $\Gamma(H\rightarrow YN)/\Gamma_{\rm tot}\approx0.7$ for $m_H\sim2m_\Lambda$. The relative branching ratios for the $H\rightarrow YN$ modes, neglecting phase space factors, are found to be[13]

$$\begin{aligned}&\Gamma(\Sigma^-p)\approx\Gamma(\Sigma^0n)\ ,\\&\Gamma(\Lambda n)\approx\tfrac{1}{3}\Gamma(\Sigma^0n)\ ,\end{aligned}\tag{14}$$

whereas the usual octet rule for the weak Hamiltonian would lead to[13]

$$\Gamma(\Sigma^-p):\Gamma(\Sigma^0n):\Gamma(\Lambda n)=5:3:2\ .\tag{15}$$

In either case, the Σ^-p mode features prominently in the decay of the H. For an H produced at rest in the c.m. system for 15 GeV/nucleon ^{32}S incident on a ^{197}Au target, the decay length of the H is about 1.5 m for $\tau_H\approx0.5\times10^{-8}$ sec. This is a favorable situation for detection. The characteristic signature of the Σ^-p decay mode is a "V" originating more than 1 m from the target. The Σ^- will travel about 10 cm before decaying into a π^- and an (unobserved) neutron. Most Λ's will decay closer to the target. The decay of the K_L, with $\tau=5.2\times10^{-8}$ sec, will also contribute "V's," but they will be distinguishable from H decay by detailed analysis.

In the range $2055<m_H<2134$ MeV/c^2, the decay mode $H\rightarrow\Lambda n$ prevails. Here the decay length is 100–500 m. A Λ decay this far from the production target would be highly unusual. Due to the long decay length, the techniques described in the next paragraph may also be applicable.

For $1879<m_H<2055$ MeV/c^2, only the mode $H\rightarrow nn$ with $\Delta S=2$ is available. Here $\tau_H>3\times10^5$ sec,[13] so the H is stable for the purposes of laboratory experiments. A possible approach consists of looking for nuclear fragments of abnormal charge to mass ratio Q/M, for example the systems consisting of an H attached to a nuclear core consisting of Z protons and N neutrons. Writing $m_H=(2+x)m_n$ with $0<x<2(m_\Lambda/m_n-1)$, we have $(A=N+Z)$

$$\frac{Q}{M}\approx\frac{Z}{A+2+x}\ ,\tag{16}$$

where we have neglected the binding energy and expressed M in units of the neutron mass m_n. In Table II, we display the range of Q/M values for various choices of nuclear core. Normal $N=Z$ nuclei have $Q/M\approx\frac{1}{2}$, whereas H nuclei have $Q/M<\frac{1}{3}$. Note that some neutron-rich metastable nuclear species (e.g., ^{8}He, with $Q/M\approx\frac{1}{4}$) also have small Q/M values, but these would also be rather infrequently produced and could be distinguished from H nuclei by a measurement of M.

In the preceding discussion, we have assumed that the H will bind to a nuclear core, but it is unclear whether the light systems listed in Table II are actually stable against strong decay. We can offer no detailed model for the binding mechanism, except to note that exchange of an $I=0$ scalar two-pion system (the σ) between the H and the core will yield an attractive potential. In addition, there will be an attractive contribution coming from second-order pion exchange, i.e., $HN\xrightarrow{\pi}H_1'N\xrightarrow{\pi}HN$, where H_1' is the flavor octet 1^+, $I=1$ first excited state of the H.[2,15] This is analogous to the chains $NN\xrightarrow{\pi}\Delta N\xrightarrow{\pi}NN$ or $\Lambda N\xrightarrow{\pi}\Sigma N\xrightarrow{\pi}\Lambda N$, which contribute attraction to the N- nucleus or Λ-nucleus potentials, respectively. The contribution of this mechanism to the H well depth has been discussed in Ref. 52. Note that since the H is much more massive than the nucleon, only a rather shallow attractive potential is needed for it to bind to a nucleus. For a square well potential of depth V_0 and radius $R=r_0A^{1/3}$, the requirement for binding the s wave is $V_0\geq\pi^2/8\mu R^2$, where

$$\mu=m_H/(1+m_H/Am_n)$$

is the reduced mass. For $r_0\approx1.4$ fm, the minimum value of V_0 varies from 16 MeV for $A=2$ to less than 3 MeV for $A\geq11$. Thus, a small well depth suffices to bind the H.

The production of H nuclei may proceed via a second-order coalescence mechanism in which $\{d,t,\alpha\}$ clusters are formed, subsequently fusing with the H to form $\{^4_HH,^5_HH,^6_H\mathrm{He}\}$. Bando *et al.*[53] have performed such calculations for $\Lambda\Lambda$ hypernuclear formation. There will be a substantial penalty factor for such a higher-order fusion

TABLE II. Charge to mass ratios of bound nuclear systems containing an H.

System	Q/M range
$H+p(^3_H\mathrm{H})$	0.296–0.33
$H+d(^4_H\mathrm{H})$	0.11–0.125
$H+t(^5_H\mathrm{H})$	0.186–0.2
$H+\alpha(^6_H\mathrm{He})$	0.313–0.33

process. We estimate a rate of order 10^{-4} to 10^{-5} for ${}^{4}_{H}H$ formation, relative to that for the H itself. Thus such objects will be rare, but their signature of an anomalously low Q/M ratio is quite distinctive.

Another rather speculative possibility is that of an $H^2=(HH)_{L=0}$ bound state $[J^\pi=0^+,\ I=0]$, an $S=-4$ analog to the α particle. Only a small amount of attraction is required to bind such an object. Some of this would be provided by second-order pion exchange $(HH\xrightarrow{\pi}H_1'H_1'\xrightarrow{\pi}HH$, analogous to $NN\rightarrow\Delta\Delta\rightarrow NN)$. The binding of a state with $B=4$, $S=-4$ has been considered by Issinskii *et al.*[54] in the context of the SU(3) Skyrme model. They find a binding energy of 15–20 MeV with respect to the $H+H$ threshold.

In heavy-ion collisions, such an object could be formed by a second-order fusion process, with a rate estimated at $10^{-6}-10^{-7}$ per central collision at AGS energies. The weak decays of the HH system could provide spectacular experimental signatures, for instance $H^2\rightarrow 4\Lambda\rightarrow 4p+4\pi^-$ (an eight-prong event) or $H^2\rightarrow\Sigma^- pH\rightarrow\Sigma^-+3p+2\pi^-$ (a six-prong event).

An alternative method for detecting the H is diffractive dissociation, a suggestion attributed to Bjorken in the review of Guaraldo.[31] Mechanisms for this process are shown in Fig. 7. For example, the H could decay virtually into two baryons, and then one of the baryons interacts with a nucleon in the target, and transfers the necessary four momentum to go on shell. The top two graphs in Fig. 7 are of this type. Alternatively, an Hp collision could result in the excitation of the H to an excited state H^*, which subsequently decays into two Λ's. This is illustrated in the bottom graph of Fig. 7. The most favorable channel for experimental detection would be ΞN, for two reasons: (i) the $H\rightarrow\Xi N$ recoupling coefficient is larger than for $\Lambda\Lambda$ or $\Sigma\Sigma$ (see Sec. II), and (ii) the final state NN amplitude t is larger than for YN scattering. One would then look for the decay chain

$$Hp\rightarrow pp\Xi^-,\quad \Xi^-\rightarrow\pi^-\Lambda,\quad \Lambda\rightarrow\pi^- p\ . \tag{17}$$

The final state would contain five charged particles $(ppp\pi^-\pi^-)$. The same final state occurs for $Hp\rightarrow p\Lambda\Lambda$, but in this case there are two V's. The dissociation mechanism should work well if the m_H is near $2m_\Lambda$, but the cross section will drop off rapidly as the H becomes strongly bound. The dissociation of the H into $\Sigma^+\Sigma^-$ or $\Sigma^0\Sigma^0$ is not favorable for study, because of the presence of two or more neutral particles in the final state.

The threshold momentum for H dissociation on a hydrogen target is shown in Fig. 8, as a function of m_H. The average lab momentum $\langle p_H\rangle$ of the H emitted from a typical heavy-ion fireball is well above threshold. To see this, we note that the total mean lab energy $\langle E_H\rangle$ is

$$\langle E_H\rangle=\gamma[m_H+\tfrac{3}{2}T+\beta(3m_HT)^{1/2}\cos\theta]\ , \tag{18}$$

where

$$\gamma=\cosh y\ ,$$

$$\beta=\frac{(\gamma^2-1)^{1/2}}{\gamma}\ ,$$

and y is the fireball rapidity, θ the c.m. emission angle of

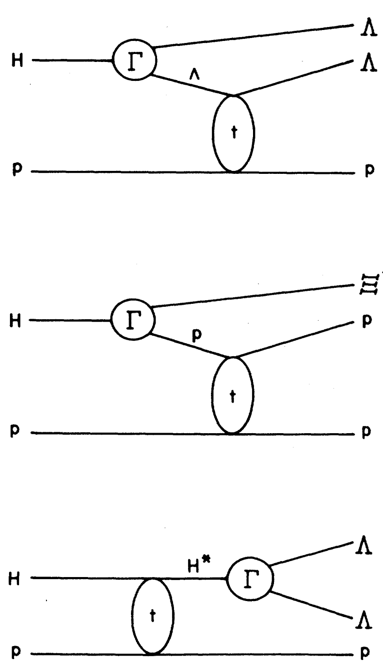

FIG. 7. Diagrams for H dibaryon dissociation.

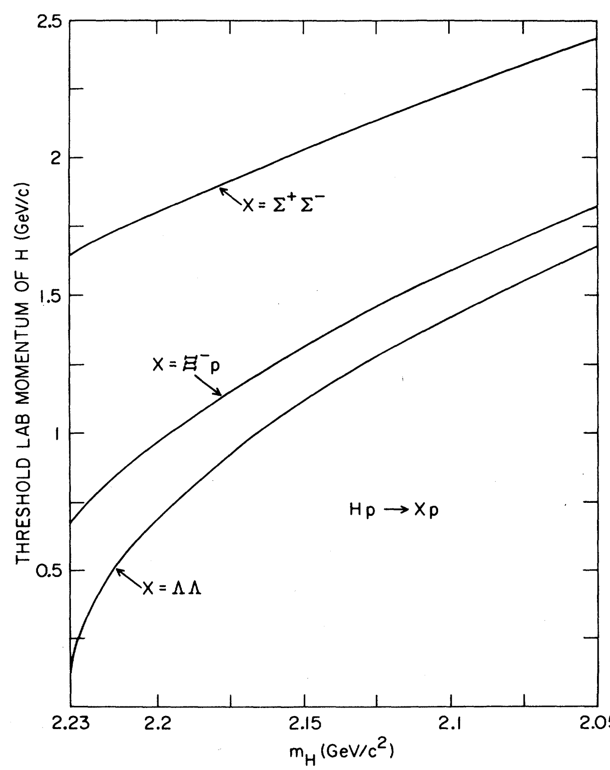

FIG. 8. Threshold momentum for H dissociation on a proton target into $\Lambda\Lambda$, $\Xi^- p$, or $\Sigma^+\Sigma^-$, as a function of the H mass.

the H. For $^{16}O+^{197}Au$ collisions at 15 GeV/nucleon we have $\gamma=1.83$, $\beta=0.84$. The maximum lab emission angle for the H is about 17.8°, corresponding to $\theta=120°$. For this case, we have $\langle p_H \rangle \approx \frac{3}{2} m_H$, while for $\theta=0°$ emission, the result is $\langle p_H \rangle \approx \sqrt{8} m_H$. In all cases, $\langle p_H \rangle$ is more than 3 GeV/c, so we are far above the thresholds indicated in Fig. 8. Note that within the fireball, the H-nucleon relative momenta are relatively small, and hence one is usually below the thresholds shown in Fig. 8. Thus, the H will not often dissociate while escaping the fireball.

For large $\langle p_H \rangle$, the main contribution to the cross section comes from the dissociation of the H into baryons of momenta around $\langle p_H \rangle /2$, with low relative momentum, followed by a Λp, $\Xi^- p$, or pp final state interaction which serves to put the process on shell. A very rough estimate yields

$$\sigma \approx \begin{cases} \frac{1}{4}\sigma_{pp} \approx 10 \text{ mb} & \text{for } Hp \to \Xi^- pp \\ \frac{1}{8}\sigma_{\Lambda p} \approx 3 \text{ mb} & \text{for } Hp \to \Lambda\Lambda p \end{cases} \tag{19}$$

assuming $\sigma_{pp} \approx 40$ mb and $\sigma_{\Lambda p}/\sigma_{pp} \approx \frac{2}{3}$.

The presence of a large neutron flux could constitute an important background which might mask the signal due to H dissociation. The cross section for the process $np \to 5$ prongs varies from 0.2–0.6 mb in the 3–4 GeV/c region,[55] which is much smaller than Eq. (19). However, N_n/N_H is very large, so a detector would need to distinguish between the $3p2\pi^-$ final state (the signal) and events containing combinations like $2p\pi^+2\pi^-$, which could arise from np interactions.

VI. CONCLUSIONS

The quest for dibaryon resonances is an exciting one. Only in the strange sector are such objects likely to be long lived or even stable with respect to strong decay. The production of such objects requires an intense source of strange baryons. This could be supplied by a high momentum kaon beam, with the (K^-,K^+) reaction used to transfer two units of strangeness to a baryon. An alternative, which we have explored here, is to exploit the copious production of strangeness in high energy heavy-ion collisions. In collisions at Brookhaven AGS energies (15 GeV/nucleon), we predict that H's will be produced at the level of 10^{-2} per central collision (~ 10 mb cross section). In contrast to the (K^-,K^+) reaction, where one detects one or two particles recoiling against the H in a two or three-body final state, one must detect the H directly when it is formed in the high multiplicity environment of a heavy-ion collision. If the H lifetime τ_H is less than 10^{-10} sec or so, direct detection will be difficult, since the decay occurs close to the production vertex. The very short lifetime also limits the possibility of detecting stable objects containing a charmed quark, for instance the pentaquark $\bar{c}suud$.[1] However, if $\tau_H \geq 10^{-9}$ sec, as predicted in the quark model,[13] several strategies for detecting the H are viable. These include the detection of weak decay products $(\Sigma^- p)$, diffraction dissociation of the H by a nuclear target, and the search for H-nucleus bound states of anomalous charge and/or mass ratio. The choice of an optimum experiment is governed by the mass of the H, which has not been reliably predicted. Several complementary experiments are thus required.

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