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ABSTRACT

We investigate the properties of multi-strange baryonic systems, comparing conventional many-Λ hypernuclei, where the strange quarks are localized in individual hyperons, to "strangelets" or chunks of strange matter, which involve delocalized quarks which roam in a single large bag. Mass formulae and strong/weak decay modes for such objects are discussed, as well as the prospects for producing multi-strange systems in relativistic heavy ion collisions. For production, we consider two extremes, one based on the coalescence model and another which assumes the formation of quark-gluon-plasma. We mention the experimental searches which are underway or planned, using heavy ion beams.

1. Introduction and Motivation

The study of multi-strange hadronic matter has a number of strong motivations. From the nuclear physics point of view, one would like to investigate the rôle of the strangeness degree of freedom in a nuclear many-body system. From the observed properties of bound states of nucleons and hyperons (Λ, Σ, Ξ), one can derive properties of effective baryon-baryon (NN, ΛN, $\Lambda\Lambda$, ...) interactions, and search for remnants of SU(3) symmetry in such interactions. Based on reasonable extrapolations from the observed properties of Λ and $\Lambda\Lambda$ hypernuclei, one anticipates the existence of a broad array of bound multi-strange nuclear systems, whose binding energy per baryon E_B/A and density ρ are comparable to those of ordinary nuclei ($E_B/A \simeq -8$ MeV, $\rho \simeq \rho_0 \simeq 1/6$ fm^{-3}). Such relatively dilute and weakly bound systems can support a significant influx of strangeness S; shell model estimates indicate that objects with strangeness as large as $|S|/A \approx 1/3$ should be stable with respect to strong baryon emission.

In addition to conventional multi-hyperonic matter, there is the exciting possibility that there may exist another branch of the binding energy curve of strange matter corresponding to more deeply bound systems of larger strangeness ($|S|/A \sim 0.8 - 1$) and higher density ($\rho \sim 2\rho_0$). The existence and properties of such "strangelets", or droplets of strange quark matter, were explored in depth by

Witten[1] and Farhi and Jaffe[2]. Unlike weakly bound systems of nucleons and hyperons, strangelets correspond to a single large bag containing comparable numbers of u, d and s quarks. In such a system, the strange (s) quarks are not localized within an individual hyperon, but are free to enjoy excursions over the entire volume of the bag.

In order to make reliable predictions for the masses of strangelets, we must exercise considerable control over quantum chromodynamics (QCD) in the non-perturbative (confinement) regime. At present, we do not possess this control, and the theoretical predictions discussed here should be regarded as educated guesses, whose validity is to be decided ultimately by experiment. As an example, the MIT Bag Model and a certain version of the SU(3) soliton model can both be adjusted to acceptably reproduce the masses of $A = 1$, $S = 0, -1, -2, -3$ baryons, yet these two treatments of non-perturbative QCD yield strikingly different predictions for the level order of dibaryons ($A = 2$) with multiple strangeness.

The possible existence of stable strange quark matter has significant ramifications in other areas, for instance astrophysics and cosmology. We mention the rôle that strange quark matter could play in the early universe, in the cores of neutron stars, and in nucleosynthesis. We do not treat these exciting questions here, but refer the reader to the Proceedings of the Aarhus Workshop[3], where an extensive list of references is to be found. In the present paper, we focus on mass formulae for strange matter, strong and weak decay properties, and production of multi-strange objects in relativistic heavy ion collisions.

2. Multi-strange Hypernuclei

In this section, we consider conventional bound systems of nucleons and hyperons, with the goal of estimating the maximum amount of strangeness that such a system could support before becoming unbound. For quantitative estimates, we rely on some recent work by the Frankfurt group[4,5].

We first provide a capsule summary of the information which is available concerning the bound states of $S = -1, -2$ hypernuclei. From (π^+, K^+) and (K^-, π^-) reactions on nuclear targets, which deposit a hyperon in the nuclear medium[6], we can establish the systematics of Λ single particle binding energies. For each orbital angular momentum L, the Λ binding is a smooth function of A, displaying no shell effects. In the dilute nuclear system, the Λ behaves as a distinguishable baryon, and there is no need to introduce any significant Pauli repulsion due to the antisymmetrization of the non-strange quarks in the Λ with similar (u, d) quarks in the nucleus. If this data is analyzed in a mean field (Hartree) approach[7], an attractive Λ-nucleus potential is found, with a well depth V_Λ roughly 1/2 as large as that for a nucleon, i.e.

$$V_\Lambda \simeq 30 \text{ MeV} \sim \frac{1}{2} V_N \tag{1}$$

The one–body Λ–nucleus spin–orbit potential is also known[6] to be at least an order of magnitude weaker than the corresponding $\vec{L}\cdot\vec{\sigma}$ interaction for a nucleon, and thus for our purposes it can be assumed that Λ single particle levels with $J = L \pm 1/2$ are effectively degenerate ($p_{1/2,3/2}$, $d_{3/2,5/2}$, etc.).

For a given A, a Λ hyperon can exist in several states $\{L, J\}$ which are stable with respect to strong decay. There is some (controversial) evidence[8] for quasi–particle states of heavier hyperons (Σ, Ξ) in a nucleus. Although these excitations can undergo strong decay ($\Sigma N \to \Lambda N$, $\Xi N \to \Lambda\Lambda$), they could be relatively narrow (width $\Gamma \lesssim 5$ MeV) due to binding effects and approximate selection rules[8].

There is some limited evidence for the doubly strange ($S = -2$) hypernuclei ${}^{6}_{\Lambda\Lambda}\mathrm{He}(\alpha + \Lambda + \Lambda)$ and ${}^{10}_{\Lambda\Lambda}\mathrm{Be}$ in early emulsion experiments and more recent results from Aoki *et al.*[9] at KEK using a 1.66 GeV/c K^- beam to tag the production of a Ξ^- hyperon ($K^- + p \to K^+ + \Xi^-$), which subsequently interacts in emulsion to form a $\Lambda\Lambda$ hypernucleus: an event was seen which most likely[10] corresponds to ${}^{13}_{\Lambda\Lambda}\mathrm{B}$ production. The analysis of these data yields an attractive $\Lambda\Lambda$ matrix element

$$\langle V_{\Lambda\Lambda}(r)\rangle \approx -4 \text{ to } -5 \text{ MeV} \tag{2}$$

which is to be compared to values of -2 to -3 MeV for ΛN and -6 to -7 MeV for NN matrix elements in the ${}^1\mathrm{S}_0$ state. Although the ${}^1\mathrm{S}_0$ nn system is known to be unbound, the corresponding two–body $\Lambda\Lambda$ system could be weakly bound, because of the larger Λ mass. Such a "quasi–molecular" $\Lambda\Lambda$ state near threshold, bound by long and medium range meson exchange forces, is to be distinguished from a six quark H dibaryon[11]: the latter, if deeply bound, would be an SU(3) singlet with a relatively small $\Lambda\Lambda$ component.

Given that $S = -1, -2$ hypernuclei are known to be bound, we can ask the question: How many units of strangeness can be added to a nuclear system before it becomes unstable with respect to direct Λ emission? This question has been addressed in mean field approximation[4,5], using interactions which are consistent with the data on ordinary nuclei and Λ hypernuclei. As an example, consider a multiply strange nucleus ${}^{208}_{n\Lambda}\mathrm{Pb}$, in which n neutrons have been replaced by Λ's. The results of Rufa *et al.*[4] for the binding energy per particle of this system may be parametrized in the form

$$\frac{E_B}{A} \approx -9 \text{ MeV} + 48\left(\frac{S - S_{\min}}{A}\right)^2$$
$$E_B = E - (A - |S|)\, m_N - |S| m_\Lambda \tag{3}$$

where $S_{\min} \approx -30$ or $|S_{\min}|/A \approx 1/7$. Thus the maximum binding per particle occurs for a substantial strangeness content $|S| \sim |S_{\min}|$ and the system resists Λ emission for even larger values of $|S|/A$. The increased binding for $|S| \neq 0$ is understood as an effect of the Pauli principle at the hadron level: neutrons in valence orbits of large L are replaced by Λ's in more deeply bound levels of lower L which are not Pauli blocked. The same idea at the quark level recurs in our later discussion of strangelets.

To understand the restriction on $|S|/A$ for stability against Λ emission, we consider the requirements of the Pauli principle for Λ's. For $A = 6$, we know that ^{6}He, ${}^6_\Lambda$He and ${}^{6}_{\Lambda\Lambda}$He are bound. In ${}^{6}_{\Lambda\Lambda}$He, the two Λ's occupy the $1s_{1/2}$ shell model orbital, with a ^{1}S$_0$ relative two-body state. If we try to add a third Λ, the Pauli principle requires that it must be injected into the p state, which only becomes bound near $A = 12$. Thus ${}^{6}_{3\Lambda}$He is unbound, and we are restricted to $|S|/A \lesssim 1/3$. As another example, consider an $A = 16$ core. Here we expect that 8Λ's $(2s_{1/2} + 2p_{1/2} + 4p_{3/2})$ could be added in bound states, whereas $L = 2$ Λ orbitals are unbound. Thus systems up to ${}^{24}_{8\Lambda}$O may be bound, or $|S|/A \leq 1/3$.

To summarize: on the basis of reasonable extrapolations from known properties of $S = 0, -1, -2$ many-baryon systems, one expects to find a wide range of bound hypernuclei with $|S|/A \lesssim 1/3$. Such objects are bound by conventional long and medium range attraction due to meson exchange. The binding of such objects is small, of order 10 MeV/A or less, and their density remains of order $\rho \sim \rho_0$ for all $|S|$, so they represent dilute systems in which the strange quarks are largely localized in individual Λ hyperons (although Σ and Ξ-like components could be admixed into the wave function at some level). For small $|S|$, the r.m.s. radius of the Λ's is smaller than that of the nucleons, but for larger $|S|$, a "Λ halo" is formed, so the strangeness is more concentrated in the surface of the nucleus[4,5].

3. Multi–strange Few–body Systems

We mention a few possibilities for rather light systems with multiple strangeness, since these could be produced in detectable numbers in relativistic heavy ion collisions. For instance, the dibaryon $H = (ssuudd)^{I=0}_{0^+}$ of Jaffe[11] can be produced in reactions[12] like $\Xi^- + d \rightarrow H + n$ or $K^- + {}^3\text{He} \rightarrow K^+ + H + n$ or with observable rates in heavy ion[13] encounters at high energy. If the H is stable, there is also the possibility[14] of bound states of the H and a nuclear core, for instance ${}^4_H\text{H} = H + d$, or bound states of two or more H's, for example an $H^2 = (HH)_{L=0}$ configuration.

The $\Sigma^-\Sigma^-$ system ($I = 2$, ^{1}S$_0$) provides a good example of the dramatic model dependence of non–perturbative approaches to QCD. This system is unbound in the bag model[11], since the net quark interaction is repulsive in this channel. On the other hand, a version of the SU(3) soliton model[15,16] leads to strongly bound $\Sigma^-\Sigma^-$, $\Xi^-\Xi^-$ and even $\Omega^-\Omega^-$ states. The binding could be sufficiently strong as to suppress the weak decay $\Sigma^-\Sigma^- \rightarrow \Sigma^- n\pi^-$; the state would then decay via a leptonic mode, with a rather long lifetime $\geq 10^{-8}$ sec. Both the bag model and the soliton model are adjusted to reproduce the observed masses of strange baryons, yet they differ markedly in their predictions for $A = 2$ systems.

At the hadron level, the members of the SU(3) baryon octet, $\{N, \Lambda, \Sigma, \Xi\}$ are distinguishable particles, so one can prepare a variety of light systems which might possibly be stable against strong decay[5]. These involve ^{1}S$_0$ pairs of baryons, all occupying $s_{1/2}$ orbitals in the many-body system. For ^{1}S$_0$ pairs of maximum isospin (nn, $\Lambda\Lambda$, $\Sigma^-\Sigma^-$, $\Xi^-\Xi^-$), all pseudoscalar (π, η) and scalar (σ, δ) meson exchange contributions are coherently attractive. Thus composite "Noah's Ark" systems such as $\{nn\,\Lambda\Lambda\,\Sigma^-\Sigma^-\,\Xi^-\Xi^-\}$ with $J^\pi = 0^+$, $I = 4$, $S = -8$ afford interesting possibilities for binding. Note that the absence of protons enables us to evade the strong conversion processes $\Sigma^- p \longrightarrow \Lambda n$ and $\Xi^- p \longrightarrow \Lambda\Lambda$.

4. Strange Quark Matter: Mass Formula

As discussed in Section 2, a conventional many–baryon system can support the injection of numerous Λ's (up to $|S|/A \lesssim 1/3$) and still remain stable. Witten[1] and Farhi and Jaffe[2] proposed that systems of even larger strangeness might be stable, in which the quarks are loaded into a single large bag, rather than being localized in single baryons. They consider a degenerate Fermi gas of quarks at temperature $T = 0$, in which flavor equilibrium is maintained by the weak processes $s, d \leftrightarrow u + e^- + \bar{\nu}_e$, and $u + d \leftrightarrow s + u$. For a bulk system at equilibrium, the chemical potentials are related by $\mu_d = \mu_s = \mu_u + \mu_e$. Because of Pauli blocking at the quark level, it is energetically favorable to add strange quarks to a system in which a number of levels are already occupied by u or d quarks. From the thermodynamic potentials Ω_i for non–interacting quarks

$$\Omega_i = \begin{cases} -\mu_i^4/4\pi^2 & \text{for } i = u, d \\ -\frac{1}{4\pi^2}\left[\mu_s\left(\mu_s^2 - m_s^2\right)^{1/2}\left(\mu_s^2 - \frac{5m_s^2}{2}\right) + \frac{3m_s^4}{2}\ln\left(\frac{\mu_s + \left(\mu_s^2 - m_s^2\right)^{1/2}}{m_s}\right)\right] & \text{for } i = s \end{cases} \tag{4}$$

we can derive the quark number density n_i and the energy density ϵ:

$$n_i = -\partial\Omega_i/\partial\mu_i,\ \epsilon = \sum_i (\Omega_i + \mu_i n_i) + B \tag{5}$$

For bulk ($A \longrightarrow \infty$) quark matter, the system is electrically neutral:

$$\frac{2}{3}n_u - \frac{1}{3}n_d - \frac{1}{3}n_s - n_e = 0 \tag{6}$$

For finite A, the electrons reside outside the system, which then acquires a positive charge Z.

Using Eqs. (4–6), one can show that the energy per baryon E/A can be less than the nucleon mass m_N for certain choices of the bag pressure B and the strange quark mass m_s. This is the regime of stable strange quark matter (SQM). The region of stability in the $\{B, m_s\}$ plane changes smoothly when quark–quark interactions via gluon exchange with coupling constant α_c are included perturbatively[2]. The typical density of a chunk of SQM is $\rho \approx (1-2)\rho_0$, where $\rho_0 \simeq 1/6\ \text{fm}^{-3}$ is the density of non–strange nuclear matter.

For ordinary nuclei, the total energy $E(A, Z)$ for baryon number A and charge Z can be described by the Weizsäcker mass formula

$$E(A, Z) = m_N A - a_V A + a_S A^{2/3} + a_C Z^2/A^{1/3} + a_{\rm symm}\left(Z - \frac{A}{2}\right)^2/A - \delta \qquad (7)$$

where volume (a_V), surface (a_S), Coulomb (a_C), symmetry ($a_{\rm symm}$) and even–odd (δ) energies are included. For SQM, Berger and Jaffe[17] have developed a mass formula of similar structure which also incorporates the strangeness degree of freedom:

$$E(A, S, Z) = \epsilon_0 A + \epsilon_1 A^{2/3} + \frac{\epsilon_2}{A}(S - S_{\rm min})^2 + \left(\frac{\epsilon_3}{A} + \frac{\epsilon_4}{A^{1/3}}\right)(Z - Z_{\rm min})^2 \qquad (8)$$

The parameters (ϵ_i, $S_{\rm min}$, $Z_{\rm min}$) are not independent, but are determined in terms of $\{B, m_s, \alpha_c\}$. Comparing with Eq. (7), we can identify ϵ_0 with $m_N - a_V$; for ordinary nuclei, $a_V \simeq 15$ MeV (binding energy per particle of nuclear matter), so $\epsilon_0 \simeq 925$ MeV for nuclei. Typically, we have $|S_{\rm min}|/A \simeq 0.8 - 1$, $Z_{\rm min}/A \simeq 0.1$, i.e. large strangeness and relatively small charge.

As an example, for the choice $\epsilon_0 = 900$ MeV, $m_s = 150$ MeV, we find for $A = 208$, $Z = 82$ the result

$$\frac{E}{A} - m_N = -17 + 86\left(\frac{|S|}{A} - 0.8\right)^2 \qquad (9)$$

This should be contrasted with the results of Rufa *et al.*[4](see Eq. (3)) for the hypernucleus ${}^{208}_{n\Lambda}$Pb:

$$\frac{E}{A} - m_N \approx -8 + 160\frac{|S|}{A} + 70\left(\frac{|S|}{A}\right)^2 \qquad (10)$$

These two forms of the binding energy of multi–strange matter are depicted in Fig. 1. We note that hypernuclei are always more massive than $S = 0$ nuclei of the same A, i.e. the binding energy of a Λ in a heavy nucleus approaches the well depth $V_\Lambda \simeq 30$ MeV, which is much less than the mass difference $m_\Lambda - m_N \simeq 176$ MeV. For small $|S|$, the energy of SQM lies above that for a hypernucleus, whereas for large $|S|$, SQM lies lower, perhaps even below $E = A m_N$. In Fig. 1, the SQM curve will move up if ϵ_0 is increased. Note that the binding energy per particle of SQM near $S = S_{\rm min}$ can be very large when measured with respect to the threshold $(A - |S|)m_N + |S|m_\Lambda$ for dissolution of the system into a gas of N's and Λ's. In Fig. 1, this binding energy for $|S| = 0.8\,A$ is about 158 MeV per particle, much larger than the values of order 10 MeV for multi–strange hypernuclei.

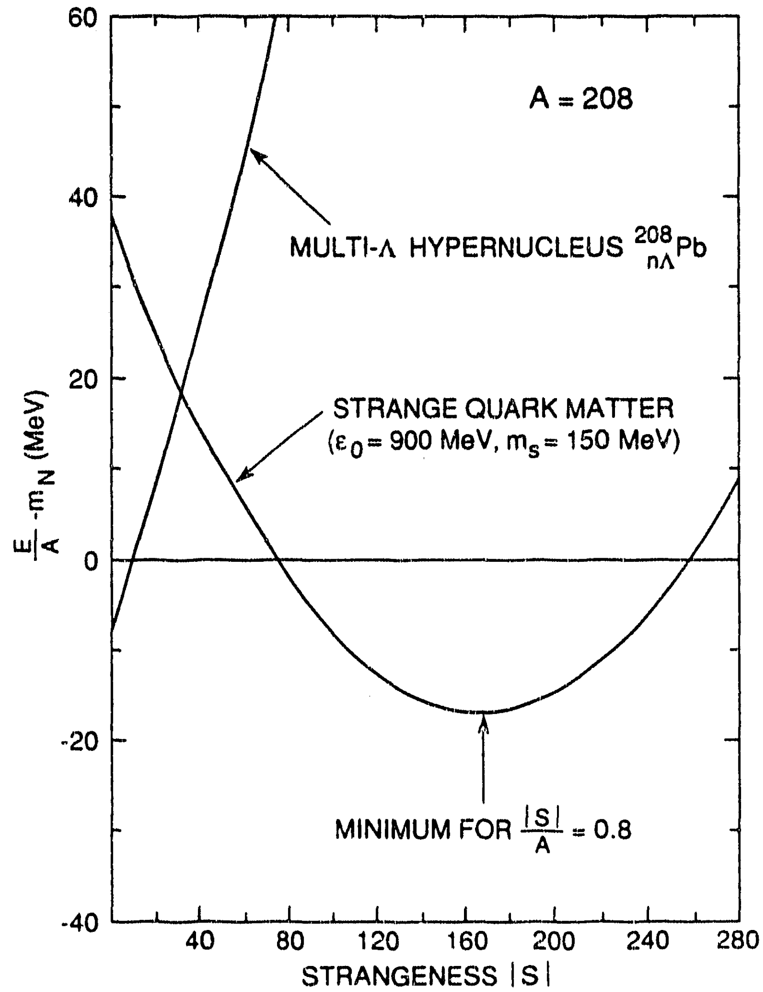

Figure 1: Energy per particle E/A relative to the nucleon mass m_N as a function of strangeness $|S|$ for a multi–Λ hypernucleus and for a droplet of strange quark matter, both with $Z = 82$, $A = 208$. The hypernuclear binding was estimated in mean field theory[4] and the strangelet energy was obtained from the Berger–Jaffe mass formula with $\epsilon = 900$ MeV, $m_s = 150$ MeV.

5. Strange Quark Matter: Strong and Weak Decays

Given the mass formula of Eq. (8), one can delineate the regions of stability of SQM with respect to various strong or weak decays. For instance the Q–values (energy release) for strong neutron, proton or pion emission are given by

$$\begin{aligned} Q_S^n &= E(A,S,Z) - E(A-1,S,Z) - m_n \\ Q_S^p &= E(A,S,Z) - E(A-1,S,Z-1) - m_p \\ Q_S^{\pi^-} &= E(A,S,Z) - E(A,S,Z+1) - m_{\pi^-} \end{aligned} \tag{11}$$

with similar expressions for weak decay (Q_W^{n,p,π^-}). By subtraction, we obtain

$$Q_W^{n,p} - Q_S^{n,p} = \frac{\epsilon_2}{A-1}\left[-1 + 2\left(|S| - |S_{\min}|\right)\right] \quad (12a)$$

$$Q_S^{\pi^-} < 0 \text{ for } Z \geq Z_{\min} - A\,m_{\pi^-}/2\left(\epsilon_3 + \epsilon_4 A^{2/3}\right) \quad (12b)$$

where $|S_{\min}|$ refers to the A-1 baryon system. Thus for $|S| < |S_{\min}| + 1/2$, we have $Q_S^{n,p} > Q_W^{n,p}$ (strong decay region), while for $|S| > |S_{\min}| + 1/2$, we obtain $Q_W^{n,p} > Q_S^{n,p}$ (weak decay region). Thus SQM of small $|S|$ decays strongly by nucleon emission, while for large $|S|$, weak nucleon emission takes over. For some values of $|S|$, both weak and strong emission may occur. Further, strong proton emission limits Z values above $Z_{\min}$, while (12b) constrains the values $Z < 0$ which are stable against strong π^- decay. Thus various strong emission processes limit the possible excursions of (Z, S) from $(Z_{\min}, S_{\min})$.

An interesting quantity, relevant to experiments, is the minimum baryon number $A_{\min}$ for which SQM is stable with respect to both weak and strong nucleon emission. For $Z \approx Z_{\min}$ and $A_{\min} \gg 1$, we find

$$A_{\min} \simeq \left(\frac{2\epsilon_1}{3\left(m_N - \epsilon_0\right)}\right)^3 \quad (13)$$

As ϵ_0 approaches m_N, $A_{\min} \to \infty$, and no stable objects exist. Conversely, if a heavy ion collision experiment is capable of detecting multi–strange clusters of $A \leq A_{\max}$, and none are found, one can use (13) to constrain the parameters of the SQM mass formula. Of course, such a constraint depends strongly on the production model.

Hypernuclei of strangeness $S = -1, -2$ decay weakly with lifetimes of order $\tau \approx \tau_\Lambda \approx 3 \times 10^{-10}$ sec. For $A > 5$, the free decay $\Lambda \to N\pi$ is strongly Pauli–blocked, since the recoiling nucleon carries a momentum of only 100 MeV/c. The primary decay mode is $\Lambda N \to NN$, involving the emission of two nucleons into the continuum. In a strangelet, if the binding energy E_B/A (see Eq. (3)) exceeds $(m_\Lambda - m_N)/2$, this two–nucleon mode has no phase space, and the weak emission of a single nucleon takes over (the second nucleon remains bound). There have been several attempts to estimate strange quark weak decay rates, both at high temperature[18] and for strangelets[19] ($T = 0$). Pauli blocking also plays a dramatic rôle at the quark level, leading to a suppression of the decay rate by several orders of magnitude: Koch[19] estimates a lifetime of order 10^{-5} to 10^{-6} sec for the non-leptonic process $u + s \to d + u$.

There is also some sparse information on the weak (single) neutron emission from hypernuclei. From the data of Coremans *et al.*[20], one obtains

$$\tau\left({}^5_\Lambda\text{He} \to n + \alpha\right) \approx 50 - 80\tau_\Lambda \sim 1.4 - 2 \times 10^{-8} \text{ sec} \quad (14)$$

A theoretical approach to single nucleon weak decay has been formulated by Filimonov and Potashev[21], who use a transition operator of the form $a + b\vec{\sigma} \cdot \vec{q}$ to obtain a decay rate

$$\Gamma = 1/\tau \approx \text{ const} \times |\psi(q)|^2 \left(a^2 + b^2 q^2\right) q \tag{15}$$

where $\psi(q)$ is the Λ wave function and q is approximately the neutron momentum. For ${}^5_\Lambda\text{He} \to n + \alpha$, $q \simeq 520$ MeV/c, a rather substantial value. For a strangelet, which is typically much more deeply bound, q will be smaller. A formula of the type (15) may be useful to extrapolate to the regime of strangelets.

6. Production Mechanisms

We now consider mechanisms for the production of multi-strange objects in hadronic collisions. For $S = -1, -2$, kaon beams offer an effective means of producing hypernuclei: basic reactions are $K^- n \to \pi^- \Lambda$ or $K^- p \to K^+ \Xi^-$, for instance. For $|S| \geq 3$, relativistic heavy ion beams offer the only feasible way to produce multi-strange clusters in the laboratory. With heavy ions, each NN collision produces strange particles independently via the $NN \to NYK$, $NNK\overline{K}$ reactions. It is already established at AGS (15 GeV/A) and CERN (60, 200 GeV/A) energies that hyperons are abundantly created in central collisions (for Si+Au at the AGS, some 4–5 Λ's + Σ's typically emerge; for Au+Au collisions, of order 15 hyperons are anticipated). Thus a central high energy heavy ion collision provides a bath of strange particles, which can serve as the building blocks of strangelets or other composite objects of multiple strangeness.

We now outline two different estimates of the production rate of strangelets in heavy ion collisions. The first is based on the coalescence model[22], in which composite particles are formed at a late stage of the collision process, at "freeze-out". This may be viewed as a conservative lower limit for strangelet production. The second method presupposes the formation of a substantial droplet of quark–gluon–plasma (QGP), which then hadronizes, occasionally yielding a multi-strange object. This is an optimistic scenario.

In the coalescence picture, we start with some non-strange cluster, say the ^{4}He nucleus and add more baryons, both strange and non-strange. The number of clusters $N(A, S)$ per collision is roughly given by

$$N(A, S) \simeq P_1^{|S|} P_2^{A-4} N\left({}^4\text{He}\right) \tag{16}$$

where P_1 and P_2 are penalty factors for converting a nucleon to a Λ or adding a nucleon, respectively. In the thermal model, P_2 is given by

$$P_2 \simeq \rho_p \lambda_T^3 \quad ; \quad \lambda_T = \left(2\pi/m_p T\right)^{1/2} \tag{17}$$

where ρ_p is the proton density at freezeout, λ_T is the thermal wavelength, and T is the temperature. At BEVALAC energies[23](0.4–2 GeV/A), we find that $P_2 \simeq 0.2$ reproduces the observed ${}^4\text{He}/p$ ratio in the coalescence model. Extrapolating to the more elevated temperatures prevailing at higher energies, and estimating P_1 from the observed Λ/p ratio, we obtain

$$P_1 \approx P_2 \approx 1/10 \tag{18}$$

at AGS energies (11 GeV/A for Au beams), and predict a $^4\text{He}/p$ ratio of 2×10^{-3}, or $N(^4\text{He}) \simeq 1/4$. In an AGS experiment with sensitivity $\epsilon = 2.5 \times 10^{-n}$ for detection of composite fragments, we are able to sample a region of the $\{S, A\}$ plane bounded by

$$|S| + A \leq 3 + n \tag{19}$$

Since $n \leq 11$ for all the proposed experiments (see next section), only rather light strangelets (up to $S = -6$, $A = 8$ or so) could be detected if the coalescence estimates are correct.

A much more optimistic scenario emerges if quark–gluon–plasma (QGP) is formed in a heavy ion collision. The formation of strangelets during the hadronization of QGP has been considered by several authors[24–28]. As the initial quark–gluon droplet cools down, the system enters a mixed phase, in which the s quarks are distilled into the QGP, and the $\bar{s}$'s are concentrated in the hadron phase and then radiated as K^0 or K^+ mesons. Greiner *et al.*[25–27] have argued that strangelets with typical baryon number $A = 10 - 30$ and negative charge $Z \approx -0.1A$ could be formed by this strangeness distillation process. Crawford *et al.*[28] have provided rough estimates of the production probability P of strangelets per heavy ion collision event. They write

$$P\,(A, Z) = P_{\text{QGP}} \cdot P_A \cdot \sum_S P\,(A, Z, S) \cdot P_{\text{cool}} \tag{20}$$

where P_{QGP} is the probability for QGP formation, $P_A \approx A/2A_{\text{beam}}$ is the probability of producing a cluster of baryon number A, and $P_{\text{cool}} \approx c/A$, with $c \approx 6 \times 10^{-2}$ (15 GeV/A) or 10^{-2} (200 GeV/A), is the probability that a cluster cools down to its ground state. In (20), we sum over all values of S for which the strangelet is sufficiently long–lived ($\tau > 10$ ns) to be measurable, using a Berger–Jaffe mass formula[17]. Finally, we have

$$\begin{aligned} P\,(A, Z, S) &= P\,(n_u)\; P\,(n_s) \\ P\,(n_{u,s}) &= e^{-\bar{n}_{u,s}} \frac{(\bar{n}_{u,s})^{n_{u,s}}}{n_{u,s}!} \end{aligned} \tag{21}$$

The small charge ($Z \approx 0.1A$) and large strangeness ($|S| \approx 0.8A$) of the strangelet are created by a Poisson distribution of fluctuations (Eq. (21)), starting from an initial droplet of large initial charge $Z_{\text{in}} \approx A/2$, with $n_u = Z_{\text{in}} - Z$, and small strangeness $\bar{n}_s \approx (0.1 - 0.2)A$ with $n_s = |S|$. For typical parameter choices ($P_{\text{QGP}} = 0.1$, $\bar{n}_s = 0.1A$, $A_{\text{beam}} = 200$, $\epsilon_0 = 900$ MeV, $m_s = 150$ MeV), one obtains[28]

$$P\,(A, Z) = \begin{cases} 2 \times 10^{-9} & (A = 10,\ Z = -1) \\ 7 \times 10^{-9} & (A = 10,\ Z = 1) \\ 8 \times 10^{-11} & (A = 20,\ Z = 1) \end{cases} \tag{22}$$

These predictions are orders of magnitude larger than the coalescence estimates of Eq. (16). In particular, we note that the penalty factor for adding ten baryons (Eq. (22)) is only 10^{-2}, to be compared to 10^{-10} from Eq. (16). This is because $P_A \cdot P_{cool}$ is assumed to be independent of A, and the decrease of $P(A, Z)$ with increasing A comes only through the Poisson distribution $P(n_u)$. If QGP is formed, strangelets with A as large as 20–30 may be produced with observable rates in heavy ion collisions. In fact, if stable strangelets exist at all, they would provide an excellent signature of QGP formation!

7. Experiments with Heavy Ion Beams

Very sensitive heavy isotope searches at Tandem accelerators have not yielded any indication of the existence of very long–lived anomalous objects. Blackman and Jaffe[29] used these results to constrain the choice of the parameters (ϵ_0, m_s) in the strangelet mass formula. In the range $0.1 \leq Z/A \leq 0.3$, experiment E814 at the AGS[30] obtained limits of 10^{-4} to 5×10^{-5} on production of anomalous particles. In the past year or two, several new experiments were proposed for the AGS (E864 of Sandweiss *et al.*, E878 of Crawford *et al.*, E882 of Price *et al.* and E886 of Imai *et al.*) and at CERN (P268 of Pretzl *et al.*, Letter of Intent of Schlein *et al.*). These proposals use a variety of techniques, and are sensitive to different regions of lifetime and Z/A. The most sensitive experiment proposed to date is E864, which is capable of detecting strangelets (or other unusual objects) at a level of a few $\times 10^{-11}$ per Au+Au event. These experiments will open up an important new era in the search for strange quark matter.

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