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ASPECTS OF STRANGE MATTER

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MASTER

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Aspects of Strange Matter

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Abstract. We discuss the stability of multiply strange baryonic systems, in the context of a mean field approach obtained from an underlying set of phenomenological meson–baryon interactions. The coupling parameters which determine the conventional $\sigma + \omega$ mean fields (Hartree potentials) seen by varius baryon species (N, Λ, Ξ) in the many-body system are constrained by reproducing the trend of observed binding energies of single particle (N, Λ, Ξ) states, as well as the energy per particle and density of non–strange nuclear matter. We also consider additional scalar (σ^*) and vector (ϕ) fields which couple strongly to strange baryons. The couplings of these fields are adjusted to produce strong hyperon–hyperon interactions, as suggested by the data on $\Lambda\Lambda$ hypernuclei. Extrapolating this approach to systems of large strangeness S, we find a broad class of objects composed of neutrons, protons, Λ's and Ξ's, which are stable against strong decay. In these systems, the presence of filled Λ orbitals blocks the strong decay $\Xi N \rightarrow \Lambda\Lambda$, leading to a strangeness fraction $f_s = |S|/A \approx 1$, density $\rho \approx (2-3)\rho_0$, and charge fraction f_q in the range $-0.1 < q/A < 0.1$, comparable to that of hypothetical stable strange quark matter ("strangelets"), but with a low binding energy per particle $E_B/A \approx -10$ to -20 MeV. Such weakly bound multi–strange objects can be stable for very large A, unlike ordinary nuclei, since the Coulomb repulsion generated by the protons is largely cancelled by the presence of a comparable number of Ξ's, leading to a small net charge (positive or negative) of order $A^{1/3}$.

INTRODUCTION

In addition to the usual atomic nuclei, containing neutrons and protons, one can imagine many–body systems which incorporate the hyperons, namely Λ, Σ and Ξ. We are thus led to ask the following question: What is the domain of stability (against strong decay) of systems of SU(3) octet baryons $\{n, p, \Lambda, \Sigma^0, \Sigma^{\pm}, \Xi^-, \Xi^0\}$? For such systems, the usual variables of baryon number A and charge per baryon q/A are supplemented by the strangeness fraction $|S|/A$ (zero for ordinary nuclei). Thus we are essentially exploring the strangeness degree of freedom in multi–baryon systems.

Our strategy is as follows: we extend the conventional shell model picture of nuclei to include any number of hyperons. The N, Λ, Σ and Ξ are treated as distinguishable particles and thus each of these may occupy the same set of shell model orbitals $\{n\ell j\}$. The calculations were conducted in an extended version of the relativistic mean field (RMF) approach [1,2]. In addition to the usual scalar (σ) and vector (ω) mesons of the Walecka model [1], we include a second scalar–vector pair (σ^*, ϕ) which is responsible for most of the hyperon–hyperon (YY) interaction. The model is described in detail in two papers, [3,4] which also contain references to earlier work.

STABILITY CONSIDERATIONS

There is substantial evidence for the existence of single Λ hypernuclei [5] and even a few emulsion events which have been interpreted as $\Lambda\Lambda$ hypernuclei [6]. The results of a mean–field analysis [7] of the Λ–nucleus data yields a well depth for the Λ of about

$$U_\Lambda \simeq 27 \pm 1 \text{ MeV}, \tag{1}$$

whereas an analysis of the few existing Ξ–nucleus events gives

$$U_\Xi \simeq 24 \pm 4 \text{ MeV}. \tag{2}$$

These well depths are comparable to the energy release $Q(\Xi N \to \Lambda\Lambda)$ in the free space $\Xi N \to \Lambda\Lambda$ strong conversion process. We have

$$Q(\Xi N \to \Lambda\Lambda) = \begin{cases} 28 & \text{MeV for } \Xi^- p \to \Lambda\Lambda \\ 23 & \text{MeV for } \Xi^0 n \to \Lambda\Lambda \end{cases} \tag{3}$$

Due to the rough equality of $U_{\Lambda,\Xi}$ and $Q(\Xi N \to \Lambda\Lambda)$, it becomes possible to overcome the free space decay $\Xi N \to \Lambda\Lambda$ by binding effects. In particular, $\Xi N \to \Lambda\Lambda$ can be Pauli–blocked if a certain number of Λ orbits are occupied. This is the crucial mechanism in stabilizing the Ξ in a strange many–body system. On the other hand, we find that the Σ cannot be stabilized against strong decay in this manner, since $Q(\Sigma N \to \Lambda N) = 75 - 80$ MeV is much larger than the well depths (U_Σ = 15–20 MeV from the data on Σ^-–atoms).

As a prototype for this argument, let us consider the ${}_{\Xi^0\Lambda\Lambda}^{\;\;\;\;7}\mathrm{He}$ system (${}^4\mathrm{He}+2\Lambda+\Xi^0$), in which all the baryons are in the $1s_{1/2}$ shell model state. Since the $s^\Lambda_{1/2}$ shell is fully occupied and the $p^\Lambda_{1/2,3/2}$ shells are unbound in such a light system, the $\Xi N \to \Lambda\Lambda$ reaction must eject the two Λ's into the continuum. The strong decay with the smallest Q value is then

$${}_{\Xi^0\Lambda\Lambda}^{\;\;\;\;7}\mathrm{He} \to 2\Lambda + {}_{\Lambda\Lambda}^{\;\;5}\mathrm{He} \tag{4}$$

for which

$$\begin{aligned} Q \approx&\ \ Q_{0n} - B_n({}^4\mathrm{He}) - B_{\Xi^0} + 2(B_\Lambda({}^4_\Lambda\mathrm{He}) - B_\Lambda({}^5_\Lambda\mathrm{He})) \\ \approx&\ \ 1.2 \text{ MeV} - B_{\Xi^0}. \end{aligned} \tag{5}$$

Thus if $B_{\Xi^0} > 1.2$ MeV, the Ξ^0 is stabilized against strong decay. This condition is achieved in our calculations. The mechanism of stabilization is clear from Eq. (5): the binding energy of a neutron in ${}^4\mathrm{He}(B_n({}^4\mathrm{He}) = 20.6$ MeV) cancels out most of the energy release $Q_{0n} = m_{\Xi^0} + m_n - 2m_\Lambda = 23.2$ MeV of Ξ conversion in free space. Our conclusion is that ${}_{\Xi^0\Lambda\Lambda}^{\;\;\;\;7}\mathrm{He}$ is likely to be the lightest stable multi–strange system containing a Ξ. Production rates for this object in heavy ion collision at Brookhaven AGS energies were estimated by Baltz *et al* [8].

SKETCH OF THE MEAN FIELD METHOD

We now briefly outline the standard relativistic mean field approach [1,2]. More details on our extensions of the standard model may be found in the thesis of J. Schaffner [9]. The relativistic σ+ω mean field model starts from the Lagrangian (for nucleon systems)

$$\begin{aligned}\mathcal{L}_N = &\bar{\psi}_N(i\gamma^\nu\partial_\nu - m_N)\psi_N + \frac{1}{2}\partial_\nu\sigma\partial^\nu\sigma - U(\sigma)\\ &- \frac{1}{4}G_{\mu\nu}G^{\mu\nu} + \frac{1}{2}m_\omega^2 V_\mu V^\mu - \frac{1}{4}\mathbf{B}_{\mu\nu}\mathbf{B}^{\mu\nu} + \frac{1}{2}m_\rho^2\mathbf{R}_\mu\mathbf{R}^\mu - \frac{1}{4}F_{\mu\nu}F^{\mu\nu}\\ &- g_{sN}\bar{\psi}_N\psi_N\sigma - g_{vN}\bar{\psi}_N\gamma_\mu\psi_N\mathbf{V}^\mu - \frac{1}{2}g_{rN}\bar{\psi}_N\gamma_\mu\tau_N\psi_N\mathbf{R}^\mu\\ &- \frac{1}{2}e\bar{\psi}_N(1+\tau_{0,N})\gamma_\mu\psi_N A^\mu \end{aligned} \tag{6}$$

which describes nucleons interacting via a scalar (σ), vector (V^μ), isovector ($\mathbf{R}^\mu$) and an electromagnetic (A^μ) potentials. The fields are defined by

$$\begin{aligned} F_{\mu\nu} &= \partial_\mu A_\nu - \partial_\nu A_\mu,\\ G_{\mu\nu} &= \partial_\mu V_\nu - \partial_\nu V_\mu,\\ \mathbf{B}_{\mu\nu} &= \partial_\mu\mathbf{R}_\nu - \partial_\nu\mathbf{R}_\mu \end{aligned} \tag{7}$$

We choose the parameter set $PL-Z$ of Reinhard [10], which yields reasonable values for nuclear matter ($\rho_0 = 0.153$ fm^{-3}, $E_B/A = -15.94$ MeV and compressibility $K = 294.3$ MeV). The addition of Λ's and Ξ's is described by adding the following terms to the nucleon Lagrangian $\mathcal{L}_N$:

$$\mathcal{L}_\Lambda = \bar{\psi}_\Lambda(i\gamma^\nu\partial_\nu - m_\Lambda)\psi_\Lambda - g_{s\Lambda}\bar{\psi}_\Lambda\psi_\Lambda\sigma - g_{v\Lambda}\bar{\psi}_\Lambda\gamma_\mu\psi_\Lambda V^\mu.$$

$$\begin{aligned}\mathcal{L}_\Xi = &\bar{\psi}_\Xi(i\gamma^\nu\partial_\nu - m_\Xi)\psi_\Xi\\ &- g_{s\Xi}\bar{\psi}_\Xi\psi_\Xi\sigma - g_{v\Xi}\bar{\psi}_\Xi\gamma_\mu\psi_\Xi V^\mu\\ &- \frac{1}{2}g_{r\Xi}\bar{\psi}_\Xi\tau_\Xi\gamma_\mu\psi_\Xi\mathbf{R}^\mu + \frac{1}{2}e\bar{\psi}_\Xi(1-\tau_{0,\Xi})\gamma_\mu\psi_\Xi A^\mu,\end{aligned} \tag{8}$$

The parameters in $\mathcal{L}_\Lambda$ are chosen to reproduce the Λ well depth ~ 27 MeV in nuclear matter and to obtain a small value for the Λ–nucleus spin–orbit potential. For the Ξ, we reproduce a well depth of 28 MeV. For simplicity, we impose SU(6) constraints

$$\frac{g_{v\Lambda}}{g_{vN}} = \frac{2}{3}, \quad \frac{g_{v\Xi}}{g_{vN}} = \frac{1}{3} \tag{9}$$

The $\sigma+\omega$ version of our model (referred to as Model 1) yields a small value $\Delta B_{\Lambda\Lambda} \leq 1$ MeV for the strength of the $\Lambda\Lambda$ interaction, whereas the empirical value is about 4.5 MeV [11], i.e., a rather strong attraction. We remedy this situation by introducing a new YY interaction mediated by a second scalar–vector pair (σ^*, ϕ) which are assumed to couple only to hyperons. The additional terms in the Lagrangian are then

$$\mathcal{L}_{YY} = -\frac{1}{4}S_{\mu\nu}S^{\mu\nu} + \frac{1}{2}m_\phi^2\phi_\mu\phi^\mu + \frac{1}{2}(\partial_\nu\sigma^*\partial^\nu\sigma^* - m_{\sigma^*}^2\sigma^{*2}) - g_{s*\Lambda}\bar{\psi}_\Lambda\psi_\Lambda\sigma^* - g_{\phi\Lambda}\bar{\psi}_\Lambda\gamma_\mu\psi_\Lambda\phi^\mu - g_{s*\Xi}\bar{\psi}_\Xi\psi_\Xi\sigma^* - g_{\phi\Xi}\bar{\psi}_\Xi\gamma_\mu\psi_\Xi\phi^\mu, \quad (10)$$

Here, $S_{\mu\nu} = \partial_\mu\phi_\nu - \partial_\nu\phi_\mu$. The coupling constants to the strange vector meson ϕ are assumed to satisfy the SU(6) relations

$$\frac{g_{\phi\Lambda}}{g_{vN}} = -\frac{\sqrt{2}}{3} \quad , \quad \frac{g_{\phi\Xi}}{g_{vN}} = -\frac{2\sqrt{2}}{3} \quad (11)$$

By fitting the σ^* strength to theoretical estimates of Λ and Ξ well depths in hyperon matter [4], we obtain

$$\frac{g_{s*\Lambda}}{g_{sN}} = 0.69 \quad , \quad \frac{g_{s*\Xi}}{g_{sN}} = 1.25 \quad (12)$$

The extended $\sigma+\omega+\sigma^*+\phi$ mean field model is referred to as Model 2. It gives a much larger value $\Delta B_{\Lambda\Lambda} = 3$ MeV than Model 1. With a model adjusted to approximate all known properties of $S = -1$ and $S = -2$ hypernuclei, we thus have a vehicle for extrapolation to systems of arbitrary (S, A).

RESULTS FOR STRANGE HADRONIC MATTER

Our procedure [4] is generally to start from closed shell proton and neutron configurations, mostly doubly magic nuclear cores, and fill up successively the Λ, Ξ^0 and Ξ^- single–particle (s.p.) states in the self consistent mean field formed by using the resulting attractive–scalar and repulsive–vector potentials. By adding many hyperons (equivalent to augmenting the meson fields), the normal sequence of nucleon s.p. energies and the nucleon magic numbers change appreciably. In particular, the nucleon spin–orbit splitting gets bigger and this gives rise to a new shell ordering in strange nuclear systems of large $|S|$, which depends on the occupation numbers of hyperons. For example, consider a system with a ^{208}Pb core, to which 70 Λ, 18 Ξ^0 and 70 Ξ^- are added. In Model 1, we have ($E_B/A = -12.9$ MeV). The nucleon magic number 20 becomes 28. For protons, the magic number 82 becomes 70. We also note that the nucleon potential becomes considerably deeper in the presence of many hyperons. For instance, the $1s_{1/2}$ neutron state, originally bound at -61.2 MeV in a nucleon bath, is shifted to -83.4 MeV in the presence of hyperons. The corresponding $1s_{1/2}$ proton level is shifted from -50.5 MeV to -88.7 MeV. The proton levels sink below the neutron levels in the presence of hyperons, since the nucleon symmetry potential, which in the absence of Coulomb effects would have made protons more bound than neutrons in ordinary nuclei, is now stronger than the residual Coulomb proton potential due to the 12 "excess" protons (82 protons minus 70 negative cascades). Hence, the last proton level here is bound by 18 MeV more than the last neutron level, so that the onset of $\Xi N \to \Lambda\Lambda$ conversion occurs in the $\Xi^0 n$ channel, thus keeping the number of Ξ^0's considerably below that of Ξ^-'s.

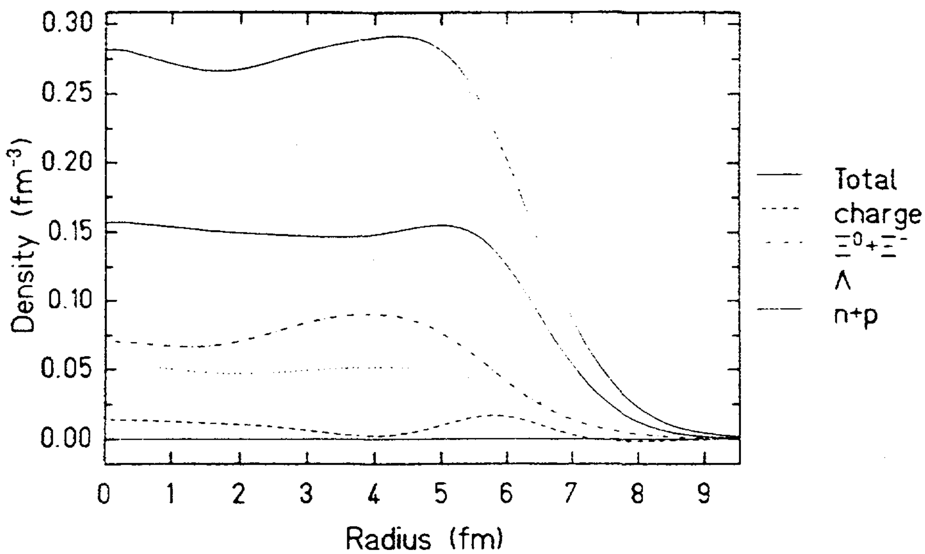

Figure 1: Vector density as a function of radius r for the ^{208}Pb + $70\Lambda + 18\Xi^0 + 70\Xi^-$ bound state in Model 1. The total density and the nucleon component thereof are shown as solid lines, while the Ξ, Λ, and charge densities appear as dot–dashed, dotted, and dashed curves, respectively [4].

The radial distributions of the vector densities for this $A = 366$, $S = -246$ system (^{208}Pb + 70Λ + $18\Xi^0$ + $70\Xi^-$) in Model 1 are displayed in Fig. 1. Although the nucleon density is close to that of ordinary nuclear matter ($\rho_0 \approx 0.16$ fm^{-3}), the total baryon density is close to $2\rho_0$. This higher density originates in the fact that the hyperons are distinguishable from each other and from nucleons, so each of the five species $\{n, p, \Lambda, \Xi^0, \Xi^-\}$ can occupy the same set of shell model orbitals. The density plots of Fig. 1 show that the Λ density readily leaks outside of the nuclear interior because of the relatively small binding of the last few levels. Comparing Ξ levels with nucleon levels, the Ξ's generally are less bound. Therefore, the total charge density changes from positive to negative as the radius r increases beyond 7 fm.

The SU(3) octet baryons are treated as distinguishable point particles in our RMF calculations [4]. As the density increases, this approximation must break down at some level, since a quark in a system of strongly overlapping composite particles need no longer be localized in a single hadron. One might therefore expect some effects of indistinguishability at the quark level, for instance, the effective "Pauli repulsion" generated by the requirements of antisymmetrization for quarks. These densities in Model 2 are not as smooth as those shown for Model 1 in Fig. 1, particularly the cascade densities which exhibit a double hump structure. The total vector density reaches peak values as high as $\rho \approx 3\rho_0$. The increased hyperon densities, and in particular the potentials generated by the σ^* and ϕ mesons in Model 2, cause the hyperon s.p. states to be considerably more bound.

In Model 2, we obtain a vast array of stable configurations containing only hyperons, bound by as much as 8.5 MeV/A. These systems, the lightest of which has $A = 6$, $S = -10$ $(2\Lambda + 2\Xi^0 + 2\Xi^-)$, are composites of all three hyperons $\{\Lambda, \Xi^0, \Xi^-\}$. Systems with only two species, $\{\Xi^0, \Xi^-\}$ to be specific, were found to be only marginally bound in some instances, with $E_B/A \approx -2$ MeV. In Model 1, on the other hand, the YY interaction is weak, and all purely hyperonic states are unstable.

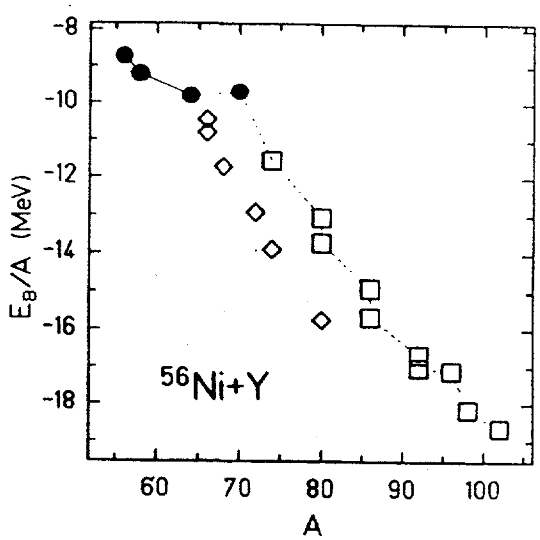

Figure 2: E_B/A *vs.* A for $^{56}\mathrm{Ni} + Y$ composites in Model 2. The solid circles correspond to $n_\Lambda = 0, 2, 8, 14$ with $n_\Xi = 0$. The other sequences (diamonds and squares) are obtained by adding Ξ's to systems with $n_\Lambda = 8, 14$. Each sequence terminates at the maximum A for which the system is stable with respect to strong baryon emission [4].

We now discuss some typical results for the binding energy per particle E_B/A of strange hadronic matter, as a function of A, $f_s = |S|/A$, and $f_q =$

q/A where $q = n_p - n_{\Xi^-}$ is the total charge. We also explore the degree of correlation of f_s and f_q. Our procedure to find metastable combinations of nucleons and hyperons is as follows: we start from magic normal nuclei (such as ^{16}O, ^{56}Ni, ^{132}Sn, ^{208}Pb, ...), add every magic combination of Λ's, Ξ^-'s, and Ξ^0's, and look for the four possible strong reactions ($\Lambda\Lambda \leftrightarrow N\Xi$) and their Q–values, respectively. If all Q–values are negative, the reactions are Pauli–blocked and the whole "nucleus" is metastable.

For systems consisting of a ^{56}Ni core plus a number of hyperons (denoted collectively as Y), the behavior of the binding energy per particle E_B/A as a function of A is shown in Fig. 2 for Model 2. The circles correspond to systems containing only nucleons and Λ's ($n_\Lambda = 0, 2, 8, 14$). For the next closed Λ shell, at $n_\Lambda = 18$, our calculation for $A = 74$ yields no stable $\{^{56}\text{Ni} + 18\Lambda\}$ configuration, since the $\Lambda\Lambda \rightarrow \Xi N$ conversion destabilizes it. The diamonds and squares correspond to some of the stable systems with Ξ hyperons.

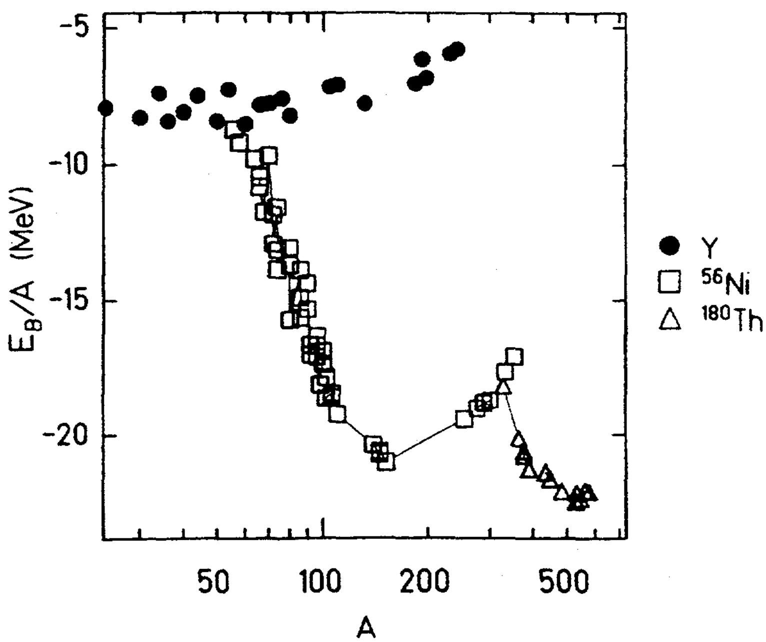

Figure 3: E_B/A *vs.* A for stable multi–strange states based on ^{56}Ni and ^{180}Th cores, for Model 2. Stable configurations Y consisting only of hyperons are indicated as black dots [4].

In Fig. 3, we plot E_B/A *vs.* A for systems of strange hadronic matter for Model 2. Several nuclear cores are used, including the doubly magic superheavy ^{310}G ($Z = 126$, $N = 184$) and the unstable core ^{180}Th ($Z = 90$, $N = 90$). The binding energy curves display a roughly parabolic shape on the average, on which is superimposed a jagged structure arising from shell effects. For Model 1, the minima lie at about $E_B/A \simeq -12$ to -13 MeV, and the region of stability extends to much higher A than for ordinary nuclei. The most important element in this increased binding energy and stability is the greatly diminished Coulomb repulsion, due to the addition of the Ξ^-'s. For Model 2 (Fig. 3), the minima lie at a larger binding of around $E_B/A \approx -21$ to -22 MeV.

The "gluelike" effect of hyperons is most evident in the ^{180}Th sequence. As an ordinary nucleus, ^{180}Th is far from the valley of stability (^{232}Th), and it readily emits nucleons. However, with a sizable infusion of strangeness, one can stabilize the ^{180}Th core, obtaining a region of stability extending from $A \simeq 330$ ($f_s = 0.73$, $f_q = 0.17$) to $A \approx 578$ ($f_s = 1.19$, $f_q = -0.08$) as shown by the triangles in Fig. 3. The simplest example of this effect is seen in $^6_\Lambda$He; this system is stable against nucleon emission, although the ^{5}He core decays strongly to $n + {}^4$He.

In Fig. 3, we note the existence of stable objects composed only of hyperons (designated as Y, not to be confused with the ordinary Ytterbium nucleus). We have only selected species bound by more than 5 MeV/A with maximal binding for each A. These systems, with typical binding energies $E_B/A \approx -5$ to -8.5 MeV, require a strong YY interaction (Model 2) to achieve stability. The most stable configuration has $A = 60$ $(14\Lambda + 28\Xi^0 + 18\Xi^-)$ and $E_B/A = -8.5$ MeV. The baryon number A of Y states cannot become arbitrarily large, since the Coulomb repulsion due to the Ξ^-'s, in the absence of protons, will lead to fission for large A, as for ordinary nuclei. In fact, these purely hyperonic systems may become particle unstable earlier than implied by nuclear fission, since the $\Lambda\Lambda \to \Xi N$ conversion prevents the use of more Λ's (than 34 in our Model 2) as A increases.

The balance of Ξ^0 and Ξ^- hyperons depends on the isospin dependent $(\vec{\tau}_1 \cdot \vec{\tau}_2)$ part of the $\Xi\Xi$ potential. If the isospin dependence is strong, the number of Ξ^0 and Ξ^- will remain about equal. However, this isospin potential for a Ξ is estimated to be small, based on the couplings of Nijmegen Model D [12]. In this case, the Coulomb potential dominates the symmetry (isospin) potential, and it becomes possible to obtain systems with $n_{\Xi^-} > n_{\Xi^0}$. In fact, we can have $n_{\Xi^-} = n_p$ without losing stability. In this case (zero net charge), we can obtain objects which are stable in the bulk limit ($A \to \infty$). The stability of strange hadronic matter in bulk may have some implications for the equation of state of neutron stars at high density.

If we start with a nuclear core A, we can typically obtain stable systems with baryon number up to $2A$ or more. The results for Model 2, involving ^{56}Ni and ^{180}Th cores, are shown in Fig. 4. For Model 1, the minimum lies near $E_B/A \approx -13.2$ MeV ($f_s \approx 0.6$), whereas for Model 2, larger binding is obtained at the minimum ($E_B/A \approx -22.4$ MeV) with much larger strangeness ($f_s \approx 1.1$). We also note the very large values $f_s \approx 1.8$ at the

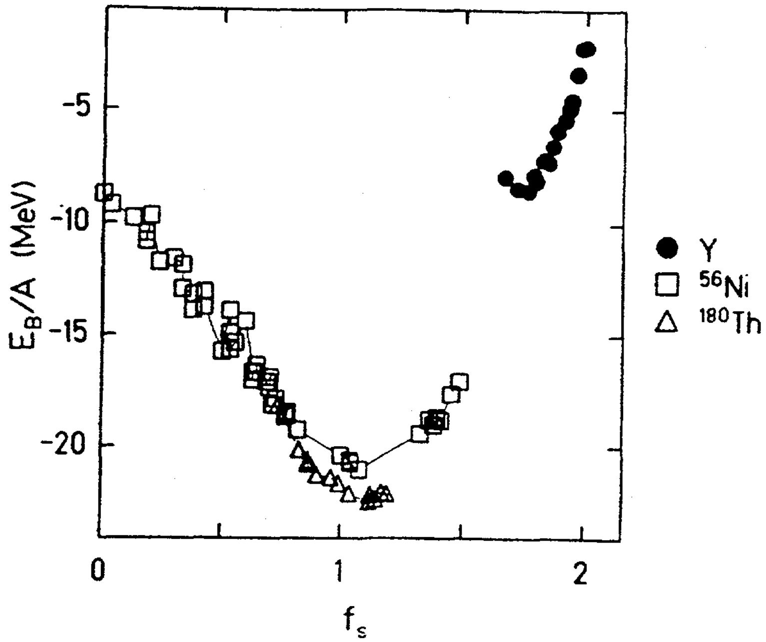

Figure 4: E_B/A *vs.* f_s for configurations of hyperons and ^{56}Ni or ^{180}Th cores, in Model 2. Purely hyperonic systems are shown as black dots [4].

minimum ($E_B/A \approx -8.5$ MeV) for the stable hyperon states Y in Model 2. The Y species displayed in Fig. 4 correspond to those which, for a given bin of f_s, are mostly bound. Hyperon matter has the features that the strangeness fraction $f_s = |S|/A$ always exceeds unity, and the charge fraction $f_q = q/A$ is always negative ($f_s = 5/3$, $f_q = -1/3$ for equal numbers of Ξ^0, Ξ^- and Λ). These values of strangeness and charge for bound hyperon systems are even more extreme than those for hypothetical stable strange quark matter [13–15]. The point $f_s = 2$ ($f_q = -1/2$) in Fig. 4 corresponds to a configuration with $n_\Lambda = 0$, $n_{\Xi^0} = n_{\Xi^-} = 8$.

Our results suggest a scaling property, namely that the f_s and f_q dependences are approximately independent of the sequence considered. A strong correlation, approximately linear, is seen to relate f_q and f_s, as shown in Fig. 5. For $N = Z$ cores, we have

$$f_q \approx \frac{1}{2}(1 - f_s) \tag{13}$$

This same relation characterizes isospin saturated strange quark matter. A similar linear relation is seen for $N > Z$ cores, with f_q normalized to Z/A at $f_s = 0$, and also for purely hyperonic matter.

The parallels of strange hadronic matter and strange quark matter are quite intriguing. The values of the density ρ, as well as f_s and f_q, are

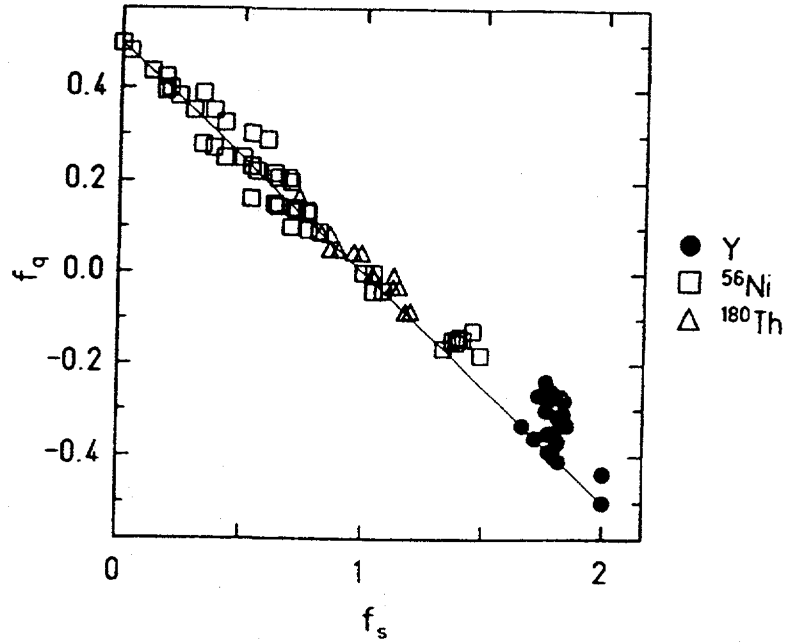

Figure 5: f_q vs f_s for Model 2. The solid line represents the linear correlation $f_q = (1 - f_s)/2$; from Reference [4].

quite similar for these two forms of matter. The difference lies in the values of E_B/A. In the RMF theory, the strengths of the mean fields are constrained by data on $S = -1, -2$ hypernuclei, and binding energies greater than $E_B/A \approx -20$ MeV or so are unlikely. Thus although the system can be stabilized against strong baryon emission, weak decays will always occur, with a typical time scale of 10^{-10} sec or less.

FINAL REMARKS

We have investigated the properties of multi–strange hadronic matter in the framework of the relativistic mean field theory. The essential phenomenological constraints on such a description consist of the observed binding energies of single Λ hypernuclei, as well as a few $S = -2$ systems. The conventional scalar (σ) plus vector (ω) theory (Model 1) can reproduce the Λ–hypernuclear data, but cannot account for the strongly attractive $\Lambda\Lambda$ interaction suggested by the analysis of the observed ${}^{6}_{\Lambda\Lambda}$He, ${}^{10}_{\Lambda\Lambda}$Be and ${}^{13}_{\Lambda\Lambda}$Be systems [6]. Accordingly, we have extended the model to include a second scalar (σ^*) – vector (ϕ) pair of meson fields, which couple strongly to hyperons, but not to nucleons. The σ^* and ϕ couplings are adjusted to produce the desired strong $\Lambda\Lambda$ interaction, and also imply a strong attraction for other hyperon pairs ($\Lambda\Xi$, $\Xi\Xi$). With $\sigma + \omega + \sigma^* + \phi$ meson fields (Model 2), the region of stability of strange hadronic matter is considerably broadened, extending to strangeness fractions f_s of order unity, and small or even negative net charges. In addition, systems of pure hyperon matter are predicted to be bound in Model 2. The smallest of these objects may be the $\Lambda\Lambda\Xi^0\Xi^0\Xi^-\Xi^-$ composite, although the mean field calculation is not reliable for such a light system.

An essential feature of our calculation is the stability with respect to the strong processes $\Xi N \leftrightarrow \Lambda\Lambda$. The presence of Λ's in bound shell model orbitals can stabilize the Ξ due to Pauli blocking. If too many Λ's are added, however, the decay $\Lambda\Lambda \rightarrow \Xi N$ occcurs, populating Ξ orbits. Thus multi–Λ systems cannot be considered in isolation: the Ξ^0 and Ξ^- are necessary constituents of multi–strange matter. The Σ's, on the other hand, are unlikely to be stabilized in the mean field picture, since the hyperon well depths remain smaller than the energy release $Q \approx 75 - 85$ MeV in the free space $\Sigma N \rightarrow \Lambda N$ reaction.

The existence of multi–strange nuclei and the possibility of stability in the bulk matter limit, are assured in both Models 1 and 2 over a wide range of choices for the coupling constants. We find that the main features of binding, such as the A, f_s and f_q dependences of E_B/A, hardly change with these variations, with the maximum binding for a given nuclear core changing by less than 1 MeV.

The presence of Ξ's in stable configurations increases the density and strangeness, and lowers the charge. The resulting conglomerates have $\rho \simeq (2-3)\rho_0$, $f_s \sim 1$, $f_q \sim 0$, rather similar to the properties of hypothetical stable states of strange quark matter. However, the binding energy per baryon is only of order $E_B/A \approx -10$ to -20 MeV. Within the dictates of the mean field theory, which is anchored to the phenomenology of $S = -1$ and -2 systems, the large bindings ($E_B/A \approx -100$ MeV or more) required to stabilize the system against weak decay are not available.

The mean field calculations discussed here represent a generalization of the nuclear shell model to include the strangeness degree of freedom in a self-consistent fashion. The SU(3) octet baryons $\{n,\, p,\, \Lambda,\, \Sigma^0,\, \Sigma^\pm,\, \Xi^0,\, \Xi^-\}$ are distinguishable at the hadron level, and are allowed to occupy the shell model orbitals $\{n\,\ell\,j\}$ independently. Strange hadronic matter exhibits a series of "magic numbers", just as in ordinary nuclei, for which the system is particularly stable. In calculations of strange quark matter, shell effects also occur [16], reflecting the operation of the Pauli exclusion principle at the quark level. Note that shell effects occur in strange quark matter even in the absence of quark–quark interactions, which are usually ignored or treated perturbatively. In mean field theory, the effects of interactions are included to all orders. The "magic numbers" are dependent on the model assumed for the strong interactions (they differ in Models 1 and 2).

Our mean field calculations represent a reasonable extrapolation from ordinary nuclear and hypernuclear physics to the domain of systems with large strangeness content. We are confident that there exists a vast array of stable objects composed of $\{n, p, \Lambda, \Xi^0, \Xi^-\}$ baryons, of arbitrarily large baryon number A, high strangeness content ($f_s \sim 1$), small net charge ($-0.1 < f_q < 0.1$), and rather weak binding ($E_B/A \approx -10$ to -20 MeV). There could conceivably exist a second branch of strange matter ("strangelets") which occupies a similar regime of stability in $\{A, f_s, f_q, \rho\}$, but with much larger values of E_B/A. There is the intriguing possibility that the weakly bound systems predicted in mean field theory could serve as "doorway states" for producing strangelets. For example, if a six quark H dibaryon [17] or a deeply bound strangelet $X(A = 7, S = -5, q = 1)$ exists with mass $M \simeq 7m_N$, one might expect to observe the strong decays

$$\begin{aligned} {}^{6}_{\Lambda\Lambda}\mathrm{He} &\to {}^{4}\mathrm{He} + H \\ {}^{7}_{\Xi^0\Lambda\Lambda}\mathrm{He} &\to K^+ + X \end{aligned} \tag{14}$$

rather than the sequential weak decays of the bound hyperons. The observation of weak decays would give limits on the mass range of H or X. Only the lightest multi–strange objects predicted here are likely to be produced with measurable rates in relativistic heavy ion collisions. Some conservative estimates based on the coalescence model are given in ref. [8]. If quark–gluon plasma is formed in a central heavy ion collision, considerably larger rates for the production of multi–strange objects can be anticipated [18]. The lightest stable object containing a Ξ is likely to be ${}^{7}_{\Xi^0\Lambda\Lambda}\mathrm{He}$. This system, the prototype of the Pauli blocking mechanism for stabilizing Ξ's in a many–baryon bound state, merits a search in a high sensitivity heavy ion experiment.

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