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Strangelets and Strange Quark Matter

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I summarize the properties of finite lumps of strange quark matter (strangelets) with emphasis on the two scenarios of producing strange matter in relativistic heavy ion collisions. As an outlook, I discuss the possibility of short-lived strange composites and charmed matter.

1. INTRODUCTION – PRODUCING STRANGE MATTER IN THE LABORATORY

As it has been discussed widely during this conference, strangeness opens a new dimensions to nuclear physics. Insofar, systems with strangeness number $S = -1, -2$ have been discussed. Here we want to examine the unknown domain of finite nuclear systems with $S < -2$. There have been speculations about the existence of finite systems of strange quark matter (strangelets) and strange hadronic matter. Here we will focus on the former objects and recent progress in this field [1] as the latter ones were discussed at the last hypernuclear meeting in detail [2]. How can one produce such strangeness-rich systems? Hadron beams enable only to explore systems up to $S = -2$. Nevertheless, relativistic truly heavy-ion collisions constitute a prolific source of strangeness as dozens of hyperons are produced on a single central event. In principle, strangelets can be produced via two different scenarios: by a coalescence of hyperons or by a distillation of a quark-gluon plasma.

The coalescence model for strangelet production in heavy-ion collisions has been put forward by Carl Dover [3]. The formation of a quark-gluon plasma is not needed in this scenario. Hyperons coalesce during the late stage of the collision forming a doorway state for strangelet production. For example, the H dibaryon would be formed by the coalescence of two Λ's or a Ξ with a nucleon which transform to a dibaryon with the same quantum numbers. The production rates are proportional to two penalty factors, one for adding a baryon number and the second one for adding one unit of strangeness to the clusters

$$P \propto q^{|A|} \cdot \lambda^{|S|} , \quad q = \frac{N_d}{N_N} , \quad \lambda = \frac{N_Y}{N_N} \quad . \tag{1}$$

Here N_i are the numbers of produced particles of the species i per collision. The penalty factors can be extracted from experiment or estimated from phase-space arguments using a cascade model [4,5]. At the AGS one finds at an energy of 14 AGeV $q \approx 0.03$, $\lambda \approx 0.13$,

*Feodor-Lynen fellow of the Alexander-von-Humboldt Stiftung, Germany while at the SPS at 200 AGeV it is $q \approx 0.0075$, $\lambda \approx 0.35$. Composites with a high baryon number are more suppressed than those with a high strangeness fraction. This feature gets more pronounced at higher bombarding energy. Assuming a sensitivity of 10^{-8}, clusters up to a baryon number of $A \leq 6$ can be produced via coalescence at the AGS, for SPS it is $A \leq 4$. Hence, very low mass numbers are expected to be seen.

On the other hand, the distillation process provides the possibility for producing strangelets up to $A = 20 - 30$ [6]. This assumes the formation of a quark-gluon plasma with a high net baryon number. Recently, it was shown that this scenarios could also hold for a quasi-baryon free region as possibly encountered at the future heavy-ion colliders RHIC and LHC due to large baryon number fluctuations [7]. The main idea is, that strange-antistrange quark pairs are abundantly produced inside a quark-gluon plasma. The antistrange quark leaves the plasma and forms a kaon with a surrounding light quarks which is ensured for a baryon-rich regime. This enriches the plasma with a finite amount of strangeness. Cooling by particle evaporation further enhances this effect and produces finally a cold metastable strangelet. Strangeness fractions of $f_S = |S|/A \geq 1$ are easily reached during the distillation process.

We will now discuss the properties of strangelets relevant for heavy-ion physics, i.e. for low baryon numbers.

2. PROPERTIES OF STRANGE MATTER

2.1. Strangelets in its ground state

The first speculation about the possible existence of collapsed nuclei, was done by Bodmer [8]. He argued that another form of matter might be more stable than ordinary nuclei. Transition to one form to the other is suppressed by a barrier so that it does not occur during the lifetime of the universe. He discussed three different forms: the nucleon model with large positive charge and a high density ($R \approx 0.5$ fm), the general baryon model with small positive charge and large strangeness, and the quark model with small positive charge and hypercharge for large baryon numbers. The second form is now called strange hadronic matter [9], while the latter one strangelets [10]. Note, that the paper lacks detailed calculation, as the MIT bag model [11] as well as the Walecka model [12] were only available a few years later.

Let us discuss briefly the MIT bag model. There are five different terms which contribute to the total mass of the bag: the volume term proportional to the bag pressure constant B, the purely phenomenological zero point energy, the kinetic energy (which is a sum over single-particle energies), and the two terms coming from the color magnetic and the color electric interaction between the quarks. The dominant term is the volume term, so that the bag pressure of bag parameter B more or less determines the mass of a strangelet. The parameters can be fixed to hadron masses. Nevertheless, the value of the bag parameter can not be fixed unambiguously. The original value of $B^{1/4} = 145$ MeV [11] has to be contrasted with $B^{1/4} = 235$ MeV from a fit to charmonium levels [13] which is compatible with estimates from QCD sum rules [14]. Furthermore, the coupling constant of the one-gluon exchange extracted from the fit to hadron masses are so large, that the pressure for massless quarks gets negative in bulk. This demonstrates that the bag model is an effective approach with effective parameters which will change when going to larger baryon numbers and/or strangeness. Therefore, we will study the properties of strange matter for a variety of parameters.

For a bag parameter of $B^{1/4} = 145$ MeV, strange quark matter is absolutely stable (this is Witten's scenario [15]). For bag parameters larger than $B \geq 210$ MeV strange quark matter is unstable. In between, it is metastable and can decay by weak interactions. The strange quark as a new degree of freedom lowers the Fermi energy and creates a global minimum located at $f_S \approx 0.7$. This means that heavy strangelets are slightly positively charged [16]. The masses of strangelets can be estimated from the MIT bag model when neglecting the color exchange contributions. One finds that finite size effects make light strangelets with $A \leq 20$ metastable even when strange quark matter is absolutely stable [17]. Moreover, there exists a broad range of charges with a quite similar binding energy. Note that the color magnetic and color electric potentials are neglected in these calculations due to their complicated group structure. The lightest strangelets have been studied in the full bag model including these two terms up to $A = 6$ [18]. It turns out that all strangelet candidates for this mass range are not bound with the well-known exception of the H dibaryon. Hence we will focus on strangelets with masses around $A = 6 - 30$ in the following.

2.2. Timescales – strong and weak decay of strangelets

Strangelets will not be in their ground state when being produced in a heavy-ion collision. Suppose, a strangelet is created in the hot and dense matter with some arbitrary strangeness, charge and baryon number. Strong interactions, the distillation process and particle evaporation will alter the composition of a strangelet on a timescale of a few hundred fm after the collision. The strangelet, if surviving this, will cool down until it reaches the domain of weak interactions. Weak hadronic decay by hadron emission takes place between $\tau = 10^{-5} - 10^{-10}$ s. Strangelets stable against strong interactions but decaying by the weak hadronic decay will be dubbed short-lived. Strangelets stable against weak hadronic decay will then be subject to the weak leptonic decay happening on a timescale of $\tau = 10^{-4} - 10^{-5}$ s. They are called long-lived in the following. The weak leptonic decay is suppressed by phase space as it involves a three-body final state. Note, that most experiments are able to detect strangelets up to a lifetime of 50 ns [19]. Recently, also the domain of short-lived strange matter is accessible by new experiments as presented at this conference [20,21].

Figure 1 shows the energy per baryon for strangelets up to $A = 40$. The line connecting the nucleon and hyperon masses stands for a free mixture of baryons. Strangelets with a strangeness fraction of $f_s < f_s^{crit}$ are lying above this line. Therefore, they are unstable and decay to nucleons and Λ's. The tangent drawn from the nucleon mass stands for a mixture of nucleons and strangelets. Strangelets in the intermediate range of $f_s^{crit} < f_s < f_s^{mix}$ will 'move' to the tangent point increasing their strangeness fraction to f_s^{mix} by emitting nucleons as this is energetically favourable. Hence, a strangelet will be highly charged with strangeness when surviving the first 100 fm of the heavy-ion collision. This is the distillation process as discussed in a dynamical approach [6]. Note that this chain of argumentation applies whenever strange quark matter is metastable and when there is a minimum in the energy per baryon at a finite f_s! Strangelets with a high strangeness fraction are most likely negatively charged, as for isospin symmetric matter $Z/A = (1 -$

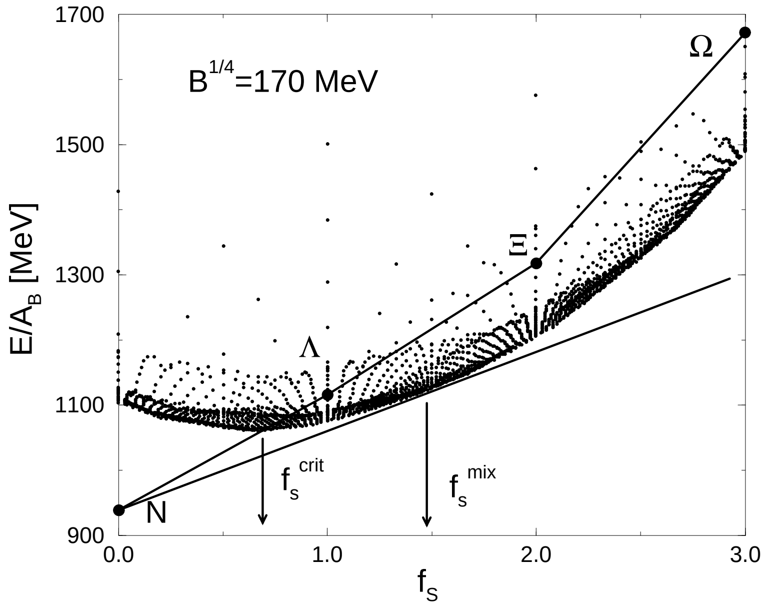

Figure 1. Energy per baryon for strangelets up to $A = 40$ using shell-mode filling for a bag parameter of $B^{1/4} = 170$ MeV. Note the pronounced shell effects.

$f_s)/2$. This gives $Z/A < -0.2$ for short-lived strangelets with $f_s \geq f_s^{mix} = 1.4$. Note also that shell effects are quite pronounced, at the order of 100 MeV/A, which we will discuss in the next subsection in more detail.

Next, strangelets will be subject to weak hadronic decay. The two major reactions are the weak nucleon decay ($Q \to Q' + N$) and the weak pion decay ($Q \to Q' + \pi$) by changing one unit of strangeness as discussed in [10]. The change in the strangeness fraction for the weak nucleon decay is $\Delta f_s = (f_s - 1)/(A - 1)$. For a strangelet with $f_s > 1$ this means that the weak nucleon decay enhances the strangeness fraction to even higher values. This is only true for bulk matter. Finite size effects and pockets in the energy per baryon will certainly alter this conclusion. We want to point out that the weak nucleon decay can only carry away positive charge. Only the weak neutron decay accompanied by an emission of a π^- carries away negative charge ($Q \to Q' + n + \pi^-$) which is suppressed by phase space. Hence, the weak nucleon decay will also drive a strangelet to negative charge.

2.3. Shell effects - the valley of stability

As we have seen, shell effects seems to be quite pronounced. Figure 2 shows the single particle energy of the quark alpha which has six up, down, and strange quarks. The lowest lying state is a 1s1/2, which can be occupied by six quarks due to the color degree of freedom. For the quark alpha, this state is filled up completely for all three quark species and one can call it triple magic then. The next states are the 1p3/2 and the 1p1/2.

Note that the level splitting is enormous between the 1s and the 1p shells (more than 100 MeV!) but also the spin-orbit splitting is huge, about 70 MeV! Therefore, the magic numbers for strangelets (6, 18, and 24) will be much more pronounced than for ordinary nuclei. The figure also shows that the single particle energy for nonstrange and strange quarks are getting quite similar for the higher lying states as the strange quark mass gets negligible compared to the kinetic energy. This will allow for filling up the strange quark levels without loosing stability, if the effective strange quark mass is not too high.

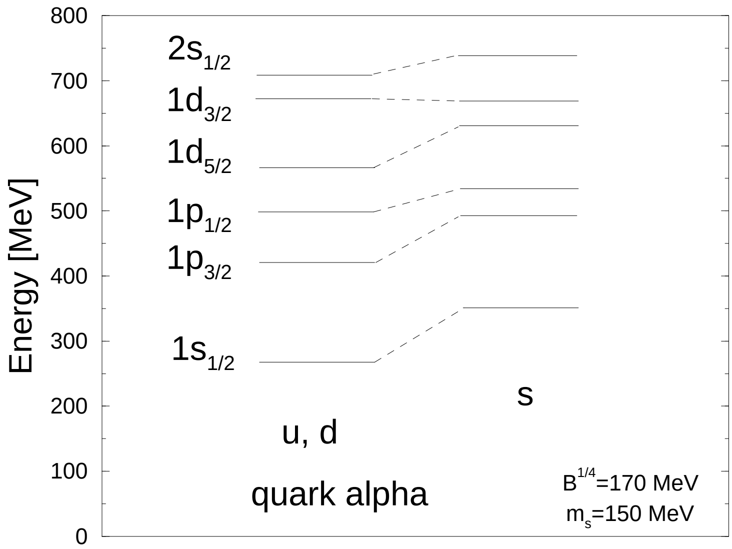

Figure 2. The shell levels of the quark alpha for a bag parameter of $B^{1/4} = 170$ MeV.

Magic strangelets with closed shells (triple magic) will appear irrespectively of the bag parameter chosen. Magic strangelets with a high amount of strange quarks and negative charge, as they most likely will survive a heavy ion collision, are for example for $N_u = N_d = 6$, $N_s = 18, 24$ and $N_u = 6$, $N_d = 18$, $N_s = 18, 24$. It is interesting to look at the masses and charges of these four strangelets: $A = 10$ with $Z = -4$, $A = 12$ with $Z = -6$, $A = 14$ with $Z = -8$, and $A = 16$ with $Z = -10$. These candidates form a valley of stability for strangelets and are highly negatively charged.

A detailed calculation in the MIT bag model including shell effects and all the hadronic decay channels has been done in [1]. It was shown that shell effects are important for long-lived candidates which are stable against weak hadronic decay. The above listed strangelet candidates turn out to be long-lived for most of the parametrizations used. Figure 3 shows the result of this investigation. Different bag parameters have been chosen but the valley of stability remains. For bag parameters of $B^{1/4} = 180$ MeV or larger, no long-lived candidates have been found at all. In these case, strange quark matter starts getting unstable against hadron emission already in bulk. The candidates will live on a timescale

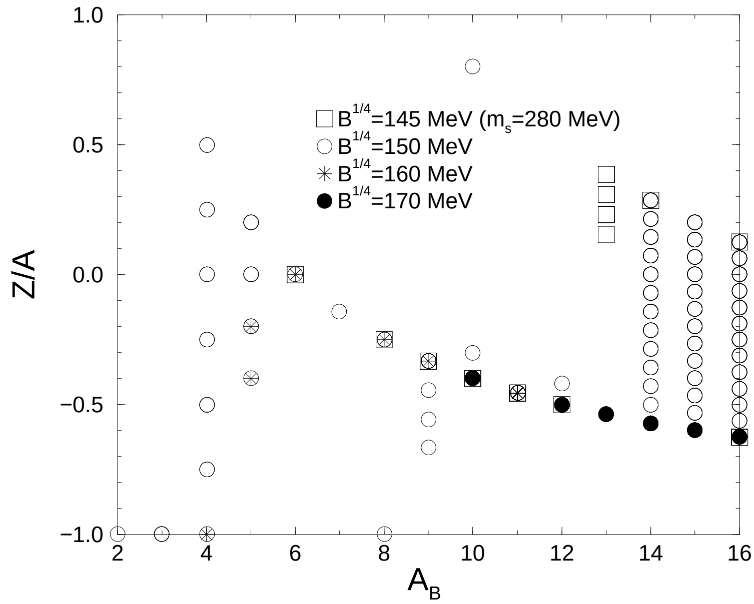

Figure 3. Long-lived strangelet candidates and their masses and charge-to-mass rations. The valley of stability is clearly visible.

of the weak leptonic decay $\tau = 10^{-4} - 10^{-5}$ s and have a charge-to-mass ratio close to an antideuteron $Z/A \approx -1/2$.

We conclude, that

- if the relativistic shell model is used, and
- if strange quark matter is at least metastable
- and has a local minimum at some strangeness fraction

then a valley of stability appears at $A = 10 - 16$ with $Z \leq -4$. If that is correct, then it has important impacts for the present searches for strangelets in heavy-ion collisions. Presently, the experiments focus on objects with a small positive charge which is only true for heavy strangelets. Nevertheless, experiments like E864 are also able to detect these highly charged strangelet candidates with their present setup [19].

3. OUTLOOK

3.1. Charm Matter

Insofar, we discussed the properties of quark matter including the strange quark. What happens when going to the next heavier quark, the charm quark with a mass of $m_c \approx 1.5$ GeV? Compared to the lighter quarks, the kinetic energy of the charm quark will be dominated by its mass. Moreover, the charm quark is heavier than the Fermi levels, therefore absolute stable charm matter most likely does not exist and will decay by weak interactions on the timescale of the lifetime of charmed hadrons, $\tau \approx 10^{-12}$ s. Nevertheless, there have been speculations about the existence of the Pentaquark with 4 light quarks (u-,d-, or s-quarks) and one anticharm quark [22]. Due to color magnetic interactions between the light quarks, the Pentaquark might be bound. Most recently, charmlets have been studied in a modified bag model [23]. The charm quark feels a strong attractive potential from the color electric term, while the light quarks are bound by color magnetic forces. This two effects enable charm matter with strange quarks to be bound even for the case of $B^{1/4} = 235$ MeV where pure strange quark matter is unstable.

Estimates for the production of charmed matter in heavy-ion collisions have been done by Carl Dover [24]. He estimates in a coalescence model that every 10^5 event at SPS might be able to produce a Pentaquark state. He also demonstrates, that charm quarks are abundantly produced at the heavy-ion collider RHIC. Coalescence estimates in the same spirit done in [23] shows that charmlets with $A + |C| \leq 4$ might be produced at every 10^6 event at RHIC. This opens the exciting perspective of probing the nuclear chart into a fourth dimensions besides baryon number, isospin and strangeness!

3.2. Dihyperons

Recently, it became possible to detect in heavy-ion collisions at the AGS (the EOS experiments E895, E910 [21]) as well as at SPS (NA49) directly weakly decaying short-lived hadrons by their decay topology in time projection chambers. In addition, one experiment (E896) is especially designed to search for short-lived strange matter [20]. They are able to detect short-lived $S < -2$ systems in the laboratory for the first time. Let us discuss in the following strange composites with $A = 2$.

A possible bound state of $(\Xi^0 \mathrm{p})_b$ can decay by weak nonmesonic decay to a Λ and a proton. The decay topology is exactly the same for the decay of a Ξ^- which has been seen in the TPC's already but with the opposite charge. A bound state of $(\Xi^0\Lambda)_b$ might decay to two Λ's. A study of $\Lambda\Lambda$ correlations, which is under investigation [21], or backtracking techniques will reveal this exotic state. There exists a lot of other candidates which will have decay patters distinct from ordinary hadrons. As E896 is also able to see neutrons [20], a decay pattern of $(\Lambda\Xi^-)_b \to \Sigma^- + \Lambda$ will be measurable in future runs.

The production rates for dihyperons are quite high. Coalescence estimates using the values given in this paper gives 0.2 and 0.03 dihyperons per central AuAu collision at the AGS for $S = -2$ and $S = -3$, respectively. At the SPS one gets 0.4 and 0.1 per event for $S = -2$ and $S = -3$, respectively.

Even if dihyperons are not bound, correlation measurements of hyperons will reveal important information about the hyperon-hyperon interactions. Note that the three double Λ hypernuclear events constitutes the only information about the hyperon-hyperon interaction so far! Any information about this would be of uttermost importance for our knowledge about the interaction of baryons and the underlying symmetry (for example, does it fulfill approximately SU(3) symmetry ?). Carl Dover studied the hyperon-hyperon interaction using the Nijmegen model D [25]. He found that the potentials between the hyperons is strongly attractive. Nevertheless, the existence or nonexistence of bound dihyperon states depends crucial on the value of the cutoff radius which is unknown. Experiments looking into this hitherto unexplored domain of nuclear physics are therefore impatiently awaited.

Acknowledgments

This work is dedicated to late Carl B. Dover. As the outlook into the domain of multiply strange nuclear systems and into the domain of charm matter has demonstrated, Carl Dover has done it already and paved the way to new physics for future generations. I owe him a most fruitful collaboration and many helpful and illuminating conversations from which I learned a lot. J.S.B. acknowledges support by the Alexander-von-Humboldt Stiftung with a Feodor-Lynen fellowship. This work is supported in part by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Nuclear Physics Division of the U.S. Department of Energy under Contract No. DE-AC03-76SF00098.

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