Stable quark matter in cosmic rays ?

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Stable lumps of quark matter may be present in cosmic rays at a flux level, which can be detected by high precision cosmic ray experiments sensitive to anomalous "nuclei" with high mass-to-charge ratio. The properties of these lumps, called strangelets, are described, and so is the production and propagation of strangelets in cosmic rays. Two experiments underway which are sensitive to a strangelet flux in the predicted range are briefly described. Finally it is summarized how strangelets circumvent the acceleration problem encountered by conventional candidates for ultra-high energy cosmic rays and move the Greisen-Zatsepin-Kuzmin cutoff to energies well above the observed maximum energies.

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1. Introduction

Iron and nickel nuclei are normally assumed to be the most stable form of hadronic matter at zero external pressure. In principle, this should be testable from the basic theory of strong inter-actions, Quantum Chromo Dynamics (QCD), but in practice this is impossible in the foreseeable future. QCD is not suited for finite density calculations and calculations involving many degrees of freedom, so much of our theoretical knowledge about dense matter is based on phenomenological model calculations that try to incorporate and parametrize some of the main features of the strong interactions, such as confinement and asymptotic freedom. In some of these studies, notably stud-ies based on the MIT bag model, it has been shown, that there is a significant range of parameters for which a three-flavor quark phase with roughly equal numbers of up, down and strange quarks (called strange quark matter in bulk, and strangelets in small lumps) has lower energy than a nu-cleus with the same baryon number. Thus, strangelets rather than nuclei may be the ground state of hadronic matter [1, 2, 3, 4, 5, 6].

A range of questions immediately occur if strangelets are more stable than nuclei:

1. Why are we here—why don't our nuclei decay?

2. Can strangelets be created in the laboratory?

3. Can strange quark matter be found in space?

The answers to all three questions rely on the properties of strangelets summarized in the following. The answers to the questions (some of which will be further explored below) are briefly:

1. Very low baryon number strangelets are likely to be unstable, even if strange matter in bulk is absolutely stable, and for intermediate baryon numbers the transformation of a nucleus requires an improbable high-order weak interaction to simultaneously transform of order *A* up and down quarks into strange quarks, where *A* is the baryon number. This explains the stability of our nuclei. Strangelets have a positive electric charge, and therefore repel nuclei from the surface, thus minimizing the risk of growth via absorption of nuclei [3, 4, 7, 8, 9, 10].

2. In principle strangelets might be formed in ultrarelativistic heavy-ion collisions by coales-cence or through a distillation mechanism leading to strangeness enhancement. However, the available baryon number is low, which makes it difficult to cross the low-*A* stability cutoff, and furthermore the environment is hot, so the process has rightly been compared to the cre-ation of ice cubes in a furnace. Experiments have been performed, but with negative results [11, 12, 13].

3. The most promising place to search for strange quark matter is in the cosmos. Originally strange quark matter "nuggets" surviving from the cosmological quark-hadron phase transi-tion were suggested as a possible and even natural candidate for the cosmological dark matter [3], but later studies showed, that this was probably unlikely due to the high temperature en-vironment that would lead to evaporation of the nuggets [14, 15, 16, 17, 18, 19]. Recent lattice results indicating, that the quark-hadron transition at low chemical potential (like the

early Universe) is not a first order phase transition seems to rule out the possibility of cosmological strange quark matter. But strange quark matter, if stable, is almost unavoidable in dense stellar objects. Pulsars and compact x-ray sources will contain strange stars containing bulk strange quark matter [3, 20, 21], and a number of observational signatures have been suggested to distinguish these from compact stars made of ordinary hadronic matter. This is an interesting story in its own right (see [5, 6] for reviews), but the focus of the present presentation is the possibility of directly observing debris from collisions of strange stars in the form of strangelets in cosmic rays.

In the following I will summarize the main properties of strangelets, describe the propagation of strangelets in cosmic rays, estimate the flux to be expected in cosmic ray detectors, and discuss the possible detection in a couple of future experiments. Finally I will argue that strangelets may also be relevant for the puzzle of ultra-high energy cosmic rays.

2. Strangelet properties

Neglecting quark Cooper pairing (to which we return in a moment), and approximating the up and down quark current masses to zero, the mass of strangelets in the MIT bag model [22] depends on three model parameters, namely the strange quark mass, m_s, the bag constant, B, and the strong fine-structure constant, α_s. In most treatments, the strong coupling and therefore α_s is set equal to zero. This is clearly not correct at the relevant densities, but it has been shown [4], that a nonzero value can be mimicked very well by a scaling of the bag constant, and taking $\alpha_s = 0$ simplifies calculations. A value of B between $B^{1/4} = 145\text{MeV}$ and $B^{1/4} = 165\text{MeV}$ allows bulk strange quark matter to be absolutely stable for not-too high strange quark mass. In fact, for $B^{1/4} < 145\text{MeV}$ even 2-flavor quark matter becomes absolutely stable, having energy per baryon lower than 930MeV, which is ruled out by the stability of nuclei. Such a transformation would not require any weak conversion of quark flavor, whereas a similar decay into strangelets requires almost one-third of the quarks transformed simultaneously into strange quarks, a highly suppressed process.

Masses of strangelets can be calculated in the MIT bag model by explicitly solving the Dirac equation with bag model boundary conditions, filling up, down, and strange quark energy levels in the way which minimizes the mass for a given A. Apart from details of closed shells the outcome of such results can be reproduced fully in a multiple reflection expansion including volume, surface, and curvature terms [4, 23, 24, 25, 26, 27].

For a strange quark mass of $m_s = 150\text{MeV}$ and a bag constant $B^{1/4} = 145\text{MeV}$, the mass of a strangelet in the lowest energy ground state as a function of baryon number is given by [26, 27]

$$M(A) = 874\text{MeV}A + 77\text{MeV}A^{2/3} + 232\text{MeV}A^{1/3}. \tag{2.1}$$

Thus, for these parameters bulk quark matter is bound by $(930 - 874)\text{MeV} = 56\text{MeV}$ per baryon. Notice that the mass is composed of a bulk term proportional to A, a surface term proportional to surface area or $A^{2/3}$, and a non-negligible curvature term proportional to radius or $A^{1/3}$. For this specific choice of parameters, only strangelets with baryon number $A > 23$ are stable relative to nuclei. Increasing the strange quark mass and/or the bag constant moves the stability cut to higher A. Compared to mass formulae for nuclei, the most striking features are the presence of

a significant curvature energy and the lack of a significant Coulomb energy. Lack of the latter also means the absence of a minimum in M/A as a function of A.

The lack of a significant Coulomb energy is due to the fortuitous cancellation of charge $+2e/3$ up quarks and charge $-e/3$ down and strange quarks in strange quark matter with equal numbers of the three quark flavors. Because of the non-zero s-quark mass the cancellation is not perfect. Typical strangelets have slightly fewer strange quarks compared to up and down, and therefore the net charge is slightly positive. A typical model result (to be compared to $Z \approx 0.5A$ for nuclei) is [28]

$$Z = 0.1 \left(\frac{m_s}{150\text{MeV}} \right)^2 A \tag{2.2}$$

$$Z = 8 \left(\frac{m_s}{150\text{MeV}} \right)^2 A^{1/3} \tag{2.3}$$

for $A \ll 700$ and $A \gg 700$ respectively (the slower growth for higher A is a consequence of charge screening).

Thus a unique experimental signature of strangelets is an unusually high mass-to-charge ratio compared to nuclei.

In recent years it has been shown, that quark matter at asymptotically high density has an interesting property called color superconductivity [29, 30]. Even the weakest attraction (and such attractions exist in QCD) leads quarks of different colors and flavors to form pairs, much like Cooper pairs in a superconductor, except that the binding in QCD is caused by a direct attraction channel rather than via indirect phonon interaction. The binding energy of a pair can be very large, ranging from MeV to over 100MeV. In general these systems are called color superconductors. If all colors and flavors pair in an equal manner one talks about color-flavor locking.

While the phenomenon of color-flavor locking seems generic in the infinite density limit, the properties of strange quark matter at densities of order or somewhat higher than nuclear matter density is at the focus of much current research and discussion. This is the density regime of relevance for strangelets, strange stars, and for quark matter cores in hybrid stars (the analogs of neutron stars if quark matter is metastable so that it forms above a certain density in compact star interiors). An additional binding energy per baryon of $(10 - 100)\text{MeV}$ is not unrealistic in these systems, thus significantly increasing the likelihood of absolutely stable strange quark matter and strangelets.

Since Cooper pairing involves quarks with equal (but opposite) momenta, the natural ground state of a color-flavor locked system has equal Fermi momenta for up, down, and strange quarks. Equal Fermi momenta implies equal number densities, and therefore a total net quark charge of zero for a bulk system [31]. However, a finite size strangelet has a surface suppression of massive strange quarks relative to the almost massless ups and downs (massive particle wave functions are suppressed at a surface), so that the total charge of a color-flavor locked strangelet is positive and proportional to area [32]:

$$Z = 0.3 \left(\frac{m_s}{150\text{MeV}} \right) A^{2/3}. \tag{2.4}$$

This phenomenon persists even for very large bags, such as strange stars, so color-flavor locked strange stars would also have a positive quark charge [32, 33, 34].

3. Astrophysical strangelet production

Strangelets may be produced when two strange stars in a binary system collide due to loss of orbital energy in the form of gravitational radiation. A strange star–black hole collision may also release lumps of quark matter. If strange quark matter is absolutely stable all compact stars are likely to be strange stars [35, 36], and therefore the galactic coalescence rate will be the one for double neutron star binaries recently updated in [37] based on available observations of binary pulsars to be $83.0^{+209.1}_{-66.1}$ Myr^{-1} at a 95% confidence interval, thus of order one collision in our Galaxy every 3,000–60,000 years (but see [38] for a somewhat lower rate estimate).

Each of these events involve a phase of tidal disruption of the stars as they approach each other before the final collision. During this stage small fractions of the total mass may be released from the binary system in the form of strange quark matter. No realistic simulation of such a collision involving two strange stars has been performed to date. Newtonian and semirelativistic simulations of the inspiral of strange stars and black holes do exist [39, 40, 41], but the physics is too different from the strange star-strange star collision to be of much guidance. Simulations of binary neutron star collisions, depending on orbital and other parameters, lead to the release of anywhere from $10^{-5} - 10^{-2} M_\odot$, where $M_\odot$ denotes the solar mass, corresponding to a total mass release in the Galaxy of anywhere from $10^{-10} - 3 \times 10^{-6} M_\odot$ per year with the collision rate above. Given the high stiffness of the equation of state for strange quark matter, strange star-strange star collisions should probably be expected to lie in the low end of the mass release range, so the canonical input for the following calculations is a galactic production rate of

$$\dot{M} = 10^{-10} M_\odot \text{yr}^{-1}. \tag{3.1}$$

All strangelets released are assumed to have a single baryon number, A. This is clearly a huge oversimplification, but there is no way of calculating the actual mass spectrum to be expected. As demonstrated in [42] the quark matter lumps originally released by tidal forces are macroscopic in size; when estimated from a balance between quark matter surface tension and tidal forces a typical fragment size is

$$A \approx 4 \times 10^{38} \sigma_{20} a_{30}^3, \tag{3.2}$$

where $\sigma_{20} \approx 1$ is the surface tension in units of $20 \, \text{MeV/fm}^2$ and a_{30} is the distance between the stars in units of 30 km. But subsequent collisions will lead to fragmentation, and under the assumption that the collision energy is mainly used to compensate for the extra surface energy involved in making smaller strangelets, it was shown [42] that a significant fraction of the mass released from binary strange star collisions might ultimately be in the form of strangelets with $A \approx 10^2 - 10^4$, though these values are strongly parameter dependent.

The total flux results derived for cosmic ray strangelets below are mostly such that values for some given A are valid as a lower limit for the flux for a fixed total strangelet mass injection if strangelets are actually distributed with baryon numbers below A.

4. Strangelet flux in cosmic rays [1]

Apart from an unusually high A/Z-ratio compared to nuclei, strangelets behave in many ways like ordinary cosmic ray nuclei. For example, the most likely acceleration mechanism would be Fermi acceleration in supernova shocks resulting in a rigidity spectrum at the source which is a powerlaw in rigidity as described below. Due to the high strangelet rigidity, R, at fixed velocity, v ($R \equiv pc/Ze = Am_0c^2\gamma(\beta)\beta/Ze$, where p is the strangelet momentum, Am_0 is the strangelet rest mass, $\beta \equiv v/c$, and $\gamma \equiv (1-\beta^2)^{-1/2}$) strangelets are more efficiently injected into an accelerating shock than are nuclei with $A/Z \approx 2$ (c.f. discussion of nuclei in [44]), and most strangelets passed by a supernova shock will take part in Fermi acceleration.

The time scales for strangelet acceleration, energy loss, spallation and escape from the Galaxy are all short compared to the age of the Milky Way Galaxy. This makes it reasonable to assume that cosmic ray strangelets are described by a steady state distribution given as a solution to a propagation equation of the form

$$\frac{dN}{dt} = 0 \tag{4.1}$$

where $N(E,x,t)dE$ is the number density of strangelets at position x and time t with energy in the range $[E, E+dE]$.

Given a solution for $N(E)$ the corresponding flux in the "average" interstellar medium with energies from $[E, E+dE]$ is given by

$$F_E(E)dE = \frac{\beta c}{4\pi}N(E)dE, \tag{4.2}$$

and the corresponding flux in terms of rigidity is

$$F_R(R)dR = Ze\beta F_E(E(R))dR \tag{4.3}$$

(using $dE/dR = Ze\beta$).

Like other charged cosmic ray particles strangelets are influenced by the solar wind when entering the inner parts of the Solar System. The detailed interactions are complicated, but as demonstrated for nuclei in [45], a good fit to the solar modulation of the cosmic ray spectrum can be given in terms of a potential model, where the charged particle climbs an electrostatic potential of order $\Phi = 500$ MeV (the value changes by a factor of less than 2 during the 11 (22) year solar cycle). This effectively reduces the cosmic ray energy by $|Z|\Phi$ relative to the value in interstellar space, and at the same time the flux is reduced by the relative reduction in particle momentum squared, so that the modulated spectrum is

$$F_{\text{mod}}(E) = \left(\frac{R(E)}{R(E+|Z|\Phi)}\right)^2 F_E(E+|Z|\Phi). \tag{4.4}$$

Solar modulation significantly suppresses the flux of charged cosmic rays at energies below a few GeV and effectively works like a smooth cutoff in flux below kinetic energy of order $|Z|\Phi$. Since

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[1]Section 4 closely follows Ref. [43], but the treatment has been slightly simplified by excluding less important terms in the propagation equation.

strangelets have a high mass-to-charge ratio they are nonrelativistic at these energies, which correspond to rigidities of $R_{\mathrm{SM}} \approx (A/Z)^{1/2}\Phi_{500}^{1/2}$ *GV*, where $\Phi_{500} = \Phi/(500\mathrm{MeV})$.

For cosmic rays to reach the Earth or an Earth-orbiting detector like the Alpha Magnetic Spectrometer on the International Space Station, the rigidities have to exceed the geomagnetic cutoff rigidity, which is a function of detector position, and for an orbiting observatory like AMS the value varies from 1–15 GV as a function of time. The geomagnetic cutoff rigidity for low mass strangelets is comparable to or higher than the solar modulation cutoff, whereas high mass strangelets experience solar modulation already at rigidities above the geomagnetic cutoff.

For a non-magnetic body like the Moon, there is no corresponding cutoff, and the total flux is given by F_{mod}. This makes the lunar surface an interesting laboratory for a strangelet search (c.f. Section 5.2; see also Ref. [46]).

Given the significant uncertainty in input parameters a simple but physically transparent model for strangelet propagation was chosen in [43]. $\frac{dN}{dt}$ is given by the following sum of a source term from supernova acceleration, a diffusion term, loss terms due to escape from the Galaxy, energy loss, and spallation,

$$\frac{dN}{dt} = \frac{\partial N}{\partial t}\big|_{\mathrm{source}} + D\nabla^2 N + \frac{\partial N}{\partial t}\big|_{\mathrm{escape}} + \frac{\partial}{\partial E}[b(E)N] + \frac{\partial N}{\partial t}\big|_{\mathrm{spallation}}. \tag{4.5}$$

The individual terms will be defined and discussed in the following. Further terms describing, e.g. decay and reacceleration of strangelets due to passage of new supernova shock waves, are discussed in [43], but will be neglected here for simplicity.

The strangelet spectrum after acceleration in supernova shocks is assumed to be a standard powerlaw in rigidity with index -2.2 as derived from observations of ordinary cosmic rays. The minimal rigidity is assumed to be given by the speed of a typical supernova shock wave, $\beta_{SN} \approx 0.005$, so $R_{\min} = \gamma(\beta_{SN})\beta_{SN}Am_0c^2/Ze \approx 5\mathrm{MV}A/Z$. The maximal rigidity from acceleration in supernova shocks, $R_{\max}$, is of order 10^6 GV, but the actual number is irrelevant since $R_{\max} \gg R_{\min}$ and the rigidity spectrum steeply declining. For a total production rate of $\dot{M} = 10^{-10}M_\odot\mathrm{yr}^{-1}$ of baryon number A strangelets spread evenly in an effective galactic volume V, the total source term is

$$G(R) = 1.2\frac{\dot{M}}{VAm_0R_{\min}}\left(\frac{R}{R_{\min}}\right)^{-2.2} \tag{4.6}$$

or in terms of energy (using $dE/dR = Ze\beta$ and $G(R)dR = G(E)dE$)

$$\frac{\partial N}{\partial t}\big|_{\mathrm{source}} \equiv G(E) = \frac{G(R(E))}{Ze\beta}. \tag{4.7}$$

The terms $D\nabla^2 N + \frac{\partial N}{\partial t}\big|_{\mathrm{escape}}$, where D is the diffusion coefficient, describe cosmic ray diffusion in real space and eventual escape from the confining magnetic field of the Galaxy. Charged cosmic rays are spiralling along field lines in the weak galactic magnetic field, but due to the very irregular structure of the field, the particles scatter on magnetic "impurities", and the motion is best described in terms of diffusion. From studies of nuclei it is known, that cosmic rays are confined to move in a region significantly larger than the galactic disk, where most stars and interstellar matter is located. A typical value for the effective galactic volume confining cosmic rays is

$V = 1000 \text{ kpc}^3$. The standard leaky box approximation assumes $D = 0$ and $\frac{\partial N}{\partial t}|_{\text{escape}} = -N/\tau_{\text{escape}}$, where $\tau_{\text{escape}}(A, Z, E)$ is the average escape time from an otherwise homogeneous distribution in the galactic volume, V. From studies of cosmic ray nuclei the escape time is known as a function of rigidity, R, as

$$\tau_{\text{escape}} = \frac{8.09 \times 10^6 \text{y}}{n\beta} \left(\frac{R}{R_0} \right)^{\delta}. \tag{4.8}$$

where $R_0 = 4.7 \text{GV}$, $\delta = 0.8$ for $R < R_0$, and $\delta = -0.6$ for $R > R_0$. n denotes the average hydrogen number density per cubic centimeter ($n \approx 0.5$ when averaging over denser regions in the galactic plane and dilute regions in the magnetic halo).

The term in the propagation equation $\frac{\partial}{\partial E}[b(E)N]$, describes the influence of energy loss processes. The energy loss rate $b(E) \equiv -dE/dt$ can be treated as a sum of ionization losses (from interaction with neutral hydrogen atoms and molecules), Coulomb losses (from interaction with ionized hydrogen), and pion production losses from inelastic collisions at high relativistic γ-factor (threshold at $\gamma = 1.3$). The various contributions are described in [43]. At speeds close to the speed of light the ionization loss is simply proportional to nZ^2.

Like nuclei strangelets have a roughly geometrical cross section proportional to $A^{2/3}$ for spallation in collisions with interstellar matter (mainly hydrogen). The corresponding spallation time scale is taken to be

$$\tau_{\text{spallation}} = \frac{2 \times 10^7 \text{y}}{n\beta} A^{-2/3}. \tag{4.9}$$

At low kinetic energy the cross section can vary somewhat due to resonances etc, but such complications will be neglected here, since the detailed physics is unknown in the case of strangelets. We have also neglected the slight reduction in geometrical area of strangelets relative to nuclei due to their slightly larger density. The largest uncertainty in the treatment of spallation is the fact that strangelets (like nuclei) are not always completely destroyed in a spallation reaction. In addition to nucleons and nuclei smaller strangelets may result from this type of reaction, but we are ignorant of the physics to an extent where it is impossible to include this effective feed-down to lower A in a meaningful manner. Therefore, spallation is assumed to be a process destroying strangelets, i.e.

$$\frac{\partial N}{\partial t}|_{\text{spallation}} = -\frac{N}{\tau_{\text{spallation}}}. \tag{4.10}$$

This leads to an overall underestimate of the strangelet flux.

Strangelet energies are redistributed according to the propagation equation. Some leave the Galaxy or are destroyed by spallation. Occasionally strangelets get a new kick from a passing supernova shock, and in a first approximation they regain the source term relative distribution of rigidity. The time scale between supernova shock waves passing a given position in interstellar space is of order $\tau_{SN} \approx 10^7$ y. This scale is comparable to or larger than the time scales for energy loss, spallation, and escape from the Galaxy, so reacceleration of cosmic ray strangelets has only a moderate influence on the steady state distribution. By adding energy (on average) to the particles it actually reduces the total flux of strangelets somewhat because higher energies make destruction and escape more likely (see Ref. [43] for details).

Introducing the terms discussed above the steady state equation $dN/dt = 0$ leads to the following differential equation for $N(E)$

$$b(E)\frac{dN}{dE} = \frac{N(E)}{\tau(E,A,Z)} - G(E),$$ (4.11)

where

$$1/\tau(E,A,Z) = 1/\tau_{\text{escape}} + 1/\tau_{\text{spallation}} + 1/\tau_{\text{loss}},$$ (4.12)

with $1/\tau_{\text{loss}} \equiv -db(E)/dE$. When energy loss can be neglected, $b(E) \approx 0$ and $|\tau_{\text{loss}}| \approx \infty$. In this limit the spectrum is simply given by

$$N(E) \approx G(E)\tau(E,A,Z),$$

and $1/\tau(E,A,Z) \approx 1/\tau_{\text{escape}} + 1/\tau_{\text{spallation}}$.

The general solution of the propagation equation requires numerical integration, but several limits can be treated analytically and provide a physical understanding of the full numerical solutions found in [43]. The special cases (disregarding solar modulation and geomagnetic cutoff) can be divided according to the relative importance of the different time scales, τ_{escape}, $\tau_{\text{spallation}}$, and $|\tau_{\text{loss}}|$ (the energy loss time scale is negative at low and high energies, describing a net increase in number of particles).

When one of the time scales is significantly smaller than the others at a given energy (rigidity), the corresponding process dominates the physics. The relative importance of the processes depends on strangelet properties A, Z, on the density of interstellar hydrogen n (though most processes have the same n-dependence), and of course on the strangelet energy, E (or rigidity, R). Energy loss dominates at low energy, spallation at intermediate E, and escape from the Galaxy at the highest energies.

The energy loss domination is important only at energies below the solar modulation cutoff, so I refer the reader to [43] for details.

At intermediate energies the spectrum is determined by the strangelet spallation time, $\tau \approx \tau_{\text{spallation}}$, and the energy distribution is approximately given by

$$N(E) \approx G(E)\tau_{\text{spallation}}(n,\beta,A)$$ (4.13)

with $\tau_{\text{spallation}} = 2 \times 10^7 \text{yr} n^{-1} \beta^{-1} A^{-2/3}$.

This gives the approximate result

$$F_R(R) = 2.34 \times 10^5 \text{m}^{-2}\text{yr}^{-1}\text{sterad}^{-1} \text{ GV}^{-1} A^{-0.467} Z^{-1.2} R_{\text{GV}}^{-2.2} \Lambda,$$ (4.14)

or for the total flux above rigidity R

$$F(>R) = 1.95 \times 10^5 \text{m}^{-2}\text{yr}^{-1}\text{sterad}^{-1} A^{-0.467} Z^{-1.2} R_{\text{GV}}^{-1.2} \Lambda.$$ (4.15)

The results scale in proportion to

$$\Lambda = \left(\frac{\beta_{\text{SN}}}{0.005}\right)^{1.2} \left(\frac{0.5\text{cm}^{-3}}{n}\right) \left(\frac{\dot{M}}{10^{-10}M_\odot\text{yr}^{-1}}\right) \left(\frac{1000\text{kpc}^3}{V}\right) \left(\frac{930\text{MeV}}{m_0c^2}\right).$$ (4.16)

Notice that the differential strangelet spectrum keeps the source term slope, $G(R) \propto R^{-2.2}$.

The intermediate energy domain is replaced by the high energy domain when $\tau_{\text{escape}} \leq \tau_{\text{spallation}}$. Except for very low A this happens when $R > 1.0\text{GV}A^{1.11}$, or $E > 1.0\text{GeV}ZA^{1.11}\beta^{-1}$. At high energies the spectrum is determined by the confinement time of strangelets in the Galaxy, $\tau \approx \tau_{\text{escape}}$, and the propagation equation leads to

$$N(E) \approx G(E)\tau_{\text{escape}}(n, \beta, R). \tag{4.17}$$

For semirelativistic or relativistic strangelets with $\beta \approx 1$, $\tau_{\text{escape}} \propto R^{-0.6}$, so the spectrum is steepened from the source term $R^{-2.2}$ to $R^{-2.8}$.

This gives the approximate result

$$F_R(R) = 2.40 \times 10^5 \text{m}^{-2}\text{yr}^{-1}\text{sterad}^{-1} \text{ GV}^{-1}A^{0.2}Z^{-1.2}R_{\text{GV}}^{-2.8}\Lambda, \tag{4.18}$$

and for the total flux above rigidity R

$$F(>R) = 1.33 \times 10^5 \text{m}^{-2}\text{yr}^{-1}\text{sterad}^{-1}A^{0.2}Z^{-1.2}R_{\text{GV}}^{-1.8}\Lambda. \tag{4.19}$$

The astrophysical input parameters in the present calculations are uncertain at the order of magnitude level, so approximate relations for the total flux of strangelets hitting the Earth or Moon accurate within a factor of 2 (for fixed input parameters) are useful. As indicated above solar modulation effectively cuts off the strangelet flux at rigidities of order $R_{\text{SM}} \approx (A/Z)^{1/2}\Phi_{500}^{1/2}\text{GV}$, which is in the part of the spectrum where the strangelet flux is governed by spallation. The total flux hitting the Moon or Earth is therefore roughly given by

$$F_{\text{total}} \approx 2 \times 10^5 \text{m}^{-2}\text{yr}^{-1}\text{sterad}^{-1}A^{-0.467}Z^{-1.2}\max[R_{\text{SM}}, R_{\text{GC}}]^{-1.2}\Lambda, \tag{4.20}$$

depending on whether solar modulation or geomagnetic cutoff dominates. In the case of solar modulation domination (always relevant for the Moon, and relevant for AMS as long as $R_{\text{SM}} > R_{\text{GC}}$) one obtains

$$F_{\text{total}} \approx 2 \times 10^5 \text{m}^{-2}\text{yr}^{-1}\text{sterad}^{-1}A^{-1.067}Z^{-0.6}\Phi_{500}^{-0.6}\Lambda. \tag{4.21}$$

For strangelets obeying the CFL mass-charge relation $Z = 0.3A^{2/3}$ this becomes

$$F_{\text{total}} \approx 4 \times 10^5 \text{m}^{-2}\text{yr}^{-1} \text{ sterad}^{-1}A^{-1.467}\Phi_{500}^{-0.6}\Lambda \tag{4.22}$$

$$\approx 2.8 \times 10^4 \text{m}^{-2}\text{yr}^{-1}\text{sterad}^{-1}Z^{-2.2}\Phi_{500}^{-0.6}\Lambda, \tag{4.23}$$

which reproduces the numerical results of Ref. [43] to within 20% for $Z > 10$ and to within a factor of a few even for small Z, where the assumptions of nonrelativistic strangelets and spallation domination both are at the limit of being valid.

The strangelet flux for Λ of order unity is high enough to be of interest for various upcoming experimental searches, and at the same time small enough to agree with previous searches which have given upper limits or shown marginal evidence for signatures consistent with strangelets (see [11] for an overview). As stressed several times above, many parameters are uncertain at the order of magnitude level. The scaling with these parameters is indicated where relevant. In particular this is true for the overall normalization of the strangelet flux as expressed via the parameter Λ (Eq. (4.16)), whereas the relative behavior of the differential flux is less uncertain.

Apart from the uncertainty in parameters within the picture discussed here, one cannot rule out the possibility that some of the basic assumptions need to be modified. In addition to strangelet production in strange star collisions, it has been suggested that a (possibly small) flux of strangelets may be a direct outcome of the Type II supernova explosions [47], where strange stars form. The treatment above does not include such additional strangelet production mechanisms, and therefore the flux predictions are conservative. Another assumption that leads to a conservative lower limit on the flux is that spallation is assumed to destroy strangelets completely. At least for low-energy collisions one would expect that fragments of strangelets would survive, but the input physics for performing a realistic strangelet spallation study is not sufficiently well known, and therefore the conservative assumption of complete destruction was made. A numerical simulation of strangelet propagation in Ref. [48] assumed stripping of nucleons rather than complete destruction of strangelets in interstellar collisions and studied two specific sets of values for the mass and energy of strangelets injected into the interstellar medium. Several other assumptions made in that numerical study differ from those of the current investigation, so a direct comparison of the results is not possible, except that both studies are consistent with the possibility of a significant, measurable strangelet flux in our part of the Galaxy.

But ultimately the question of whether strangelets exist in cosmic rays is an experimental issue.

5. Experiments underway

Several experiments have searched for strangelets in cosmic rays and/or have had their data reanalyzed to provide limits on the strangelet abundance. While some interesting events have been found that are consistent with the predictions for strangelets, none of these have been claimed as real discoveries. Interpreted as flux limits rather than detections these results are consistent with the flux estimates given above. For discussions see [11].

If the interesting events were actual measurements, two experiments that are currently underway (with the author as a humble theoretician participating in collaborations with very capable experimentalists!) will reach sensitivities, that would provide real statistics. These experiments are the Alpha Magnetic Spectrometer (AMS-02) on the International Space Station, and the Lunar Soil Strangelet Search (LSSS) at the Wright Nuclear Structure Laboratory at Yale.

5.1 AMS-02

The Alpha Magnetic Spectrometer (AMS) is a high profile space-based particle physics experiment involving approximately 500 physicists from more than 50 institutions in 16 countries, led by Nobel laureate Samuel Ting of the Massachusetts Institute of Technology (MIT). The AMS flew in space in June of 1998 aboard the Space Shuttle Discovery [49], and it is currently scheduled to fly to the International Space Station (ISS) on Utilization Flight # 4.1 (UF4.1) with launch in 2008. Once on the ISS, the AMS-02 will remain active for at least three years before it is returned to Earth. AMS-02 will provide data with unprecedented accuracy on cosmic ray electrons, positrons, protons, nuclei, anti-nuclei and gammas in the GV-TV range in order to probe issues such as antimatter, dark matter, cosmic ray formation and propagation. And in addition it will be uniquely suited to discover strangelets characterized by extreme rigidities for a given velocity compared to nuclei [11]. AMS-02 will have excellent charge resolution up to $Z \approx 30$, and should therefore be

able to probe a large mass range for strangelets. A reanalysis of data from the 1998 AMS-01 mission has given hints of some interesting events, such as one with $Z = 2, A = 16$ [50], but with the larger AMS-02 detector running for 3 years or more, real statistics is achievable.

5.2 LSSS

The Lunar Soil Strangelet Search (LSSS) is a search for $Z = 8$ strangelets using the tandem accelerator at the Wright Nuclear Structure Laboratory at Yale. The experiment involves a dozen people from Yale, MIT, and Århus, led by Jack Sandweiss of Yale. The experiment which is currently in its calibration phase, studies a sample of 15 grams of lunar soil from Apollo 11 (sample number AH10084). The expectation is to reach a sensitivity of 10^{-17} over a wide mass range, sufficient to provide detection according to Equation (4.21) if strangelets have been trapped in the lunar surface layer, which has an effective cosmic ray exposure time of around 500 million years.

6. Moving the GZK-cutoff–Strangelets as ultra-high energy cosmic rays

One of the most interesting puzzles in cosmic ray physics is the apparent existence of cosmic rays with energies well beyond 10^{19}eV, with measured energies as high as 3×10^{20}eV. Briefly speaking there are two puzzles [51, 52]:

1. It is almost impossible to find an astrophysical mechanism capable of accelerating cosmic rays to these energies.

2. Even if acceleration happens, ultra-high energy cosmic rays have a relatively short mean-free path for interactions with the cosmological microwave background photons, and only cosmic rays from fairly nearby (unknown) sources would be able to reach us with the high energies measured.

Interestingly, strangelets circumvent both of these problems, and therefore provide a possible mechanism for cosmic rays beyond the socalled Greisen-Zatsepin-Kuzmin (GZK) cutoff. The benefits of strangelets with respect to the two problems are [53]:

1. All astrophysical cosmic ray "accelerators" involve electromagnetic fields, and the maximal energy that can be transmitted to a charged particle is proportional to the charge, Z. For instance some mechanisms can provide a maximal rigidity, R_{MAX}, proportional to the product of the magnetic field and size of the region where acceleration takes place, and for relativistic particles this corresponds to a maximal energy $E_{\text{MAX}} = Z R_{\text{MAX}}$. The charge of massive strangelets has no upper bound in contrast to nuclei, so highly charged strangelets are capable of reaching energies much higher than those of cosmic ray protons or nuclei using the same "accelerator" [53].

2. The GZK-cutoff is a consequence of ultrarelativistic cosmic ray projectiles hitting a 2.7K background photon with a Lorentz-factor γ large enough to boost the 7×10^{-4}eV photon to energies beyond the threshold of a process leading to significant projectile energy loss, such as photo-pion production, photo-pair production, or photo-disintegration. The threshold for

such a process has a fixed energy, $E_{\text{Threshold}}$, in the frame of the cosmic ray (e.g., $E_{\text{Threshold}} \approx$ 10MeV for photo-disintegration of a nucleus or a strangelet), corresponding to a Lorentz-factor $\gamma_{\text{Threshold}} = E_{\text{Threshold}}/E_{2.7\text{K}}$ ($\approx 10^{10}$ for photo-disintegration). But this corresponds to a cosmic ray total energy

$$E_{\text{Total}} = \gamma_{\text{Threshold}}Am_0c^2. \tag{6.1}$$

Thus, E_{Total} ($\approx 10^{19}A$ eV at the photo-disintegration threshold) is proportional to the rest mass of the cosmic ray, and since strangelets can have much higher A-values than nuclei, this pushes the GZK-cutoff energy to values well beyond the current observational limits for ultra-high energy cosmic rays [53, 54].

A testable prediction of the strangelet scenario for ultra-high energy cosmic rays is that strangelets at a given energy will be more isotropically distributed than protons and nuclei because the gyro-radius in the galactic magnetic field is proportional to Z^{-1}, so the arrival direction for high-Z strangelets points back to the source to a lesser extent than for low-Z candidates which have gyro-radii comparable to the size of the Galaxy.

7. Conclusion

The total strangelet flux reaching the Moon or a detector in Earth orbit is in a regime that could be within experimental reach, and therefore provide a crucial test of the hypothesis of absolutely stable strange quark matter. Experiments are underway which are sensitive to the high mass-to-charge signature expected for such events within the flux-range predicted. Strangelets may also provide an interesting explanation of ultra-high energy cosmic ray events.

Acknowledgments

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