Review of speculative "disaster scenarios" at RHIC

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This paper discusses speculative disaster scenarios inspired by hypothetical new fundamental processes that might occur in high-energy relativistic heavy-ion collisions. The authors estimate the parameters relevant to black-hole production and find that they are absurdly small. They show that other accelerator and (especially) cosmic-ray environments have already provided far more auspicious opportunities for transition to a new vacuum state, so that existing observations provide stringent bounds. The possibility of producing a dangerous strangelet is discussed in most detail. The authors argue that four separate requirements are necessary for this to occur: existence of large stable strangelets, metastability of intermediate size strangelets, negative charge for strangelets along the stability line, and production of intermediate size strangelets in the heavy ion environment. Both theoretical and experimental reasons why each of these appears unlikely are discussed. In particular, the authors know of no plausible suggestion for why the third or especially the fourth might be true. Given minimal physical assumptions, the continued existence of the Moon, in the form we know it, despite billions of years of cosmic-ray exposure, provides powerful empirical evidence against the possibility of dangerous strangelet production.

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I. INTRODUCTION

Fears have been expressed that heavy-ion collisions at the Relativistic Heavy Ion Collider (RHIC), which Brookhaven National Laboratory (BNL) is now commissioning, might initiate a catastrophic process with profound implications for health and safety. In this paper we explore the physical basis for speculative disaster scenarios at RHIC.[1]

Concerns have been raised in three general categories: first, formation of a black hole or gravitational singularity that accretes ordinary matter; second, initiation of a transition to a lower vacuum state; and third, formation of a stable "strangelet" that accretes ordinary matter. We have reviewed the scientific literature, evaluated recent correspondence, and undertaken additional calculations where necessary, to evaluate the scientific basis of these safety concerns.

Our conclusion is that the candidate mechanisms for catastrophe scenarios at RHIC are firmly excluded by compelling arguments based on well-established physical laws. In addition, where the data exist, a conservative analysis of existing empirical evidence excludes the possibility of a dangerous event at RHIC at a very high level of confidence. Accordingly, we see no reason to delay the commissioning of RHIC on account of these safety concerns.

Considerable attention has been focused on the possibility of placing a bound on the probability of a dangerous event at RHIC by making a "worst case" analysis of certain cosmic-ray data (Dar *et al.*, 1999). We believe it is reasonable to assume that the laws of physics will not suddenly break down in bizarre ways when entering a regime that actually differs only slightly and in apparently inessential ways from regimes already well explored. We will review the work that has been done on

[1]This paper is a revision and adaptation of a report commissioned by Dr. John Marburger, Director of BNL. Dr. Marburger orginally charged our committee to review the issues and "to reduce to a single comprehensive report the arguments that address the safety of each of the speculative 'disaster scenarios'."

empirical bounds and point out where and how the laws of physics must be bent in order to avoid very firm bounds on the probability of a dangerous event at RHIC. No limit is possible if one allows arbitrarily poor physics assumptions in pursuit of a worst case scenario.

Some of the expressed anxiety seems to be based on a misunderstanding of the nature of high-energy collisions: It is necessary to distinguish carefully between total energy and energy density. The total center-of-mass energy ($E_{\mathrm{c.m.}}$) of gold-gold collisions at RHIC will exceed that of any existing accelerator. But $E_{\mathrm{c.m.}}$ is surely not the right measure of the capacity of a collision to trigger exotic new phenomena. If it were, a batter striking a Major League fastball would be performing a far more dangerous experiment than any contemplated at a high-energy accelerator. To be effective in triggering exotic new phenomena, energy must be concentrated in a very small volume.

A better measure of effectiveness is the center-of-mass energy of the elementary constituents within the colliding objects. In the case of nuclei, the elementary constituents are mainly quarks and gluons, with small admixtures of virtual photons, electrons, and other elementary particles. Using the Fermilab Tevatron and the LEP collider at the European Center for Nuclear Research (CERN), collisions of these elementary particles with energies exceeding what will occur at RHIC have already been extensively studied.

What is truly novel about heavy-ion colliders compared to other accelerator environments is the volume over which high-energy densities can be achieved and the number of quarks involved. In a central gold-gold collision, hundreds of quarks collide at high energies. Black holes and vacuum instability are generic concerns that have been raised, and ought to be considered, each time a new facility opens up a new high-energy frontier. The fact that RHIC accelerates heavy ions rather than individual hadrons or leptons makes for somewhat different circumstances. Nevertheless there are simple, convincing arguments that neither poses any significant threat. The strangelet scenario is special to the heavy-ion environment. It could have been raised before the commissioning of the Alternating Gradient Synchrotron (AGS) or CERN heavy-ion programs. Indeed, we believe the probability of a dangerous event, though still immeasureably small, is greater at AGS or CERN energies than at RHIC. In light of its special role at RHIC, we pay most attention to the strangelet scenario.

In the remainder of this Introduction we give brief, nontechnical summaries of our principal conclusions regarding the three potential dangers. In the body of the paper which follows we consider each problem in as much detail as seems appropriate. First, in Sec. II we present a summary of cosmic-ray data necessary to make empirical estimates regarding vacuum decay and strangelets. Sections III, IV, and V are devoted to gravitational singularities, vacuum decay, and strangelets, respectively.

When we make quantitative estimates of possible dangerous events at RHIC, we will quote our results as a probability, $\mathfrak{p}$, of a single dangerous event over the lifetime of RHIC (assumed to encompass approximately 2×10^{11} gold-gold collisions over a ten-year lifetime at full luminosity). We do not attempt to decide what is an acceptable upper limit on $\mathfrak{p}$, nor do we attempt a ''risk analysis,'' weighing the probability of an adverse event against the severity of its consequences. Ultimately, we rely on compelling physics arguments which, we believe, exclude a dangerous event beyond any reasonable level of concern.[2]

A. Gravitational singularities

Exotic gravitational effects may occur at immense densities. Conservative dimensionless measures of the strength of gravity give 10^{-22} for classical effects and 10^{-34} for quantum effects in the RHIC environment, in units where 1 represents gravitational effects as strong as the nuclear force. The theoretical basis for these estimates is presented in Sec. III. In fact RHIC collisions are expected to be less effective at raising the density of nuclear matter than collisions at lower energies where the ''stopping power'' is greater and existing accelerators have already probed larger effective energies. In no case has any phenomenon suggestive of gravitational clumping, let alone gravitational collapse or the production of a singularity, been observed.

B. Vacuum instability

Physicists have grown quite accustomed to the idea that empty space—what we ordinarily call ''vacuum''—is in reality a highly structured medium, that can exist in various states or phases, roughly analogous to the liquid or solid phases of water. This idea plays an important role in the standard model. Although certainly nothing in our existing knowledge of the laws of Nature demands it, several physicists have speculated on the possibility that our contemporary vacuum is only metastable, and that a sufficiently violent disturbance might trigger its decay into something quite different (Kobzarev, Okun, and Voloshin, 1974; Callan and Coleman, 1977; Coleman, 1977; Frampton, 1977). A transition of this kind would propagate outward from its source throughout the universe at the speed of light, and would be catastrophic.

We know that our world is already in the correct (stable) vacuum for quantum chromodynamics. Our knowledge of fundamental interactions at higher energies, and in particular of the interactions responsible for electroweak symmetry breaking, is much less complete. While theory strongly suggests that any possibility for triggering vacuum instability requires substantially larger energy densities than RHIC will provide, it is difficult to give a compelling, unequivocal bound based on theoretical considerations alone.

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[2]We thank A. Kent for correspondence on the subject of risk analysis.

Fortunately in this case we do not have to rely solely on theory; there is ample empirical evidence based on cosmic-ray data. Cosmic rays have been colliding throughout the history of the universe, and if such a transition were possible it would have been triggered long ago. Motivated by the RHIC proposal, in 1983 Hut and Rees (1984) calculated the total number of collisions of various types that have occurred in our past light cone—whose effects we would have experienced. Even though cosmic-ray collisions of heavy ions at RHIC energies are relatively rare, Hut and Rees found approximately 10^{47} comparable collisions have occurred in our past light cone. Experimenters expect about 2×10^{11} heavy-ion collisions in the lifetime of RHIC. Thus on empirical grounds alone, the probability of a vacuum transition at RHIC is bounded by 2×10^{-36}. We can rest assured that RHIC will not drive a transition from our vacuum to another. We review and update the arguments of Hut and Rees in Sec. IV after introducing the necessary cosmic-ray data in Sec. II.

C. Strangelets

Theorists have speculated that a form of quark matter, known as ''strange matter'' because it contains many strange quarks, might be more stable than ordinary nuclei. Hypothetical small lumps of strange matter, having atomic masses comparable to ordinary nuclei have been dubbed ''strangelets.'' Strange matter may exist in the cores of neutron stars, where it is stabilized by intense pressure.

For strange matter to pose a hazard at a heavy-ion collider, four conditions would have to be met:

- Strange matter would have to be absolutely stable in bulk at zero external pressure. If strange matter is not stable, it will not form spontaneously.
- Strangelets would have to be at least metastable for very small atomic mass, for only very small strangelets can conceivably be created in heavy-ion collisions.
- It must be possible to produce such a small, metastable strangelet in a heavy-ion collision.
- The stable composition of a strangelet must be *negatively* charged. Positively charged strangelets pose no threat whatsoever.

Each of these conditions is considered unlikely by experts in the field, for the following reasons:

- At present, despite vigorous searches, there is no evidence whatsoever for stable strange matter anywhere in the Universe.
- On rather general grounds, theory suggests that strange matter becomes unstable in small lumps due to surface effects. Strangelets small enough to be produced in heavy-ion collisions are not expected to be stable enough to be dangerous.
- It is overwhelmingly likely that the most stable configuration of strange matter has positive electric charge.
- Theory suggests that heavy-ion collisions (and hadron-hadron collisions in general) are a poor way to produce strangelets. Furthermore, it suggests that the production probability is lower at RHIC than at lower energy heavy-ion facilities like the AGS and CERN. Models and data from lower energy heavy-ion colliders indicate that the probability of producing a strangelet decreases very rapidly with the strangelet's atomic mass.
- A negatively charged strangelet with a given baryon number is much more difficult to produce than a positively charged strangelet with the same baryon number because it must contain proportionately more strange quarks.

To our knowledge, possible catastrophic consequences of strangelet formation have not been studied in detail before.[3] Although the underlying theory (quantum chromodynamics, or QCD) is fully established, our ability to use it to predict complex phenomena is imperfect. A reasonable, conservative attitude is that theoretical arguments based on QCD can be trusted when they suggest a safety margin of many orders of magnitude. The hypothetical chain of events that might lead to a catastrophe at RHIC requires several independent, robust theoretical arguments to be wrong simultaneously. Thus theoretical considerations alone would allow us to exclude any safety problem at RHIC confidently.

However, one need not use theoretical arguments alone. We have considered the implications of natural ''experiments'' elsewhere in the Universe, where cosmic-ray induced heavy-ion collisions have been occurring for a long time. Recent satellite based experiments have given us very good information about the abundance of heavy elements in cosmic rays, making it possible to obtain a reliable estimate of the rate of such collisions. We know of two domains where empirical evidence tells us that cosmic-ray collisions have not produced strangelets with disasterous consequences: first, the surface of the Moon, which has been impacted by cosmic rays for billions of years, and second, interstellar space, where the products of cosmic-ray collisions are swept up into the clouds from which new stars are formed. In each case the effects of a long-lived, dangerous strangelet would be obvious, so dangerous strangelet production can be bounded below some limit. For example, we know for certain that iron nuclei with energy in excess of 10 GeV/nucleon (equivalent to AGS energies) collide with iron nuclei on the surface of the Moon approximately 6×10^{10} times per second. Over the five-billion year life of the Moon approximately 10^{28} such collisions have occurred. None has produced a dangerous strangelet which came to rest on the lunar surface, for if it had, the Moon would have been converted to strange matter. Similarly, we know that the vast number of heavy-ion collisions in interstellar space have not

[3]The paper by Dar, DeRujula, and Heinz appeared after the completion of the bulk of our work and addresses only a subset of the basic issues.

created a dangerous strangelet that lived long enough to be swept up into a star (Dar *et al.*, 1999). A dangerous strangelet would trigger the conversion of its host star into strange matter, an event that would resemble a supernova. The present rate of supernovae—a few per millennium per galaxy—translate into a strong upper limit on the probability of long-lived dangerous strangelet production at RHIC.

To translate each of these results into a bound on $\mathfrak{p}$, it is necessary to model some aspects of strangelet production, propagation, and decay. By making sufficiently unlikely assumptions about the properties of strangelets, it is possible to render both of these empirical bounds irrelevant to RHIC. Dar *et al.* (1999) construct just such a model in order to discard the lunar limits: They assume that strangelets are produced only in gold-gold collisions, only at or above RHIC energies, and only at rest in the center of mass. We are skeptical of all these assumptions. If they are accepted, however, lunar persistence provides no useful limit. Others, in turn, have pointed out that the astrophysical limits of Dar *et al.*, 1999 can be avoided if the dangerous strangelet is metastable and decays by baryon emission with a lifetime longer than $\sim 10^{-7}$ s. In this case strangelets produced in the interstellar medium decay away before they can trigger the death of stars, but a negatively charged strangelet produced at RHIC could live long enough to cause catastrophic results. Under these conditions the Dar, De Rujula, and Heinz bound evaporates.

We wish to stress once again that we do not consider these empirical analyses central to the argument for safety at RHIC. The arguments which are invoked to destroy the empirical bounds from cosmic rays, if valid, would not make dangerous strangelet production at RHIC more likely. Even if the bounds from lunar and astrophysical arguments are set aside, we believe that basic physics considerations rule out the possibility of dangerous strangelet production at RHIC.

II. HEAVY NUCLEI IN COSMIC RAYS

Cosmic-ray processes accurately reproduce the conditions planned for RHIC. Cosmic rays are known to include heavy nuclei and to reach extremely high energies. Hut and Rees (1984) pioneered the use of cosmic-ray data in their study of decay of a false vacuum. Dar, De Rujula, and Heinz (1999) have recently used similar arguments to study strangelet production in heavy-ion collisions. Here we summarize data on heavy nuclei (iron and beyond) in cosmic rays and carry out some simple estimates of particular processes which will figure in our discussion of strange matter. In some instances we use observations directly; elsewhere reasonable extrapolation allows us to model behavior where no empirical data are available.

We are interested in cosmic-ray collisions which simulate RHIC and lower energy heavy-ion facilities like the AGS. Equivalent stationary target energies range from 10 GeV/nucleon at the AGS to 20 TeV/nucleon corresponding to the center-of-mass energy of 100 GeV/nucleon at RHIC. The flux of cosmic rays has been measured accurately up to total energies of order 10^{20} eV.[4] Many measurements of the abundance of ultraheavy nuclei in cosmic rays at GeV/nucleon energies are summarized in Binns (1988). These measurements are dominated by energies near the lower energy cutoff of 1.5 GeV/nucleon. More extensive measurements have been made of the flux of nuclei in the iron-nickel ($Z = 26$–28) group and lighter. Data on iron are available up to energies of order 2 TeV/nucleon (Swordy *et al.*, 1993). However, we know of no direct measurements of the flux of nuclei heavier than the iron-nickel group at energies above 10 GeV/nucleon.

Thus data on iron are available over almost the entire energy range we need. For nuclei heavier than iron, data are available close to AGS energies, but not in the 100 GeV/nucleon–20 TeV/nucleon domain. For ultraheavy nuclei at very high energies, we extrapolate existing data to higher energies using two standard scaling laws, which agree excellently with available data.

- At energies of interest to us, the flux of every species which has been measured shows a simple power-law spectrum $dF/dE \propto E^{-\gamma}$ with $\gamma \approx 2.5$–2.7. Swordy *et al.* (1993) found this behavior for oxygen, magnesium, and silicon as well as hydrogen, helium, and iron. The same power law is observed at high energies where data are dominated by hydrogen (Wiebel-Sooth and Biermann, 1998).[5]
- At all energies where they have been measured, the relative abundance of nuclear species in cosmic rays reflects their abundance in our solar system. (See, for example, Fig. 6 in Binns, 1988.) Exceptions to this rule seem to be less than an order of magnitude. If anything, heavy nuclei are expected to be relatively more abundant in high-energy cosmic rays.

In light of these facts we adopt the standard idealization that the A (baryon number or atomic mass) and E (energy per nucleon) dependence of the flux of primary cosmic rays factors at GeV/nucleon–TeV/nucleon energies:

$$\frac{dF}{dE} = \Gamma(A,E_0)(E_0/E)^{\gamma}, \tag{1}$$

where E_0 is some reference energy. To be conservative we will usually take $\gamma = 2.7$. The total flux at energies above some energy E is given by

$$F(A,E) = \int_E^{\infty} dE' \frac{dF}{dE'} = \frac{E}{\gamma - 1}\frac{dF}{dE} = \frac{E}{\gamma - 1}\Gamma(A,E). \tag{2}$$

The units of dF/dE are $\{\mathrm{sr,s,m^2,GeV}\}^{-1}$. The flux of cosmic rays is very large in these units. For example, for iron at 10 GeV/nucleon, according to Swordy *et al.* (1993)

[4]For a review of cosmic rays and references to the original literature, see Wiebel-Sooth and Biermann (1998).

[5]At energies above 10^{15} eV the power γ changes abruptly. This occurs above the energies of interest to us.

$$\frac{dF}{dE}(\mathrm{Fe},\ 10\ \mathrm{GeV}) \equiv \Gamma(\mathrm{Fe},\ 10\ \mathrm{GeV}) \approx 4\times 10^{-3}\{\mathrm{sr\,s\,m^2\ GeV}\}^{-1}. \quad (3)$$

Combining all nuclei with $Z>70$ into our definition of ''gold,'' we find an abundance of $\sim 10^{-5}$ relative to iron.[6]

We are interested in cosmic-ray initiated heavy-ion collisions which have occurred where we can observe their consequences. Three particular examples will figure in our subsequent considerations: (a) Cosmic-ray collisions with nuclei on the surface of planetoids that lack an atmosphere, like the Moon; (b) Cosmic-ray collisions in interstellar space resulting in strangelet production at rest with respect to the galaxy; (c) The integrated number of cosmic-ray collisions in our past light cone.

A. Cosmic-ray impacts on the Moon

First we consider cosmic rays impinging on the surface of a planetoid similar to the Moon. The number of impacts per second with energy greater than E on the surface of the planet is given by $8\pi^2 R^2 F(A,E)$, where we measure R in units of R_{Moon},

$$\frac{dN(A,E)}{dt} = 2\times 10^{14}\frac{\Gamma(A,E)}{\gamma-1}E\left(\frac{R}{R_{\mathrm{Moon}}}\right)^2. \quad (4)$$

For convenience, we use iron with $E=10$ GeV/nucleon as our reference. From Eqs. (2)–(4) we find

$$\frac{dN(A,E)}{dt} \approx 5\times 10^{12}\frac{\Gamma(A,\ 10\ \mathrm{GeV})}{\Gamma(\mathrm{Fe},\ 10\ \mathrm{GeV})}\left(\frac{10\ \mathrm{GeV}}{E}\right)^{1.7}\left(\frac{R}{R_{\mathrm{Moon}}}\right)^2. \quad (5)$$

This large instantaneous rate makes it possible to obtain useful limits from cosmic-ray collisions with nuclei on the lunar surface.

B. Cosmic-ray collisions in space

Following Dar *et al.* (1999), we consider collisions of cosmic rays in which the center-of-mass velocity is less than $v_{\mathrm{crit}}=0.1$ in units of c. With this v_{crit} strangelets produced at rest in the center of mass will have high probability of slowing down without undergoing nuclear collisions which would destroy them. The flux given in Eq. (1) is associated with a density, $dn/dE = (4\pi/c)(dF/dE)$. The rate per unit volume for collisions of cosmic rays with energy per nucleon greater than E in which all components of the center-of-mass velocity are less than v_{crit} is given by

$$R(E) = 2c\sigma f_\theta \int_E^\infty dE_1 \int_{(1-v_{\mathrm{crit}})E_1}^{(1+v_{\mathrm{crit}})E_1} dE_2 \frac{dn}{dE_1}\frac{dn}{dE_2}, \quad (6)$$

where $\sigma = 0.18A^{2/3}$ b is the geometric cross section, and $f_\theta = 4v_{\mathrm{crit}}^2$ is a geometric factor measuring the fraction of collisions in which the transverse velocity is less than v_{crit}. Substituting from Eq. (1), and normalizing to iron-iron collisions at $E=10$ GeV/nucleon, we obtain

$$R(E,A) = 10^{-45}\left(\frac{10\ \mathrm{GeV}}{E}\right)^{3.4}\left(\frac{\Gamma(A)}{\Gamma(\mathrm{Fe})}\right)^2 \times \left(\frac{A}{56}\right)^{2/3}\ \mathrm{cm^{-3}\,s^{-1}}. \quad (7)$$

Although this rate appears very small, these collisions have been occurring over very large volumes for billions of years.

C. Cosmic-ray collisions in our past light cone

Finally we update the calculation of Hut and Rees of the total number of high-energy collisions of cosmic rays in our past light cone. The number of such collisions for cosmic rays with energy greater than E is given by

$$N \sim 10^{47}\left(\frac{\Gamma(\mathrm{A})}{\Gamma(\mathrm{Fe})}\right)^2\left(\frac{56}{A}\right)^{2.7}\left(\frac{100\ \mathrm{GeV}}{E}\right)^{3.4}, \quad (8)$$

where we have normalized to iron at $E=100$ GeV/nucleon. The difference between the extremely small coefficient in Eq. (7) and the extremely large coefficient in Eq. (8) reflects integration over our past light cone, i.e., over the volume and age of the universe ($VT \approx C^3T^4 \approx 3\times 10^{101}$ s cm^3, where $T \approx 10^{10}$ y is the Hubble time).

III. STRENGTH OF GRAVITATIONAL EFFECTS

Two possible sources of novel gravitational effects might in principle be activated in collisions at RHIC. The first type is connected with classical gravity, the second type with quantum gravity.

To estimate the quantitative significance of classical gravity, an appropriate parameter is

$$k_{\mathrm{cl}} \equiv \frac{2GM}{Rc^2} \quad (9)$$

for a spherical concentration of mass M inside a region of linear dimension R, where G is Newton's constant and c is the speed of light. It is when $k_{\mathrm{cl}}\to 1$ that the escape velocity from the surface at R, calculated in Newtonian gravity, becomes equal to the speed of light. The same parameter, $2GM/c^2$, appears in the general relativistic line element

$$ds^2 = c^2dt^2\left(1-\frac{2GM}{rc^2}\right) - \frac{dr^2}{1-(2GM/rc^2)} - r^2d^2\Omega \quad (10)$$

[6]Estimates range from 10^{-5} (Binns, 1988) to as high as 10^{-4} (Swordy, 1999). To be conservative, we chose a value on the low side.

outside a spherical concentration of mass M. In this language, it is when $k_{\rm cl}=1$ that a horizon appears at R, and the body is described as a black hole.

Now for RHIC we obtain a very conservative upper bound on $k_{\rm cl}$ by supposing that all the initial energy of the collision becomes concentrated in a region characterized by the Lorentz-contracted nuclei with a Lorentz contraction factor of 10^{-2}. We are being extremely conservative by choosing the largest possible mass and the smallest possible distance scale defined by the collision, and also by ignoring the effect of the electric charge and the momentum of the constituents, which will resist any tendency to gravitational collapse. Thus our result will provide a bound upon, not an estimate of, the parameters that might be required to have a realistic shot at producing black holes.

With $M=10^4\,{\rm GeV}/c^2$ and $R=10^{-2}\times10^{-13}\,{\rm cm}$, we arrive at $k_{\rm cl}=10^{-22}$. The outlandishly small value of this overgenerous estimate makes it pointless to attempt refinements.

To estimate the quantitative significance of quantum gravity, we consider the probability to emit the quantum of gravity, a graviton. It is governed by

$$k_{\rm qu}\equiv\frac{GE^2}{\hbar c^5}, \tag{11}$$

where $\hbar$ is Planck's constant and E is the total center-of-mass energy of collision. For collisions between elementary particles at RHIC, we should put $E\approx200\,{\rm GeV}$. This yields $k_{\rm qu}\approx10^{-34}$. Once again, the tiny value of $k_{\rm qu}$ makes it pointless to attempt refinements of this rough estimate. Of course higher energy accelerators than RHIC achieve larger values of $k_{\rm qu}$, but for the foreseeable future values even remotely approaching unity are a pipe dream.

IV. DECAY OF THE FALSE VACUUM

Hut and Rees (1984) first examined the question of vacuum stability in the context of cosmic-ray collisions in 1983. Kobzarev, Okun, and Voloshin (1974) had shown that the transition to the true vacuum, once initiated, would propagate outward at the speed of light. Thus our existence is evidence that no such transition occurred in our past light cone at least since the time of decoupling. Hut and Rees then estimated the total number of cosmic-ray collisions in the RHIC energy regime which have occurred in our past light cone. They used data on cosmic-ray fluxes that have subsequently been confirmed and updated. Not knowing which would be more effective at triggering a transition, Hut and Rees looked at both proton-proton collisions and collisions of heavy nuclei. Cosmic-ray data on proton fluxes go up to energies of order $10^{20}\,{\rm eV}$ (Wiebel-Sooth and Biermann, 1998). They concluded that proton-proton collisions with a center-of-mass energy exceeding 10^8 TeV have occurred so frequently in our past light cone that even such astonishingly high-energy collisions can be considered safe.

For heavy ions, Hut and Rees derived an estimate of the number of cosmic-ray collisions in our past light cone. We have updated their result in Eq. (7), and normalized it so that the coefficient 10^{47} equals the number of iron-iron collisions at a center-of-mass energy exceeding 100 GeV/nucleon. The abundance of iron in cosmic rays has now been measured up to energies of order 2 TeV/nucleon (Swordy *et al.*, 1993) and agrees with the estimate used by Hut and Rees. This result translates into a bound of 2×10^{-36} on $\mathfrak{p}$, the probability that (in this case) an iron-iron collision at RHIC energies would trigger a transition to a different vacuum state. While we do not have direct measurements of the fractional abundance of elements heavier than iron in cosmic rays of energy of order 100 GeV/nucleon, we do have good measurements at lower energies, where they track quite well with the abundances measured on earth and in the solar system. For ''gold'' (defined as $Z>70$) at lower energies $\Gamma({\rm Au})/\Gamma({\rm Fe})\approx10^{-5}$, leading to a bound, $\mathfrak{p}<2\times10^{-26}$ on the probability that a gold-gold collision at RHIC would lead to a vacuum transition. Even if this estimate were off by many orders of magnitude, we would still rest assured that RHIC will not drive a transition from our vacuum to another.

Since the situation has not changed significantly since the work of Hut and Rees, we do not treat this scenario in more detail here. The interested reader should consult Hut's 1984 paper (Hut and Rees, 1984) for further details.

V. STRANGELETS AND STRANGE MATTER

The scientific issues surrounding the possible creation of a negatively charged, stable strangelet are complicated. Also, it appears that if such an object did exist and could be produced at RHIC, it might indeed be dangerous. Therefore we wish to give this scenario careful consideration.

This section is organized as follows. First we give a pedagogical introduction to the properties of strangelets and strange matter. Second we discuss the mechanisms that have been proposed for producing a strangelet in heavy-ion collisions. We examine these mechanisms and conclude that strangelet production at RHIC is extremely unlikely. Nevertheless, we go on to discuss what might occur if a stable, negatively charged strangelet could be produced at RHIC. In light of the possible consequences of production of a stable negatively charged strangelet, we shall refer to such an object as a ''dangerous'' strangelet.

We then turn to the cosmic-ray data. We obtain strong bounds on the dangerous strangelet production probability at RHIC from physically reasonable assumptions. We also describe the ways in which these bounds can be evaded by adopting a sequence of specially crafted assumptions about the behavior of strangelets, which we consider physically unmotivated. It is important to remember, however, that evading the bounds does not make dangerous strangelet production more likely.

A. A primer on strangelets and strange matter

Strange matter is the name given to quark matter at zero temperature in equilibrium with respect to the weak interactions. At and below ordinary nuclear densities, and at low temperatures, quarks are confined to the interiors of the hadrons they compose.

It is thought that any collection of nucleons or nuclei brought to high enough temperature or pressure,[7] will make a transition to a state where the quarks are no longer confined into individual hadrons. At high temperature the material is thought to become what is called a quark-gluon plasma. The defining property of this state is that it can be accurately described as a gas of nearly freely moving quarks and gluons. One main goal of RHIC is to provide experimental evidence for the existence of this state, and to study its properties. At high pressure and low temperature the material is expected to exhibit quite different physical properties. In this regime, it is called quark matter. Quarks obey the Pauli exclusion principle—no two quarks can occupy the same state. As quark matter is compressed, the exclusion principle forces quarks into higher and higher energy states.

Given enough time (see below), the weak interactions will come into play, to reduce this energy. Ordinary matter is made of up (u) and down (d) quarks, which are the lightest species (or "flavors") of quarks. The strange quark (s) is somewhat heavier. Under ordinary conditions when an s quark is created, it decays into u and d quarks by means of the weak interactions. In quark matter the opposite can occur. u and d quarks, forced to occupy very energetic states, will convert into s quarks. Examples of weak interaction processes that can accomplish this are strangeness changing weak scattering, $u+d\rightarrow s+u$, and weak semileptonic decay, $u\rightarrow s+e+\overline{\nu}_e$. These reactions occur rapidly on a natural time scale $\sim 10^{-14}$ s. When the weak interactions finish optimizing the flavor composition of quark matter, there will be a finite density of strange quarks—hence the name "strange matter."

The most likely location for the formation of strange matter is deep within neutron stars, where the mammoth pressures generated by the overlayers of neutrons may be sufficient to drive the core into a quark matter state. When first formed, the quark matter at the core of a neutron star would be nonstrange, since it was formed from neutrons. Once formed, however, the quark matter core would rapidly equilibrate into strange matter, if such matter has lower free energy at high external pressure.

Initially, the nonstrange quark matter core and the overlaying layer of neutrons were in equilibrium. Since the strange matter core has lower free energy than the overlaying neutrons, its formation disrupts the equilibrium. Neutrons at the interface are absorbed into the strange matter core, which grows, eating its way outward toward the surface. There are two possibilities. If strange matter has lower internal energy than nuclear matter even at zero external pressure, the strange matter will eat its way out essentially to the surface of the star. On the other hand, if below some nonzero pressure, strange matter no longer has lower energy than nuclear matter, the conversion will stop. Even in the second case a significant fraction of the star could be converted to strange matter. The "burning" of a neutron star as it converts to strange matter has been studied in detail (Alcock *et al.*, 1991).[8] It is not thought to disrupt the star explosively, because the free-energy difference between strange matter and nuclear matter is small compared to the gravitational binding energy.

In 1984, E. Witten suggested that perhaps strange matter has lower mass than nuclear matter even at zero external pressure (Witten, 1984). Remarkably, the stability of ordinary nuclei does not rule this out. A small lump of strange matter, a "strangelet," could conceivably have lower energy than a nucleus with the same number of quarks. Despite the possible energy gain, the nucleus could not readily decay into the strangelet, because it would require many weak interactions to occur simultaneously, in order to create all the requisite strange quarks at the same time. Indeed, we know that changing one quark (or a few) in a nucleus into an s quark(s)—making a so-called hypernucleus—will raise rather than lower the energy.

Witten's paper sparked a great deal of interest in the physics and astrophysics of strange quark matter. Astrophysicists have examined neutron stars both theoretically and observationally, looking for signs of quark matter. Much interest centers around the fact that a strange matter star could be considerably smaller than a neutron star, since it is bound principally by the strong interactions, not gravity. A small quark star could have a shorter rotation period than a neutron star and be seen as a submillisecond pulsar. At this time there is no evidence for such objects and no other astrophysical evidence for stable strange matter, although astrophysicists continue to search and speculate.[8]

Strange matter is governed by QCD. At extremely high densities the forces between quarks become weaker (a manifestation of asymptotic freedom) and one can perform quantitatively reliable calculations with known techniques. The density of strange matter at zero external pressure is not high enough to justify the use of these techniques. Nevertheless the success of the ordinary quark model of hadrons leads us to anticipate that simple models which include both confinement and perturbative QCD provide us good qualitative guidance as to the properties of strange matter (Farhi and Jaffe, 1984).

Such rough calculations cannot answer the delicate question of whether or not strange matter is bound at zero external pressure reliably. Stability seems unlikely, but not impossible.

[7]For theoretical purposes a better variable is chemical potential, instead of pressure. But either can be used.

[8]For a review and extensive references, see Madsen (1998).

Some important qualitative aspects of strange matter dynamics that figure in the subsequent analysis are as follows.

a. Binding systematics[9]

The overall energy scale of strange matter is determined by the confinement scale in QCD which can be parametrized by the "bag constant." Gluon exchange interactions between quarks provide important corrections. Calculations indicate that gluon interactions in quark matter are, on average, repulsive, and tend to destabilize it. To obtain stable strange matter it is necessary to reduce the value of the bag constant below traditionally favored values (Farhi and Jaffe, 1984; Madsen, 1998). This is the reason we describe stability at zero external pressure as "unlikely."

b. Charge and flavor composition[10]

If strange matter contained equal numbers of u, d, and s quarks it would be electrically neutral. Since s quarks are heavier than u and d quarks, Fermi gas kinematics (ignoring interactions) would dictate that they are suppressed, giving strange matter a positive charge per unit baryon number, $Z/A>0$.

If this kinematic suppression were the only consequence of the strange quark mass, strange matter and strangelets would certainly have positive electric charge. In a bulk sample of quark matter this positive quark charge would be shielded by a Fermi gas of electrons electrostatically bound to the strange matter, as we discuss further below. Energy due to the exchange of gluons complicates matters. As previously mentioned, perturbation theory suggests this energy is repulsive, and tends to unbind quark matter. However, gluon interactions weaken as quark masses are increased, so the gluonic repulsion is smaller between s-s, s-u, or s-d pairs than between u and d quarks. As a result, the population of s quarks in strange matter is higher than expected on the basis of the exclusion principle alone. If, in a model calculation, the strength of gluon interactions is increased, there comes a point where strange quarks dominate. Then the electric charge on strange matter becomes negative.

Increasing the strength of gluon interactions pushes the charge of quark matter negative. However, it also unbinds it. Unreasonably low values of the bag constant are necessary to compensate for the large repulsive gluonic interaction energy.[11] For this reason we consider a negative charge on strange matter to be extremely unlikely.

[9]Farhi and Jaffe, 1984.

[10]Farhi and Jaffe, 1984.

[11]Some early studies that suggested negatively charged strange matter for broad ranges of parameters were based on incorrect applications of perturbative QCD.

c. Finite-size effects[12]

If it were stable, strange matter would provide a rich new kind of "strange" nuclear physics (De Rujula and Glashow, 1984; Farhi and Jaffe, 1984; Berger and Jaffe, 1987). Unlike nuclei, strangelets would not undergo fission when their baryon number grows large. Nuclear fission is driven by the mismatch between the exclusion principle's preference for equal numbers of protons and neutrons and electrostatics' preference for zero charge. In strange matter there is little mismatch: $u\approx d\approx s$ coincides with approximately zero charge.

On the other hand strangelets, like nuclei, become less stable at low baryon number. Iron is the most stable nucleus. Lighter nuclei are made less stable by surface effects. Surface energy is a robust characteristic of degenerate fermion systems. Estimates suggest that strange matter, too, has a significant surface energy, which would destabilize small strangelets (Farhi and Jaffe, 1984; Berger and Jaffe, 1987; Madsen, 1994). The surface tension which makes light nuclei and water droplets roughly spherical is a well-known manifestation of positive surface energy. The exact value of A below which strangelets would not be stable is impossible to pin down precisely, but small values of A (e.g., less than 10–30) are not favored.

Some very small nuclei are very stable. The classic example is ^{4}He. The reasons for helium's stability are very well understood. A similar phenomenon almost certainly does not occur for strangelets. The pattern of masses for strangelets made of 18 or fewer quarks can be estimated rather reliably (Farhi and Jaffe, 1984). Gluon interactions are, on average, destabilizing. They are most attractive for six quarks, where they still fail to produce a stable strange hadron. The most bound object is probably the H, composed of $uuddss$ (Jaffe, 1977). It is unclear whether this system is stable enough to be detected. On empirical grounds, it is certainly not lighter than the nonstrange nucleus made of six quarks—the deuteron. For $2<A\leqslant 6$, QCD strongly suggests complete instability of any strangelets. Larger strangelets, with baryon numbers up to of order 100, have been modelled by filling modes in a bag (Greiner *et al.*, 1988; Gilson and Jaffe, 1993; Madsen, 1994). These admittedly crude studies indicate the possible existence of metastable states, but none are sufficiently long lived to play a role in catastrophic scenarios at a heavy-ion collider. Thus, even if it were stable in bulk, strange matter would be unlikely to be stable in small aggregates.

d. Strangelet radioactivity and metastability[13]

If strange matter is stable in bulk and finite-size effects destabilize small strangelets, then there will likely be a range of A over which strangelets are metastable and decay by various radioactive processes. The lighter a

[12]Farhi and Jaffe, 1984; Berger and Jaffe, 1987; Madsen, 1994; Gilson and Jaffe, 1993.

[13]Farhi and Jaffe, 1984; Berger and Jaffe, 1987.

strangelet, the more unstable and shorter lived it would be. Two qualitatively different kinds of radioactivity concern us: baryon emission and lepton or photon emission.

- *Baryon emission*. It might be energetically favorable for a small strangelet to emit baryons (neutrons, protons, or α particles, in particular), and reduce its baryon number. Such decays are likely to be very rapid. Strong baryon emission would have a typical strong interaction lifetime of order 10^{-23} s. α decay, which can be very slow for nuclei, would be very rapid for a negatively charged strangelet on account of the absence of a Coulomb barrier. Weak baryon emission would be important for some light strangelets that must adjust their strangeness in order to decay. The lifetime for weak baryon emission can be approximated by

$$\tau^{-1} \approx \frac{Q}{4\pi} \sin^2 \theta_c G_F^2 \mu^4, \tag{12}$$

where G_F is Fermi's constant ($G_F = 10^{-5} M_p^{-2}$), $\sin\theta_c$ is Cabibbo's angle, Q is the Q value of the decay, and μ is the quark chemical potential in strange matter. Reasonable choices for these parameters put τ below 10^{-8} s. Baryon emission leaves a small strangelet smaller still, and less stable. Strangelets unstable against baryon emission quickly decay away to conventional hadrons.

- *Lepton or photon emission*. A strangelet which is stable against baryon emission would adjust its flavor through a variety of weak processes until it reached a state of minimum energy. The underlying quark processes include electron or positron emission, $(d \text{ or } s) \rightarrow u e^- \overline{\nu}_e$, $u \rightarrow (d \text{ or } s) e^+ \nu_e$, electron capture, $u e^- \rightarrow (d \text{ or } s) \nu_e$, and weak radiative strangeness changing scattering, $ud \rightarrow su\gamma$. These processes are much slower than baryon emission because they typically have three-body final states, initial state wave-function factors, or other suppression factors. Rates would depend on details of strangelet structure which cannot be estimated without a detailed model. We would expect lifetimes to vary as widely as the β decay and electron capture lifetimes of ordinary nuclei, which range from microseconds to longer than the age of the universe.
- *Systematics of stability*. The only studies of strangelet radioactivity were done in the context of a rather primitive model (Gilson and Jaffe, 1993). Even then, some features emerge that would have significant implications for the disaster scenarios which concern us. Specifically:

—Even if the asymptotic value of Z/A were negative, there probably would exist absolutely stable strangelets with positive charge. Production of such a species would terminate the growth of a dangerous strangelet (see below). The opposite case (a negatively charged strangelet in a world where Z/A is asymptotically positive) would not present a hazard.

—Calculations indicate that the lightest (meta)stable strangelet can occur at a value of $A \equiv A_{\min}$ well below the onset of general stability, with no further stable species until some $A' \gg A_{\min}$. This phenomenon occurs in conventional nuclear physics at the upper end of the periodic table, where occasional (meta)stable nuclei exist in regimes of general instability. In this case a dangerous strangelet could not grow by absorbing matter.

Even though these features of strangelet stability could stop the growth of a negatively charged strangelet produced at RHIC, we cannot use them to argue for the safety of RHIC because we do not know how to model them accurately.

For the sake of definiteness, we will refer to any strangelet with a lifetime long enough to be produced at RHIC, come to rest, and be captured in matter as ''metastable.'' To summarize, strangelets which decay by baryon emission have lifetimes which are generally too short to be metastable. Thus any strangelets which eventually evaporate away do so very quickly. On the other hand, strangelets which decay by lepton or photon emission could be quite long lived.

B. Searches for strange matter

In addition to the astrophysical searches reviewed in Madsen and Haensel (1991) and Madsen (1998), experimental physicists have searched unsuccessfully for stable or quasistable strangelets over the past 15 years. Searches fall in two principal categories: (a) searches for stable strangelets in matter; (b) attempts to produce strangelets at accelerators.

Stable matter searches look for stable stangelets created sometime in the history of our Galaxy, either in cosmic-ray collisions or as by products of neutron star interactions. Due to its low charge-to-mass ratio, a stable light strangelet would look like an ultraheavy isotope of an otherwise normal element. For example, a strangelet with $A \approx 100$ might have $Z = 7$. Chemically, it would behave like an exotic isotope of nitrogen, ^{100}N(!). Searches for ultraheavy isotopes place extremely strong limits on such objects (Nitz *et al.*, 1986; Hemmick *et al.*, 1990). The failure of these searches is relevant to our considerations because it further reduces the likelihood that strange matter is stable in bulk at zero external pressure (Blackman and Jaffe, 1989).

Accelerator searches assume only that strangelets can be produced in accelerators and live long enough to reach detectors. Experiments to search for strangelets have been carried out at the Brookhaven National Laboratory Alternating Gradient Synchrotron (AGS) and at the CERN Super Proton Accelerator (SPS). At the AGS the beam species and energy were gold at an energy of 11.5 GeV/nucleon (Xu, 1999). At the CERN SPS the beam was lead at an energy of 158 GeV/nucleon (Klingenberg *et al.*, 1996). Experiments (with less sensitivity) were also done at CERN with sulfur beams at an energy of 200 GeV/nucleon (Borer *et al.*, 1994). In all of these experiments the targets were made of heavy elements (lead, platinum, and tungsten).

All of the experiments were sensitive to strangelets of both positive and negative electric charge. All of the experiments triggered on the low value of Z/A characteristic of strangelets. The experiments were sensitive to values of $|Z/A| \lesssim 0.3$, masses from 5 GeV/c^2 to 100 GeV/c^2, and lifetimes longer than 50 ns $(5\times10^{-8}\ \text{s})$.

None of the experiments detected strangelet signals. Limits were therefore set on the possible production rates of strangelets with the stated properties. The limits achieved were approximately less than one strangelet in 10^9 collisions at the AGS and from one strangelet per 10^7-10^9 collisions at CERN energies, depending on the precise properties of the strangelet.

Of course the limits obtained from previous strangelet searches cannot be used to argue that experiments at RHIC are safe because the total luminosity of earlier searches would not place a decisive limit on the probability of negative strangelet production at RHIC. However, attempts to understand possible strangelet production mechanisms in these experiments figure importantly in our consideration of dangerous strangelet production at RHIC.

C. Strangelet production in heavy-ion collisions

The lack of a plausible mechanism whereby hypothetical dangerous strangelets might be produced is one of the weakest links in the catastrophe scenario at a heavy-ion collider. Before discussing production mechanisms in detail, it is worthwhile to summarize some of the very basic considerations that make dangerous strangelet production appear difficult.

- Strangelets are cold, dense systems. Like nuclei, they are bound by tens of MeV (if they are bound at all). Heavy-ion collisions are hot. If thermal equilibrium is attained, temperatures are of order 100 MeV or more. The second law of thermodynamics fights against the condensation of a system an order of magnitude colder than the surrounding medium. It has been compared to producing an ice cube in a furnace.
- $q\bar{q}$ pairs, including $s\bar{s}$ pairs, are most prevalent in the central rapidity region in heavy-ion collisions. Baryon chemical potential is highest in the nuclear fragmentation regions. To produce a strangelet one needs both high chemical potential and many s quarks made as $s\bar{s}$ pairs. But the two occur in different regions.
- Strangelets include many strange quarks. The more negative the strangelet charge, the more strange quarks. For example, a strangelet with $A=20$ and $Z=4$ would include 12 s quarks if the number of u and d quarks are equal (as expected). However, a strangelet with $A=20$ and $Z=-1$ would have to contain 22 s quarks (again assuming $u=d$). The more strange quarks, the harder it is to produce a strangelet. Thus dangerous strangelets are much harder to make than benign $(Z>0)$ strangelets.
- As we have previously discussed, the smaller the strangelet, the less likely it is to be stable or even metastable. The last several items make it clear that the larger the strangelet, the less likely it is to be produced in a heavy-ion collision.

We find that these arguments, though qualitative, are quite convincing. Especially, they strongly suggest that strangelet production is even more unlikely at RHIC than at lower-energy facilities (e.g., AGS and CERN) where experiments have already been performed.

Unfortunately, the very unlikelihood of production makes it difficult to make a reasonable model for how it might occur, or to make a quantitative estimate.

Two mechanisms have been proposed for strangelet production in high-energy heavy-ion collisions: (a) coalescence and (b) strangeness distillation. The coalescence process is well known in heavy-ion collisions and many references relate to it. A recent study which summarizes data at the AGS energies has been reported (Nagle, 1999). The strangeness distillation process was first proposed by Heinz *et al.* (1987) and Greiner *et al.* (1987).

The coalescence process has been carefully studied at AGS energies (Nagle, 1999). The coalescence model is most easily summarized in terms of a penalty factor for coalescing an additional unit of baryon number and/or strangeness onto an existing clump. By fitting data, Nagle (1999) finds a penalty factor of 0.02 per added baryon. The additional penalty for adding strangeness has been estimated at 0.2; however, the data of Nagle (1999) suggests that it might be as small as 0.03. The model was originally intended to estimate the probability of producing nuclei and hypernuclei from the coalescence of the appropriate number and types of baryons. When it is used to estimate strangelet production, it is assumed that the transition from hadrons to quarks occurs with unit probability. This is certainly a gross overestimate, since wholesale reorganization of the quark wave functions is necessary to accomplish this transition. By ignoring this factor we obtain a very generous *overestimate* of the strangelet production probability. Given that the probability of producing a deuteron in the collision is about unity, this suggests that the yield of a strangelet with, for example $A=20$, $Z=-1$, and $S=22$ is about one strangelet per 10^{46} collisions (taking the strangeness penalty factor as 0.2). This would lead to a probability $\mathfrak{p}\approx2\times10^{-35}$ for producing such a strangelet at RHIC. The difficulty of producing a (meta)stable, negatively charged strangelet (if it exists) is one of the principal reasons we believe there is no safety problem at RHIC.

In addition, the coalescence factors are expected to decrease as the collision energy increases. This is because the produced particles are more energetic, and therefore less likely to be produced within the narrow range of relative momentum required to form a coalesced state. If one compares the coalescence yields at the Bevalac, the AGS, and the CERN experiments, this expectation is dramatically confirmed. From the point of view of coalescence, the most favorable energy for strangelet production is below that of the AGS.

Closely related to the coalescence model is the thermal model, in which it is assumed that particle production reflects an equilibrium state assumed to exist until the fireball cools and collisions cease. In this model the ''free'' parameters are the temperature and the baryon chemical potential at freeze-out (Braun-Munzinger *et al.*, 1995). Applying this model to the AGS experimental situation gives a reasonably good account of particle ratios, and indicates a freeze-out temperature of 140 MeV and a baryon chemical potential of 540 MeV. With these parameters the model can predict the production probability of strangelets with any given baryon number, charge, and strangeness. Braun-Munzinger and Stachel (1995) have carried out detailed calculations for the AGS case and find very small production. For example, the yield of a strangelet with $A=20$, $Z=2$, and $S=16$ is $\sim 2\times 10^{-27}$ per central collision. Since central collisions are about 0.2 of all collisions this translates into a yield of one strangelet (with these parameters) in 2×10^{27} collisions if such a strangelet were stable and if we scale without change from AGS to RHIC energy. The yield of a negatively charged strangelet would be much smaller still.

As the collision energy increases, this model predicts higher temperatures and smaller baryon chemical potentials. The result is that in this model strangelet production is predicted to decrease quickly with total center-of-mass energy in this model. The thermal model clearly favors an energy even lower than the AGS for the optimum for producing strangelets, should they exist.

The strangeness distillation mechanism is considerably more speculative. It assumes that a quark-gluon plasma (QGP) is produced in the collision and that the QGP is baryon rich. It further assumes that the dominant cooling mechanism for the QGP is evaporation from its surface. Since it is baryon rich, there is a greater chance for an $\bar{s}$ quark to find a u or d quark to form a kaon with positive strangeness than for an s quark to find a $\bar{u}$ or $\bar{d}$ quark to form a kaon with negative strangeness. The QGP thus cools to a system containing excess s quarks, which ultimately becomes a strangelet.

This mechanism requires a collision energy sufficient to form a QGP. RHIC should be high enough. Many heavy-ion physicists believe that even the fixed target CERN experiments have reached a sufficient energy and are in fact forming a QGP. If this is the case, the failure of the CERN experiments to find strangelets argues against either the existence of this mechanism or the existence of strangelets. A substantial body of evidence supports the view that a QGP is formed at CERN energies, but a truly definitive conclusion is not possible at present. In any case, fits to data from the AGS and CERN, and theoretical models suggest that the baryon density at central rapidity, where a QGP can be formed, will decrease at RHIC. Moreover, there is considerable evidence that the systems formed in CERN heavy-ion collisions do not cool by slow evaporation from the surface but rather by rapid, approximately adiabatic expansion, as is also expected theoretically. Altogether, the strangeness distillation mechanism seems very unlikely to be effective for producing strangelets at RHIC.

In summary, extrapolation from particle production mechanisms that describe existing heavy-ion collision data suggests that strangelets with baryon number large enough to be stable cannot be produced. With one exception, all production models we know of predict that strangelet production peaks at low energies, much lower than RHIC and perhaps even lower than the AGS. The one exception is the hypothetical strangeness distillation mechanism. However, available data and good physics arguments suggest that this mechanism does not apply to actual heavy-ion collisions.

D. Catastrophe at RHIC?

What is the scenario in which strangelet production at RHIC leads to catastrophe? The culprit would be a stable (or long-lived, metastable) negatively charged strangelet produced at RHIC. It would have to be a light representative of a generic form of strange matter with *negative* electric charge in bulk. It would have to live long enough to slow down and come to rest in matter. Note that the term ''metastable'' is used rather loosely in the strangelet literature. Sometimes it is used to refer to strangelets that live a few orders of magnitude longer than strong interaction time scales. As mentioned above, we use ''metastable'' to refer to a lifetime long enough to traverse the detector, slow down and stop in the shielding. Since strangelets produced at high rapidity are likely to be destroyed by subsequent collisions, we assume a production velocity below $v_{\mathrm{crit}}=0.1c$ (Dar *et al.*, 1999). Hence it requires a lifetime greater than $\sim 10^{-7}$ s in order to satisfy our definition of metastable.

Once brought to rest, a negative metastable strangelet would be captured quickly by an ordinary nucleus in the environment. Cascading quickly down into the lowest Bohr orbit, it would react with the nucleus, and could absorb several nucleons to form a larger strangelet. The reaction would be exothermic. After this reaction its electric charge would be positive. However, if the energetically preferred charge were negative, the strangelet would likely capture electrons until it once again had negative charge. At this point the nuclear capture and reaction would repeat. Since there is no upper limit to the baryon number of a strangelet, the process of nuclear capture and weak electron capture would continue.

There are several ways that this growth might terminate without catastrophic consequences: First, as mentioned earlier, a stable positively charged species might be formed at some point in the growth process. This object would be shielded by electrons and would not absorb any more matter. Second (also mentioned before), the lightest metastable strangelet might be isolated from other stable strangelets by many units in baryon number.[14] Third, the energy released in the capture process might fragment the strangelet into smaller, unstable objects. Unfortunately, we do not know enough about QCD either to confirm or exclude these possibilities.

A strangelet growing by absorbing ordinary matter would have an electric charge very close to zero. If its electric charge were negative, it would quickly absorb (positively charged) ordinary matter until the electric charge became positive. At that point absorption would cease until electron capture again made the quark charge negative. As soon as the quark charge became negative the strangelet would absorb a nucleus. Thus the growing strangelet's electric charge would fluctuate about zero as it alternately absorbed nuclei and captured electrons. Even though the typical time for a single quark to capture an electron might be quite long, the number of participating quarks grows linearly with A, so the baryon number of the strangelet would grow exponentially with time, at least until the energy released in the process began to vaporize surrounding material and drive it away from the growing strangelet. This process would continue until all available material had been converted to strange matter. We know of no absolute barrier to the rapid growth of a dangerous strangelet, were such an object hypothetically to exist and be produced. This is why we have considered these hypotheses in detail to assure ourselves beyond any reasonable doubt that they are not genuine possibilities.

We should emphasize that production of a strangelet with positive charge would pose no hazard whatsoever. It would immediately capture electrons forming an exotic ''strangelet-atom'' whose chemical properties would be determined by the number of electrons. The strange ''nucleus'' at its core would be shielded from further nuclear interactions in exactly the same way that ordinary nuclei are shielded from exothermic nuclear fusion. We see no reason to expect enhanced fusion processes involving atoms with strangelets at their core. It has been suggested that an atom with a strangelet at its core would undergo fusion reactions with light elements in the environment and, like a negatively charged strangelet, grow without limit (Wagner, 1999). This will not occur. First, the strength and range of the strong interactions between a strangelet-atom and an ordinary atom are determined by well-known, long-range properties of the nuclear force which are exactly the same for strangelets as for nuclei. Second, fusion is suppressed by a barrier penetration factor proportional to the product of the charge on the strangelet times the charge on the nucleus, $f \propto e^{-Z_1 Z_2 K}$. The most favorable case would be a strangelet of charge one fusing with hydrogen. Hydrogen-hydrogen fusion at room temperature is so rare that it is a subject of intense debate whether it has ever been observed. Even if strangelet-atom-hydrogen fusion were enhanced by some unknown and unexpected mechanism, the suppression factor that appears in the exponent would be doubled as soon as the strangelet had acquired a second unit of charge. As the strangelet's charge grows each successive fusion would be breathtakingly more suppressed.

To provide a concrete example, we have calculated the rate of fusion of a thermalized (room temperature) strangelet with baryon number 396 (the baryon number present in the entire Au-Au collision) and $Z=6$, with hydrogen. Using standard and well tested nuclear reaction theory, we find a fusion rate of $\sim 10^{-2\times 10^5}\,\mathrm{s}^{-1}$.

On theoretical grounds alone, as discussed above, we believe creation of a dangerous strangelet at RHIC can be firmly excluded. We now turn to the important empirical evidence from cosmic rays.

E. Cosmic-ray data relevant to the strangelet scenario

It is clear that cosmic rays have been carrying out RHIC-like ''experiments'' throughout the Universe since time out of mind. Here we choose some specific conditions and summarize briefly the arguments that place restrictions on dangerous strangelet production at RHIC. We have made estimates based on cosmic-ray collisions with the Moon. We also review the astrophysical estimates in a recent paper by Dar, De Rujula, and Heinz (1999).

In order to extract bounds from cosmic-ray data, it is necessary to model the rapidity distribution of strangelets. It will turn out that the most important distinguishing features of a production mechanism are how it behaves at central and extreme values of the rapidity. Inclusive hadronic processes generally fall like a power of the rapidity near the limits of phase space. In light of this, we see no reason for strangelet production to be exponentially suppressed at Y_{min} and Y_{max}. On the other hand, long-standing theoretical ideas and phenomenology suggest the emergence of a ''central plateau'' away from the kinematic limits of rapidity, along which physics is independent of the rapidity. Insofar as these ideas are correct, a singularity at central rapidity would violate the principle of relativity.

So for our first model we assume a power-law dependence at the kinematic limits of rapidity, and an exponential fall off away from the target fragmentation region, where the baryon chemical potential decreases. By convention we take $y=0$ to be the kinematic limit and we model the strangelet production near $y=0$ by

$$\left.\frac{d\Pi}{dy}\right|_{\mathrm{BG}} = Npy^{a}e^{-by}, \tag{13}$$

where a and b are parameters, N is a normalization constant chosen so that p is half the total strangelet production probability per collision (the other half comes near the other rapidity limit). The subscript ''BG'' stands for ''best guess.''

Dar et al. (1999) have made an extreme model of strangelet production, where production is completely confined to central rapidity. We know of no physical

[14] A similar barrier (the absence of a stable nucleus with $A=8$) prevents two α particles from fusing in stellar interiors.

motivation for this assumption. On the contrary, what we know about particle production in heavy-ion collisions argues against such a model. Their model can be approximated by a δ function at central rapidity,

$$\left.\frac{d\Pi}{dy}\right|_{\mathrm{DDH}} = p\,\delta(y - Y/2), \tag{14}$$

where Y is the total rapidity interval and DDH is the Dar, De Rujula, and Heinz bound. Although we find such a model impossible to justify on any theoretical grounds, we will use this rapidity distribution when we review the work of Dar *et al.* (1999).

The limits from cosmic-ray considerations depend on the assumed rapidity distribution of strangelet production, in the following respect. If strangelets are produced in the nuclear fragmentation regions, then cosmic-ray collisions with stationary nuclei on the surface of the Moon provide more than adequate limits on dangerous strangelet production at RHIC. On the other hand, if strangelets were produced only at zero rapidity in the center of mass, then strangelets produced on the Moon would not survive the stopping process. Under this hypothetical—and we believe, quite unrealistic—assumption the persistence of the Moon provides no useful limit on strangelet production.

Dar, De Rujula, and Heinz introduce a parameter p as a simple way to compare limits obtained in different processes (Dar *et al.*, 1999). p measures the probability to make a strangelet in a single collision with speed low enough to survive the stopping process at RHIC. p is related to the parameter $\mathfrak{p}$ which we introduced earlier by $\mathfrak{p} = 2\times10^{11}p$. We will analyze cosmic-ray data in terms of p and relate the results to $\mathfrak{p}$ when necessary. We assume that p is independent of the atomic mass of the colliding ions, at least for iron and gold. We also assume p is the same for RHIC and AGS energies. A single choice of p simplifies our presentation. We will discuss the qualitative differences between AGS and RHIC energies and between collisions of different nuclear species where they arise. Of course our aim is to bound $\mathfrak{p}$ far below unity.

We begin with our neighbor, the Moon, because we know the environment well and know the Moon is not made of strange matter.[15] The Moon has a rocky surface rich in iron. Using the data from Sec. II it is easy to calculate the rate of collisions between specific heavy ions on the lunar surface.

Consider a cosmic-ray nucleus A colliding with a nucleus A' with fractional abundance $f_{A'}$ in the lunar soil. The total number of collisions at energies greater than E over the five-billion year lifetime of the Moon [from Eq. (5)] is[16]

$$N(A,E)|_{\mathrm{Moon}} \approx 8\times10^{29} f_{A'} \frac{\Gamma(A,\ 10\ \mathrm{GeV})}{\Gamma(\mathrm{Fe},\ 10\ \mathrm{GeV})} \times \left(\frac{10\ \mathrm{GeV}}{E}\right)^{1.7}. \tag{15}$$

Using iron, $f_{\mathrm{Fe}} = 0.012$,[17] and the cosmic-ray abundance of iron and ''gold,'' we can calculate the number of dangerous strangelets which would have been created on the surface of the moon in several cases of interest as a function of p.

I. *Dangerous strangelet production in lunar iron-iron collisions at AGS energies.*

Taking $E = 10$ GeV/nucleon and $f_{\mathrm{Fe}} = 0.012$ we obtain $N_{\mathrm{Moon}}(\text{Fe-Fe, AGS}) \approx 10^{28}p$ for the number of dangerous strangelets produced on the surface of the Moon in terms of the probability to produce one in a single collision at RHIC (p).

II. *Dangerous strangelet production in lunar iron-iron collisions at RHIC energies.*

Scaling E to 20 TeV/nucleon, we find $N_{\mathrm{Moon}}(\text{Fe-Fe, RHIC}) \approx 2\times10^{22}p$.

III. *Dangerous strangelet production in lunar ''gold''-iron collisions at AGS energies.*

The penalty of demanding ''gold'' is a factor of 10^{-5} in cosmic-ray flux, so $N_{\mathrm{Moon}}(\text{Au-Fe, AGS}) \approx 10^{23}p$.

IV. *Dangerous strangelet production in lunar ''gold''-iron collisions at RHIC energies.*

Scaling E to 20 TeV/nucleon, we find $N_{\mathrm{Moon}}(\text{Au-Fe, RHIC}) \approx 2\times10^{17}p$.

The Moon does not provide useful limits for targets less abundant than iron.

The total number of collisions on the surface of the Moon is huge compared to the number anticipated at RHIC. However, strangelets produced with even relatively low rapidity in the lunar rest frame do not survive subsequent collisions with nuclei in the lunar soil. Dar, De Rujula, and Heinz model the survival probability by assuming that strangelets with $v_{\mathrm{crit}} < 0.1c$ survive and all others are torn apart (Dar *et al.*, 1999). Here, we assume

[15] Collisions of cosmic rays with the outer envelopes of stars, gaseous planets, or even terrestrial planets with atmospheres like the Earth and Venus, lead overwhelmingly to collisions with light nuclei like hydrogen, helium, etc. This is not a likely way to make strange matter.

[16] Equation (15) was obtained by multiplying Eq. (5) by $\sim 15\times10^{16}$, the number of seconds in five-billion years, and by the fractional abundance, $f_{A'}$. In addition, the collision cross section varies with A and A' like $(A^{1/3} + A'^{1/3})^2$. Since the dominant constituents of the Moon are lighter than iron, the probability of a cosmic ray interacting with iron (or gold) is higher than measured by its fractional abundance alone. We ignore the A dependence of the cross section because it is small, it increases the strength of our bounds, and it complicates our equations.

[17] Measurements of the iron abundance on the Moon exist for six different Apollo landing sites. FeO abundance by weight ranges from 4.2 to 17.2%. To be conservative we took 4% which, converted to an abundance by number, yields 1.2%. See Taylor (1982).

a geometric strangelet dissociation cross section which is independent of energy, and use standard methods to calculate a survival probability. Our results agree with those of Dar, De Rujula, and Heinz to within a factor of 2 for all cases of interest. Consider a strangelet with atomic mass A, charge Z, and rapidity y in the lunar rest frame. Its survival probability is

$$P(y,A,Z)=\exp[-n\sigma(A)\lambda(y,Z,A)]$$
$$=\exp\left[-4.85\left(1+\frac{1}{3}A^{1/3}\right)^2(\cosh y-1)A/Z^2\right]. \tag{16}$$

Here n is the density of lunar soil (assuming silicon, $n=0.5\times10^{23}\ \mathrm{cm}^{-3}$), $\sigma(A)$ is the geometric cross section for the strangelet to collide with a silicon nucleus, $\sigma(A)=0.4(1+\frac{1}{3}A^{1/3})^2$ b, and $\lambda(y,Z,A)$ is the stopping distance calculated assuming that the strangelet loses energy only by ionization, $\lambda(y,Z,A)=242(\cosh y-1)A/Z^2$ cm.

For a representative dangerous strangelet, e.g., $A=20$, $Z=-1$, the suppression factor in Eq. (16) is very large, $P(y,20,-1)=\exp[-350(\cosh y-1)]$, so only strangelets with $y\approx0$ survive. For the rapidity distribution, Eq. (14), chosen by Dar, De Rujula, and Heinz, all dangerous strangelets produced at RHIC would survive stopping, but no strangelet would survive stopping on the Moon. The more realistic production mechanism of Eq. (13) yields lunar suppression factors of 3×10^{-3}, 10^{-4}, 2×10^{-6}, and 5×10^{-8} when the parameter a (which controls the small y behavior of dN/dy) is chosen as 1, 2, 3, and 4.[18] However, this mechanism also reduces the probability that a strangelet produced at RHIC will survive the stopping process. The survival probabilities are 8×10^{-3}, 8×10^{-3}, 10^{-2}, and 2×10^{-2}, for $a=1,2,3,4$, respectively. Thus the effective lunar suppression factors are as follows: an enhancement of 3 for $a=1$, no suppression for $a=2$, suppression by 2×10^{-4} for $a=3$, and by 3×10^{-6} for $a=4$. Choosing a suppression factor of 10^{-6} we obtain survival probabilities of $10^{22}p$ for case I (iron-iron at AGS energies), $2\times10^{16}p$ for case II (iron-iron at RHIC energies), $10^{17}p$ for case III ("gold"-iron at AGS energies), and $2\times10^{11}p$ for case IV ("gold"-iron at RHIC energies).

To compare with other estimates we convert these results to bounds on $\mathfrak{p}$, the probability of producing a dangerous strangelet at RHIC which survives the stopping process. The fact that the Moon has not been converted to strange matter over its lifetime bounds $\mathfrak{p}$ by $\mathfrak{p}<2\times10^{-11}$, 10^{-5}, 2×10^{-6}, and 1 for cases I–IV, respectively. Since we believe strangelet production to be more likely at AGS energies than at RHIC, and believe iron to be a reasonable "heavy nucleus," we take the limit from case I very seriously. However, if one insists on recreating exactly the circumstances at RHIC and insists on the worst case rapidity distribution, then lunar limits are not applicable.

Dar, De Rujula, and Heinz explore the consequences of dangerous strangelet production in nucleus-nucleus collisions in interstellar space. They adopt "worst case" assumptions at several points. In particular, they demand RHIC energies and ultraheavy nuclei (gold rather than iron), and they assume that a dangerous strangelet is produced only at zero rapidity in the center of mass. Given these restrictive conditions they compute the rate at which strangelets are produced at rest relative to the galaxy. Taking an energy of 100 GeV/nucleon and an abundance relative to iron of 10^{-5} in Eq. (7),[19] we reproduce their result, $R(100\,\mathrm{GeV},\mathrm{Au})\approx10^{-58}$. Multiplying by the age of the galaxy ($T_0=$ten-billion years) and by the probability, p, of dangerous strangelet production, we find the number of dangerous strangelets produced per cm^3 in the galaxy,

$$N(100\ \ \mathrm{GeV},\mathrm{Au})=T_0pR(100\ \ \mathrm{GeV},\mathrm{Au})$$
$$=10^{-41}p\ \ \mathrm{cm}^{-3}. \tag{17}$$

Dar, De Rujula, and Heinz estimate that the material contained in a volume of $10^{57}\,\mathrm{cm}^3$ is swept up in the formation of a "typical star," so that the probability of a dangerous strangelet ending up in a star is approximately $P_\star\approx10^{16}p$. They then go on to argue that the subsequent destruction of the star would be detectable as a supernovalike event. Based on $P_\star$ and the observed rate of supernovas, Dar, De Rujula, and Heinz limit p to be less than 10^{-19}. This corresponds to a limit of 2×10^{-8} on $\mathfrak{p}$, the probability of producing a dangerous strangelet during the life of RHIC. Actually, we believe that Dar, De Rujula, and Heinz have been too conservative. Good physics arguments indicate that lower energy collisions are more likely to create strangelets, and iron is nearly as good a "heavy" ion as gold. If we scale down E from RHIC energies (100 GeV/nucleon) to AGS energies (4.5 GeV/nucleon) we gain a factor of 4×10^4 from the $E^{-3.4}$ dependence in Eq. (7). If we replace gold by iron we gain a factor of 10^{10}. So the bound on dangerous strangelet production during the RHIC lifetime is more nearly $\mathfrak{p}<10^{-21}$.

Finally, we point out the implications of strangelet metastability for these arguments. Dar, De Rujula, and Heinz have implicitly assumed that the dangerous strangelet produced in interstellar space lives long enough to be swept up into a protostellar nebula. Suppose, instead, that the dangerous strangelet was only metastable, and that it decays away by baryon emission with a lifetime greater than 10^{-7} s but much less than the millions of years necessary to form a star. In this case a dangerous strangelet produced at RHIC would

[18]These estimates apply to $A=20$, $Z=-1$. Larger A are more suppressed, but we do not consider production of a negatively charged strangelet with A much larger than 20 to be credible. Larger Z reduces the suppression.

[19]Dar, De Rujula, and Heinz assume an $E^{-2.6}$ decay of the cosmic-ray spectrum and take $\Gamma(\mathrm{Au})/\Gamma(\mathrm{Fe})\approx3\times10^{-5}$, slightly different from our choices.

have time to stop in matter, stabilize, and begin to grow. However, a strangelet formed in interstellar space would decay harmlessly into baryons, etc.[20]

We have estimated baryon emission lifetimes for strangelets. A lifetime of 10^{-7} s is near the upper limit of our estimates. Since the strangelet production cross section is likely to fall so quickly with A and S, the strangelet most likely to be created at RHIC would be the least stable and would likely decay on time scales much shorter than 10^{-7} s by strong baryon emission. A strangelet heavy enough to have a baryon emission lifetime of order 10^{-7} s would be much harder to produce at RHIC. Still, the astrophysical argument of Dar, De Rujula, and Heinz is compromised by the possibility of producing a metastable strangelet with a long enough baryon emission lifetime. Note, however, that instability to decays which do not change baryon number (and therefore do not lead the strangelet to evaporate) is irrelevant. Also, note that metastability does not compromise the lunar arguments: a metastable strangelet produced in the lunar rest frame would have just as much time to react as one produced at RHIC.

This discussion shows the pitfalls of pursuing the worst case approach to the analysis of empirical limits. The rapidity distribution necessary to wipe out lunar limits is bizarre. The metastability scenario necessary to wipe out the astrophysical limits seems less unphysical, but still highly contrived. Compelling arguments assure us that RHIC is safe. Nevertheless, a worst case analysis, based on arguments which bend, if not break, the laws of physics, leads to a situation where there is no totally satisfactory, totally empirical limit on the probability of producing a dangerous strangelet at RHIC.

In summary, we have relied on basic physics principles to tell us that it is extremely unlikely that negatively charged strange matter is stable, that if it is stable in bulk, it is unlikely to be stable in small droplets, and that even small strangelets are impossibly difficult to produce at RHIC. In addition, empirical arguments using the best physics guidance available, as opposed to worst case assumptions, together with data on cosmic-ray fluxes, bound the probability of dangerous strangelet production at RHIC to be negligibly small.

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