

Can Cosmic Strangelets Reach the Earth?

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The mechanism for the propagation of strangelets with low baryon number through the atmosphere of the Earth has been explored. It has been shown that, under suitable initial conditions, such strangelets may indeed reach depths near mountain altitudes with mass numbers and charges close to the observed values in cosmic ray experiments.

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The existence of strange quark matter (SQM), containing a large amount of strangeness, had been postulated by various authors quite a few years ago. In a seminal work in 1984, Witten [1] proposed that SQM with roughly equal numbers of up, down, and strange quarks could be the *true* ground state of quantum chromodynamics (QCD), the accepted theory of strong interactions. While only SQM with very large baryon numbers was initially thought to be favorable (in terms of stability), later calculations have shown [2–6] that small lumps of SQM can also be stable. The occurrence of stable (or metastable) lumps of SQM, referred to in the literature as strangelets, would lead to many rich consequences; for a recent review, see [7].

Strangelets may arise from various scenarios; they could be formed in highly energetic nuclear collisions associated with the formation of quark-gluon plasma [3,5] or they might be of cosmological origin, as remnants of the cosmic QCD phase transition [8]. Collisions of strange stars could also lead to the formation of strangelets which could contribute to the cosmic ray flux [7]. In heavy ion collisions, strangelets with atomic number A up to 20–30 may be formed [5], the stability of which depend rather sensitively on the parameter values (like the Bag constant) and an underlying shell-like structure. For larger strangelets ($A > 40$), the stability appears to be more robust [6,7]. We confine our attention to these larger strangelets, which may not be readily formed in laboratory heavy ion collisions but could be of cosmic origin. A discerning property of such strangelets would be an unusual charge to mass ratio ($Z/A \ll 1$) [7].

The obvious place to look for such strangelets would be in the cosmic ray flux. In this context, it may be recalled that there have been intermittent reports in the literature [9–13] about the detection of exotic cosmic ray events, with unusually low charge to mass ratios; some of these events are tabulated in Table I. Although it appears natural to identify these events with strangelets, no consensus has yet emerged, primarily because of the ambiguities associated with the mechanism of propagation of strangelets through the terrestrial atmosphere. For example, if a strangelet arriving at the top of the atmosphere has a baryon number $A \sim 1000$, there would be a seri-

ous problem with its penetrability through the atmosphere, as the *exotic* events are observed at quite low altitudes. One could assume that their geometric cross sections are very small. Alternately, one could conjecture, *a la* Wilk *et al.* [14–16] and others [17], that although the initial mass of the strangelet is very large, it decreases rapidly due to collisions with air molecules, until the mass reaches a critical value m_{crit} below which the strangelet simply evaporates into neutrons.

The difficulties associated with this interpretation are twofold. First, one has to take account of the fact that, unlike ordinary nuclear fragments which tend to break up in collisions, strangelets can become more strongly bound if they absorb matter [1]. Second, since a strangelet has a net electric charge, it experiences an ever increasing geomagnetic field, which considerably lengthens its path before reaching a certain altitude. This implies many more interactions with the nuclei of the atmospheric atoms, as a result of which the strangelet would “evaporate” much before the desired depth is reached.

These difficulties can be naturally overcome in a different scenario, proposed recently by the present authors [18], in which the stability of the strangelet plays a very important role. In this model, an initially small strangelet, during its travel through the Earth’s atmosphere, picks up mass, rather than losing it, from the atmospheric atoms. Such a situation may prevail unless the propagation velocity of the strangelet is so high that in a collision with the atmospheric nucleons, the excitation energy would exceed the binding energy. We have estimated that for our case (initial A larger than 40), this upper limit on the velocity comes out to be above $0.7c$. (We disregard the possibility of fissionlike fragmentation of the strangelets.) The

TABLE I. Mass and charge obtained from cosmic ray experiments.

Event	Mass	Charge
Counter experiment [9]	$A \sim 350\text{--}450$	14
Exotic track [10]	$A \sim 460$	20
Price’s event [11]	$A > 1000$	46
Balloon experiments [12,13]	$A \sim 370$	14

equation governing the rate of change of mass with respect to distance traveled is given by

$$\frac{dm_s}{dh} = \frac{f \times N \times m_n}{\lambda} \quad (1)$$

and the equation of motion reads

$$\frac{d\vec{v}}{dt} = -\vec{g} + \frac{q}{m_s} (\vec{v} \times \vec{B}) - \frac{\vec{v}}{m_s} \frac{dm_s}{dt}. \quad (2)$$

In the above equation, m_s and \vec{v} represent the instantaneous mass and the velocity of the strangelet, q represents the charge, and λ represents the mean free path of the strangelet in the atmosphere. The factor f determines the fraction of neutrons that are actually absorbed out of the incident neutrons (N). In this case, λ is both a function of h (which determines the density of air molecules) and m_s (which is related to the interaction cross section). The initial velocity has to be bigger than a threshold value, so that a strangelet of a given initial mass and charge can arrive at an altitude ~ 25 km from the sea level, surmounting the geomagnetic barrier. The upper limit of 25 km is chosen primarily to economize on the computation time and is *a fortiori* justified since the density of the atmosphere above this height is almost negligible for our purpose. The variation of atmospheric density with height has been described by a parametric fit to the data given in the standard reference of Kaye and Laby [19].

According to the above analysis, a strangelet with an initial mass of ~ 64 amu and charge ~ 2 evolves to a mass ~ 340 amu or so, by the end of its journey, an altitude ~ 3.6 km above the sea level (typically the height of a north-east Himalayan peak in India, like *Sandakphu*, at a geomagnetic latitude $\sim 30^\circ$ N). This mass is quite close to the few available data (see Table I) and seems to support the interpretation that exotic cosmic ray events with very small Z/A ratios could result from SQM droplets. However, it was assumed in [18] that only neutrons are absorbed preferentially over the protons from the nuclei of the atmospheric atoms (i.e., charge of the strangelet remains constant), the protons being Coulomb repelled. It should nonetheless be realized that in the earlier phase of the journey, when the relative velocity between the strangelet and the air molecule is large, some protons will indeed be absorbed, albeit with a lower cross section than that for neutron capture. As the strangelet builds up in mass as well as in charge, the Coulomb barrier at the surface of the strangelet gets steeper and the relative velocity also gets further reduced. This will slow down the charge transfer process and ultimately inhibit it. Also, one cannot avoid the issue of loss of energy of the strangelet through ionization of the surrounding media. As we shall see, the ionization losses, which become quite significant at comparatively low altitudes, actually provide a lower limit to the height at which the strangelets can be detected successfully.

In this Letter, we therefore try to explore, in a relativistic setting, the consequences of the absorption of protons by the strangelets traversing the terrestrial atmosphere. The equation of motion (2) can be generalized to a relativistic form in a straightforward manner:

$$\gamma m_s \frac{d\vec{v}}{dt} = -m_s \vec{g} + q(\vec{v} \times \vec{B}) - \gamma \vec{v} \left(\frac{dm_{s_n}}{dt} + \frac{dm_{s_p}}{dt} \right) - m_s \vec{v} \frac{d\gamma}{dt} - \frac{f(v)}{\sqrt{3}} \vec{v}, \quad (3)$$

where γ is the Lorentz factor. The third term takes care of the deceleration of the strangelet due to the absorption of neutrons as well as protons, where proton absorption is related to neutron absorption as

$$\frac{dm_{s_p}}{dt} = \frac{\sigma_p}{\sigma_n} \frac{dm_{s_n}}{dt} \equiv f_{pn} \frac{dm_{s_n}}{dt}, \quad (4)$$

where σ_p and σ_n are the cross sections for neutron and proton absorption, respectively. Treating, classically, the proton of energy E as a free charged particle of unit charge in the repulsive Coulomb field of the strangelet, we can easily estimate the minimum separation r_{\min} along the trajectory to be given by

$$\frac{(mv_o b)^2}{2mr_{\min}^2} + U(r_{\min}) = E,$$

where $U(r)$ represents the potential energy of the proton due to the Coulomb field of the strangelet; v_o is the relative speed with which the N_2 nuclei (and hence, its constituent protons) approach the strangelet, and b is the impact parameter. Assuming that charge transfer can take place when $r_{\min} \leq R_s$ (the radius of the strangelet), the corresponding value of b ($\equiv b_c$) is $b_c^2 = R_s^2 [1 - U(R_s)/E]$, so that the proton capture cross section (σ_p) by a strangelet of atomic number Z_s is $\sigma_p = \pi b_c^2 = \pi R_s^2 [1 - \frac{Z_s e^2}{4\pi \epsilon_0 R_s} / E]$.

In contrast, the scattering cross section for neutrons (σ_n) is just $\pi(r_n + R_s)^2$ and hence

$$f_{pn} = \frac{R_s^2}{(r_n + R_s)^2} \left(1 - \frac{1}{E} \frac{Z_s e^2}{4\pi \epsilon_0 R_s} \right). \quad (5)$$

Finally, the last term of Eq. (3) accounts for the ionization loss. The expression for $f(v)$ is given by [20]

$$f(v) = -\frac{dE}{dx} = \frac{Z_s^2 e^4 n Z_{\text{med}}}{4\pi \epsilon_0^2 m_e v^2} \ln \left(\frac{b_{\max}}{b_{\min}} \right). \quad (6)$$

Here, n represents the number density of the atmospheric atoms at a particular altitude, Z_{med} is the number of electrons per atom of N_2 which can be ionized, m_e is the mass of the electron, and b_{\max} and b_{\min} are the maximum and minimum values of the impact parameter. At large velocities, expression (6) reduces to, with I denoting the average ionizing energy,

$$f(v) = \frac{Z_s^2 e^4 n Z_{\text{med}}}{4\pi \epsilon_0^2 m_e v^2} \ln \left(\gamma^2 \frac{2m_s v^2}{I} - \beta^2 \right). \quad (7)$$

However, when the velocity of the strangelet falls below a critical value $v \leq 2Z_s v_0$ ($v_0 = 2.2 \times 10^6$ m/s is

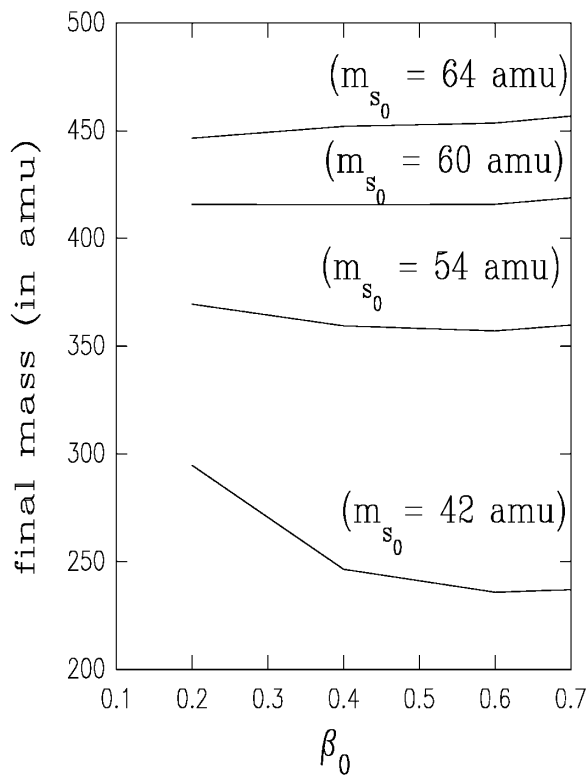


FIG. 1. Variation of final masses with initial β for different initial masses.

the speed of the electron in the first Bohr orbit), electron capture becomes significant which can be accounted for by the replacement $Z_s \rightarrow Z_s \frac{1}{3} \frac{v}{v_0}$ [20,21].

Equation (3) was solved by the 4th order Runge-Kutta method with different sets of initial mass, charge, and β . It may be mentioned at this point that the first term in Eq. (3) is not important in magnitude, as is to be expected. We have nonetheless included it for numerical stability. This serves to define the downward vertical direction in the vector algorithm, especially for very small initial velocities.

Figure 1 shows the final masses (for initial masses 42, 54, 60, and 64 amu and a fixed initial charge 2) as a function of initial β . It is seen, especially for smaller initial masses, that the final mass decreases at first with increasing initial β and then begins to increase again after a critical β is reached and this critical value of β shifts to the left for larger initial mass. This can be ascribed to higher speeds leading to larger mass increments, whose effect would be more pronounced for lower initial masses.

Let us now consider a representative set of data with initial mass 64 amu and charge 2 for detailed discussion. The results for $\beta_0 = 0.6$ are shown in Figs. 2 and 3, where the variation of speed (β) and the energy of the strangelet with altitude are depicted. The sharp change seen at ~ 13 km corresponds to the onset of electron capture, which is handled phenomenologically through the effective Z_s . The insets in Figs. 2 and 3 show a zoomed-up view of the respective quantities near the end point

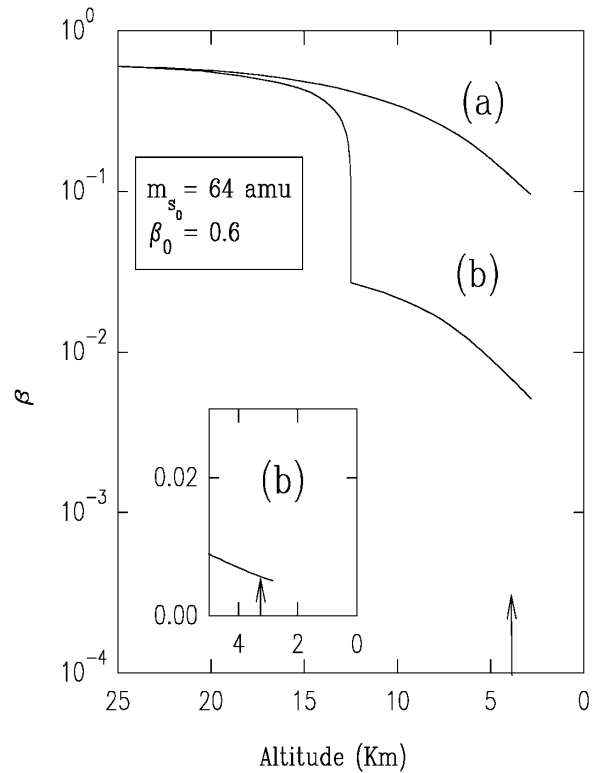


FIG. 2. Variation of final β with altitude (a) for constant charge and without ionization loss and (b) including proton absorption, as well as ionization loss. The inset shows a zoomed view of the graph at mountain altitude.

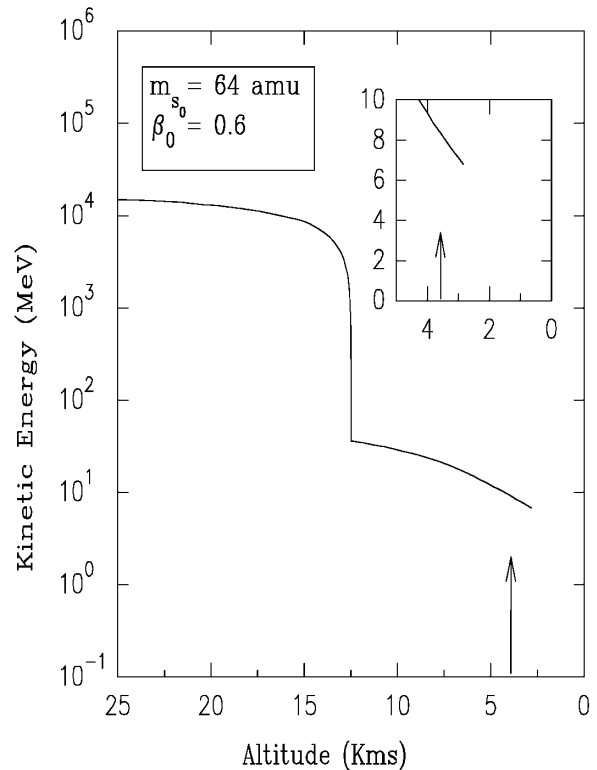


FIG. 3. Variation of kinetic energy with altitude. Inset as in Fig. 2.

TABLE II. The final values, denoted with suffix l , are tabulated along with initial β (β_0).

β_0	m_{s_0}	m_l (amu)	q_l	$\beta_l \times (10^{-3})$	e_l (MeV)
0.2	42	294.7	3	2.8	1.05
	54	369.4	4	3.0	1.55
	60	415.8	4	3.0	1.80
	64	446.5	5	3.1	1.98
0.4	42	246.4	6	4.9	2.84
	54	359.5	8	4.7	3.73
	60	415.6	8	4.7	4.25
	64	452.0	9	4.6	4.63
0.6	42	235.8	10	7.4	5.97
	54	357.1	12	6.6	7.15
	60	416.0	13	6.4	7.87
	64	453.6	14	6.3	8.39
0.7	42	236.4	12	8.6	8.16
	54	359.1	14	7.6	9.59
	60	418.3	15	7.3	10.46
	64	456.3	16	7.2	11.11

of the journey. It is apparent from the figures that the ionization term reduces the overall energy and speed considerably from the nondissipative situation [18]. However, the zoomed-up insets in Figs. 2 and 3 show that the strangelets may have enough energy to be detectable at an altitude of 3.6 km from sea level. For example, for the values of the initial quantities m_{s_0} and β_0 shown here, the strangelet is left with a kinetic energy ~ 8.5 MeV (corresponding to $\frac{dE}{dx} = 2.35$ (MeV/mg)/cm² in a solid state nuclear track detector (SSNTD) like CR-39), which, although small, is just above the threshold of detection $(\frac{dE}{dx})_{\text{crit}} \sim 1-2$ (MeV/mg)/cm² for $\beta < 10^{-2}$ in CR-39 for the present configuration. Below this height, the possibility of their detection with passive detectors like SSNTD reduces to almost zero.

Table II lists the final values of the quantities mass, charge, β , and the energy of the strangelet at the end of the journey for different initial velocities. A comparison between Tables I and II shows that the final masses and charges are very similar to the ones found in cosmic ray events.

In conclusion, we have presented a model for the propagation of cosmic strangelets of none-too-large size through the terrestrial atmosphere and shown that when a proper account of charge and mass transfer as well as ionization loss is taken, they may indeed reach mountain altitudes, so that a ground based large detector experiment would have a good chance of detecting them.

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