Pulsar scintillation patterns and strangelets

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We propose that interstellar extreme scattering events, usually observed as pulsar scintillations, may be caused by a coherent agent rather than the usually assumed turbulence of \( \text{H}_2 \) clouds. We find that the penetration of a flux of ionizing, positively charged strangelets or quark nuggets into a dense interstellar hydrogen cloud may produce ionization trails. Depending on the specific nature and energy of the incoming droplets, diffusive propagation or even capture in the cloud are possible. As a result, enhanced electron densities may form and constitute a lens-like scattering screen for radio pulsars and possibly for quasars.

A variety of scintillation phenomena observed from pulsars and quasars require interstellar scattering screens that contain compact regions of high electron density. These include quasar Extreme Scattering Events (ESE) ¹ ², pulsar parabolic arcs ³ ⁴ and Galactic Center scattering of OH maser sources ⁵ ⁶. Many of these phenomena, in particular the ESEs, require enhanced electron density regions of A.U. size (\( \sim 1.5 \times 10^{13} \) cm). The overpressure \( P/k_B \approx 10^6 - 10^8 \) K cm\(^{-3} \), as estimated from typical temperatures ⁷ \( T \approx 10^4 \) K and particle densities \( n \approx 10^2 - 10^4 \) cm\(^{-3} \), is difficult to explain in any conventional scattering screens embedded in dense molecular \( \text{H}_2 \) clouds. The only plausible environment where such pressures might be attained would be in the dense cloud cores (cf. the compact ionized cloud model developed by Walker ⁸ ⁹). The required source properties in this latter model require a significant mass in extremely dense cold gas clumps to source the ionized gas. The stability of such cold clumps is questionable, although exotic models have been proposed ¹⁰.

There is some direct evidence that the ionized clouds are highly elongated ¹¹. It is an already well-known fact that plasma lenses result from electron (over-) under-densities. Electron over-densities result in a faster phase velocity, corresponding to a concave (divergent) optical lens, while under-densities are associated with a convergent lens ¹². Quantitative lens models for the ESEs ⁷ ¹³ estimate electron column densities \( N_e \approx 10^{15} \) cm\(^{-2} \).

Here we are interested in estimating the effect of the enhancement of the electron column density on A.U. scales as a result of the formation of ionized trails in \( \text{H}_2 \) (molecular) and HI (atomic) hydrogen clouds caused by an external agent. We propose a mechanism that provides an alternative to postulating the existence of controversial clumps of dense molecular hydrogen, and naturally generates pervasive ionized trails in dense interstellar cloud cores. Our model invokes strangelets ¹² (also known as nuclearites), finite droplets of quark matter with non-zero strangeness fraction and slightly charged. They are currently being searched on earth ¹³ as final products in heavy ion collisions, with the ALICE experiment at the LHC, in the CDMSII, under the form of light ionizing particles, or in the AMS-02 mission. It is expected that these quark droplets can be naturally generated by a series of different astrophysical events where a nucleon-quark deconfinement transition may take place, e.g. neutron star (NS) collisions, NS or black hole combined binary mergers or in NS to quark star (QS) conversions. This latter process may be induced by internal heating due to dark matter annihilations ¹⁴ (under the assumption of a Majorana particle candidate) and even leave observable traces in the pulsar distribution ¹⁵ or in the emission of very short gamma ray bursts (GRBs) with typical time scales \( T_{90} \approx 0.1 \) s ¹⁶ detectable with modern projected missions ¹⁷. Due to the large gravitational and nuclear binding energies released in the transition process, a mass ejection episode is expected, possibly seeding the interstellar medium with a fraction of strangeness-carrying lumps of matter formed during the phase of nucleon deconfinement. Energetics show that the measured short GRBs isotropic equivalent photon emission value \( E_{\gamma,iso} \approx 10^{48} - 10^{49} \) erg is compatible with relativistic mass ejecta \( M_{ej} \approx 10^{-4} M_{\odot} \) able to consistently produce observable gamma rays. It has been actually proposed that quark matter droplets might partially populate cosmic ray (CR) primaries (see e.g. ¹⁸–²⁰).

In this context, let us consider a cloud of mixed \( \text{H}_2 \) and HI. Typical ionization reactions are of the form, \( \text{X} + \text{H}_2 \rightarrow \text{X} + \text{H}_2^+ + e^- \), or \( \text{X} + \text{HI} \rightarrow \text{X} + \text{H}^+ + e^- \), where X is the incoming charged strangelet. In addition, electron capture by the positively charged strangelet could be, in principle, possible ²¹. The energy needed to ionize a hydrogen atom (molecule), initially in the ground
state is I(HI) = 13.6 eV (I(H_2) = 15.6 eV). In astrophysical CGS units a more practical conversion factor 1 eV = 1.62 \times 10^{-12} \text{ erg} is used. The dimensions of the molecular cloud (MC) vary but the denser regions with \( n \sim 10^{4+5} \text{ cm}^{-3} \) are typically less than \( \sim 1 \text{ pc} \). Assuming for the cloud core \( R_C \sim 0.1 \text{ pc} \sim 3 \times 10^{17} \text{ cm} \), then the core volume is \( V_C \sim 4\pi R_C^3/3 \approx 1.1 \times 10^{53} \text{ cm}^3 \). Taking in the cloud core \( n_0 \sim 10^4 \text{ cm}^{-3} \), the number of particle species (HI, protons, electrons) is accordingly given by \( n_0 V_C \approx 1.1 \times 10^{57} \). In addition, non-vanishing magnetic fields are expected in the MC and can be parametrized as \( B \sim 100(\frac{n_0}{10^4 \text{ cm}^{-3}})^{1/2} \mu \text{G} \). In general, other ionizing agents, such as CRs can cause ionization of \( \text{H}_2 \) or HI with a rate \( \zeta^{\text{HI}} \sim 10^{-15} \text{s}^{-1} \) or \( \zeta^{\text{H}_2} \sim 10^{-17} \text{s}^{-1} \). A more exotic flux of potentially ionizing electrically charged strangelets through the cloud will depend, on the one hand, on the time and spatial distribution of their astrophysical sources (besides a possible primordial background) and, on the other hand, on their peculiar nature.

We will consider a simplified model with emission of lumps of baryonic number \( A \) and mass \( m_A \lesssim A m_N \), where \( m_N \) is the nucleon mass. The detailed mass formula can be obtained from existing calculations \([21]\). The strangelet number production rate in the astrophysical \( \text{ith}-\text{process} \) will be given by \( dN_{A,i}/dt = \eta_i M_{A,i}/m_A \). \( M_{A,i} \) is the mass rate and \( \eta_i \) is the efficiency of strangelet ejection for the \( \text{ith}-\text{process} \) involved, respectively. For example, for NS collisions it is expected that \( M_{A,j} \sim (10^{-5} - 10^{-2}) M_\odot \) while in a NS merger the mass ejection is \( M_{A,j} \sim (5 \times 10^{-4} - 7 \times 10^{-3}) M_\odot \) for equal-mass binaries with total mass \( m = 2.7 M_\odot \). In a NS to QS transition capable of emitting a GRB, it is expected that \( M_{A,j} \lesssim 10^{-4} M_\odot \). According to \([15]\), the rate of these transitions is \( R_{\text{NS-QS}} \sim (8 \times 10^{-4} - 3 \times 10^{-3}) R_{\text{SNII}} \) being \( R_{\text{SNII}} \sim 10^{-2} \text{ yr}^{-1} \) the rate of type II supernovae in our galaxy.

Having enumerated the possible processes that may constitute a source of the strangelet flux, for the rest of this work for practical purposes we will consider a generic source where the deconfinement transition can take place with galactic appearance rate \( R = \dot{R} \sim 10^{-5} \text{ yr}^{-1} \) and ejected mass \( M_{A,j} = M_{A,j} \sim 10^{-5} M_\odot \). The ejected mass rate is then \( \dot{M}_{A,j} = R M_{A,j} \sim 10^{-10} M_\odot \text{ yr}^{-1} \). Since the efficiency of strangelet production depends on the so far unknown details of the engine model, we will consider that only a small fraction \( \eta \sim 10^{-2} \) is ejected under exotic form \([14]\). Nevertheless, strangelets should be emitted with \( A \)-values larger than a critical stability value \([20]\), \( A > A_{\text{min}} \approx 10^2 - 10^3 \) so that they can possibly decay to the lightest energetically stable \( A_{\text{min}} \)-species.

The peculiar nature of the quark droplets is highly uncertain but it is usually assumed that their charge is small and distributed positively on the surface \([18]\). For ordinary strangelets \( Z \approx A^{1/3} \) \([27]\) while for CFL (colour-flavour-locked) strangelets \( Z \approx 0.3 A^{2/3} \). Even smaller charge-to-mass ratios are energetically possible for intermediate masses, \( A \sim 10^2 - 10^{18} \), assuming a strong coupling \( \alpha_s = 0.9 \) \([28]\). For example, for \( A \approx 10^9 \), \( Z/A \approx 10^{-4} \) while larger strangelets \( A \approx 10^{18} \) have \( Z/A \approx 10^{-7} - 10^{-5} \) and even \( Z/A < 0 \). In this work we constrain strangelets to have \( Z \geq 1 \).

To estimate the strangelet production number rate we are going to consider scintillation from a galactic emitting generic source located at a distance \( d_{\text{SO}} \lesssim 10 \text{ kpc} \) from an observer. As an example, pulsar simulations yield a spatial distribution peaking about galactocentric radii \( \sim 5 \text{ kpc} \) and vanishing beyond \( \sim 12 \text{ kpc} \) \([22]\). We will also discuss the effect from a possibly nearby pulsating source \([30] \) \([31]\) at \( d_{\text{SO}} \sim 600 \text{ pc} \). Then, the A-sized strangelet number galactic production rate at the generic source is given by

\[
\frac{dN_A}{dt} = 2 \times 10^{15} \left( \frac{\eta}{0.01} \right) \left( \frac{M_{A,j}}{10^{10} M_\odot \text{ yr}^{-1}} \right) \frac{f_S(Z, \beta)}{A} \text{ yr}^{-1}.
\]

(1)

It is expected that a possible emission distribution function at the source \( f_S(Z, \beta) \) can modulate this rate. We will not consider this refinement here, and in what follows we will assume \( f_S(Z, \beta) \sim 1 \).

Due to the fact that strangelets are electrically charged they will diffuse in the magnetized medium and the effective distance traveled over the rectilinear distance \( d_s \), is obtained as \( l(d_s) = d_s^2 c/(2 D) \) in a corresponding diffusive time \( t_{\text{diff}} \sim l(d_s)/c \). For the diffusion coefficient, \( D \), in the galactic halo we take \([22] \) \([D(E) = 1.33 \times 10^{28} H_{\text{kepc}}[E/(3Z \text{ GeV})]^{1/2} \text{ cm}^2 \text{ s}^{-1}, \text{where } H_{\text{kepc}} = H/(1 \text{ kpc}) \) is its height. In the MC larger values of the magnetic field are assumed and following \([32]\), we take \( D_{\text{MC}}(E) \approx 1.7 \times 10^{27} [E/(Z \text{ GeV})]^{1/2} B/10 \mu G \text{ cm}^{-1/2} \text{ s}^{-1} \) with an averaged value over the MC of \( B \sim 10 \mu G \).

Typically, the strangelet ejection energy at a transitioning source allows Lorentz factors bounded by a saturation value, \( \gamma \lesssim \gamma_{\text{sat}} \approx 20 - 1000 \) \([14]\). Correspondingly, the kinetic energy is \( T \sim (\gamma - 1) A \text{ GeV}/c^2 \) so that we will assume \( T \sim A T_0 \sim A \text{ TeV} \) droplets with a moderate \( A \gtrsim A_{\text{min}} \) and \( Z \) charge. In such scenario and in the universe lifetime, \( \tau_{\gamma}, N_{\gamma} \sim \tau_{\gamma} \sim 10^5 \text{ s} \) sources could be expected at \( d_{\text{SO}} \lesssim 10 \text{ kpc} \). In that case the unscreened diffusive flux is \( F_0 \sim 0.01 N_{\gamma} \) and

\[
F_0 \approx \frac{dN_A}{dt} \frac{1}{4 \pi l(d_{\text{SO}})^2} \sim 2.3 \times 10^{16} Z^{-2/3} A^{-1/3} \text{ cm}^{-2} \text{ sr}^{-1} \text{ s}^{-1}.
\]

(2)

We must note, however, that some sources may be closer than the assumed \( 10 \text{ kpc} \). In the case that \( d_{\text{SO}} \sim 600 \text{ pc} \) \([29]\) then there are additional volumes \( V_{600 \text{ pc}} \approx 2.1 \times 10^{-4} \) and distance \( \left( \frac{100 \text{ kpc}}{600 \text{ pc}} \right)^2 \approx 7.7 \times 10^4 \) factors yielding an estimate one order of magnitude larger than
previous value (we have roughly assumed averaged source distribution in the halo).

As a comparison, for strangelets being currently searched in neutrino telescopes on earth, there is a lower limit $A \gtrsim 10^{13}$, our estimates yield in that case $F \sim 1.1 \times 10^{-15}$ cm$^{-2}$ sr$^{-1}$ s$^{-1}$. We will restrict to $Z = 1$, $A \sim 10^9$ droplets but larger $A$ are allowed if they are less energetic remaining below observational CR bounds $\sim 10^{20}$ eV. If the source is nearby by chance there are presently competitive limits from experiments such as ANTARES [34], who report a testing capability flux $F_{\text{ANTARES}} \sim 2 \times 10^{-14}$ cm$^{-2}$ sr$^{-1}$ s$^{-1}$ for $A \gtrsim 10^{13}$, or IceCube-22, who report $F_{\text{IceCube-22}} \sim 10^{-18}$ cm$^{-2}$ sr$^{-1}$ s$^{-1}$ for $A \gtrsim 10^{17}$ nuclearites [35]. It is however uncertain whether such large-A droplets can arrive on earth without suffering spallation or decay processes [20].

Once the nuclearites are produced it may happen that their diffusive trajectories intersect with a MC. These are accumulated in the midplane of the galactic disk although we suppose that may also be present at higher latitudes. If the Larmor radius is comparable to the typical coherence length of the galactic magnetic field, $r_L \lesssim l_c$, with $l_c = 10 - 100$ pc [30], then a strangelet will suffer an accumulated deflection that can be estimated from random walk approximation as,

$$\theta(T) \simeq 5.4^2 \left( \frac{l_c}{10 \text{ pc}} \right) \left( \frac{d_{\text{SC}}}{10 \text{ kpc}} \right)^{1/2} \left( \frac{Z/A}{10^{-7}} \right) \left( \frac{B}{1 \mu G} \right) \left( \frac{1 \text{ TeV}}{T_3} \right).$$

Thus, propagation effects in the galaxy or MC for moderate $Z/A$ will not make possible to determine the line-of-sight direction to a distant emitting source.

In this scenario one can estimate $N_A^C$, the effective number of ionizing $A$-sized strangelets traversing the MC core in its lifetime, $\tau_C \sim 10^7$ yr, as $N_A^C \simeq \eta_0 \pi R_c^2 \frac{c}{v} \frac{d_{\text{SC}}}{d_{\text{SO}}} F dt$. Since beaming is observed in some very short GRBs (expected in this scenario [16]) with opening angle $\theta_i \sim 1^\circ - 30^\circ$, an efficiency factor $\eta_0 \simeq \frac{1}{1 - \cos \theta_i}$ has been introduced. Due to the deflection effects we expect typically $\eta_0 \simeq 0.1 - 1$. Using values $d_{\text{SC}} \sim d_{\text{SO}} \sim 10$ kpc then

$$N_A^C \simeq 2 \times 10^{38} A^{-1/3} Z^{-2/3} \left( \frac{\eta_0}{0.1} \right) \left( \frac{R_c}{0.1 \text{ pc}} \right) \left( \frac{d_{\text{SC}}}{d_{\text{SO}}} \right)^2 \left( \frac{\tau_C}{10^5 \text{ yr}} \right).$$

If closer sources are considered [20] then the MC would be illuminated, effectively, by an order of magnitude larger amount of particles as discussed previously and even higher if the source was inside the MC [20].

As the slightly charged and heavy strangelet traverses the MC with a diffusive behaviour, continuously loses kinetic energy. Strangelets of net effective charge $Z_1$ may alter the cloud ionization fraction either by ionizing hydrogen or capturing electrons so that they change its own incident state of charge, as measured for massive ions on gas [31]. For low velocities ($\beta = v/c$) this may be important as it diminishes the ionizing power from the bare charge $Z$ to an effective $Z_1 = Z(1 - e^{-(0.95/\gamma Z^{2/3})})$ [38].

Interaction in the cloud may arise due to a variety of processes [28][21]. The strangelet (kinetic) energy loss per unit length due to ionization and pion production can be obtained from the stopping power $\frac{dT}{dx} = \frac{dT}{dt} \frac{1}{v}$. At lower energies, ionization is the most important process, and for an incoming particle with effective charge $Z_1$ and velocity $\beta$ the energy loss rate reads [39]

$$\frac{dT}{dt} = -1.82 \times 10^{-3} \left( \frac{n}{10^{4} \text{ cm}^{-3}} \right) \left[ Z_1^2 \Psi(Z, \beta) + \Xi(A, \beta) \right] \text{eV/s,}$$

where

$$\Psi(Z, \beta) = \left[ 1 + 0.0185 \ln(\beta) \beta(\beta - \beta_0) \right] 2 \beta^2,$$

and

$$\Xi(A, \beta) = 0.72 A^{0.53} \gamma^{1.28} \theta(\gamma - 1.3),$$

$\theta(x)$ is the Heaviside function, $\gamma^{-1} = \sqrt{1 - \beta^2}$ and $\beta_0 = 0.01$ is the electron orbital velocity. This function is shown in Fig. 1 versus the strangelet kinetic energy per baryon number for an $A = 10^3$ case. Different amounts of charge $Z/A = 5 \times 10^{-3}$ (solid), ordinary (dotted) and CFL (dot-dashed line) strangelets and a CR proton case (dashed line) have been considered. For relativistic strangelets most of the energy is kinetic and they show an enhanced ionizing power with respect to the linear case and a change in slope around $T/A \sim 500$ MeV, due to pion production, for larger energies.

To estimate the hydrogen ionization rate by strangelets, we consider the number density of old NSs in our galaxy assuming $N_8 \sim 10^8$ NSs in a halo radius $\sim 10$ kpc. This yields $n_{\text{NSs}} \sim (10^{18} \text{pc})^3 \sim 10^{-4}$ pc$^{-3}$. Note that MC density is $n_{\text{MC}} \sim 10^{-6}$ pc$^{-3}$. In the lifetime of the MC from a rate of conversions $R \sim 10^{-5}$ yr$^{-1}$ the conversion fraction is $f_{\text{conv}} = \tau_C R/N_8 \sim 10^{-6}$. Therefore conversions are rare non-repeating events as naturally arises in this scenario.

![FIG. 1. Logarithm of the kinetic energy loss rate in the cloud as a function of the logarithm of the incoming $(A = 10^3)$ strangelet kinetic energy per baryon number. Several charge states are shown and a CR proton case with $Z = 1$ is depicted for comparison.](image-url)
Having estimated the incoming strangelet flux at distance $d_{SC}$ and the effective number of strangelets traversing the core, $N_{S}^{\chi}$, the injection rate, $S$, must take into account the fewer conversions in the MC lifetime $S \sim F \frac{N_{S}^{\chi}}{\tau_{\text{rec}}}$ with $\tau_{\text{rec}} = R \tau_{C}$.

As the rate depends on the nature and energetics of the droplet it thus influences the strangelet density in the MC. For example, taking a diffusion escape time for a few TeV/A strangelet entering the MC, $t_{\text{diff}} \sim l(R_{C})/c \sim 0.027 (Z/A)^{1/2}$ yr. In that case the strangelet number density in the MC is $n_{A} \approx S t_{\text{diff}} R_{C}^{-1} \sim 6.5 \times 10^{-26} Z^{-1/6} A^{-5/6} n_{e} \text{cm}^{-3}$. As we have seen in Fig. 2, some lumps may be captured in the MC lifetime and in that case $n_{A}$ would be a factor $\tau_{C}/t_{\text{diff}}$ larger. For moderate A lumps, these values are much lower than the CR density at the low-energy break in the spectrum of Galactic CRs at $\sim 1$ GeV, computed from the CR flux $F_{CR} \sim 10^{4} m^{-2} s^{-1} sr^{-1}$ or $n_{CR} \sim 3 \times 10^{-11} \text{cm}^{-3}$ and at $\sim 1$ TeV, $F_{CR} \sim 10^{-1} m^{-2} s^{-1} sr^{-1}$, $n_{CR} \sim 3 \times 10^{-16} \text{cm}^{-3}$.

The hydrogen ionization rate due to strangelets $\zeta_{H}^{H}$ averaged over the MC, is now estimated to be $\zeta_{H}^{H} = |dI/dt| n_{A}(I)^{-1} n_{e}^{-1}$, per H atom. Assuming ($I$) $\approx 36$ eV per H ionization and heavy strangelets with $\gamma \sim 10^{3}$ then $\zeta_{A}^{H} = 1.2 \times 10^{-25} [Z^{1.8} A^{-5/6} + 5 \times 10^{3} A^{-1/3} Z^{-1/6}] s^{-1}$. The strangelet mean free path or average range can be calculated as $R \sim f dT \left(\frac{dT}{dx}\right)^{-1}$. If strangelets have charge $Z > 1$ their range will be smaller than those previously discussed. In Ref. 2 we show the logarithm of the strangelet range in the cloud (in A.U.) as a function of the logarithm of the incoming kinetic energy per baryon number. We consider a proton CR case $Z = 1$, ordinary and CFL strangelets with $A = 10^{2}$ and $A = 10^{6}$ respectively, and larger strangelets with $A = 10^{13}$ and $A = 10^{18}$ with $Z/A = 10^{-7}$. As seen the cloud can effectively stop the very heavy nuggets while this is not possible for the smaller droplets even at lower kinetic energies.

Once the cloud core is populated with trails of HII (ionized hydrogen), this region contains $H^{+}$ ions that tend to recombine radiatively as $H^{+} + e^{-} \rightarrow H + h\nu$. The balance between ionization and recombination determines the ionized fraction. To estimate the possibility of a net ionization in the HI cloud one must compare the ionization and recombination times, $\tau_{\text{ioniz}}$ and $\tau_{\text{rec}}$, respectively, so that $\tau_{\text{ioniz}} \lesssim \tau_{\text{rec}}$. If recombination is slower, then there would be a net amount of ionized HI in the cloud to cause an electron density enhancement. In addition, there might be a fraction of net charge due to the accumulation of positively charged lumps stopped by the cloud. The recombination time, $\tau_{\text{rec}}$, is calculated in the on the spot approximation as $\tau_{\text{rec}} = 1/n_{e} \alpha_{2} = 3.85 \times 10^{2} n_{e}^{-1} T_{4}^{-0.8} s$ where we take $T_{4} = T/10^{4} K$ and $n_{e}$ is the electron number density. We estimate the average ionization time from the average number of ionizations in the cloud lifetime as $\tau_{\text{ioniz}} \approx \tau_{C}/N_{A}$. Since $\tau_{\text{ioniz}} < \tau_{\text{rec}}$ then it is indeed possible to obtain net ionization. The ionization of interest for the ESE occurs along the trajectory, that spans the MC, but only over a trajectory width given by the product of hydrogen recombination time and sound speed, $c_{s} \sim 13 T_{4}^{1/2} \text{km s}^{-1}$, namely at $T = T_{4}$ $h \sim c_{s} \alpha^{-1} n_{e}^{-1} \sim 33.3/n_{A} \text{A.U.}$, where $n_{A} = 10^{4} \text{cm}^{-3}$. Therefore the local ionization rate is increased by a factor $\sim (R_{C}/h)^{2} \sim 3.6 \times 10^{9} n_{A}^{2}$ at TeV energies. The condition $t_{\text{rec}}^{-1} = \alpha_{2} n_{e} = \zeta_{H}^{H}$ yields an electron density of $n_{e} \sim 8.4 \times 10^{-4} A^{-1/3} Z^{-1/6} n_{A} \text{cm}^{-3}$. Typical ESE events have been argued to be explained by localized and dense plasma structures of length comparable to trail widths of $h \sim 1.7 \times 10^{4} - 10^{3} \text{km}^{-1}$ for densities $n_{A} = 10 - 1$ typical in dense cores. The enhancement in the column density coming from the ionization produced by the incoming particles is estimated as $\Delta n_{e} = \int n_{e}(s) ds$. Then for a typical heavy species with energies of TeV that is stopped in the cloud the accumulated enhancement over the cloud lifetime for $n_{A} = 1 - 10$ is $\Delta n_{e} \sim n_{e} R_{C} \sim 8.4 \times 10^{-4} A^{-1/3} Z^{-1/6} n_{A} 3 \times 10^{17} \text{cm}^{-2} 2.5 \times 10^{14+15} A^{-1/3} Z^{-1/6} \text{cm}^{-2}$ as required by these scattering events.

We have shown that pulsar scintillations may be caused by an ionization agent constituted by positively charged lumps of strange matter. Typical obtained over-densities in electron column densities are compatible with those of quantitative ESE models.Compact A. U. sized regions are involved in the geometry of these events. Since these quark nuggets remain to be experimentally discovered further work is needed to experimentally confirm this theoretical scenario.

FIG. 2. Logarithm of the strangelet range in the cloud as a function of the logarithm of the incoming kinetic energy per baryon number. See text for details.
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