Asymmetric neutrino propagation in newly born magnetized strange stars: GRB and kicks

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Abstract

We consider the early cooling evolution of strongly magnetized strange stars in a CFL phase with high gap $\Delta \gtrsim 100$ MeV. We demonstrate how this model may explain main features of the gamma-ray burst phenomena and also yield a strong star kick. The mechanism is based on beaming of neutrino emission along the magnetic vortex lines. We show that for sufficiently high initial temperatures $T_0 \sim 30$ to 60 MeV and surface magnetic fields $B_s \sim 10^{15}$ to $10^{17}$ G, the energy release within the narrow beam is up to $10^{52}$ erg with a magnetic field dependent time scale between $10^{-2}$ s (for a smaller magnetic field) to 10 s. The above mechanism together with the parity violation of the neutrino-producing weak interaction processes in a magnetic field allow for the strange star kick. The higher the magnetic field the larger is the stars kick velocity. These velocities may cover the same range as observed pulsar kick velocities.

1 Introduction

Compact stellar objects are subdivided into three groups: usual neutron stars (having only hadronic matter interiors and an extended crust), hybrid stars (with quark core, hadronic shell and a crust similar to that in neutron stars) and strange stars (hypothetical self-bound objects composed of $u$, $d$ and $s$ quark matter with only a thin crust if any). The possibility of strange stars has been extensively discussed in the literature, see [1] and Refs. therein.

Gamma ray bursts (GRB) are among the most intriguing phenomena in the Universe, see [2] and Refs. therein. If the energy is emitted isotropically, an energy release of the order of $10^{53}$ to $10^{54}$ erg is needed to power the GRB. However, there is now compelling evidence that the gamma ray emission is not isotropic, but displays a jet-like geometry. In the case of beaming a smaller energy, of the order of $10^{51}$ erg, would be sufficient to power the GRB [3]. Although typically one may speak about a short GRB with a time scale $t \sim 10^{-3}$ to $10^{-1}$ s and a long GRB with a time scale $t \sim 1$ to 100 s, energies, durations, time scales involved in a burst, pulse shape structures, sub-burst numbers, etc, vary so much that it is hard to specify a typical GRB.

Many scenarios for the GRB have been proposed: the “collapsar”, or “hypernova” model linking the GRB with ultra-bright type Ibc supernovae (hypernovae) and the subsequent black hole formation.
neutron star mergers \cite{3}, or accretion of matter onto a black hole; strange star collisions \cite{6}; the model \cite{7} assuming a large surface magnetic field up to $10^{16}$ G of a millisecond pulsar produced in the collapse of a $10^9$ G white dwarf, with a powerful pulsar wind, as the source of the GRB; the model \cite{8} suggesting a $ee^+$ plasma wind between heated neutron stars in close binary systems as consequence of the $\nu\bar{\nu}$ annihilation; the model \cite{10} of a steadily accreting $\sim 10^6$ G white dwarf collapse to a millisecond pulsar with a $10^{17}$ G interior toroidal field, causing the GRB; an (isotropic) first order phase transition \cite{11} of a pure hadronic compact star to a strange star (see also \cite{12}); asymmetric core combustion \cite{13} in neutron stars due to the influence of the magnetic field generating an acceleration of the flame in the polar direction, etc.

Pulsars are rotating neutron stars with rather high magnetic fields causing observable radio dipole signals. Most of the known pulsars are born in the neighborhood of the galactic plane and move away from it with natal kick velocities which are typically higher than those of their progenitors \cite{14}. This implies that the birth process of pulsars also produces their high velocities and thus cannot be entirely isotropic \cite{15}. Up to now, the mechanisms driving this asymmetry are far from clear. There are several hypotheses, ranging from asymmetric supernova explosions \cite{17}, neutron star instabilities \cite{18,19} and magneto-rotational effects \cite{20,21} to the model of an electromagnetic or neutrino rocket \cite{17}. It is also not clarified whether the distribution of pulsar kick velocities is bimodal with a low-velocity component of $v \leq 100$ km/s (20 \% of the known objects) and a high-velocity component of $v \geq 500$ km/s (80\%) as suggested by \cite{22} or whether it can be explained by a one-component distribution \cite{10}. Most of the models are capable of explaining kick velocities of $v \sim 100$ km/s, but it is a nontrivial problem to explain the highest measured pulsar velocities around 1600 km/s.

In this work we assume that a strongly magnetized strange star has been formed in the color superconducting color-flavor-locked (CFL) state with a large gap $\Delta(T)$, $\Delta(T = 0) \lesssim 100$ MeV, cf. \cite{23}.\footnote{Our results are also relevant for the mCFL phase \cite{24}. We only need all quarks to be gapped with large gaps} as the result of a phase transition. The latter could be caused by the accretion of the matter from a companion star, the neutron star angular momentum decrease owing to the gravitational and electromagnetic radiation after the supernova event had occurred, or during the collapse of a magnetized white dwarf to the neutron star state, the proto-neutron star collapse to the new stable state, or by some other reason. We also can deal with a hybrid star instead of the strange star. However, cf. \cite{10}, there is experimental evidence that the GRB carries only a tiny baryon load of mass $\lesssim 10^{-4}M_\odot$. Therefore the hadronic shell and the crust should be rather thin or there should be a special reason for a low baryon loading. We suggest to explain the beaming by the presence of a strong magnetic field. \textit{Thus the beaming and the low baryon loading stimulate us to conjecture about a magnetized strange star, which has a tiny hadron shell, if any, and only a thin crust.} Varying the value of the surface magnetic field $B_s$, we search for an optimal configuration to explain the GRB characteristics and to estimate a range of velocities of the star kicks. We use units $\hbar = c = 1$.

\section{Vortices in magnetic field, neutrino beaming and GRB}

\textbf{Vortex structures in the CFL phase of the strange star in the magnetic field.} Let us further assume that the bulk of the star is in the CFL phase with no hadronic shell and a thin crust generating a surface magnetic field which may achieve the values $B_s \sim 10^{12} - 10^{16}$ G. Then, one may expect by flux conservation that the inner magnetic field (at the crust-core interface) $B_{in,s} = B_s(n_{in}/n_s)^{2/3}$ can reach the values $\gtrsim 10^{14} - 10^{18}$ G, where $n_s$ and $n_{in}$ are the corresponding densities at the surface and in the core, respectively. Below we will vary $B_{in,s}$ within the interval $10^{15} \lesssim B_{in,s} \lesssim 10^{17}$ G. From the BCS relation the critical temperature of the superconductivity is $T_c = 0.57 \Delta(T = 0)$ and the temperature dependent gap can be parameterized as \cite{22} $\Delta(T) \simeq \Delta(T = 0)[1 - (T/T_c)^{3.4}]^{0.53}$. For initial temperatures $T_0 \simeq 30 - 60$ MeV and the gap $\Delta(T = 0) \gtrsim 100$ MeV we have $T_0 < T_c$. Thus the
strange star is formed in the superconducting phase. Refs. [20, 27, 28, 29] have shown, that typical values of the Ginzburg-Landau parameter for color superconductors are \( \kappa_{CFL} = \lambda_\xi \approx 1 \). The coherence length \( \xi_{CFL} \) and penetration depth \( \lambda_{CFL} \) for CFL matter can be estimated in the weak coupling limit as

\[
\xi_{CFL} \simeq 0.3 \left( \frac{100 \text{ MeV}}{T_c} \right) \left( 1 - \left( \frac{T}{T_c} \right)^{3.4} \right)^{-0.53} \text{ fm}, \quad \lambda_{CFL} \simeq 2 \left( \frac{3\sqrt{2}}{\sqrt{3} g^2} \right) \left( \frac{300 \text{ MeV}}{\mu_q} \right) \left( 1 - \left( \frac{T}{T_c} \right)^{3.4} \right)^{-0.53} \text{ fm},
\]

where we used the above \( \Delta(T) \) dependence. In our case the strong coupling constant \( \alpha_s = g^2/4\pi \approx 1 \), and the quark chemical potential is \( \mu_q \geq 350 \text{ MeV} \). The critical Ginzburg-Landau parameter is \( \kappa_{CFL} \simeq 1/\sqrt{2} \) distinguishes between type I (\( \kappa_{GL} < \kappa_{CFL} \)) and type II (\( \kappa_{GL} > \kappa_{CFL} \)) superconductors. Then for \( T_c > T_c^{1-II} \approx 15 \div 20 \text{ MeV} \) we deal with a type II superconductor. The latter entails the existence of a mixed phase (normal vortices embedded in superconducting matter) for \( T < T_c \) in a broad interval of magnetic fields \( B_{c1} < B < B_{c2} \). The value \( B_{c2} \) is \( \approx B_c \kappa_{GL} \), with \( B_c \gtrsim 10^{18} \text{ G} \) for \( \Delta \gtrsim 100 \text{ MeV} \). Actually, even if \( B < B_{c1} = B_c/\kappa_{GL} \), the magnetic field, in spite of the Meissner effect, cannot be expelled from the star interior within a pulsar lifetime [30], thus being concentrated in vortices aligned along the magnetic axis (feasibly being parallel to the rotation axis). Therefore, the mixed state exists for all \( B < B_{c2} \). Making below only rough estimates we will assume the simplest Abrikosov vortex structure of the mixed phase.

Due to magnetic flux conservation the number of vortices in the strange star interior is \( N_{vo} = \pi B_{in,s} R^2/\Phi_q \), where \( \Phi_q = 6\Phi_0 \) with the magnetic flux quantum \( \Phi_0 = 2 \cdot 10^{-7} \text{ G cm}^2 \) and the star radius \( R \approx 10 \text{ km} \) [20, 27, 28]. Superconductivity is expelled from the vortex interior, \( r \ll \xi \), where the matter is in the state of a strongly magnetized quark-gluon plasma. The magnetic field decreases on the scale of the penetration depth \( \lambda_{CFL} \) as \( B \sim B_0 \exp(-r/\lambda_{CFL}) \), for \( r > \lambda_{CFL} \), cf. [31, 26, 27, 28]. Typical values of the magnetic field in the vortex center are \( B_0 \sim \Phi_q/(\pi \lambda_{CFL}^2) \sim 10^{18} - 10^{19} \text{ G} \), depending on \( \lambda_{CFL} \).

In hadronic matter and in metals, the pairing gap is rather small and the coherence length \( \xi \propto 1/\Delta(T) \) is larger than the Debye screening length. Therefore, charge neutrality is easily fulfilled for \( r \ll \xi \). In the quark matter case we have opposite situation [32, 33]. For both, quarks and electrons, \( \lambda_D = \nu \lambda_{CFL} \) with typical values of \( \nu \sim 3 - 5 \). The presence of the Coulomb field diminishes the difference between the chemical potentials of quarks of different species since one does not need to fulfill the charge neutrality condition locally at distances \( r \ll \lambda_D \) but rather globally at \( r \gtrsim \lambda_D \). This implies that the cylindrical volume within distances \( r < \lambda_D \) from the vortex center is possibly filled by a 2SC+s phase plus a normal phase [34], also characterized by a still rather high magnetic field \( B \sim 10^{17} \text{ G} \). In case of the two-flavor superconducting (2SC) phase quarks of two colors, e.g., green and red, are paired with a large gap, whereas blue quarks are unpaired or paired with only a small gap \( \lesssim 1 \text{ MeV} \) [23]. The 2SC+s phase emits neutrinos with a still high rate of about 1/3 of the emissivity of normal quark matter. In this phase at densities under consideration typical values of the electron fraction \( Y_e = n_e/n_b \), where \( n_e \) and \( n_b \) are electron and baryon densities, are \( Y_e^{vo} \sim 10^{-2} \) (for \( r > \lambda_D \)). The electron fraction in the CFL phase at \( T = 0 \) vanishes [23], but at finite temperatures \( Y_e^{CFL}(T) \neq 0 \) due to thermal \( ee^+ \) excitations.

The typical distance between vortices is \( d = (\Phi_q/B_{in,s})^{1/2} \). The region near the magnetic vortex, where the magnetic field decreases by up to two orders of magnitude compared to the central value \( B_0 \), has the volume \( V_{vo} \sim 2\pi R (\nu \lambda_{CFL})^2 \sim 10^{21} \text{ fm}^3 \). The fraction of the total vortex volume \( N_{vo} V_{vo}/V \) (\( N_{vo} \) is the total number of vortices, \( V \) is the star volume) varies from several \( 10^{-7} \) to several \( 10^{-3} \) for \( B_{in,s} \) from \( \sim 10^{13} \text{ G} \) to \( 10^{17} \text{ G} \), respectively. Therefore the condition \( d \gg \lambda_D \) is safely fulfilled.

Goldstone transport. The CFL (and mCFL) phase is also characterized by the presence of (almost) massless weakly interacting Goldstone excitations (like phonons). For \( T \gg T_G^{coup} \approx 2 - 3 \text{ MeV} \) of our interest the mean free path of Goldstones is much shorter than the star radius \( \sim 10 \text{ km} \), cf. [35], and the typical time of their transport to the surface is very large. In the presence of a strong magnetic field generating vortices Goldstones cannot pass from CFL matter to the 2SC+s matter and
to the normal quark matter phase of the vortex core, where they cannot exist. Thus, the mean free path of Goldstone excitations $\lambda_{GB}^{\nu}$ in presence of vortices is still much shorter than in a homogeneous CFL strange star matter. For magnetic fields of our interest one finds $\lambda_{GB}^{\nu} \sim d^2/\xi \approx 6 \times (10^2 - 10^6)$ fm. Thus Goldstones may efficiently equilibrate the temperature between the vortex interior and exterior regions.

**Neutrino radiation and beaming.** The emissivity of the quark direct Urca (QDU) process in normal quark matter in absence of the magnetic field is very high \[36\]

$$\epsilon^{QDU}_\nu \simeq 2 \times 10^{26} \alpha_s u Y_e^{1/3} T_9^6 \text{ erg cm}^{-3} \text{ s}^{-1}. \quad (2)$$

Correspondingly, the neutrino mean free path in normal quark matter becomes very short

$$\lambda_{QDU}^{\nu} = 3 \times 10^6 \alpha_s^{-1} u^{-1} Y_e^{-1/3} T_9^{-2} \text{ cm}. \quad (3)$$

For $T \ll p_{Fe}$ one $Y_e = p_{Fe}/p_{Fb}$, where $p_{Fi} = (3\pi^2 n_i)^{1/3}$, $i = e, p$ and $u = n_b/n_0$ is the compression.

In superfluid matter the suppression factor of the emissivity is roughly \[37\]

$$\zeta^{DU}_\nu \simeq \exp(-\Delta(T)/T), \quad T < T_c, \quad (4)$$

and the enhancement factor of the mean free path is then $\zeta^{DU}_\nu$. In the CFL phase $p_{Fe} = 0$ but there exist thermal $ee^\pm$ excitations. Then $p_{Fe}$ should be replaced by $T$ and we may use above expressions (including appropriate suppression and enhancement factors) with $Y_e^{CFL} \simeq T^3/(3\pi^2)n_b$. The emissivity remains virtually unaltered in the presence of a magnetic field when many Landau levels are occupied \[38\], and it is increased by 10 – 100 times if only the first Landau level is filled. However, the latter occurs only for $B \gtrsim 10^{18}$ G. With the field values which we have assumed this condition is satisfied only in the region of the vortex cores, $r \lesssim \xi$, of a tiny volume. In the remaining part of the vortex volume $V_{vo}$ ($\xi < r < \nu \lambda_D$, $\xi \ll \nu \lambda_D$) one can roughly neglect the influence of the magnetic field on eq. (2). Also multiplying (2) by a factor $1/3 \leq s < 1$ we may estimate the emissivity of the 2SC+s+normal phase since due to the pairing in the 2SC phase the contribution of the red and green quarks is exponentially suppressed according \[41\]. Also in the region between the vortices occupied by CFL phase with the volume $V - N_{vo}V_{vo}$, the emissivity of the direct Urca process is suppressed by \[41\]. The neutrino mean free path increased by this factor exceeds the radius of the star for temperatures below $T_{opac}^{CFL} \sim 20$ MeV (for $\Delta(T = 0) \sim 150$ MeV), whereas for the normal quark matter one gets $T_{opac}^{norm} \lesssim$ few MeV, cf. \[41\]. For higher gaps the value $T_{opac}^{CFL}$ still increases.

For $T > T_{opac}^{CFL}$ the heat transport is governed by the transport equation, $C_V^\nu \dot{T} = \kappa \Delta T$, $C_V$ is the total specific heat of the matter and $\kappa = \kappa_\nu + \kappa_G$ is the total heat conductivity. Since Goldstones (phonons) are efficiently captured by vortices one has $\kappa_\nu \gg \kappa_G$, and thus $\kappa \approx \kappa_\nu$. A typical time scale for the energy transport to the surface is $t_{tr} \sim R^2/C_V/\kappa$. Since $\kappa_\nu \sim C_V \lambda_\nu$, the energy transport time is $t_{tr} \sim R^2/(\lambda_\nu)$. Using the above value $\lambda_{QDU}^{\nu}$ for $T \sim 20 - 40$ MeV, we get $t_{tr} \sim 10^{-4} - 10^{-1}$ s (estimate is done for $\Delta(T = 0) \sim 150$ MeV). For that time the star essentially loses its initial thermal energy. If initially (for $T_0 = 40$ MeV, $\Delta \gtrsim 100$ MeV) the thermal energy is mainly concentrated in massless excitations, the energy falls by a factor $2^4$ for $T \sim T_{opac}^{CFL} \sim 20$ MeV. After the transport time $t_{tr}$ the CFL regions become transparent for neutrinos.

For $T < T_{opac}^{CFL}$, the prominent difference between the neutrino mean free path in the 2SC+s+normal matter on the one hand and the CFL quark matter on the other hand leads to anisotropic neutrino emission via the direct (no rescatterings) QDU reaction from the star within a cone of the temperature-dependent opening angle $\theta_\nu(T) \sim \lambda_\nu(T)/R$ around the magnetic axis, where $\lambda_\nu(T) = \lambda_\nu(T) V/(N_{vo}V_{vo})$. Thus after $t > t_{tr}$ the star begins to radiate the remaining energy $\sim 10^{52}$ erg within a narrow beam cone.
Cooling evolution. Numerical results. The cooling evolution for $t > t_{tr}$ within the opening angle can be described by inverting the solution of

$$t = - \int_{T_{CFL}^{\text{spec}}}^{T(t)} dT' \frac{C_V(T')}{L(T')}$$

(5)

where

$$C_V(T) = \left\{ [1 - N_{\nu_o}V_{\nu_o}/V] \zeta_{DU} + sN_{\nu_o}V_{\nu_o}/V \right\} C_V^0(T) + C_V^{\text{ex}}(T),$$

(6)

$C_V^0(T) \simeq 10^{39} u^{2/3}(R/10\text{km})^3 T_9 \text{ erg/K}$ is the specific heat of the normal quark matter and $C_V^{\text{ex}}(T) \simeq 3 \cdot 10^{32} g^{\text{ex}}(R/10\text{km})^3 T_9^3 \text{ erg/K}$ is the contribution of (almost) massless Goldstone excitations and electrons and positrons (for the effective value of the degeneracy factor $g^{\text{ex}} \simeq 11$, we have assumed $v_G \simeq 1/\sqrt{3}$ for the velocity of Goldstones, the maximum Fermi velocity for relativistic fermi liquids. If $v_G$ were $\simeq 1$, we would get $g^{\text{ex}} \simeq 3$). The luminosity is

$$L(T) = \left[ 1 - \cos \theta_\nu(T) \right] L_0(T) \simeq L_0(T) \theta_\nu^2(T)/2$$

(7)

that takes into account the effect of the neutrino beaming. Here the isotropic luminosity is

$$L_0(T) = N_{\nu_o}V_{\nu_o} s \epsilon_{\nu}^{QDU}(T, Y_{e^{\nu}}^{\nu}) + (V - N_{\nu_o}V_{\nu_o}) \epsilon_{\nu}^{QDU}(T, Y_{e}^{\text{CFL}}(T)) \zeta_{DU}.$$  

(8)

We take the electron fraction $Y_{e^{\nu}} \simeq 10^{-2}$ for the 2SC+s + normal phase and a thermally produced electron fraction $Y_{e}^{\text{CFL}}(T)$ for the CFL phase.

Here the Goldstones play an essential role since they rapidly (in a microscopic time) equilibrate the temperature between the vortices (of volume $N_{\nu_o}V_{\nu_o}$) and their surrounding. In Eq. (5) we neglected this very short time assuming an instantaneous heat transport from the hotter CFL regions of vortex exteriors to the regions of the vortex volume which then cool down by the direct neutrino radiation within the beaming angle.

Results for the time evolution of neutrino luminosity, temperature and the neutrino beaming angle are shown in Fig. [4]. They demonstrate the effect of the cooling delay due to neutrino beaming in a strong magnetic field. For gaps $\sim 100 - 200$ MeV and $B_{in,s} \sim 10^{15} \text{G}$ we get $\theta_\nu \sim 10$ grad and the energy $\sim 10^{32}$ erg that remained in the star after the neutrino transport era has passed is radiated in the beam during several $\sim 10^{-2} - 10^{-1}$ s. Decreasing further the magnetic field one would get increase of the opening angle $\theta_\nu$ angle finally yielding inappropriate smearing of the beaming. For $B_{in,s} \sim 10^{17} \text{G}$ we obtain a narrow beam, $\theta_\nu \sim 0.1 - 1$ grad, and the initial star energy is radiated in the beam during several seconds.

Neutrino energy conversion to photons, GRB. The neutrino/antineutrino collisions produce $e^+e^-$ pairs outside the star, which efficiently convert to photons. We use the result of [5] for the $\nu\bar{\nu} \to e^+e^-$ conversion rate:

$$\dot{E}_{e^+} \sim 5 \cdot 10^{32} T_9^2(R/10\text{km})^3 \text{ erg s}^{-1}$$

(9)

to estimate that for a surface temperature $T_s \sim 10 - 30$ MeV most of the initial thermal energy can be converted in $e^+$ for an appropriate short time step $t \sim 10^2 - 10^{-2}$ s. The advantage of the strange star model is that the strange star has a thin hadron shell and a tiny crust, if any. In this case we may assume that the surface temperature is of the same order of magnitude as the internal temperature. Note that in the presence of beaming already $E_\gamma \sim$ several $\cdot 10^{51}$ erg produced on a time scale $t \sim 10^{-3} - 10^3$ s could be sufficient to explain GRB. Thus, above estimates are very optimistic for the GRB model dealing with a magnetized quark star, if $T_s \gtrsim 10$ MeV.

Moreover, $e^+$ pairs are accelerated in the strong magnetic field of the strange star exterior, increasing the hard component of the X ray spectrum, in coincidence with experimental findings.
3 Strange star kicks from parity violation

Range of kick velocities. Parity non-conservation in the weak interaction neutrino induced direct processes in presence of a magnetic field leads to a violation of reflection symmetry since the neutrino flux is a polar vector while the magnetic field is an axial one. The magnitude of this asymmetry has been estimated considering different reactions in neutron stars as the modified Urca reaction \[39\], see also \[40\]; \(\beta\) decays on the pion condensate \[41\]; formation and breaking of nucleon pairs \[42\]; coherent neutrino electron scattering on nuclei \[43\]; neutrino-polarized neutron elastic scattering; polarized electron capture \[44\]; the \(\nu \rightarrow \nu ee^+\) process \[45\]; the reaction \(\nu \bar{\nu} \rightarrow ee^+\); etc. Some of these reactions produce an asymmetry factor up to \(A_\nu \sim 10^{-4} B_{14}\), where \(B_{14} = B/10^{14}\) G. In our case of the strange star in the CFL phase with a strong magnetic field the direct QDU process and formation and breaking of quark pairs are relevant as well as the reaction \(\nu \bar{\nu} \rightarrow ee^+\) near the star surface. Here we assume that the QDU processes for \(T < T_{\text{CFL}}\) within the beaming angle (in accordance with the no go theorem \[43, 46\]) produce an asymmetry factor of the same order of magnitude as has been estimated in the literature. This \(A_\nu\) factor then leads to a net momentum transfer from the neutrino flux to the star of mass \(M\) resulting in a time-dependent kick velocity

\[
v(t) = \frac{A_\nu}{M} \int_{t(T_{\text{CFL}})}^{t} dt' L(T),
\]

which saturates at the magnetic field dependent asymptotic value

\[
v \simeq 1.7 \cdot B_{14} (M/M_\odot)^{-1} \text{ km s}^{-1}, \quad \text{for} \quad A_\nu \simeq 10^{-4} B_{14},
\]

as soon as the beaming ceases. Numerical results are shown in Fig. 2 for a typical strange star with \(M = 1.4 M_\odot\) and \(R = 10\) km. In order to obtain large star kicks with \(v \sim 10^3\) km s\(^{-1}\), magnetic fields \(B_{\text{in,s}} \sim 10^{17}\) G and large gaps \(\Delta \sim 200\) MeV are required.
4 Conclusions

We have shown that in the presence of a strong magnetic field $B_{\text{in,s}} \sim 10^{15} - 10^{17}$ G, the early cooling evolution of strange stars with a color superconducting quark matter core in the CFL phase right after a short transport era is characterized by anisotropic neutrino emission, collimated within a beaming angle $\theta_\nu$ around the magnetic axis. Initial temperatures of the order of $T_0 \sim 30 - 60$ MeV allow for the energy release $\sim 10^{53}$ erg. An energy $\sim 10^{52}$ erg is released within a narrow beam on time scales from $10^{-2}$ s to 10 s, appropriate for the phenomenology of both short and long GRB. Exploring the beaming mechanism one may obtain a wide range of strange star kick velocities up to $10^3$ km/s, in dependence on the magnetic field, pairing gap, radius, mass and burst duration of the star. The range of strange star kick velocities is thereby in a qualitative agreement with recent observational data on pulsar kick velocity distribution. All our estimates are very rough and essentially vary with the parameters of the model. Thus we just have shown a principal possibility of the application of the model to the GRB and kick description. A more systematic investigation, together with a consistent modeling of the compact star structure will be developed along the lines described in this contribution [47].

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