

Magnetic field instability in a neutron star driven by the electroweak electron-nucleon interaction versus the chiral magnetic effect

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We show that the Standard Model electroweak interaction of ultrarelativistic electrons with nucleons (the eN interaction) in a neutron star (NS) permeated by a seed large-scale helical magnetic field provides its growth up to $\gtrsim 10^{15}$ G during a time comparable with the ages of young magnetars $\sim 10^4$ yr. The magnetic field instability originates from the parity violation in the eN interaction entering the generalized Dirac equation for right and left massless electrons in an external uniform magnetic field. We calculate the average electric current given by the solution of the modified Dirac equation containing an extra current for right and left electrons (positrons), which turns out to be directed along the magnetic field. Such a current includes both a changing chiral imbalance of electrons and the eN potential given by a constant neutron density in a NS. Then we derive the system of the kinetic equations for the chiral imbalance and the magnetic helicity which accounts for the eN interaction. By solving this system, we show that a sizable chiral imbalance arising in a neutron protostar due to the Urca process $e_{\bar{L}} + p \rightarrow N + \nu_{eL}$ diminishes very rapidly because of a huge chirality-flip rate. Thus the eN term prevails over the chiral effect, providing a huge growth of the magnetic helicity and the helical magnetic field.

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Some neutron stars, called magnetars, having magnetic fields $B \sim 10^{15} - 10^{16}$ G, can be considered the strongest magnets in our Universe [1]. Despite the existence of various models for the generation of such strong fields, based, e.g., on the turbulent dynamo [2], the origin of magnetic fields in magnetars is still an open problem. Recently, in Ref. [3] the authors tried to apply the chiral magnetic effect [4,5], adapted successfully for the QCD plasma [6], to tackle the problem of magnetic fields in magnetars. The approach of Ref. [3] implies the chiral kinetic theory, where the Vlasov equation is modified when adding the Berry curvature term to the Lorentz force [7].

The fate of such a chiral plasma instability is based on the Adler anomaly in QED with the nonconservation of the pseudovector current for massless fermions $\bar{\psi}\gamma_{\mu}\gamma_5\psi$ in external electromagnetic fields. Since this current is the difference of right j_{μ}^R and left j_{μ}^L currents, the assumption of a seed imbalance between their densities given by the difference of chemical potentials, $(n_R - n_L) \sim \mu_5 = (\mu_R - \mu_L)/2 \neq 0$, where $n_{R,L}$ are the densities of right and left fermions (electrons) and $\mu_{R,L}$ are their chemical potentials, could lead to the magnetic field instability we study here adding electroweak interactions in the Standard Model (SM).

The same effect (while without weak interactions) was used in Ref. [8] to study the self-consistent evolution of

the magnetic helicity in the hot plasma of the early Universe driven by the change of the lepton asymmetry $\sim \mu_5$. In Ref. [8] it was shown that such an asymmetry diminishes, $\mu_5 \rightarrow 0$, due to the growth of the chirality-flip rate in the cooling Universe through electron-electron (ee) collisions, $\Gamma_f \sim \alpha_{\text{em}}^2 (m_e/3T)^2$, where $\alpha_{\text{em}} = e^2/4\pi \approx 1/137$ is the fine structure constant, m_e is the electron mass, and T is the plasma temperature.

This negative result encouraged the appearance of Ref. [9], where another mechanism for the generation of magnetic fields was proposed. It is based on the parity violation in electroweak plasma resulting in the nonzero Chern-Simons (CS) term Π_2 that enters the antisymmetric part of the photon polarization operator in the plasma of massless particles. Here we adopt the notation for the CS term from Ref. [9]. In Ref. [10], a similar CS term $\Pi_2^{(\nu l)}$, based on the neutrino interactions with charged leptons, was calculated. Based on this calculation, the magnetic field instability driven by neutrino asymmetries was revealed. This instability is implemented in different media such as the hot plasma of the early Universe and a supernova (SN) with a seed magnetic field.

The amplification of a seed magnetic field during the SN burst driven by a nonzero electron neutrino asymmetry $\Delta n_{\nu_e} \neq 0$ which enters the CS term $\Pi_2^{(\nu e)}$ was suggested in Ref. [10] to explain the generation of the strongest magnetic fields in magnetars. Note that after the SN burst, a cooling neutron star (NS) as the

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corresponding SN remnant equally emits neutrinos and antineutrinos. Thus, the neutrino asymmetry vanishes. The inclusion of the electroweak ee interaction with a stable fraction of degenerate electrons $n_e \approx \text{const}$ instead of the νe interaction with vanishing neutrino asymmetry $\Delta n_{\nu_e} \rightarrow 0$ makes no sense since the corresponding parity violating CS term $\Pi_2^{(ee)}$ tends to zero in the static limit $\omega \rightarrow 0$ for an electron gas, $\Pi_2^{(ee)} \rightarrow 0$, as found in Ref. [11].

In the present work we suggest taking into account the electroweak electron-nucleon (eN) interaction providing a longtime acting source of the magnetic field instability that plays a role of a CS term in the pseudovector electron current $\mathbf{J}_5 = \Pi_2^{(eN)} \mathbf{B}$. Instead of the Matsubara technique used in Refs. [10,11], here we calculate the total electric current in SM (as an additive to the standard ohmic current) solving the Dirac equation for the massless right and left electrons (positrons) in a magnetic field.

We start the derivation of the aforementioned CS term by solving the Dirac equation for a massless electron in the magnetic field $\mathbf{B} = (0, 0, B)$ accounting for the electroweak eN interaction in a NS. This equation reads

$$[\gamma^\mu (i\partial_\mu + eA_\mu) - \gamma^0 (V_L P_L + V_R P_R)] \psi_e = 0, \quad (1)$$

where $\gamma^\mu = (\gamma^0, \boldsymbol{\gamma})$ are the Dirac matrices, $A^\mu = (0, 0, Bx, 0)$ is the vector potential, $P_{L,R} = (1 \mp \gamma^5)/2$ are the chiral projection operators, $\gamma^5 = i\gamma^0 \gamma^1 \gamma^2 \gamma^3$, and $e > 0$ is the absolute value of the electron charge.

In Eq. (1) we assume that there are no macroscopic fluid (nucleon) currents in a NS. The effective potentials $V_{L,R}$ in Eq. (1) are given by the SM Lagrangian of the eN interaction via neutral currents in the Fermi approximation (see, e.g., Ref. [12]),

$$\begin{aligned} \mathcal{L} = & \sqrt{2} G_F \bar{\psi}_e \gamma_\mu (g_L^{(e)} P_L + g_R^{(e)} P_R) \psi_e \\ & \times [\bar{\psi}_n \gamma^\mu \psi_n - (1 - 4\xi) \bar{\psi}_p \gamma^\mu \psi_p], \end{aligned} \quad (2)$$

where $G_F \approx 1.17 \times 10^{-5} \text{ GeV}^{-2}$ is the Fermi constant $g_L^{(e)} = -1/2 + \xi$ and $g_R^{(e)} = \xi$ are the standard coupling constants in SM with the Weinberg parameter $\xi = \sin^2 \theta_W \approx 0.23$, and $\psi_{n,p}$ are the neutron and proton wave functions. We reduced the total eN Lagrangian in Ref. [12] to Eq. (2), omitting the axial nucleon currents $\sim \bar{\psi}_{n,p} \gamma^\mu \gamma^5 \psi_{n,p}$ that are irrelevant to our problem.

Taking the statistical averaging $\langle \dots \rangle$ in Eq. (2) over the equilibrium (Fermi) distributions of nucleons in a neutron star and recalling that macroscopic nucleon currents are absent, i.e., $\langle \bar{\psi}_{n,p} \boldsymbol{\gamma} \psi_{n,p} \rangle = 0$, we get the following definition of $V_{R,L}$ to be used in Eq. (1):

$$\begin{aligned} V_L &= -\frac{G_F}{\sqrt{2}} [n_n - n_p (1 - 4\xi)] (2\xi - 1), \\ V_R &= -\frac{G_F}{\sqrt{2}} [n_n - n_p (1 - 4\xi)] 2\xi, \end{aligned} \quad (3)$$

where $n_{n,p} = \langle \psi_{n,p}^\dagger \psi_{n,p} \rangle$ are the number densities of neutrons and protons.

Let us decompose ψ_e in the chiral projections as $\psi_e = \psi_L + \psi_R$, where $\psi_{L,R} = P_{L,R} \psi_e$. Then, using Eq. (1) we get that $\psi_{L,R} = e^{-iE_{L,R}t + ip_y y + ip_z z} \psi_{L,R}(x)$, where

$$\begin{aligned} \psi_{L,R}^{(n)}(x) &= \frac{1}{4\pi \sqrt{E_{L,R} - V_{L,R}}} \begin{pmatrix} \sqrt{E_{L,R} - V_{L,R} \mp p_z} u_{n-1} \\ \mp i \sqrt{E_{L,R} - V_{L,R} \pm p_z} u_n \\ \mp \sqrt{E_{L,R} - V_{L,R} \mp p_z} u_{n-1} \\ i \sqrt{E_{L,R} - V_{L,R} \pm p_z} u_n \end{pmatrix}, \\ \psi_{L,R}^{(0)}(x) &= \frac{1}{2\pi \sqrt{2}} \begin{pmatrix} 0 \\ u_0 \\ 0 \\ \mp u_0 \end{pmatrix}. \end{aligned} \quad (4)$$

Here $\psi_{L,R}^{(n)}$ corresponds to $n = 1, 2, \dots$, $\psi_{L,R}^{(0)}$ to $n = 0$, $\eta = \sqrt{eB}x + p_y/\sqrt{eB}$, $u_n(\eta) = (eB/\pi)^{1/4} \exp(-\eta^2/2) \times H_n(\eta)/\sqrt{2^n n!}$, and $H_n(\eta)$ is the Hermite polynomial. The upper signs in Eq. (4) stay for ψ_L and the lower ones for ψ_R . To derive Eq. (4) we use the γ matrices in the Dirac representation as in Ref. [13]. The energy levels $E_{L,R}$ in Eq. (4) can be obtained from the following expression:

$$(E_{L,R} - V_{L,R})^2 = p_z^2 + 2eBn. \quad (5)$$

The normalization factors in Eq. (4) correspond to

$$\int (\psi_{L,R})_{np,p_z}^\dagger (\psi_{L,R})_{n'p'_z} d^3x = \delta_{nn'} \delta(p_y - p'_y) \delta(p_z - p'_z), \quad (6)$$

since the chiral projections $\psi_{L,R}$ are independent. It is worth mentioning that a more general solution of Eq. (1), which accounts for the nonzero electron mass, was found in Ref. [14].

The spinors in Eq. (4) are then used to calculate the induced electric current which has a nonzero projection on the z axis $\sim \bar{\psi}_e \boldsymbol{\gamma}^3 \psi_e$. Analogously to Ref. [4], one shows that the averaged current gets the contribution from the main Landau level $n = 0$ only. It should be noted that massless particles have a strong correlation between their momentum and helicity. Thus, at $n = 0$, left electrons have $p_z > 0$, whereas right ones have $p_z < 0$.

Making the statistical average with the Fermi-Dirac distribution of the left and right electrons (positrons)

$f_{e,\bar{e}}(E) = [\exp(\beta(E \mp \mu_{L,R}) + 1)]^{-1}$, where $\beta = 1/T$ is the reciprocal temperature, $\mu_{L,R}$ are their chemical potentials, and the lower sign stays for positrons, then using Eq. (5), one obtains the component of the current J_z along the magnetic field in the form

$$J_z = \frac{e^2 B}{4\pi^2} \left\{ \int_{-\infty}^0 dp_z [f_e(-p_z + V_R) - f_{\bar{e}}(-p_z - V_R)] - \int_0^{+\infty} dp_z [f_e(p_z + V_L) - f_{\bar{e}}(p_z - V_L)] \right\}. \quad (7)$$

Basing on Eq. (7) and introducing vector notations, we derive the average induced current in the final form as

$$\mathbf{J} = \frac{2\alpha_{\text{em}}}{\pi} (\mu_5 + V_5) \mathbf{B}, \quad (8)$$

which is an additive to the ohmic current \mathbf{J}_{Ohm} in a standard QED plasma. It should be noted that Eq. (8) is valid for any electron temperature.

The current in Eq. (8) is proportional to α_{em} and consists of two parts: the vector term given in QED by the *pseudoscalar* coefficient $\mu_5 = (\mu_R - \mu_L)/2$ ($\mu_5 \rightarrow -\mu_5$ under spatial inversion) and the pseudovector current $\mathbf{J}_5 = (2\alpha_{\text{em}}/\pi)V_5\mathbf{B} = \Pi_2^{(eN)}\mathbf{B}$ given in SM by the *scalar* factor $V_5 = (V_L - V_R)/2$. Indeed, after statistical averaging, the interaction Lagrangian in Eq. (2) becomes

$$\mathcal{L} = \frac{1}{2}(V_L + V_R)\bar{\psi}_e\gamma_0\psi_e + \frac{1}{2}(V_R - V_L)\bar{\psi}_e\gamma_0\gamma_5\psi_e. \quad (9)$$

The factor $\bar{\psi}_e\gamma_0\gamma_5\psi_e$ in the parity violation term of Eq. (9) is the pseudoscalar with respect to the spatial inversion $P = P^{-1} = \gamma_0$ since $P\gamma_0\gamma_5P^{-1} = -\gamma_0\gamma_5$. Thus V_5 should be scalar; cf. Ref. [15]. The true pseudoscalar for both P inversion and Lorentz transformation should be $\bar{\psi}\gamma_5\psi$. It should be noted that one loses Lorentz invariance in a medium with a selected reference frame like a NS at rest.

The weak interaction coefficient in Eq. (8),

$$V_5 = \frac{G_F}{2\sqrt{2}} [n_n - (1 - 4\xi)n_p], \quad (10)$$

is of the order $V_5 \approx G_F n_n / 2\sqrt{2} = 6 \text{ eV}$ in a NS with $n_n = 1.8 \times 10^{38} \text{ cm}^{-3}$, which corresponds to $\rho_n = 3 \times 10^{14} \text{ g} \cdot \text{cm}^{-3}$ since $n_p \ll n_n$. At first glance, the electromagnetic QED term in the current in Eq. (8), $\sim \mu_5$, seems to be much bigger than the weak one in Eq. (10) [16]. However, we show below that the latter term remains a stable source of the magnetic field instability in a NS while the former one vanishes, $\mu_5 \rightarrow 0$, e.g., for helical magnetic fields with the maximum helicity, contrary to the statement in Ref. [7] that an

imbalance $\mu_5 \neq 0$ could lead to the generation of strong magnetic fields in magnetars.

The evolution of the magnetic field in the presence of the induced current in Eq. (8), proportional to the magnetic field, obeys the modified Faraday equation; cf. Ref. [10]. However, it is more convenient to study the evolution of the magnetic helicity density

$$h(t) = \frac{1}{V} \int d^3x (\mathbf{A} \cdot \mathbf{B}), \quad (11)$$

where \mathbf{A} is the 3D vector potential and V is the volume of space. Defining the helicity density spectrum $h(k, t)$ as $h(t) = \int dk h(k, t)$ and accounting for the induced current in Eq. (8), which includes both the chiral imbalance contribution $\sim \mu_5$ and the electroweak term $\sim V_5$, we get the kinetic equation for $h(k, t)$ which is the generalization of Eq. (6) in Ref. [8],

$$\frac{\partial h(k, t)}{\partial t} = -\frac{2k^2 h(k, t)}{\sigma_{\text{cond}}} + \frac{\alpha_{\text{em}}}{\pi} \left[\frac{k(\Delta\mu + 2V_5)}{\sigma_{\text{cond}}} \right] h(k, t). \quad (12)$$

Here $\Delta\mu = \mu_R - \mu_L = 2\mu_5$ and we have just assumed, as in Ref. [8], the maximal helicity field configuration; i.e., the magnetic energy density reads $\rho_B(t) = \int dk \rho_B(k, t) = (1/2) \int dk k h(k, t)$. It is worth mentioning that the sign of the $\Delta\mu$ term in Eq. (12) is opposite to that in Ref. [8] since we use the different definition of γ^5 .

We should then derive the kinetic equation which governs the chiral imbalance evolution, which is complementary to Eq. (12). Using Eq. (11) and the Maxwell equations, we get the helicity density change in the standard form,

$$\frac{dh(t)}{dt} = -\frac{2}{V} \int d^3x (\mathbf{E} \cdot \mathbf{B}), \quad (13)$$

where \mathbf{E} is the electric field. Then, accounting for the Adler anomaly for the pseudovector current in the electromagnetic fields, $\partial_\mu (j_R^\mu - j_L^\mu) = \partial_\mu (\bar{\psi}\gamma^\mu\gamma^5\psi) = (2\alpha_{\text{em}}/\pi)(\mathbf{E} \cdot \mathbf{B})$, we derive the conservation law involving $h(t)$ and $n_{R,L}$,

$$\frac{d}{dt} \left[n_R - n_L + \frac{\alpha_{\text{em}}}{\pi} h(t) \right] = 0, \quad (14)$$

which is valid in a QED plasma.

Taking into account that $n_{L,R} = \mu_{L,R}^3/3\pi^2$ and assuming that $\mu_L \sim \mu_R \sim \mu$, where μ is the chemical potential of the degenerate electron gas in a NS, which is true at least at the beginning of the imbalance in a NS, we get that $n_R - n_L \approx 2\mu_5\mu^2/\pi^2$. Eventually we obtain from Eq. (14), using the expression for $\partial h(k, t)/\partial t$ in Eq. (12), the evolution equation for μ_5 ,

$$\frac{d\mu_5}{dt} = \frac{\pi\alpha_{\text{em}}}{\mu^2\sigma_{\text{cond}}} \int dk k^2 h(k, t) - \left[\frac{2\alpha_{\text{em}}^2 \rho_B(t)}{\mu^2\sigma_{\text{cond}}} \right] (\mu_5 + V_5) - \Gamma_f \mu_5. \quad (15)$$

In Eq. (15) we added the rate of the chirality-flip processes, $\Gamma_f \simeq (m_e/\mu)^2 \nu_{\text{coll}}$, given by the Rutherford electron-proton (ep) collision frequency $\nu_{\text{coll}} = \omega_p^2/\sigma_{\text{cond}}$ without flip. Here $\omega_p = \mu\sqrt{4\alpha_{\text{em}}/3\pi}$ is the plasma frequency in a degenerate ultrarelativistic electron gas and σ_{cond} is the electric conductivity in a degenerate electron-proton plasma consisting of ultrarelativistic degenerate electrons and non-relativistic degenerate protons [17]. Note that in a degenerate electron gas ν_{coll} depends on the temperature T ; cf. Ref. [18]. This is due to the Pauli principle when all electron states with the momenta $0 \leq p \leq \mu$ are busy; i.e., ep scattering is impossible at $T = 0$.

One can see that Eq. (15) is different from the simplified kinetic approach $d\mu_5/dt = \Gamma_{\text{inst}}\mu_5 - \Gamma_f\mu_5$, where $\Gamma_{\text{inst}} = \alpha_{\text{em}}^2\mu_5$ is a factor providing the magnetic field growth used in Refs. [3,7]. The first term on the rhs of Eq. (15) can really be estimated as $\sim\alpha_{\text{em}}^2\mu_5^2$ for all “equal” parameters $\mu \sim \mu_5 \sim \sigma_{\text{cond}}$, which is not the case we rely on. The more important difference is the appearance of the second term $\sim\rho_B$ that is the backreaction from the magnetic field that diminishes an imbalance μ_5 .

Let us choose the simplest case of the monochromatic helicity density spectrum $h(k, t) = h(t)\delta(k - k_0)$, where we can vary the wave number k_0 and the magnetic field scale $\Lambda_B = k_0^{-1}$ to later find some critical regimes for the imbalance evolution $\mu_5(t)$ through Eq. (15). Using the dimensionless functions $\mathcal{M}(\tau) = (\alpha_{\text{em}}/\pi k_0)\mu_5(t)$ and $\mathcal{H}(\tau) = (\alpha_{\text{em}}^2/2k_0\mu^2)h(t)$, which depend on the dimensionless diffusion time $\tau = (2k_0^2/\sigma_{\text{cond}})t$, we can recast the self-consistent system of Eqs. (12) and (15) as

$$\begin{aligned} \frac{d\mathcal{M}}{d\tau} &= (1 - \mathcal{M} - \mathcal{V})\mathcal{H} - \mathcal{G}\mathcal{M}, \\ \frac{d\mathcal{H}}{d\tau} &= - (1 - \mathcal{M} - \mathcal{V})\mathcal{H}. \end{aligned} \quad (16)$$

Here for fixed $V_5 = 6$ eV the dimensionless parameters $\mathcal{V} = (\alpha_{\text{em}}/\pi k_0)V_5$ and $\mathcal{G} = (\sigma_{\text{cond}}/2k_0^2)\Gamma_f = (2\alpha_{\text{em}}/3\pi)(m_e/k_0)^2$ are the function of the parameter k_0 only. Note that \mathcal{G} does not depend on the conductivity σ_{cond} since the rate of the chirality flip can be estimated as $\Gamma_f \simeq (m_e/\mu)^2 \nu_{\text{coll}}^{(\text{no flip})}$, where in the magnetohydrodynamic plasma $\nu_{\text{coll}}^{(\text{no flip})} = \omega_p^2/\sigma_{\text{cond}}$ is the ep collision frequency without flip. The dimensionless diffusion time τ depends on the conductivity found in Ref. [18]:

$$\sigma_{\text{cond}} = \frac{1.6 \times 10^{28}}{(T/10^8 \text{ K})^2} \left(\frac{n_e}{10^{36} \text{ cm}^{-3}} \right)^{3/2} \text{ s}^{-1}, \quad (17)$$

which is valid for cooling NS matter consisting of degenerate nonrelativistic nucleons and ultrarelativistic degenerate electrons.

For the magnetic field scale Λ_B comparable to the NS radius $R_{\text{NS}} = 10$ km, or for the small wave number $k_0 = 1/R_{\text{NS}} = 2 \times 10^{-11}$ eV, one gets the electroweak interaction contribution in Eq. (16) as $\mathcal{V} = 7 \times 10^8$ coming from the current in Eq. (8), where we substitute the small $V_5 = 6$ eV. We choose the initial chiral imbalance as $\mu_5(0) = 1$ MeV $\ll \mu$, where for $n_e = \mu^3/3\pi^2 = 10^{36} \text{ cm}^{-3}$ in Eq. (17) the electron chemical potential is equal to $\mu = 60$ MeV. Hence at the beginning the dimensionless chiral imbalance $\mathcal{M}(0) = (\alpha_{\text{em}}/\pi k_0)\mu_5(0) \simeq 10^{14}$ is much bigger than the electroweak term \mathcal{V} . At first glance, such an inequality could be expected when comparing electromagnetic and weak interaction effects, $\mathcal{M}(0) \gg \mathcal{V} = \text{const}$. We assume also the constant temperature in a cooling NS $T = 10^8$ K [19]. Therefore the electric conductivity in Eq. (17) is also constant, $\sigma_{\text{cond}} = 10^7$ MeV.

The dimensionless chirality-flip rate,

$$\mathcal{G} = \frac{2\alpha_{\text{em}}}{3\pi} \left(\frac{m_e}{k_0} \right)^2 = 10^{30}, \quad (18)$$

is huge for the given small $k_0 = 2 \times 10^{-14}$ keV. If we change $m_e = 511$ keV $\rightarrow m_{\text{eff}} = \mu\sqrt{\alpha_{\text{em}}/2\pi}$ [20], the rate in Eq. (18) would be even bigger, diminishing μ_5 faster in the first line of Eq. (16). Finally, for the acceptable initial magnetic field $B_0 = 10^{12}$ G, the initial helicity density $h(0) = B_0^2/k_0 = 2 \times 10^{13}$ MeV³ gives $\mathcal{H}(0) = (\alpha_{\text{em}}^2/2k_0\mu^2)h(0) = 6 \times 10^{21}$.

We solved the system of the self-consistent kinetic equations in Eq. (16) numerically for the adopted \mathcal{V} and \mathcal{G} as well as the initial conditions $\mathcal{M}(0) = 10^{14}$ and $\mathcal{H}(0) = 6 \times 10^{21}$ chosen above. In Fig. 1 we plot the evolution of the chiral imbalance $\mathcal{M}(\tau)$. In the inset, one can see how a large chirality imbalance $\mu_5 \sim \mathcal{O}(\text{MeV})$ vanishes owing to the huge chirality-flip rate in Eq. (18), $\mu_5 \rightarrow 0$, corresponding to $t \sim 10^{-12}$ s during a very short time $\tau \sim 10^{-30}$. In the main plot one finds a sharp slope for \mathcal{M} somewhere at $\tau \approx 3 \times 10^{-8}$ that corresponds to the time $t \sim 8000$ yr. The obtained critical time is of the order of young magnetar ages [1]. In Fig. 2 we see that, at the same time $\tau \approx 3 \times 10^{-8}$, the magnetic helicity density \mathcal{H} grows by about 10 orders of magnitude, which corresponds to the growth of $B = \sqrt{k_0 h}$ by 5 orders of magnitude, getting just $B \simeq 10^{17}$ G if we start from the seed field $B_0 = 10^{12}$ G.

It is interesting to mention that, in Fig. 1, a *positive* primeval chirality imbalance, $\mu_5 = (\mu_R - \mu_L)/2 > 0$, which appears, e.g., due to the direct Urca process, $e_L^- + p \rightarrow n + \nu_{eL}$, becomes *negative*, $\mu_R - \mu_L < 0$. This happens due to the simultaneous growth of the helicity density h (see Fig. 2), which amplifies the negative derivative

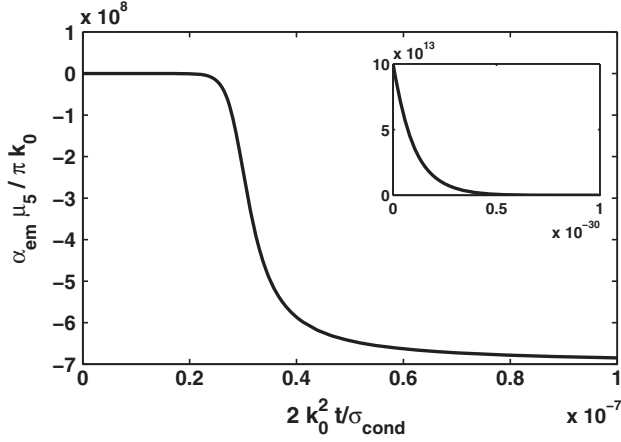


FIG. 1. The dimensionless chiral imbalance \mathcal{M} versus τ . The horizontal axis of the main plot starts at $\tau \gtrsim 10^{-30}$. (Inset) The evolution of \mathcal{M} in the initial time interval corresponding to $\tau < 10^{-30}$.

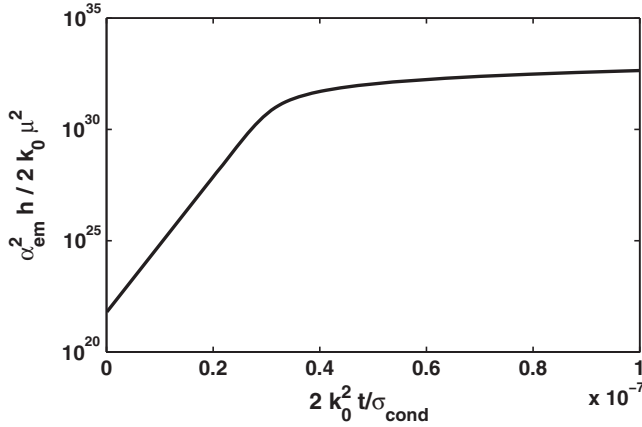


FIG. 2. The dimensionless helicity \mathcal{H} versus τ .

$d\mathcal{M}/d\tau < 0$ much more intensely than the chirality flip $\sim \mathcal{G}$. Vice versa, the attenuation of \mathcal{M} owing to the chirality flip is more important at the first stage illustrated in the inset of Fig. 1. Since $\mathcal{M} \rightarrow -\mathcal{V} = -7 \times 10^8$ ($\mu_5 \rightarrow -6$ eV; see Fig. 1), while the decreasing sum $\mathcal{V} + \mathcal{M}$ remains positive, the value of the positive derivative $d\mathcal{H}/d\tau > 0$ diminishes, or the helicity evolution simulates a saturation; see Eq. (16) and Fig. 2.

Finally we notice that helical magnetic fields determine the evolution of the chiral imbalance $\mu_5(t)$, rather than a nonzero seed $\mu_5 \neq 0$ can lead the growth of the magnetic helicity density $h = B^2/k_0$ or the magnetic field itself. This imbalance starting from a sizable value $\mu_5 \sim \mathcal{O}(\text{MeV})$ decreases down to the eN interaction term $|\mu_5| \sim V_5 \sim 6$ eV. We stress that, namely, the electro-weak interaction term $V_5 > |\mu_5|$ drives the amplification of the seed magnetic field in a NS, see the second line of Eq. (16). If one takes into account the cooling of a

neutron star, $dT/dt < 0$, a more realistic model to generate strong magnetic fields in magnetars can be developed. We plan to do that in our future work.

Of course, we considered here only the largest scale $k_0^{-1} = R_{\text{NS}} = 10$ km as the most interesting case for magnetic fields in NS. Our model is simplified both due to the choice of the maximum helicity density $kh(k, t) = 2\rho_B(k, t)$ instead of the more general inequality $kh(k, t) \leq 2\rho_B(k, t)$ [21] and owing to the choice of the monochromatic helicity density spectrum $h(k, t) = h(t)\delta(k - k_0)$. The generalization of our model, e.g., accounting for an initially nonhelical magnetic field, the continuous magnetic energy spectrum, complicates the problem. This requires solving the system of kinetic equations for the magnetic helicity density and magnetic energy density instead of the single Eq. (12) here. We also skipped the stage of a supernova collapse with nonequilibrium processes at that time, considering in our model mostly longtime intervals $\sim (10^3 - 10^4)$ yr for a thermally relaxed NS core.

We would like to mention that recently, in Ref. [22], the application of the chiral plasma instability in SN was also criticized because the chirality flip was underestimated in Ref. [3] in the approximation $(m_e/\mu)^2 \ll 1$. Instead of the tedious calculation made in Ref. [22], we can reproduce in a simpler way the flip rate Γ_f obtained by the authors in Ref. [3] and demonstrate why their derivation is invalid. Indeed, in Ref. [3] the authors incorrectly relied on the flip rate $\Gamma_f \sim \alpha_{\text{em}}^2 (m_e/\mu)^2 \mu_5$, meaning instead that the collision frequency without flip, entering the flip rate as $\Gamma_f = (m_e/\mu)^2 \nu_{\text{coll}}^{(\text{no flip})}$, is given by the common formula $\nu_{\text{coll}}^{(\text{no flip})} = \sigma n_e \approx (\alpha_{\text{em}}^2/\mu^2) \mu^3 = \alpha_{\text{em}}^2 \mu_5$. Here it was assumed that $\mu \sim \mu_5$, using the electron density $n_e \sim \mu^3$ and the Rutherford cross section for $e p$ collisions $\sigma \sim \alpha_{\text{em}}^2/\langle E \rangle^2$, where $\langle E \rangle \sim \mu$ is the mean electron energy. Such an estimate of the flip rate Γ_f is incorrect for a degenerate electron gas because the Pauli principle was not taken into account.

To resume, we have suggested here a novel mechanism for the magnetic field amplification in NS based on the eN electroweak interaction. For this purpose, in Eq. (8), we have generalized the CS term, derived in Ref. [4], to include the electroweak interaction of right and left ultrarelativistic degenerate electrons with nucleons. Then, in Eqs. (12) and (15), we have obtained a new system of kinetic equations for the evolution of the chiral imbalance and the magnetic helicity. This system accounts for the eN interaction and the backreaction. Finally we have applied our results to predict the magnetic field growth in a NS up to the values observed in magnetars.

It should be noted that our model is absolutely different from the well-known approach put forward in Ref. [2] based on a strong turbulent convection in the core of the SN and the fast dynamo operating only for a few

seconds, being driven by the high neutrino luminosity $L_\nu > 10^{52} \text{ erg} \cdot \text{s}^{-1}$ at that time. It should be noted that, in Ref. [23], it was found that protostars, which were progenitors to some magnetars, did not seem to reveal a fast rotation as required in Ref. [2]. We also refute the arguments in Ref. [3] that suggest the generation of strong magnetic fields in magnetars is based on chiral magnetic instability.

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