Antistars or antimatter cores in mirror neutron stars?

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The oscillation of the neutron $n$ into mirror neutron $n'$, its partner from dark mirror sector, can gradually transform an ordinary neutron star into a mixed star consisting in part of mirror dark matter. The implications of the reverse process taking place in the mirror neutron stars depend on the sign of baryon asymmetry in mirror sector. Namely, if it is negative, as predicted by certain baryogenesis scenarios, then $n' - \pi$ transitions create a core of our antimatter gravitationally trapped in the mirror star interior. The annihilation of accreted gas on such antimatter cores could explain the origin of our antimatter in the Universe (for reviews see e.g. 1–2).

A specific feature of this scenario is that any neutral particle (elementary as e.g. neutrinos or composite as e.g. the neutron) can have maximal mixing with its M partner, and such twin species can effectively oscillate between each other. Namely, the active–sterile $\nu - \nu'$ oscillations between the O and M neutrinos, violating the lepton number $L$ and $L'$ oscillations between the O and M neutrinos, violating the lepton number $L$ and $L'$ mixings of the O and M species (fermions, gauge bosons and Higgs fields of two sectors), implies that the two worlds should have identical microphysics. In this way, all O particles: electron $e$, proton $p$, neutron $n$, neutrinos $\nu$, photon $\gamma$ etc. must have the mass-degenerate M twins $e'$, $p'$, $n'$, $\nu'$, $\gamma'$ etc. which are sterile to the interactions of our Standard Model (SM) but have their own SM interactions of exactly the same pattern. Like O matter, during cosmological evolution the M matter should form nuclei, atoms, stars and planets which, being invisible for us in terms of O photons, can represent dark matter (DM) in the Universe (for reviews see e.g. 3–4).

A specific feature of this scenario is that any neutral particle (elementary as e.g. neutrinos or composite as e.g. the neutron) can have maximal mixing with its M partner, and such twin species can effectively oscillate between each other. Namely, the active–sterile $\nu - \nu'$ oscillations between the O and M neutrinos, violating the lepton numbers $L$ and $L'$ of both sectors by one unit, can be experimentally observed as the deficit of neutrinos. Analogously, active–sterile oscillation between the neutron $n$ and M neutron $n'$, induced by the mass mixing

$$\varepsilon \, n' n + h.c. \quad (1)$$

violates the baryon numbers $B$ and $B'$ by one unit but conserves the combination $B = B + B'$. This phenomenon can be tested via the neutron disappearance ($n \to n'$) and regeneration ($n \to n' \to n$) experiments 5–6.

As it was shown in Ref. 8, the phenomenological and cosmological bounds do not exclude $nn'$ mixing mass to be as large as $\varepsilon \sim 10^{-15}$ eV, which corresponds to few seconds for the characteristic oscillation time $\tau = 1/\varepsilon$. In fact, for free neutrons $n - n'$ oscillation is suppressed by the medium effects as the presence of matter and magnetic fields 8,9. But in deep cosmos it could proceed without suppression, with interesting implications e.g. for the cosmic rays at super-GZK energies, as far as the neutrons produced in $p\gamma \to n\pi^+$ reactions on relic photons can promptly oscillate into $n'$ and then decay in M sector, $n' \to p' \gamma^\prime$ 7,8. Several dedicated experiments searching for $n - n'$ oscillations still allow rather short oscillation times 9,10. Moreover, some of their results show deviations from null hypothesis 9,14,16 and new experiments are underway for testing these anomalies 17,18.

It is remarkable that the nuclear stability yields no limit on $\varepsilon$ since $n \to n'$ conversion in nuclei is forbidden by the energy conservation 3. However, in the neutron stars (NS) $n \to n'$ conversion in matter is energetically favored, and it can gradually transform the NS into mixed stars partially consisting of M matter. Since in super-dense nuclear medium $n - n'$ oscillation is strongly suppressed, for $\varepsilon \sim 10^{-15}$ eV the effective conversion time appears to be much larger than the universe age 8. Nevertheless, it can have observable effects for the NS which were analyzed in details in Ref. 19. Various aspects of $n - n'$ oscillation in the NS were discussed also in Refs. 20–23.

The sign of ordinary BA, $B = \text{sign}(n_b - n_{\bar{b}}) = 1$, is fixed by some baryogenesis mechanism which created the primordial excess of baryons over antibaryons. At first glance

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\* In this paper we use natural units $c = 1$ and $\hbar = 1$. 2 Notice that by naming $n'$ as mirror neutron, we implicitly extend the notion of our baryon charge $B$ to that of M species $B'$ classifying the latter via the combined charge $B = B + B'$ which is conserved by $nn'$ mixing 8. In fact, the M species with $B = 1$ and $B = -1$ for us respectively are the mirror baryons (MB) and anti-mirror baryons (AMB) (namely, $\pi'$ is anti-mirror neutron or mirror antineutron), no matter how they are qualified by the inhabitants of M world.
glance, the sign of mirror BA should be the same by M parity, B’ = sign(nB − nV) = 1. This would be the case if the identical baryogenesis mechanisms act separately in the O and M sectors. However, one can also envisage a co-baryogenesis scenario e.g. via the processes which violate B and B’ but conserve B = B + B’ (which processes can in turn be related to underlying new physics which induces nn' mixing \( \eta \)). In this case, for B being positive, the null O+M asymmetry in the Universe would imply the negative B’. The co-genesis models which induce BA’s in both sectors via the B−L and B’−L’ violating cross-interactions between the O and M particle species indeed predict B and B’ of the opposite signs \( \eta \).

Hence, depending on the sign of B’\( \eta \), the compact stars in M sector can be the mirror neutron stars (MNS) or anti-mirror neutron stars (AMNS) consisting respectively of the MB or AMB. Correspondingly, n’ → n transitions in the MNS (or \( \pi’ \rightarrow \pi’ \) in the AMNS) can create the cores of ordinary matter (or antimatter) in their interior. In both cases these cores can be detectable as hot compact sources emitting the photons in the far UV and soft X-ray ranges \( \eta \), but this electromagnetic emission cannot trace their composition (matter or antimatter). However, the presence of interstellar medium makes the difference. Namely, accretion of interstellar gas on the MNS with O matter core can only heat it and induce X-ray emission. However, in the case of AMNS, the accreted gas will annihilate on the antimatter core producing gamma-rays with energies up to a GeV or so.

Interestingly, the recent analysis \( \eta \) based on the 10-year Fermi Large Area Telescope (LAT) gamma-ray source catalog \( \eta \) has identified 14 point-like candidates emitting γ’s with a spectrum compatible to baryon-antibaryon annihilation \( \eta \), and not associated with objects belonging to established gamma-ray source classes.

In this letter we discuss the possibility whether the unusual sources of this type can be the AMNS with the antimatter cores which produce the annihilation γ-rays by accretion of interstellar gas, and whether some part of this antimatter can escape from the AMNS producing the antinuclei detected by AMS-02 in cosmic rays.

2. Neutron stars are presumably born after the supernova (SN) explosions of massive progenitor stars followed by the neutronization of their iron cores. The NS are formed with very high rotation speed and with the surface magnetic fields as large as \( 10^{9} \div 10^{15} \) Gauss. This makes many of them observable as pulsars due to their electromagnetic radiation. The NS are believed to have a compact onion-like structure: the inner core of dense nuclear liquid dominantly consisting of neutrons (with less than 10 % fraction of protons and electrons), and the outer shell consisting of heavy nuclei which form a rigid crust at the surface (for reviews, see e.g. \( \eta \)). The NS mass-radius (\( M−R \)) relations depend on the equation of state (EoS), i.e. the pressure–density relation in dense nuclear matter. The masses of the known pulsars range within \( M = (1 \div 2) M_{\odot} \), and the observation of \( 2M_{\odot} \) ones disfavors the too soft EoS. E.g. the Sly EoS \( \eta \) allows the maximum mass \( M_{\text{max}} = 2.05 M_{\odot} \), while some of the realistic EoS reviewed in Refs. \( \eta \) can afford somewhat larger masses, up to \( 2.5 M_{\odot} \) or so. All these EoS predict the NS radii in the range \( R \approx (10 \div 15) \) km.

The baryon amount \( N \) is related to the NS mass as

\[
N = \kappa M/m \simeq \left( \frac{\kappa M}{1.5 M_{\odot}} \right) \times 2 \cdot 10^{57}
\]  

where \( m \) is the nucleon mass, and \( \kappa \simeq 1 \) is the EoS dependent factor which, for a given EoS, mildly depends also on \( M \). E.g. the SLy EoS implies \( \kappa \approx 1.1 \) for the typical NS masses \( M \approx 1.5 M_{\odot} \), increasing to \( \kappa \approx 1.2 \) for heavy NS with \( M \approx 2 M_{\odot} \). \( \eta \). The deficit between the gravitational mass \( M \) and equivalent baryonic mass \( M_{b} = mN_{b} \), \( M_{b} = M − M = (\kappa − 1)M \), corresponds to the gravitational binding energy.

In principle, at extremal densities the energetically favored ground state can be strange quark matter \( \eta \). Therefore, the NS which mass can reach a critical value by accretion from the companion or by the matter fallback after the SN explosion can be transformed into hybrid stars (HS) with more onion-like structure including the quark matter core, or even entirely quark stars (QS), depending on the amount of accreted mass and on the details of the quark matter EoS. Interestingly, the bimodal distribution of the pulsar masses \( \eta \) can correspond two different types of stars, the lighter stars being the NS and heavier stars being the HS/QS. In the following, we shall concentrate on the NS, addressing the possibility of the HS/QS only in proper occasions.

3. The core-collapse of O star should produce a NS entirely consisting of ordinary nuclear matter. But once the NS is formed, then \( n \rightarrow n’ \) transitions in its liquid nuclear core can effectively produce M neutrons.

The oscillation \( n \rightarrow n’ \) in nuclear medium is described by the Schrödinger equation with effective Hamiltonian

\[
H = \begin{pmatrix}
\Delta E & \varepsilon \\
-\varepsilon & 0
\end{pmatrix}
\]  

where the off-diagonal term \( \varepsilon \) comes from \( n−n’ \) mass mixing \( \eta \), and \( \Delta E \) is the medium-induced energy splitting between \( n \) and \( n’ \) states. It can be taken as the difference of the in-medium optical potentials induced by their coherent scatterings, \( \Delta E = V − V’ \). Namely, for O neutrons we have \( V = 2\pi a_{n}b/m_{n} \), where \( m_{n} \) is the neutron mass, \( a = 1.5 \times 3 \text{ fm} \) is its scattering length, and \( n_{b} = \xi \times 0.16 \text{ fm}^{-3} \) is the baryon density, i.e. \( \xi \) is \( n_{b} \) in units of nuclear density.

\( \eta \) Modulo some tiny amount of M matter which can be accreted by the progenitor star during its lifetime, or produced via \( n−n’ \) oscillation of the neutrinos involved in certain chains of nuclear reactions at late stages of its evolution before the core-collapse.

\( \xi \) The NS liquid core is dominated by the neutrons and we shall neglect subdominant fraction of protons for simplicity. In the crust the neutrons are bound in heavy nuclei and for them \( n−n’ \) conversions is ineffective, but this fraction is also negligible.
If the M baryon density in the star is small, \( n'_b \ll n_b \), \( V' \) can be neglected and we get \( \Delta E = V = a_3 \xi \times 125 \text{MeV} \) \( ^5 \)

Hence, in a dense nuclear medium \( n \sim n' \) mixing angle is
\[
\theta = \varepsilon/V \approx (\varepsilon_{15}/a_3 \xi) \times 8 \cdot 10^{-24} \ (\varepsilon_{15} = \varepsilon \times 10^{15} \text{eV}^{-1}),
\]
and the mean probability of \( n \sim n' \) oscillation \( P_{nn'} = 2\theta^2 \) is extremely small.

The mirror neutrons are mainly produced via the scattering process \( nn \rightarrow nn' \) \( ^4 \) Its rate can be estimated as
\[
\Gamma = 2\theta^2 \nu(\sigma\nu)_{nn}, \quad \sigma = 4\pi a^2 \quad \text{is the nn elastic scattering cross-section and} \quad v = p_F/m \quad \text{is the mean relative velocity, while the factor} \quad \eta = 0.18 \quad \text{takes into account the Pauli blocking implying that in} \quad nn \rightarrow nn' \quad \text{process the final state} \quad n \quad \text{should have the momentum above the Fermi momentum} \quad n_n.
\]

\[
\varepsilon_n \text{ is extremely small.}
\]

\( H_n \) Hence, in a dense nuclear medium \( \varepsilon \) can be neglected and we get \( \Delta \)
\[
\varepsilon_n \text{ is conserved, for the fraction of } M \text{ baryons produced}
\]
\[
\frac{M}{M} \varepsilon \sim 10^{-13} \text{eV} \quad \text{or so.}
\]

Thus, by integrating the Boltzmann equations
\[
\dot{N'} = \langle \Gamma \rangle N, \quad \dot{N} = -\langle \Gamma \rangle N
\]
and observing that the overall O+M baryon number \( N + N' \) is conserved, for the fraction of M baryons produced in the NS of the age \( t < t_U \) we get:
\[
\frac{N'}{N + N'} = \frac{t}{\tau_c} = 10^{-5} \times \varepsilon_{15}^2 a_R \left( \frac{M_{1.5}}{1.5 \text{M}_\odot} \right)^{2/3} \left( \frac{t}{10 \text{Gyr}} \right)
\]

Thus, for \( \varepsilon \sim 10^{-15} \) eV the amount of M baryons in the oldest NS are comparable to their amount in planets.

Due to the Pauli blocking, the process \( nn \rightarrow nn' \) process takes place for the neutron momenta close to the Fermi surface, so that \( n' \) is typically produced with the Fermi energy \( E_F = p_{F,n}^2/2m = \xi^{2/3} \times 60 \text{MeV} \). Hence, the energy production rate per baryon is independent of \( \xi \):
\[
\dot{E} = \Gamma E_F N \sim \varepsilon_{15}^2 \left( \frac{M}{1.5 \text{M}_\odot} \right) \times 10^{31} \text{erg/s} \quad (8)
\]
which energy will be radiated away via the mirror photons and neutrinos.

In the case of HS/QS situation is somewhat different. Transition \( n \sim n' \) should be ineffective in quark matter since there are not much neutrons. In addition, quark matter is thought to be self-bound \( ^{34} \), in particular if it is in color superconducting phase \( ^{37} \), and its transition to M matter is not energetically favored. Hence in the QS, almost entirely consisting of quark matter, \( n \sim n' \) transitions play practically no role. In the hybrid stars \( n \sim n' \) transition should take place in its shell consisting of the neutron liquid. Thus, in principle the HS can also develop the M matter cores in their interior. However, the energy production rate in the HS, with less neutrons employed, can be much lower than it was estimated in Eq. (8) for the NS.

For the evolution time \( \tau_c \) being much larger than the cooling time, Eq. (8) in fact determines the NS mass loss rate due to \( n \sim n' \) conversion:
\[
\dot{M}/M \sim -\varepsilon_{15}^2 \times 10^{-16} \text{yr}^{-1} \quad (9)
\]
The mass loss changes the orbiting period in the compact NS binaries as \( P/P_b \sim -2\dot{M}/M \) \( ^{22} \). The observational data on the orbital period decay in some well-studied compact binaries (as e.g. the famous Hulse-Taylor pulsar B1913+16) yield an upper bound roughly as \( \varepsilon < 10^{-13} \text{eV} \) \( ^{19} \). Somewhat stronger bounds, \( \varepsilon < 10^{-15} \text{eV} \) or so \( ^{19} \), can be obtained from the surface temperature measurements of old pulsars, which for PSR J2114–3933, with \( T_{\text{surf}} < 4 \times 10^9 \text{K} \) \( ^{38} \), approaches \( \varepsilon < 10^{-16} \text{eV} \). But this object, being an isolated slow pulsar with characteristic age \( \tau_c \sim 3 \text{Gyr} \) and unknown mass, could in principle be quark matter dominated. In this case the above limit would be invalid. More generally, the observational determination of the NS surface temperature is influenced by environmental factors as composition of the pulsar magnetosphere and interstellar extinction, and its theoretical interpretation depends on the cooling models. Therefore, in the following we conservatively stick to benchmark values \( \varepsilon \sim 10^{-15} \text{eV} \).

M baryons produced in the NS will form a mirror core in its interior. In fact, \( n' \) produced at the first instances after the NS formation will undergo \( \beta \)-decay \( n' \rightarrow p' + e^-\bar{\nu} \) forming a plasma consisting of \( p' \) and \( e^- \) components. But very soon, in hundred years or so, the density produced M matter \( n'_b \) will exceed \( 10^{98}/\text{cm}^3 \), equivalent to the baryon density in the sun’s centre. In gravitationally captured hot M matter nuclear reactions will be ignited among \( p' \)
In any case, the fraction of nuclear reactions in hot core, it is difficult to precisely
\[ \sigma \]
and \[ \epsilon \times 10^{15} \] for \[ \epsilon \] should be significant, and the “mirror photosphere”
temperature in the NS can be estimated by equating \[ xE = 4\pi R_{ph}^2 \times \sigma T^4 \]
\[ \text{where} \]
\[ R_{ph} \text{ is the photosphere radius} \]
and \[ \sigma \] is the Stefan-Boltzmann constant:
\[ T = \epsilon^{1/2} \left( \frac{xNn}{1.5 M_{\odot}} \right)^{1/4} \left( \frac{1 \text{ km}}{R_{ph}} \right)^{1/2} \times 10^6 \text{ K} \]  
(10)

For determining the observable temperature \[ T_{\infty} \], the surface redshift effect should be taken into account.

4. Let us reverse now the situation and consider neutron stars in M sector. As we have anticipated in the introduction, the BA \[ \bar{B} \] in M world can be positive or negative, depending on the baryogenesis mechanism. Hence, if \[ \bar{B} > 0 \] all compact mirror stars should be the MNS, whereas if \[ \bar{B} < 0 \] all should be the AMNS. As far as the O and M sectors have the same microphysics by mirror parity, the EOS describing the O and M nuclear matters should be identical, and, needless to say, it should be identical for the nuclear and antinuclear matters by C invariance of strong interactions. In other words, the NS, MNS and AMNS should have the same \( M-R \) relations.

In the absence on \( n-n' \) mixing these will be dark compact objects, sort of solar mass MACHOs which can be detected by microlensing, but ordinary observer cannot distinguish between the MNS and AMNS. If \( n-n' \) mixing is switched on, then \( n'-n \) transitions will create O matter in the MNS interior, while in the AMNS \( \bar{\pi} \rightarrow \pi \) will take place forming the O antimatter. These transitions with the effective time and energy production rate given again by Eqs. (2) and (8), would form hot cores which can be visible for us as bright compact stars with high temperatures (10). Clearly, this will not allow to determine whether the core is composed of O matter or O antimatter. However, the two cases can be discriminated by the accretion of ordinary gas which, in the case of the AMNS, will annihilate with the antibaryons in its interior. Thus, detection \( \gamma \)-ray sources with a typical baryon-antibaryon annihilation spectrum (20) can be the way to determine the BA sign in mirror world.

Let us consider an AMNS of a typical mass \( M \) and radius \( R \). Transitions \( \bar{\pi} \rightarrow \pi \) will produce O antimatter in its liquid core. The production rate (antibaryon per second) can be estimated as
\[ \dot{N}_{\bar{\pi}} = \frac{N_{\bar{\pi}}}{\tau_e} \approx \epsilon^{2/15} \left( \frac{M}{M_{\odot}} \right)^{2/3} \times 3 \cdot 10^{34} \text{ s}^{-1} \]  
(11)

where the amount of AMB in the AMNS \( N_{\bar{\pi}} \) is given by Eq. (2), and the conversion time \( \tau_e \) is given by Eq. (15). For simplicity, we put \( \kappa = 1.1 \) and \( R = 12 \text{ km} \).

On the other hand, the AMNS will accrete O gas while it travels in diffuse interstellar medium (ISM). The accretion rate (baryon per second) reads
\[ \dot{N}_b \approx \frac{(2GM)^2 n_{\text{is}}}{v^3} = \frac{10^{32}}{v_{100}^3} \left( \frac{n_{\text{is}}}{1/\text{cm}^3} \right) \left( \frac{M}{M_{\odot}} \right)^2 \text{ s}^{-1} \]  
(12)

where \( n_{\text{is}} \) is the baryon density in the ISM and \( v = v_{100} \times 100 \text{ km/s} \) is the star velocity relative to the ISM which is normalized taking into account that typical kick velocities of pulsars are order 100 km/s.

The ratio between the two rates
\[ \frac{N_{\bar{\pi}}}{\dot{N}_b} \approx \epsilon^{2/15} v_{100}^3 \left( \frac{1/\text{cm}^3}{n_{\text{is}}} \right) \left( \frac{M}{M_{\odot}} \right)^{-4/3} \times 300 \]  
(13)

depends, for a given \( \epsilon \), on the AMNS mass \( M \) (which can range within \( M = (1 \div 2) M_{\odot} \) as that of the normal NS) and, more critically, on its velocity \( v \) and on the ISM density. The velocity distribution of observed pulsars seems to be bi-modal: a part of pulsars have kick velocities \( v > 100 \text{ km/s} \) and some achieving 1000 km/s, while others probably receive a very little kick – otherwise they would not be contained in globular clusters which have small escape velocities. The origin of pulsar kicks is largely unknown, but we consider that this bimodal distribution applies also to the AMNS or MNS.

For definiteness, let us take \( \epsilon \approx 10^{-15} \text{ eV} \) as a benchmark value. For the AMNS with \( v > 100 \text{ km/s} \) we have \( N_{\bar{\pi}}/\dot{N}_b \gg 1 \), i.e. the antibaryon production rate (11) is much larger than the baryon accretion rate (12). In this case the antimatter core can be formed and it will emit O photons with the energy luminosity (8). Such cores can be observed as bright point-like sources in the far UV and soft X-ray ranges. As for the accreted baryons, they should annihilate on the core surface. The annihilation photons will be produced with the rate \( L_{\gamma} = \lambda_\gamma N_b \), where \( \lambda_\gamma \approx 4 \) is the average multiplicity of \( \gamma \)-s per \( p\bar{p} \) annihilation (30). The rate of energy production is \( 2\pi N_{\bar{\pi}} \), about a half of which contributes in heating the core (in addition to (8)), and another half will be emitted from its surface in \( \gamma \)-rays. Hence, the energy flux from such a source at a distance \( d \) will be \( J \approx m_{\bar{\pi}} v_b / 4\pi d^2 \), or numerically
\[ J \approx \frac{10^{-12}}{v_{100}} \left( \frac{n_{\text{is}}}{1/\text{cm}^3} \right) \left( \frac{M}{1.5 M_{\odot}} \right)^2 \left( \frac{50 \text{ pc}}{d} \right)^2 \text{ erg cm}^{-2} \text{s}^{-1} \]  
(14)

For the AMNS travelling with \( v > 100 \text{ km/s} \) in the Milky Way (MW) this emission can be below the diffuse \( \gamma \)-background and the source may remain unresolved unless this source is closer than 50 pc or so. However, if the AMNS has less velocity, say \( v \approx 30 \text{ km/s} \), and it incidentally crosses a high density region as e.g. cold molecular cloud with \( n_{\text{is}} > 10^4 \text{ cm}^{-3} \), the observability distance can be increased up to several kpc.
On the other hand, for the slow AMNS moving in galactic discs with $v < 10 \text{ km/s}$, the antibaryons produced in its interior can be outnumbered by the accreted baryons, i.e. $N_T < N_b$. In this case the antimatter core cannot be formed and the thermal emission with temperature $T$ will be suppressed. The produced antibaryons will readily annihilate with the already accreted baryons and the $\gamma$-ray luminosity will be $L_\gamma = 1 \times N_T$. In other words, the $\gamma$-luminosity of the object will be defined by the lesser value between $N_T$ and $N_b$. Namely, if the ratio $N_T/N_b$ is less than one, then instead of Eq. 14 we would have for the energy flux

$$J \approx 10^{-12} \times \left( \frac{\varepsilon_{15}}{0.06} \right)^2 \times \left( \frac{M}{1.5 M_\odot} \right)^{2/3} \times \left( \frac{50 \text{ pc}}{d} \right)^{2} \times \frac{\text{erg}}{\text{cm}^2\text{s}}$$  (15)

where the $\varepsilon$-dependent factor is normalized to indicate that for $\varepsilon < 5 \times 10^{-17} \text{ eV}$ such a sources become too faint to be resolved at distances larger than 50 pc or so.

The search of $\gamma$-ray sources with a spectrum compatible with baryon-antibaryon annihilation was recently performed in Ref. 28. Analyzing 5787 sources included in the Fermi LAT, 14 candidates were found which were selected by applying the following criteria:

(i) extended candidates were excluded (with angular size larger than the LAT resolution at energies $E < 1 \text{ GeV}$);
(ii) sources associated with objects known from other wavelengths and belonging to established $\gamma$-ray sources were excluded, as e.g. pulsars and active galactic nuclei;
(iii) sources with significant higher energy tail above a GeV were excluded since the baryon-antibaryon annihilation $\gamma$-spectrum should end up at the nucleon mass.

Interestingly, the distribution of the sources in the sky shown in Fig. 1 of Ref. 28 very much resembles the distribution of observed pulsars. Only two candidates have galactic coordinates compatible with the MW disc, while the 11 candidates having galactic latitudes $|b| > 10^\circ$ (among which 7 candidates with $|b| > 30^\circ$) can be assigned to the MW halo. Interestingly, the sources belonging to the disc are the brightest, with the energy fluxes $J \geq 10^{-11} \text{ erg cm}^{-2} \text{ s}^{-1}$, while the ones with higher galactic latitudes become increasingly fading, and the source J2330-2445 ($b = -71, 7^\circ$) is the taintest, with $J < 2 \times 10^{-12} \text{ erg cm}^{-2} \text{ s}^{-1}$. In view of Eq. 12, this may well explained by correlation of the accretion rate with the distribution of the ISM densities, though the velocity distribution of the stars remains the key issue.

Probably it is too early to claim the discovery. These sources are all faint, close to the Fermi LAT detectability threshold, and they may well belong to a known $\gamma$-ray source classes, or maybe mimicked by imperfections of the background interstellar emission. In fact, the authors of Ref. 28 take a conservative attitude and translate their findings into an upper limit on the local fraction of such objects with respect to normal stars. The clear identification of these sources is very challenging, and requires serious multiwavelength search.

The possibility of mirror NS (or HS) being the engines of our antimatter can have interesting implications since they can produce antinuclei in the ISM. Namely, for $\varepsilon \sim 10^{-15} \text{ eV}$, transitions $\pi^+ \rightarrow \pi^-$ transition in the ANS produces about $10^{52}$ antibaryons forming the hot and dense core in which nuclear reactions should form the antinuclei. However, these antinuclei are gravitationally trapped and the question is how they can escape from the star. This possibility can be provided by the mirror neutron star mergers. In the coalescence of the two ANS, their small antimatter cores do not merge at the same instant but continue the orbiting and then explode due to the decompression producing a hot cloud of the neutron rich antinuclei. Most of these antinuclei, being stable only at the extreme densities, will decay into the lighter ones which are stable in normal conditions. Therefore, the antinuclei produced by such “slings” effect can leave the coalescence site and propagate in the outer space. In addition, they can have some additional acceleration if reasonably large magnetic fields are formed the rotating ordinary anti-core during its evolution before the merger. The AMS-02 experiment hunting for the antinuclei in the cosmic rays has reported, as the preliminary results, the detection of eight antihelium events (among which two are compatible with antihelium-4 and the rest with antihelium-3). The fraction $\sim 10^{-9}$ of antihelium with respect to measured fluxes of the helium is too high to be explained by the conventional production mechanisms. Interestingly, the rate of the NS mergers $\sim 10^3 \text{ Gpc}^{-3} \text{ yr}^{-1}$, with $\sim 10^{52}$ antibaryons produced per a merger, is nicely compatible with this fraction of antihelium. In addition, our mechanism should produce also the heavier antinuclei, and thus AMS-02 can be the place where to find fantastic animals as antical or antioxygen which would be a key for many mysteries.

5. Let us discuss our scenario in wider context, and have a pleasant walk and pleasant talk with the Walrus and the Carpenter, viewing and reviewing panorama of two parallel worlds. This picture is based on the direct product $G \times G'$ of two identical gauge groups represented by the SM (or some its extension), with the total Lagrangian

$$\mathcal{L}_{\text{tot}} = \mathcal{L} + \mathcal{L}' + \mathcal{L}_{\text{mix}}$$  (16)

where $\mathcal{L}$ describes the O matter and $\mathcal{L}'$ the dark M matter, while the mixed Lagrangian $\mathcal{L}_{\text{mix}}$ describes the possible cross-interactions between the particles of two sectors. The identical forms of $\mathcal{L}$ and $\mathcal{L}'$ is ensured by a discrete symmetry $G \leftrightarrow G'$ interchanging all O species: the fermion, gauge and Higgs fields of $G$ sector, with their M partners: the fermion, gauge and Higgs fields of $G'$ sector. In the context of extra dimensions, it can be viewed as a geometric symmetry between two parallel 3-branes on which the O and M particle species are localized.

M baryons as cosmological DM have specific cosmological implications. Although O and M components have identical microphysics, their cosmological realizations cannot be quite different. Namely, the viability of
M sector requires the following conditions \cite{40,41}:

- **Asymmetric reheating:** after inflation the O and M sectors are reheated asymmetrically, with \( T > T' \), which can naturally occur in certain models;

- **Out-of-equilibrium:** interactions between O and M particles are feeble enough in order to maintain the initial temperature asymmetry in subsequent epochs. In other words, all cross-interactions in \( \mathcal{L}_{\text{mix}} \) should remain out-of-equilibrium at any stage after inflation, at least before the Big Bang Nucleosynthesis (BBN);

- **No extra heating:** after inflaton both sectors evolve almost adiabatically and the temperature difference \( T' < T \) is not erased by entropy production in M sector due to possible 1st order phase transitions.

Namely, the BBN bounds demand \( T'/T < 0.5 \) \cite{41} while the post-recombination cosmology is yet more restrictive, requiring \( T'/T < 0.2 \) or so \cite{13}. This has interesting consequences for the primordial chemical content in M world: while O world is dominated by hydrogen, with its primordial mass fraction being 75\%, and that of helium being 25\%, M world should be helium dominated, with the mass fractions of M hydrogen and M helium being respectively 25\% and 75\% or so \cite{41}.

Along with the ordinary stars also dark M stars can be formed the Galaxy. However, since mirror world is colder, the first M stars should start to form somewhat earlier than the first (population III) stars in O sector. M stars, being helium dominated, should have somewhat different initial mass function, and their evolution should be more rapid as compared to O stars \cite{44}. Since helium is less opaque than hydrogen, the massive M stars should suffer less mass losses due to stellar winds end up their life collapsing directly into black holes (BH). The intermediate mass M stars can explode as SN and form mirror neutrons stars, while the solar mass M stars as white dwarfs can survive until present times. All these objects can constitute a dominant fraction of dark matter in the Galaxy. Namely, the galactic halo can be viewed as elliptical mirror galaxy of M stars and BH in which O matter forms the disc \cite{41}. M matter could contribute also to the disc, but the density of M stars in the disc should not exceed the density of O stars \cite{44}. All these objects can be observed via microlensing as the Machos in different mass ranges. The present limits from EROS-MACHO observations do not exclude the possibility of the galactic halo dominated by dark objects as BH with masses \( M > (10 \div 100) M_\odot \) while the fraction of dark stars with \( M \sim M_\odot \) can be \( \sim 10\% \) or so. This proportion can correspond to the abundance of the LIGO gravitational wave (GW) signals \cite{40} from the BH mergers with typical masses \( M \sim (10 \div 50) M_\odot \) or so \cite{47}. In addition, some of the peculiar LIGO events with one or both light components and no optical counterpart can be viewed as the BH-MNS or MNS-MNS mergers \cite{??}.

Now the time has come to talk of many things: of shoes and ships and sealing-wax, of cabbages and kings... there is a subtlety related to the chiral character of the fermion representations in the SM: in our weak interactions fermions are left-handed (LH) while the antifermions are right-handed (RH), and two systems would be symmetric if not CP-violating effects. The value of the BA in the Universe, and in particular its sign \( B = \text{sign}(n_\text{b} - n_\text{b'}) = 1 \), is fixed by (unknown) baryogenesis mechanism a l´a Sakharov \cite{40} that created primordial excess of baryons over antibaryons due to CP-violation in some out-of-equilibrium processes violating B (or B−L) \cite{54}. In parallel sector the situation is same, apart of an ambiguity in the CP-violation pattern distinguishing between the M fermions and M antifermions, and determining the sign of mirror BA \( B' = \text{sign}(n_\text{b} - n_\text{b'}) \). This ambiguity is related to the fact that \( G \leftrightarrow G' \) symmetry can be realized in two ways: with or without chirality change between the O and M species \cite{41}.

Namely, for three families of O fermions \( f_{L,R} \) described by \( SU(3) \times SU(2) \times U(1) \), the left-handed (LH) quarks \( q_L = (u_L, d_L) \) and leptons \( l_L = (\nu_L, e_L) \) are weak doublets while the right-handed (RH) ones \( u_R, d_R \) and \( e_R \) are singlets. The antifermion fields are obtained by complex conjugation, \( \bar{f}_{R,L} = C\gamma_0 f_{R,L}^\dagger \) and they have opposite chiralities: \( \bar{q}_L = (\bar{u}_R, \bar{d}_R) \) and \( \bar{u}_L, \bar{d}_L \) are antiquarks, and \( \bar{e}_R = (\bar{\nu}_R, \bar{\epsilon}_R) \) and \( \bar{e}_L \) are antileptons. The invariance under the transformation CP: \( f_{L,R} \rightarrow \bar{f}_{R,L}, \) i.e. symmetry between the fermions and antifermions, is violated by the complex Yukawa couplings with the Higgs doublet \( \Phi \).

As for three analogous families of \( SU(3)' \times SU(2)' \times U(1)' \) in M sector, we invert definitions and denote the species \( \bar{f}_{L,R}' \) with the LH weak interactions as M antiquarks: \( \bar{q}_L' = (\bar{u}_L', \bar{d}_L') \) and \( \bar{u}_R', \bar{d}_R' \), and M antileptons: \( \bar{e}_L' = (\bar{\nu}_L', \bar{\epsilon}_L') \). Correspondingly, the respective anti-species \( f_{R,L}' \) with the RH weak interactions we call M quarks: \( q_R' = (u_R', d_R') \) and \( u_L', d_L' \), and M leptons: \( \nu_R' = (\nu_R', \epsilon_R') \) and \( \epsilon_L' \). This is just a convention: we could name them in the opposite way. The M fermions and antifermions are equivalent modulo CP violating phases in the Yukawa couplings of the mirror Higgs doublet \( \phi' \).

Hence, one can impose a symmetry \( Z_2 : f_{L,R} \leftrightarrow f_{L,R}' \) interchanging the twin species of the same chirality, i.e. each O fermion with the corresponding M antifermion. Alternatively, we can employ \( Z_2^{LR} = Z_2 \times CP \) under exchange \( f_{L,R} \leftrightarrow f_{R,L}' \) between the O and M fermions which, in our definition, have the opposite chiralities. (Clearly, both of these transformations should be also complemented by a proper exchange between the gauge and Higgs fields of two sectors.) In the former case the M antimatter should have exactly the same CP-violating physics as our matter. In the latter case the equivalence holds between the ‘left-handed’ O matter and the ‘right-handed’ M matter which means that P parity, maximally violated in weak interactions of each sector, is in some sense restored between two sectors. In fact, this was the original motivation for introducing mirror fermions \cite{51,52} (for a historical overview see Ref. \cite{53}). But the real difference is related to CP-violation which was not yet discovered at the time of original works \cite{51}.

In the absence of CP-violating phases \( Z_2 \) and \( Z_2^{LR} \) would
be equivalent.

Both of these discrete symmetries ensure the identical form of the O and M Lagrangians $\mathcal{L}$ and $\mathcal{L}'$ in (10), modulo the issues of CP-violation. Nevertheless, in the absence of cross-interaction terms $\mathcal{L}_{\text{mix}}$ in (10), with the O and M particles interacting only gravitationally, no experiment can discriminate between the two possibilities.

On the other hand, there are no fundamental reasons for neglecting the cross-interactions in $\mathcal{L}_{\text{mix}}$ which in fact are the portals for the mirror DM detection and identification. For example, the simplest possibility is a kinetic mixing term between the O and M photons, $\varepsilon F^\mu\nu F'^{\mu\nu}$ [54]. The particles of two sectors can interact also through the gauge bosons of e.g. the common family symmetry $SU(3)_H$ [55] or common $U(1)_{B-L}$ symmetry [57]. They can also share the Peccei-Quinn symmetry and their cross-interaction can be mediated by the axion [58]. The M gas (dominated by helium, and probably containing some M nuclei of carbon-neutrogen-oxygen) can be subject of direct detection via the kinetic mixing of the O–M photons or via other portals in $\mathcal{L}_{\text{mix}}$ [59].

The most interesting interactions in $\mathcal{L}_{\text{mix}}$ are the ones that violate baryon and lepton numbers of both sectors.

In fact, the conservation of B and L in the SM is related to accidental global symmetries of the Lagrangian at the level of the renormalizable terms. However, these global symmetries can be broken by higher order operators cutoff emerging from new physics at some large energy scales. Namely, L should be violated if the neutrino have the Majorana masses, while B (or B − L) violation is needed for the baryogenesis.

In particular, the lowest dimension effective operators which violate the lepton numbers are of dimension D=5:

$$\frac{A}{M} (\ell_L \phi) ^2 + \frac{A'}{M} (\bar{\ell}_L \phi') ^2 + \frac{A}{M} (\bar{\ell}_L \phi') (\ell_L \phi) + \text{h.c.} \quad \text{(17)}$$

where $M$ is the relevant mass scale, and $A = A^T$, $A' = A'^T$ and $A$ are the matrices of generically complex Wilson coefficients in family space (the family indices as well as C-matrix are suppressed). The first term, violating L by two units ($\Delta L = 2$), after substituting the Higgs VEV $\langle \phi \rangle = v \sim 10^2$ GeV, induces the small Majorana masses to the neutrinos, $m_\nu \sim v^2 / M$ [60]. The second term ($\Delta L' = 2$) works analogously in M sector and induces the Majorana masses of M neutrinos. As for the third term, it violates both L and L' by one unit conserving the combination $L = L + L'$, and induces the O–M neutrino (active-sterile) mixing [61].

The first operator in (17) can be induced by seesaw mechanism, by introducing the heavy Majorana fermions N in weak singlet and triplet representations which are coupled to the leptons via the Yukawa terms $Y_L N \phi + \text{h.c.}$, with $Y$ being the matrix of respective coupling constants. Then the second operator is induced via the Yukawa terms $Y' \ell_L N' \phi' + \text{h.c.}$ with the analogous heavy fermions $N'$ of M sector. However, there can exist also singlet heavy fermions $N$ which are coupled with both O and M leptons: $\gamma \ell_L N \phi + \gamma' \bar{\ell}_L N' \phi' + \text{h.c.}$, and thus are messengers between two sectors. In this way, all three operators in (17) are induced by the seesaw mechanism, and for their coefficients we get:

$$A = Y Z Y^T + \gamma Z Y'^T, \quad A' = Y' Z' Y'^T + \gamma' Z Y'^T,$$

$$\bar{A} = Y' Z Y'^T \quad \text{(18)}$$

where the coefficient matrices $Z$ etc. parametrize the inverse mass matrices of $N$, $N'$ and $N$ fermions respectively as $Z/M$, $Z'/M$ and $Z/M$. Without losing generality, these matrices can be taken to be diagonal and real, while the symmetry under $N \leftrightarrow N'$ implies $Z' = Z$.

Now the question comes to the sign of $Z_2$ in M sector. Clearly, this depends on the baryogenesis models, for which two possible realizations of the discrete inter-sector symmetry, $Z_2$ and $Z_2^{LR}$, have different implications.

Let us consider the case of a baryogenesis mechanisms acting separately in O and M sectors in identical manner. For example, the BA’s can be induced by means of electroweak baryogenesis (EWB) in both sectors [41]. Then, in the case of $Z_2$, which yields the CP-violation pattern for the M antimat matter identical to that of the O matter, the BA’s B and B' induced in two sectors should have the opposite signs. Alternatively, we can consider a leptogenesis scenario related to the seesaw mechanism, due to the CP-violation in the decays of the heavy Majorana fermions $N \rightarrow \ell_{L/R} \phi$ and $N' \rightarrow \ell_{L/R} \phi'$ (and analogous decays of $N'$-fermions coupled to both sectors). Then $Z_2 \left( \ell_{L/R} \phi \right)$ implies $Y' = Y$ and $Y' = Y$ for the Yukawa couplings in (15), and so for B being positive, B' should be negative.

As for $Z_2^{LR} \left( \ell_{L/R} \phi \right)$, it implies the equivalent CP-violation between the O and M fermions, and the above mechanisms of separate baryogenesis mechanisms in this case predict the same signs of B and B'.

However, the BA of opposite signs between the O and M sectors can be generated also in the case of $Z_2^{LR}$ symmetry. This occurs e.g. in the co-leptogenesis models discussed in [62, 63] which assumes that after inflation the O and M sectors are reheated asymmetrically, with $T \gg T'$, and masses of messenger $N'$ fermions between two sectors are larger than the reheating temperature. Nevertheless, the operators in (17) mediate scattering processes as $\ell_L \phi \rightarrow \bar{\ell}_L \phi'$, $\ell_L \phi \rightarrow \bar{\ell}_R \phi'$ etc. which violate both L and L', and they are out-of-equilibrium. The invariance under $Z_2^{LR} \left( \ell_{L/R} \phi \right)$ for the Yukawa couplings in (15) implies $Y' = Y$ and $Y' = Y$, in which case the CP-violating factors in the above scattering processes are non-zero, and they induce $B$ and $B'$ of the opposite signs [26, 27]. (Interestingly, this mechanism is ineffective in the case of $Z_2$ symmetry yielding $Y' = Y$ since CP-violating factors appear to be vanishing [26, 27].) Let us remark that this mechanism implies $\Omega_\nu > \Omega_b$, which is
related to the fact that M sector is colder and the produced B’ – L’ suffers less damping \[6\]. Hence, it can naturally explain the observed cosmological fractions of the baryons and DM, \(\Omega_B/\Omega_D \approx 5\), which also makes clear who has eaten more oysters, the Walrus of the Carpenter.

Hence, depending on the baryogenesis model and the type of discrete exchange symmetry, M sector can have positive or negative BA which sign should be also correlated to the chirality of M baryon in their weak interactions. In principle, two situations could be distinguished to the chirality of M baryon in their weak interaction. In an UV complete theory they can be induced via effective energy production rate due to antimatter creations which can be rendered visible, via the baryon-antibaryon annihilation γ-rays, by accretion of ordinary matter from the ISM.

We have discussed these effects in the minimal situation which assumes that n – n’ mixing occurs only due to mass mixing \[1\]. n – n’ mixing is absent, and n and n’ are exactly degenerate in mass. However, the concept permits more variations which we briefly mention below: 

Transitional magnetic moment. In difference from the n – π system where transitional magnetic moment between n and π is forbidden by Lorentz invariance, non-diagonal magnetic moment \(\mu_{n n’}\) (or dipole electric moment) is allowed between n and n’ \[6\].\[67, 68\] and they can be effectively induced in certain models of n – n’ mixing \[22\]. In this case the n – n’ transition time will depend on the magnetic field in the NS, and it can be simply estimated by replacing \(\varepsilon\) into \(|\mu_{n n’}/B|\) in Eq. \[5\], or more concretely

\[
\varepsilon_{15} \rightarrow \varepsilon_{15}^B = \left( \frac{\mu_{n n’}}{10^{-27}\text{ eV/G}} \right) \left( \frac{B}{10^{12}\text{ G}} \right)
\]

Taking e.g. \(\mu_{n n’} \sim 10^{-27}\text{ eV/G}\), which is 16 orders of magnitude smaller than the neutron magnetic moment itself, and making replacement \[21\], we see that for a mirror magnetar (\(B \sim 10^{15}\text{ G}\)) the antimatter production rate will be \(\sim 10^{40}/\text{s}\) while for an old recycled AMNS with \(B \sim 10^8\text{ G}\) it will be \(\sim 10^{26}/\text{s}\). Thus, the former objects should be very bright while the latter can be practically invisible in annihilation γ-rays. Therefore, in this case the analysis similar to that of Ref. \[28\], would require a specific selection of the source samples which would take into account the distribution of magnetic field values in the NS.

n – n’ mixing. For a simplicity, we have considered the situation with only n – n’ mixing \[1\], induced via effective D=9 operators \[19\], which conserves \(\bar{B} = B + B’\). However, there can exist also n – \(\bar{n}\)’ mixing: there is

\[8\] We single out n – n’ mixing though generically operators \[18\] can induce also other mixings as e.g. \(\Lambda – \Lambda’\) between the O and M hyperons etc.
no fundamental reason to forbid it. However, the latter mixing, due to the SM structure, emerges from D=10 operators \[69\], and the depending on the model parameters, \( n - \bar{n}' \) mixing mass \( \varepsilon_{n\bar{n}'} \) can be much smaller than \( n - n' \) mixing mass \( \varepsilon \), but can be also comparable to it. In the latter case, with \( \varepsilon_{n\bar{n}'} \sim \varepsilon_{nn'} \), both the MNS or AMNS could produce the baryon-antibaryon annihilation \( \gamma \)-rays, without ’help’ of the ordinary gas accretion. Interestingly, in the presence of both \( \varepsilon_{n\bar{n}'} \) and \( \varepsilon_{nn'} \) with the comparable values is not conflict with the nuclear stability limits, while for the free neutron case it could allow to effectively induce \( n - \bar{n} \) oscillation with pretty large rates provided that experimental conditions are properly tuned \[69\].

\( n - n' \) mass splitting. We considered the minimal situation when \( n \) and \( n' \) have exactly the same masses in which case the experimental bounds \[8\] - [15] imply \( \varepsilon < 10^{-15} \) eV or so. In this case the time of \( n - n' \) transition \[15\] is much larger than the Universe age, and thus it should be an ongoing process in the existing neutron stars (or M neutron stars). However, much larger values of \( \varepsilon \) are allowed by the experiment if \( n \) and \( n' \) are not degenerate in mass. In particular, \( n - n' \) oscillation e.g. with \( \varepsilon \sim 10^{-10} \) eV or so can solve the neutron lifetime problem, the \( 4\sigma \) discrepancy between the neutron lifetimes measured via the bottle and beam experiments, provided that \( n \) and \( n' \) have a mass splitting \( m_{n'} - m_n \sim 100 \) neV \[70\]. In fact, mass splitting will emerge in models in which M parity is spontaneously broken \[40\] but with a rather small difference between the O and M Higgs VEVs \( \langle \phi \rangle \) and \( \langle \phi' \rangle \) \[71\]. In this case \( n - n' \) conversion time will be much smaller, \( \tau_c \sim 10^6 \) yr or so, so that the most of existing NS should be already transformed in maximally mixed stars with equal amounts and equal radii of the O and M components. Hence, half of the AMNS in this case will be our antimatter matter, and the \( \gamma \)-ray emission rate due to accretion will be given by Eq. \[13\].

Concluding, we have discussed a possibility of M world having a negative BA, so that the M neutron stars are the AMNS, and \( n' - \bar{n} \) transition in their interior can create antimatter cores. The ordinary gas accreted from the ISM annihilating on the surface of these cores give rise to \( \gamma \)-rays with the typical spectrum reducible to the baryon-antibaryon annihilation.

The alternative our mechanism can be the existence of antimatter stars (antistars) \[72\]. The commonly accepted baryogenesis mechanisms fix the value as well as the sign of the BA universally in the whole Universe. In addition, the observations rule out the existence of significant amount of antimatter on scales ranging from the solar system to galaxy and galaxy clusters, and even at very large scales comparable to the present horizon \[73\], \[74\]. However, more exotic baryogenesis mechanisms (for a review see \[77\]) can in principle allow the existence of small domains at well-tempered scales in which antimatter could survive in the form of antistars \[78\], \[79\]. In particular, the Affleck-Dine mechanism \[74\] can be extended by the coupling of the Affleck-Dine B-charged scalar field to the inflaton \[70\]. This modification, with properly tuned parameters, can give rise to large baryon overdensities at needed scales in which the stars of specific pattern (or the baryon-dense objects (BDO) as they were named in Ref. \[81\]) can be formed. In addition, in these overdensities the difference between the baryon and antibaryon amounts can be non-vanishing, and it could be positive as well as negative. Provided that part of the BDO consisting of antibaryons survive in the Milky Way (MW) halo up to present days, they can be observed as the emitters of the \( \gamma \)-annihilation \( \gamma \)-rays.

In principle, the BDO and AMNS mechanisms can be distinguished by the spectral shape of the \( \gamma \)-annihilation \( \gamma \)-rays. Namely, the proton annihilation on the surfaces of the BDO should produce \( \gamma \)-rays with typical spectrum peaked at 70 MeV or so \[30\]. In the case of the AMNS, the spectral shape will be deformed by the surface redshift effect, by a factor \( \exp\left[\phi \right] \), where \( \phi = \phi(r) \) is the gravitational potential at the surface of antimat-ter core inside the AMNS. This will rescale down the spectral shape by \( (15 - 30)\% \) depending on the AMNS mass, the EoS specifics and the radius of antimatter core \( r < R \). In addition, one has to take into account the energy blueshift of the accreted protons: in fact, at the core surface they will be semi-relativistic, with the speeds nearly approaching the speed of light .

In addition, let us recall that the AMNS can radiate substantial energy \[8\] via the photons in the far UV/soft X-ray ranges (provided that the ratio \[13\] is much larger than 1) which can be an additional tracer for their identification. In addition, in Ref. \[28\] the sources possibly associated with pulsars were excluded from the possible antistar candidates. On the other hand, it is plausible that the AMNS are also observable as ordinary pulsars, if large ordinary magnetic fields are somehow developed in their antibaryon cores. This could be realized, for example, if along with \( nn' \) mixing, there is also a kinematic mixing between the O and M photons \[54\] which effectively renders the mirror electrons and protons mini-charged (with tiny ordinary electric charges). The value of these electric mini-charges are severely restricted by the the cosmological \[52\] and experimental \[53\] bounds. Nevertheless, their existence can be effective. Since the antimatter core in the AMNS should consist of the heavy antimatter and positrons, the AMNS rotation can induce circular electric current in its antimatter core by the drag mechanism \[54\] which can be sufficient for these cores to acquire significant magnetic field, as it was discussed in Ref. \[19\]. Therefore, the AMNS could mimic ordinary pulsars, perhaps with some unusual properties. Having this in mind, maybe the pulsars should not be ex-cluded from the candidate selection provided that their \( \gamma \)-emission has no high energy tail above 1 GeV or so.

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