Overdamping of neutron-mirror-neutron transitions in neutron stars

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The neutron–to–mirror-neutron transitions in neutron stars could possibly result in significant effects. In this work, we show that collisional decoherence entails exponential relaxation in lieu of oscillations. Decoherence is a great many orders of magnitude faster than the expected oscillations. The admixture of mirror neutrons at all times remains very small with respect to the ordinary-neutron component.

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I. INTRODUCTION

The idea of the existence of a hypothetical hidden mirror sector of the Standard Model was first explicitly formulated in [1]. It took a distinct form in [2]. The subject has a long history—for the review of the emergence and evolution of the concept, see Ref. [3].

The possible existence of mirror matter (MM) and the intersection of MM and ordinary matter (OM) have been the subject of intense studies during the last few decades. The significance of MM for cosmology, astrophysics, and the search for dark matter has been discussed in a number of papers [4–11]. The MM and OM interact via gravity. Besides gravity, other tiny cross-interactions between the two sectors possibly exist and can give rise to interesting observable effects.

In particular, the neutron n can be mixed with the mirror neutron n', leading to the n-n' transitions, as was suggested in Refs. [12,13].

Since then, a number of phenomenological applications of n-n' oscillations have been discussed. These include the Greisen-Zatsepin-Kuzmin cutoff [12,13], the suppression or enhancement of oscillations due to a nonzero mirror magnetic field B' and its interplay with the ordinary field B [14], and the role of n-n' transitions in the neutron lifetime puzzle [15]. An important question was raised already in Refs. [12,13]—namely, whether n-n' transitions can proceed in nuclei. The short answer given in [12,13] was that they are forbidden by the neutron binding energy as far as n and n' are degenerate in mass (Z_2 symmetry). According to Ref. [13], similar arguments apply to the neutron stars (NSs). The last statement has undergone an important evolution [16–18]. The n-n' conversion in NSs is allowed, since the binding of neutrons in NSs is provided mainly by the gravitational pressure [18]. The hole created by the neutron conversion is filled by the neighboring n, while the created n' gravitates to the NS's center. In view of the breakthrough in observations of NSs [19], the problem of NS transformation into a mixed star (MS) partially consisting of MM is now under intense discussion [18,20–24]. The rate of this process and the observable manifestations of the MS formation are addressed in these publications.

In the present work, we concentrate on the renormalization of the n-n' transition inside the NS caused by the incoherent collisions with the surrounding neutrons. The point is that the rate of oscillations critically depends on the ambient conditions, which include the surrounding medium, the external magnetic field B, and the mirror magnetic field B'. We show that the coherence needed for n-n' transition to occur is ruined by high-frequency neutron collisions in the NS matter. A correct way to describe the quantum system in contact with the environment is to use the density matrix formalism [25,26] and the theory of open quantum systems [27,28] based on the Lindblad [29,30] and Bloch [31,32] equations. Put differently, the open quantum system cannot be described by the two-by-two Hamiltonian without the risk of obtaining dubious results. The problem of oscillations in a two-state physical system interacting with the environment was for the first time formulated and solved using the density matrix in a seminal paper by G. Feinberg and S. Weinberg [33] devoted to muonium-to-antimuonium conversion. A clear and comprehensive example of applying the above formalism to treat positronium–to–mirror-positronium oscillations is presented in Ref. [34].

The paper is structured as follows: In Sec. II, we introduce the density matrix formalism and Bloch equation. As a warming-up exercise, the standard equation for the n-n' oscillations in vacuum is derived. In Sec. III, the

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Lindblad and Bloch equations are introduced and discussed. In Sec. IV, we apply Lindblad and Bloch equations to n-n' transitions inside the NS and derive the overdamped relaxation in lieu of oscillations. In Sec. V, we discuss the intricate interplay of the collision friction parameter and the time interval between collisions. Conclusions and future research perspectives are presented in Sec. VI.

II. THE DENSITY MATRIX APPROACH TO n-n' OSCILLATIONS IN VACUUM

As a warming-up task, we solve the problem of n-n' oscillations in vacuum by means of the density matrix formalism. For the n-n' system, the density matrix reads (with the indices 1 and 2 attributed to n and n')

$$\hat{\rho}(t) = \begin{pmatrix} \psi_1\psi_1^* & \psi_1\psi_2^* \\ \psi_2\psi_1^* & \psi_2\psi_2^* \end{pmatrix} = \begin{pmatrix} \rho_{11} & \rho_{12} \\ \rho_{21} & \rho_{22} \end{pmatrix}. \tag{1}$$

Consider an isolated n-n' system in vacuum with β-decay neglected. We assume that the mirror Z_2 symmetry is exact, so that the fermion masses in both sectors are the same. The Hamiltonian has the form

$$\hat{H} = \begin{pmatrix} E & \varepsilon \\ \varepsilon & E' \end{pmatrix}, \tag{2}$$

where ε is a mixing parameter and $E - E' = d$. Currently, we do not specify the origin and the value of the parameter d. If we identify it with the energy difference between n and n' due to the superimposed magnetic field B in absence of the mirror field B', then $d = |\mu_n B| = 6 \times 10^{-12}(\frac{B}{1\,\mathrm{G}})$ eV. Laboratory experiments on n-n' are typically done for $B < 1$ G. For magnetars, $B \simeq (10^{12}$–$10^{14})$ G. For the mixing parameter, we may take the value

$$\varepsilon < 1.5 \times 10^{-18}\ \mathrm{eV}$$

from the direct experiment [35,36]. It corresponds to the oscillation time $\tau_{nn'} > 448$ s [35]. This value was obtained under the assumption of zero mirror magnetic field and perfect degeneracy of n and n' masses.

Dropping the assumption $B' = 0$ leads to the dependence of the lower limit $\tau_{nn'}$ on the strength and the angle between B and B' [14,37]. The search for n-n' oscillations accounting for the possible mirror magnetic field has been performed in Ref. [38] with the result $\tau_{nn'} > (6$–$9)$ s for $B' \simeq (10$–$20)$ μT. Clear indications on the signal-like anomalies in favor of B' manifestation have not been observed. Under the assumption $B' = 0$, this experiment yielded the result $\tau_{nn'} > 352$ s ($\varepsilon < 1.9 \times 10^{-18}$ eV). Note that values of ε as large as 10^{-15} eV [12,18] and even $\varepsilon \sim 10^{-13}$ eV [20] are discussed in the literature. Note that ε sharply depends on the mass splitting $\Delta m = m - m'$. We recall that in this article, we consider the perfectly degenerate masses of n and n'. Otherwise, Δm might have contributed to d. The detailed investigation of $\Delta m \neq 0$ case is given in Ref. [39]. We stick to the exact Z_2 symmetry in order not to make less clear our core message, and not to introduce additional complications. The above estimates and considerations allow us to conclude that $d \gg \varepsilon$ is a reliable conjecture.

The time evolution of $\hat{\rho}(t)$ is described by the von Neumann–Liouville equation

$$\frac{d\hat{\rho}}{dt} = -i[\hat{H}, \hat{\rho}]. \tag{3}$$

It yields four coupled linear differential equations:

$$\dot{\rho}_{11} = -i\varepsilon(\rho_{21} - \rho_{12}), \tag{4}$$

$$\dot{\rho}_{12} = i\varepsilon(\rho_{11} - \rho_{22}) - id\rho_{12}, \tag{5}$$

$$\dot{\rho}_{21} = -i\varepsilon(\rho_{11} - \rho_{22}) + id\rho_{21}, \tag{6}$$

$$\dot{\rho}_{22} = i\varepsilon(\rho_{21} - \rho_{12}). \tag{7}$$

The quantity of interest $\rho_{22}(t) = |\psi_{n'}(t)|^2$ may be found directly by solving this set of equations. For further purposes, it makes sense to act differently—namely, to reformulate the von Neumann–Liouville equation in terms of the Bloch 3-vector $\boldsymbol{R}$ [31,32]. It is introduced by the expansion of the density matrix over the Pauli matrices:

$$\hat{\rho} = \frac{1}{2}(1 + \boldsymbol{R} \cdot \boldsymbol{\sigma}), \tag{8}$$

$$\boldsymbol{R} = \begin{pmatrix} \rho_{12} + \rho_{21} \\ -i(\rho_{21} - \rho_{12}) \\ \rho_{11} - \rho_{22} \end{pmatrix}. \tag{9}$$

The von Neumann–Liouville equation (3) is equivalent to the equation of motion of the Bloch vector

$$\dot{\boldsymbol{R}} = \boldsymbol{V} \times \boldsymbol{R}, \tag{10}$$

where

$$\boldsymbol{V} = \begin{pmatrix} 2\varepsilon \\ 0 \\ d \end{pmatrix}. \tag{11}$$

This form of the Bloch equation was proposed by Leo Stodolsky [28]. Equation (10) describes the precession of the Bloch vector $\boldsymbol{R}$ around the "magnetic field" $\boldsymbol{V}$. According to Eq. ([10]), the length of $\boldsymbol{R}$ does not change. This means the absence of decoherence. Decoherence does not happen as soon as the system is isolated from the environment. In terms of $\boldsymbol{R}$, the system of equations (4)–(7)

reduces to the following:

$$\dot{R}_x = -dR_y, \quad \dot{R}_y = dR_x - 2\varepsilon R_z, \quad \dot{R}_z = 2\varepsilon R_y. \tag{12}$$

We solve these equations with the initial condition $\boldsymbol{R}(t=0) = (0,0,1)$, which means that the system is initially in the $|n\rangle$ state. Taking in Eq. (12) the second time derivative of R_y and using the other two equations, one arrives at the result

$$\ddot{R}_y + \mathcal{D}^2 R_y = 0, \tag{13}$$

where

$$\mathcal{D}^2 = d^2 + \varepsilon^2. \tag{14}$$

Insertion of the solution of (13) into the third equation (12) yields

$$R_z(t) = \rho_{11} - \rho_{22} = 1 - \frac{8\varepsilon^2}{\mathcal{D}^2}\sin^2\frac{\mathcal{D}}{2}t. \tag{15}$$

Invoking the condition $\rho_{11} + \rho_{22} = 1$, one obtains the well-known result

$$\rho_{22}(t) = |\psi_{n'}(t)|^2 = \frac{4\varepsilon^2}{\mathcal{D}^2}\sin^2\frac{\mathcal{D}}{2}t. \tag{16}$$

In the ultrashort time limit $\mathcal{D}t \ll 1$, Eq. (16) yields

$$\rho_{22}(t) \approx \varepsilon^2 t^2. \tag{17}$$

This law does not allow us to define the transition probability per unit time [25]. The oscillation formula (16) could have been obtained by solving the Schrödinger equation. Our purpose was to introduce the density matrix formalism which lies at the core of the *correct approach* to quantum systems in contact with the environment. The problem of n-n' oscillations in the framework of the density matrix formalism with absorption has been considered in Ref. [40].

III. LINDBLAD AND BLOCH EVOLUTION EQUATIONS

The departure from ordinary quantum mechanics is needed to describe a system that interacts with the surroundings. The Lindblad and Bloch equations are the instruments suited for this purpose. The n-n' system in a NS may be regarded as a part of a large system composed of the n-n' subsystem and the ensemble of neutrons forming the NS. The density matrix of the full system evolves according to the von Neumann–Liouville equation (3). The reduced density matrix [26,27] of the n-n' subsystem is obtained by the partial trace over the unobserved states of the environment [26]:

$$\hat{\rho}(t) = \mathrm{tr}_{\mathrm{env}}(\hat{\rho}_{\mathrm{full}}(t)). \tag{18}$$

The evolution of the reduced density matrix is described by the Lindblad equation [29,30]

$$\frac{d\hat{\rho}}{dt} = -i[\hat{H}, \hat{\rho}] + \sum_n \left[\hat{L}_n \hat{\rho} \hat{L}_n^\dagger - \frac{1}{2}\{\hat{L}_n^\dagger \hat{L}_n, \hat{\rho}\}\right]. \tag{19}$$

The first term in Eq. (19) is the usual Schrödinger term, $\hat{L}_n$ is a set of additional operators called Lindblad or jump operators, and $\{\cdot,\cdot\}$ is the anticommutator. The second term in (19) is the dissipative one responsible for the loss of information into the environment. It transforms pure states into mixed ones, inducing decoherence. A lucid pedagogical derivation of the Lindblad equation may be found in Ref. [41]. This equation has been used to describe oscillations in a wide range of physical systems, from the atomic clock [42] and ultracold atoms [43] to oscillating neutrino scattering on leptons in plasma [44], heavy quark dynamics in quark-gluon plasma [45], bottomonium regeneration [46], and Hawking radiation [47].

IV. NEUTRON–MIRROR-NEUTRON CONVERSION IN A NEUTRON STAR

First, we formulate the underlying assumptions under which we consider the n-n' conversion in the NS. They include the following:

(1) Z_2 mirror symmetry is exact, so that n and n' have equal masses.
(2) The mirror neutron does not scatter the ordinary-matter (OM) neutrons forming the NS, and there is no n'-n regeneration.
(3) The admixture of MM in the NS which may cause n'-n' scattering is neglected.
(4) β-decay in both n and n' channels is temporarily omitted.

Some comments on the above assumptions are needed. There are arguments to soften the exact Z_2 symmetry. Big bang nucleosynthesis requires asymmetric inflationary reheating in the two sectors [6,48]. In addition, the breaking of Z_2 may serve as an explanation of the neutron lifetime puzzle [15]. As mentioned previously, in order to avoid complications unrelated to the topic of this work, we assume the perfect n and n' mass degeneracy. The next remark is that by NS number density, we imply the neutron density leaving aside the admixture of protons and hyperons. The n'-n regeneration is blocked by the Pauli principle. For demonstration purposes, we choose a number density twice the value of the normal nuclear density, keeping in mind the density variation from the crust to the core.

Under the listed assumptions, the elements entering into the Lindblad equation (19) are of the following form:

$$\hat{H} = \begin{pmatrix} E - \frac{2\pi}{k} nv\mathrm{Re}f(0) & \varepsilon \\ \varepsilon & E' \end{pmatrix}, \tag{20}$$

$$\hat{L} = \sqrt{nv}\hat{F}, \qquad \hat{F} = \begin{pmatrix} f(\theta) & 0 \\ 0 & 0 \end{pmatrix}. \tag{21}$$

Here, k is the neutron momentum, n is the NS number density, v is the neutron velocity, and $f(\theta)$ is the neutron-neutron elastic scattering amplitude. The term $\frac{2\pi}{k} nv\mathrm{Re}f(0)$ corresponds to the energy shift due to forward scattering. Inserting Eqs. (20) and (21) into Eq. (19), one arrives, after some simple algebra, at a set of four coupled differential equations:

$$\dot{\rho}_{11} = -i\varepsilon(\rho_{21} - \rho_{12}), \tag{22}$$

$$\dot{\rho}_{22} = i\varepsilon(\rho_{21} - \rho_{12}), \tag{23}$$

$$\dot{\rho}_{12} = -M\rho_{12} - i(d + K)\rho_{12} + i\varepsilon(\rho_{11} - \rho_{22}), \tag{24}$$

$$\dot{\rho}_{21} = -M\rho_{21} + i(d + K)\rho_{21} - i\varepsilon(\rho_{11} - \rho_{22}). \tag{25}$$

The quantities M and K stand for

$$M = \frac{2\pi}{k} nv\,\mathrm{Im}\,f(0), \qquad K = -\frac{2\pi}{k} nv\mathrm{Re}f(0). \tag{26}$$

According to the optical theorem $M = \frac{1}{2} n\sigma v$, with σ being the nn cross section, K gives the energy due to the index of refraction. The coefficient $\frac{1}{2}$ in M comes from the active component which can scatter [28,33,34]. To our knowledge, the set of equations (22)–(25) does not have a transparent analytical solution.

Equations (22)–(25) are equivalent to the following matrix equation for the Bloch vectors $\vec{R}$:

$$\dot{\vec{R}} = \vec{V} \times \vec{R} - D_T R_T, \tag{27}$$

where

$$\vec{V} = \begin{pmatrix} 2\varepsilon \\ 0 \\ d + K \end{pmatrix}, \quad D_T = \begin{pmatrix} M & 0 \\ 0 & M \end{pmatrix}, \quad R_T = \begin{pmatrix} R_x \\ R_y \end{pmatrix}. \tag{28}$$

The difference of Eqs. (27) and (28) from Eqs. (10) and (11) is the presence of the term $D_T R_T$, which represents the quantum friction and is responsible for the loss of coherence. It leads to the destruction of the off-diagonal elements of the density matrix.

Like the Lindblad equations (22)–(25), the Bloch equations (27) and (28) can hardly be solved analytically. To proceed further, we drop the component $V_z = (d + K)$ of $\vec{V}$. It corresponds to the contributions of the external magnetic field and the index of refraction. The component V_z does not participate in decoherence. With V_z omitted, the evolution equations for R_z and R_y are

$$\dot{R}_z = 2\varepsilon R_y,$$
$$\dot{R}_y = -MR_y - 2\varepsilon R_z. \tag{29}$$

Taking in Eq. (29) the second time derivative of R_z, one arrives at the equation

$$\ddot{R}_z + M\dot{R}_z + 4\varepsilon^2 R_z = 0. \tag{30}$$

This is the equation of an oscillator with friction M. It possesses three types of solutions depending on the relative magnitude of the damping parameter M and the tunneling parameter ε. In the next section, we shall relate M to the flight time t_{nn} between collisions. The connection is not entirely simple, as might be expected. The character of the solution depends on the sign and the value of the quantity

$$\Omega^2 = \frac{M^2}{4} - 4\varepsilon^2. \tag{31}$$

If $\Omega^2 > 0$, the interaction with the environment destroys the off-diagonal elements of the density matrix. This is the overdamping regime. The opposite underdamped case $\Omega^2 < 0$ corresponds to slightly or strongly suppressed oscillations. The critical damping $\Omega^2 = 0$ implies a fine-tuning of M and ε. Finally, for $M = 0$, the system is in a free oscillation regime.

It is not difficult to prove that n-n' conversion in the NS proceeds in the overdamped mode. To be aware of this, we compare the values of M and ε. For the value of ε, we set $\varepsilon = 1.5 \times 10^{-18}$ eV $\simeq 2 \times 10^{-3}$ s^{-1} from the direct experiment [35]. As we mentioned in Sec. II, a wide interval of ε values is discussed in literature. To estimate M, we proceed as follows. We take $n = 2n_0 = 0.34$ fm^{-3} for the number density. There is scarce information regarding the nn collision cross section in the NS. We employ the data plotted in Fig. 7 of Ref. [49]. At the above density, $\sigma \simeq 30$ mb $= 3$ fm^2. The authors of Ref. [49] emphasize that the cross section in medium is not a well-defined quantity, and hence the above number should be taken with caution. The neutron velocity may be estimated as $v \simeq \frac{k_F}{m} \simeq \frac{(3\pi^2 n)^{1/3}}{m} \simeq 0.4$.

With the above numbers, one gets the following estimate for M:

$$M = \frac{1}{2} n\sigma v \simeq 0.4 \times 10^8 \text{ eV} = 0.6 \times 10^{23} \text{ s}^{-1}. \tag{32}$$

These crude estimates confirm the conclusion that $M \gg \varepsilon$, namely $M/\varepsilon \sim 10^{25}$–$10^{26}$.

The solution of Eq. (30) with the initial condition $R_z(0) = 1$ (pure $|n\rangle$ state) is given by

$$R_z(t) = e^{-\frac{Mt}{2}}\left[\frac{\frac{M}{2}+\Omega}{2\Omega}e^{\Omega t} + \frac{\frac{M}{2}-\Omega}{2\Omega}e^{-\Omega t}\right]. \quad (33)$$

The two terms in Eq. (33) behave as $\exp(-\omega_i t)$, $i = 1, 2$ with dissimilar frequencies

$$\omega_1 = \frac{M}{2} - \Omega \simeq \frac{4\varepsilon^2}{M}, \qquad \omega_2 = \frac{M}{2} + \Omega \simeq M. \quad (34)$$

Note that $\omega_1/\omega_2 \simeq 4\varepsilon^2/M^2 \ll 1$. It is clear that the large-time ($t \gg 1/M$) behavior is dominated by the first term with the eigenfrequency $\omega_1 = 4\varepsilon^2/M$. The second term can only be probed at short times, $t \ll 1/M$. Interference effects at intermediate times is the subject of a paper in preparation. At large time values, $R_z(t)$ becomes an overdamped solution describing the relaxation without oscillations with a constant decoherence rate,

$$R_z(t) \sim \exp\left(-\frac{4\varepsilon^2}{M}t\right). \quad (35)$$

This form of $R_z(t)$ is in complete agreement with the corresponding results of Refs. [28,34]. Invoking the requirement that $R_z = \rho_1 - \rho_2$, one can split $R_z(t)$ into two terms:

$$R_z(t) = \rho_1 - \rho_2 \simeq \exp\left(-\frac{4\varepsilon^2}{M}t\right) - \frac{4\varepsilon^2}{M^2}\exp\left(-\frac{4\varepsilon^2}{M}t\right). \quad (36)$$

This expression is becoming exact for $t \gg 1/M$. The first term in (36) describes the evolution of $|\psi_n(t)|^2$. The disappearance time $T = M/4\varepsilon^2$ gets tremendously prolonged due to high-frequency collisions inside the NS. The mirror neutron component $|\psi_{n'}(t)|^2$ is given by the second term and is damped by a huge factor $4\varepsilon^2/M^2$ at all times. An equation similar to (36) has been derived in Ref. [50] for the neutron-antineutron conversion in nuclei. The same type of equations describe the probability of a photon oscillating into a hidden photon in the medium [51,52]. According to Eq. (36), the n-n' transition rate $\Gamma(n\text{-}n')$ is equal to

$$\Gamma(n\text{-}n') = \frac{4\varepsilon^2}{M} \simeq 10^{-20}\ \text{yr}^{-1} \simeq 0.2 \times 10^{-51}\ \text{GeV} \quad \text{for } \varepsilon = 1.5 \times 10^{-18}\ \text{eV}. \quad (37)$$

Note that the quantity M does not come from the two-by-two Hamiltonian. It originates from the Lindblad operator and is responsible for the decoherence. In what follows, we shall discuss the relation between M and the time interval t_{nn} between collisions.

V. IMPLICATIONS

We have shown that n-n' oscillations in NSs are overdamped due to the difference in scale between the mixing parameter ε and the collision friction parameter M. The n-n' system evolves slowly compared to the timescales of its interaction with the environment. It is desirable to obtain the relation between the damping factor M and the time interval t_{nn} between collisions. We have estimated M as $M \simeq 10^{23}\ \text{s}^{-1}$ [Eq. (32)]. On dimensional grounds, it is clear that $t_{nn} \sim 1/M$. It turns out that to find the proportionality coefficient is not a trivial matter. Recall that according to Eq. (26), $M = \frac{1}{2}n\sigma v$. To relate M and t_{nn}, it is tempting to bring in the simplest expression $l = 1/(n\sigma)$ for the neutron mean free path (MFP) and to obtain $\Gamma(n\text{-}n') = 8\varepsilon^2 t_{nn}$, replacing Eq. (37). The result—namely, the coefficient 8—can hardly be considered reliable. The evaluation of the MFP in dense interacting matter is a complicated problem involving the whole machinery of many-body theory. The encountered difficulties include the choice of the potential, the account of the Pauli principle, the three-body forces, the geometry of the trajectory, etc. Despite the great efforts of several authors in various approaches (see a list of references in [49]), the problem is still pending. Therefore, the expression $l = 1/(n\sigma)$, as well as $\Gamma_{nn} = \sigma_{nn} v n = 1/t_{nn}$ used in Ref. [20] is merely correct on the dimensional grounds. Corrections induced by Pauli blocking and medium-dependent energy splitting between n and n' have been proposed in Refs. [18,21]. However, the approach based on the Schrödinger equation with a two-by-two matrix Hamiltonian cannot describe the damping of the n-n' conversion caused by the collisional decoherence. The Lindblad term (19) must be added [27–30,42]. The relaxation pattern [Eqs. (35) and (36)] sets in instead of oscillations. The transition rate in given by Eq. (37) or by $\Gamma(n\text{-}n') \sim \varepsilon^2 t_{nn}$ with the proportionality coefficient, and t_{nn} can at best be estimated in a model-dependent way.

VI. CONCLUSIONS AND OUTLOOK

In this paper, a new approach to the process of neutron–mirror-neutron transition in NSs is proposed. The research is driven by the idea that high-frequency collisions with the surrounding neutrons cause a deep change in the way the n-n' conversion proceeds. The n-n' can be regarded as a subsystem in contact with the environment. Its time evolution is described as the equation of motion of the reduced density matrix $\hat{\rho}(t)$. It is no longer the von Neumann–Liouville one but contains the new terms in the form of Lindblad operators. These terms are responsible for the interaction with the environment. Under their action, the off-diagonal elements of $\hat{\rho}(t)$ decay to zero.

This process is referred to as decoherence or density matrix collapse. The Lindblad equation may be recast into a Bloch equation for the 3-vector $\vec{R}$ on the Bloch sphere. The equation for $\vec{R}$ is the equation for the oscillator with damping parameter M, which is proportional to the rate of neutron collisions in the NS. For the NS density, this rate is ∼25 orders of magnitude higher than the predicted n-n' oscillation rate. Overdamping with no oscillations at all times takes place.

For "long" times $t \gtrsim 10^{-23}$ s, the time dependence of ordinary and mirror components reads

$$\rho_{11} = |\psi_n(t)|^2 \simeq \exp\left(-\frac{4\varepsilon^2}{M}t\right),$$
$$\rho_{22} = |\psi_{n'}(t)|^2 \simeq \frac{4\varepsilon^2}{M^2}\exp\left(-\frac{4\varepsilon^2}{M}t\right). \quad (38)$$

In order to compare with the results of other authors, it is instructive, following Refs. [18,53], to express the mixing parameter ε in terms of $\varepsilon_{15} = \frac{\varepsilon}{10^{-15}\ \mathrm{eV}}$. For $\varepsilon = 1.5 \times 10^{-18}$ eV, one has $\varepsilon_{15}^2 = 2.25 \times 10^{-6}$. With this value of ε_{15}^2 and the ratio $\xi = 2$ of the NS density to normal nuclear density $n_0 = 0.16$ fm^{-3}, Eq. (40) of Ref. [18] for the effective "starting" rate of the NS transformation yields $\Gamma(n\text{-}n') \simeq 0.4 \times 10^{-52}$ GeV, and Eq. (18) of Ref. [53] gives $\Gamma(n\text{-}n') \simeq 0.2 \times 10^{-52}$ GeV (in the last case, the NS density is the average over the NS profile).

Note that about an order of magnitude difference between the results of [18,53] and Eq. (37) is not very significant in view of the fact that all three results for $\Gamma(n\text{-}n')$ predict a conversion time larger than the Universe's age. What distinguishes our result from that of Refs. [18,53] is that we predict the MM admixture to be small at all times, while according to Refs. [18,53], the NS evolution ends up with equal amounts of OM and MM. And we do not need to say that the theoretical approaches of Refs. [18,53] and the present one are different.

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DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

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