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Gauged $B - L$ number and neutron–antineutron oscillation: long-range forces mediated by baryophotons

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Abstract Transformation of a neutron to an antineutron $n \to \tilde{n}$ has not yet been experimentally observed. In principle, it can occur with free neutrons in the vacuum or with neutrons bound inside the nuclei. In a nuclear medium the neutron and the antineutron have different potentials and for that reason n–$\tilde{n}$ conversion in nuclei is heavily suppressed. This transformation can also be suppressed for free neutrons in the presence of an environmental vector field that distinguishes the neutron from the antineutron. We consider the case of a gauge field coupled to the $B - L$ charge of the particles ($B - L$ photon), and we show that discovery of n–$\tilde{n}$ oscillation in experiment will lead to few order of magnitudes stronger limits on its coupling constant than present limits from the tests of the equivalence principle. If n–$\tilde{n}$ oscillation will be discovered via nuclear instability, but not in free neutron oscillations at a corresponding level, this would indicate the presence of such environmental fifth forces. In the latter case the $B - L$ potential can be measurable by varying the external magnetic field for achieving the resonance conditions for n–$\tilde{n}$ conversion. As for neutron–mirror neutron oscillation, such potentials should have no effect once the fifth forces are associated to a common quantum number $(B - L) - (B' - L')$ shared by the ordinary and mirror particles.

1 Introduction

The oscillation phenomenon between the neutron and antineutron was suggested in the early 1970s by Kuzmin [1]. The first theoretical scheme for n–$\tilde{n}$ oscillation was proposed by Mohapatra and Marshak [2], followed by other models as e.g. [3–5]. Experimental observation of the transformation of neutron to antineutron, $n \to \tilde{n}$, would be a manifestation of baryon number violation by two units, from $B = +1$ for the neutron to $B = -1$ for the antineutron. This will also be a demonstration that one of the Sakharov conditions [6,7] required for the generation of baryon asymmetry in the Universe is indeed realized in nature. So far $n \to \tilde{n}$ conversion was not experimentally observed, however, this does not exclude the possibility that it can be a rare/suppressed process. For a review of the present theoretical and experimental situation on n–$\tilde{n}$ oscillation; see [8–10].

The neutron–antineutron mixing which violates baryon number by two units is induced by a small Majorana mass term $\frac{1}{2}\varepsilon_{n\tilde{n}}(n^T C n + \text{h.c.}) = \frac{1}{2}\varepsilon_{n\tilde{n}}(\overline{n^c}\, n + \text{h.c.})$, where $n^c = \tilde{n} = C\overline{n}^T$ is the antineutron field. As far as the neutron is a composite particle, n–$\tilde{n}$ mixing can be induced by six-fermion effective operators involving the first family quarks. In terms of the standard model ingredients, left-handed $q = (u, d)_L$ and right-handed $u = u_R$, $d = d_R$ quarks, these $D = 9$ operators read

$$\frac{1}{\mathscr{M}^5}\big(uddudd + qqdqqd + uddqqd\big) + \text{h.c.} \qquad (1)$$

where $\mathscr{M}$ is some large mass scale of new physics beyond the standard model. These operators can have different convolutions of the Lorentz, color and weak isospin indices which are not specified, and the coupling constants in front of different terms are suppressed.[1] Models of Refs. [2–5] are just different theoretical realizations for these operators. Taking matrix elements from these operators between the neutron and antineutron states, one can estimate

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[1] Needless to say, the combination qq in the third term in (1) must be in a weak isosinglet combination, $qq = \frac{1}{2}\epsilon^{\alpha\beta} q_\alpha q_\beta = u_L d_L$, where $\alpha, \beta = 1, 2$ are the weak $SU(2)$ indices, while in the second term qq can be taken in a weak isotriplet combination as well.

$$\varepsilon_{n\tilde{n}} \sim \frac{\Lambda_{\rm QCD}^6}{\mathscr{M}^5} \sim \left(\frac{1\,\text{PeV}}{\mathscr{M}}\right)^5 \times 10^{-25}\,\text{eV}. \quad (2)$$

The Clebsch coefficients of the matrix elements $\langle\tilde{n}|uddudd|n\rangle$ for different Lorentz and color structures of operators (1) were studied in Ref. [11], but we do not concentrate here on these details and take them as $O(1)$ factors. More generally, having in mind that all quark families can be involved, operators like (1) can induce the mixing phenomena also for other neutral baryons, e.g. between the hyperon Λ and the anti-hyperon $\bar{\Lambda}$.

Concerning the bounds on the n–$\tilde{n}$ oscillation time $\tau_{n\tilde{n}} = 1/\varepsilon_{n\tilde{n}}$, the direct limit from experiment with free neutrons yields $\tau_{n\tilde{n}} > 0.86 \times 10^8$ s [12]. The nuclear stability limits, with uncertainties in the evaluation of nuclear matrix elements, translate into $\tau_{n\tilde{n}} > 1.3 \times 10^8$ s [13] and $\tau_{n\tilde{n}} > 2.7 \times 10^8$ s [14]. The latter sets the strongest upper limit on n–$\tilde{n}$ mixing, $\varepsilon_{n\tilde{n}} < 2.5 \times 10^{-24}$ eV. The future long-baseline experiment with cold neutrons at the European Spallation Source (ESS) can improve the existing limits on $\tau_{n\tilde{n}}$ by more than an order of magnitude [8], reaching the sensitivity for n–$\tilde{n}$ mixing down to $\varepsilon_{n\tilde{n}} \sim 10^{-25}$ eV, and thus can test underlying B-violating physics producing operators (1) up to the effective scales $\mathscr{M} \sim 1$ PeV.

One can consider a scenario when baryon number is broken not explicitly but spontaneously. Such a baryon symmetry can be global or local, with different physical implications. The possibility of spontaneous violation of global lepton symmetry after which the neutrinos can get non-zero Majorana masses was widely discussed in the literature. As a result, a Goldstone boson should appear in the particle spectrum, named majoron [15]. Spontaneous violation of global baryon number in connection with the Majorana mass of the neutron was first discussed in Ref. [16], in the context of the Mohapatra–Marshak model [2]. Recently the discussion was revived by one of us in Ref. [17], where also the seesaw model for the n–$\tilde{n}$ transition with low-scale spontaneous violation of baryon number was suggested. An associated Goldstone particle – a baryo-majoron – can have observable effects in n–$\tilde{n}$ transitions in nuclei or dense nuclear matter. The low-scale baryo-majoron model [17] has many analogies with the low-scale Majoron model for the neutrino masses [18]. By extending baryon number to $B - L$ symmetry, the baryo-majoron can be identified with the ordinary majoron associated with the spontaneous breaking of lepton number, with interesting implications for neutrinoless 2β decay with the majoron emission [18–20], for matter-induced effects of the neutrino decay [21–25] and for the active–sterile neutrino oscillations in the early Universe [26,27].

In this paper we discuss a situation when baryon number is related to a local gauge symmetry. The idea to describe the conservation of baryon number B and lepton number L similar to the conservation of electric charge by introducing gauge symmetries $U(1)_B$ and $U(1)_L$, i.e. in terms of baryon or lepton charges coupled to the massless vector fields of leptonic or baryonic photons with tiny coupling constants, was suggested a long time ago [28–30]. Their effects for the neutron oscillations were discussed in Refs. [31–33]. Nowadays the limits on such interactions are very stringent. The best limits on the coupling strength of baryonic and leptonic photons were obtained from the Eötvös type of experiments testing the equivalence principle [34]. Then the common sense argument used here was that coupling of such photons is many orders of magnitude weaker than the gravitational interaction between baryons or leptons and therefore such photons are likely non-existent. We will try to revise this concept.

Since baryon number B and lepton number L separately are not conserved due to non-perturbative effects, it is difficult to promote them as gauge symmetries without altering the particle content of the standard model, i.e. without introducing new exotic particles just for cancellation of electroweak anomalies. However, in the standard model $B - L$ symmetry is anomaly free and respective current is conserved at the non-perturbative level as well. Therefore, we discuss not baryonic and leptonic photons separately but the vector field b_μ associated with $U(1)_{B-L}$ gauge symmetry. Clearly, such a $B - L$ photon couples with opposite charges not only between the baryons and anti-baryons (and between the leptons and anti-leptons), but also between the baryons and leptons. Thus, $B - L$ charge of the neutral hydrogen atom is zero while the $B - L$ charge of heavier neutral atoms is determined by the number of neutrons in the nuclei, and so the regular matter built by nuclei heavier than hydrogen is $B - L$ charged.

On the other side, n–$\tilde{n}$ mixing cannot be induced without violating B and $B - L$, and thus its existence would imply that gauge symmetry $U(1)_{B-L}$, if it exists, should be spontaneously broken, which should also render the $B - L$ photon massive. However, if its gauge coupling constant is very small, it can remain extremely light, and it can mediate observable long-range forces (fifth force) between material bodies even at large astronomical distances.

2 Experimental limits on B-L photons

The baryophoton b_μ associated with $U(1)_{B-L}$ gauge symmetry is coupled as $gb_\mu Q\overline{\psi}\gamma^\mu\psi$ to a current of fermion ψ bearing a non-zero charge $B - L = Q$, where g is the gauge coupling constant.[2] Namely, with the neutron, proton, electron and neutrino it interacts as $gb_\mu(\overline{n}\gamma^\mu n + \overline{p}\gamma^\mu p -$

[2] Let us remark that, for promoting $U(1)_{B-L}$ as a gauge symmetry, along with the gauge symmetry $SU(3) \times SU(2) \times U(1)$ of the standard model, in addition to the quarks $q_L = (u, d)_L$, u_R, d_R with $Q = 1/3$ and leptons $l_L = (\nu, e)_L$, e_R with $Q = -1$, also right-handed neutrino fields, ν_R, with $Q = -1$ should be introduced per each generation for cancellation of $U(1)^3_{B-L}$ triangle anomalies.

$\overline{e}\gamma^\mu e - \overline{\nu}\gamma^\mu \nu$). As far as the existence of n–$\tilde{n}$ mixing implies violation of $U(1)_{B-L}$, baryophoton cannot remain massless.

In particular, $D = 9$ effective operators (1) are now forbidden by $U(1)_{B-L}$ symmetry. One can consider instead the higher-order effective operators involving a complex scalar field χ bearing two units of $B - L$ charge, $Q_\chi = 2$:

$$\frac{\chi^\dagger}{M^6}\,(uddudd \,+\, qqdqqd \,+\, uddqqd) \,+\, \text{h.c.} \tag{3}$$

where M is some large mass scale. These operators can be effectively obtained via renormalizable couplings, e.g. in the context of seesaw-like scheme suggested in Refs. [4,17].

At low energies operators (3) reduce to the Yukawa coupling of scalar χ with the neutron,

$$\frac{y}{\sqrt{2}}\left(\chi^\dagger n^T C n \,+\, \text{h.c.}\right) \tag{4}$$

with the coupling constant $y \sim \left(\Lambda_{\text{QCD}}/\Lambda\right)^6$. Let us assume that χ gets a non-zero vacuum expectation value (VEV) $\langle\chi\rangle = v_\chi/\sqrt{2}$ from the potential $\mathcal{V}(\chi) = \frac{\lambda}{2}\left(\chi^\dagger\chi - v_\chi^2/2\right)^2$. In the vicinity of the ground state the field χ can be expanded as $\chi = 2^{-1/2}(v_\chi + \rho)\exp(i\beta/v_\chi)$, where ρ is the massive (Higgs) mode with a mass $m_\rho = \sqrt{\lambda}\,v_\chi$ and β is the massless Goldstone mode, so that the Yukawa coupling (4) reads[3]

$$\frac{y}{2}\left(v_\chi\,\overline{n^c}\,n \,+\, \rho\,\overline{n^c}\,n \,+\, i\beta\,\overline{n^c}\gamma^5 n \,+\, \text{h.c.}\right). \tag{5}$$

Hence, the VEV v_χ induces n–$\tilde{n}$ mass mixing,

$$\varepsilon_{n\tilde{n}} = y v_\chi \sim \frac{v_\chi \Lambda_{\text{QCD}}^6}{M^6} \sim \left(\frac{v_\chi}{1\,\text{GeV}}\right)\left(\frac{100\,\text{TeV}}{M}\right)^6 \times 10^{-25}\,\text{eV}. \tag{6}$$

For a benchmark value $v_\chi = 1\,$GeV, the neutron–antineutron mixing can be within the experimental reach at the ESS, i.e. $\varepsilon_{n\tilde{n}} > 10^{-25}\,$eV, for the scale M as large as 100 TeV. A priori, the scale v_χ could be much smaller. E.g., $\varepsilon_{n\tilde{n}} = 10^{-25}\,$eV could be obtained for $v_\chi = 1\,$keV and $M = 10\,$TeV, or even for $v_\chi = 10^{-3}\,$eV and $M = 1\,$TeV. However, as we show below, for a given value of $\varepsilon_{n\tilde{n}}$ the scale v_χ is limited by the nuclear stability bounds, namely by the bounds on the neutron decay into antineutron in dense nuclear matter, a phenomenon resembling the matter-induced neutrino decay with the majoron emission [21,22].

Our reasoning goes as follows. While in vacuum these transitions are suppressed since the neutron and antineutron states are degenerate, in nuclei their energy levels are different and thus the neutron state can decay into the antineutron one with emission of ρ and β. Then antineutron promptly annihilates with other nucleons producing pions with total energy roughly equal to two nucleon masses. The decay rate per neutron can be estimated as $\Gamma(n \to \tilde{n} + \rho, \beta) = (y^2/4\pi)\Delta E_{n\tilde{n}}$ [17], where $\Delta E_{n\tilde{n}}$ is the energy gap the neutron and antineutron states in the nucleus. which is typically $\sim$100 MeV for heavier nuclei.[4]

The present limits on the n–$\tilde{n}$ transition lifetime in nuclei are at the level of 10^{32} year. Namely, the experimental limits imply $\Gamma_{n\tilde{n}}^{-1} > 7.2\times10^{31}$ yr for iron [13] and $\Gamma_{n\tilde{n}}^{-1} > 1.9\times10^{32}$ yr for oxygen [14]. Taking into account that, for both elements, $\Delta E_{n\tilde{n}} \simeq 100\,$MeV [36], the latter bounds translate into a conservative upper limit on the Yukawa coupling constant $y < 10^{-31}$ or so [17]. Thus, in view of the relation $y = \varepsilon_{n\tilde{n}}/v_\chi$, we obtain the following limit on the $U(1)_{B-L}$ symmetry breaking scale:

$$v_\chi > 10^{31}\varepsilon_{n\tilde{n}} = \left(\frac{\varepsilon_{n\tilde{n}}}{10^{-25}\,\text{eV}}\right) \times 1\,\text{MeV}. \tag{7}$$

Hence, n–$\tilde{n}$ mixing exists at the level experimentally accessible at the ESS, $\varepsilon_{n\tilde{n}} = 10^{-25}\,$eV or so, then the nuclear stability bounds will give the limit $v_\chi > 1\,$MeV.

On the other hand, the VEV $\langle\chi\rangle = v_\chi/\sqrt{2}$ spontaneously breaks $U(1)_{B-L}$ and induces the baryophoton mass. In principle, other scalars η_i with non-zero $B - L$ charges Q_i and non-zero VEVs v_i can also contribute in breaking $U(1)_{B-L}$. In this case, the baryophoton mass would be

$$M_b = 2gv, \qquad v = \left[v_\chi^2 + (Q_i/Q_\chi)^2 v_i^2\right]^{1/2} \geq v_\chi. \tag{8}$$

Thus, v_χ defines the minimal possible value of M_b for a given coupling constant g. Contributions of the other scalars η_i can only make M_b larger.

The exchange of massive baryophoton mediates a spin-independent Yukawa-type fifth force between the material bodies. The potential energy of the interaction between the test particle a with non-zero $B-L$ charge Q_a, in our case the neutron or antineutron, and an attractor A (a massive body as e.g. Earth or sun) with overall $B - L$ charge Q_A, reads

$$V_a = \alpha_{B-L}\frac{Q_a Q_A}{r}\,e^{-r/\lambda}, \tag{9}$$

[3] Here we commit some simplification (which however will not affect our results) neglecting another possible coupling $\chi^\dagger n^T\gamma^5 C n$, which could emerge along with the coupling (4) as far as in general the effective operators (3) do not respect parity and also violate CP. The corresponding 'pseudo-scalar' mass term induced by the VEV $\langle\chi\rangle$ can be rotated away [35], however both ρ and β generically will have both scalar and pseudo-scalar couplings. Unfortunately, the effects of P and CP violation in these couplings can be hardy detectable in experiments.

[4] For $U(1)_{B-L}$ being a local gauge symmetry, the majoron is 'eaten' via the Higgs mechanism. However, the decay $n \to \tilde{n} + \beta$ inside the nuclei will still occur, with $\partial_\mu\beta$ constituting the longitudinal component of the massive vector field b_μ. As for the decay $n \to \tilde{n} + \rho$, it will take place if the mass of massive (Higgs) mode ρ is less than $\Delta E_{n\tilde{n}}$.

where $\alpha_{B-L} = g^2/4\pi$, and λ is the Yukawa radius:

$$\lambda = \frac{1}{M_b} \simeq \left(\frac{10^{-49}}{\alpha_{B-L}}\right)^{1/2} \left(\frac{1\,\text{GeV}}{v}\right) \times 8.8 \cdot 10^9\,\text{cm}. \quad (10)$$

The $B - L$ charge of the massive body Q_A depends on its chemical composition, and it is determined by the overall amount of neutrons in this body: due to electric neutrality, the amount of protons and electrons should be equal and their contributions cancel each other. Hence, neglecting less than 1% corrections due to binding energies and due to difference between the neutron and proton masses, we get

$$Q_A = Y_{nA} N_{BA} \simeq Y_{nA} \frac{M_A}{m_n}, \quad (11)$$

where Y_{nA} is the average neutron fraction per baryon in the attractor and N_{BA} is its overall baryon charge, m_n being the neutron mass. In particular, $Q = 0$ for hydrogen and $Q \approx 0.5$ for a typical heavy nuclei.

Different from universal gravity, the $B - L$ force induces for the neutron ($Q_n = 1$) and antineutron ($Q_{\tilde{n}} = -1$) potential energies of different sign, $V_{\tilde{n}} = -V_n < 0$, which can be related to their gravitational potential energy $V^{\text{gr}}_{n,\tilde{n}} = -Gm_n M_A/r$, G being the Newton constant, as

$$V_{n,\tilde{n}}/V^{\text{gr}}_{n,\tilde{n}} = \pm \tilde{\alpha}\, q_A\, e^{-r/\lambda} = \pm \tilde{\alpha}\, Y_{nA}\, e^{-r/\lambda}, \quad (12)$$

where $\tilde{\alpha} = \alpha_{B-L}/Gm_n^2$ is a dimensionless parameter and $q_A = Q_A/(M_A/m_n)$ is the massive objects $B-L$ charge per neutron mass unit which according to Eq. (11) is roughly the average fraction of neutrons in the attractor body, $q_A \approx Y_{nA}$.

The results of torsion-balance tests of the weak equivalence principle from Ref. [34] can be interpreted as limits on fifth forces, and in particular, for the force mediated by the $B - L$ photon. The upper limits on the parameter $\tilde{\alpha}$ as a function of the Yukawa radius λ are given in Fig. 6 of Ref. [34]. (Let us note that in Ref. [34] this value is normalized per atomic mass unit, 1 amu$= 0.99\, m_n$, a negligible difference.) In Fig. 1 we show these limits translated directly for α_{B-L}. As we see, for the values of λ much larger than the Earth radius, $\lambda \gg R_\oplus = 6 \cdot 10^8$ cm, the upper limit on α_{B-L} becomes practically independent on λ and it corresponds to $\alpha_{B-L} < 10^{-49}$, or $\tilde{\alpha} < 1.7 \times 10^{-11}$ [34].

Now we can estimate the neutron potential energy V_n as produced due to the baryophoton field by the Earth, sun and our galaxy, relative to the corresponding gravitational potential energies which in fact are substantial. Namely, the Earth induces the gravitational potential energy for the neutron on the Earth's surface as $V^{\text{gr}}_{nE} = -Gm_n M_\oplus/R_\oplus \approx -0.66\,$eV. The sun's contribution is bigger by an order of magnitude,

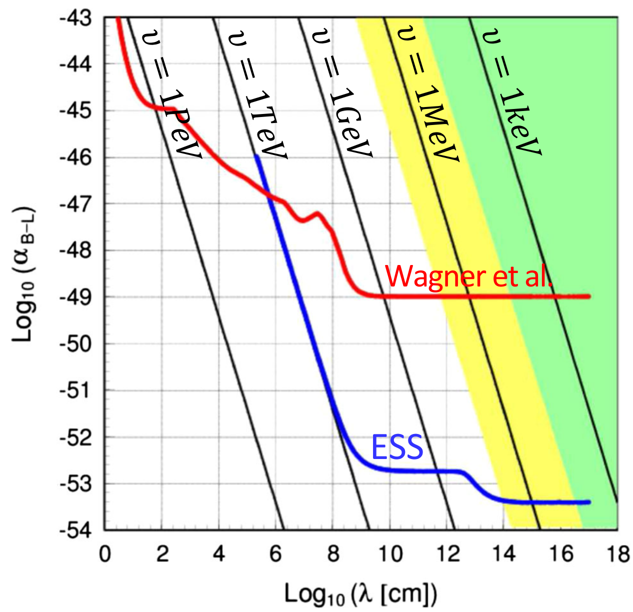

Fig. 1 Limits on α_{B-L} as a function of the Yukawa radius λ. Region above the *red curve* is experimentally excluded by the results of Wagner et al. [34]. The correlation between α_{B-L} and λ of Eq. (10) for different values of $U(1)_{B-L}$ breaking VEV v are shown for the sake of demonstration of the possible scale of the mechanism. *Light* and *dark shaded* regions are excluded by the nuclear stability limits, Eq. (15), respectively, for $\varepsilon_{n\tilde{n}} > 2.5 \times 10^{-24}$ eV and $\varepsilon_{n\tilde{n}} > 10^{-26}$ eV. The region above the *blue curve* can be excluded if the ESS experiment will discover n–$\tilde{n}$ mixing at the level $\varepsilon_{n\tilde{n}} > 10^{-25}$ eV (see Sect. 3)

$V^{\text{gr}}_{nS} = -Gm_n M_\odot/\text{AU} \approx -9.3\,$eV. Finally, the galaxy itself induces even bigger potential $V^{\text{gr}}_{nG} \sim -1\,$keV.[5]

Let us discuss first the potential V_{nE} induced by the Earth on its surface. Since the Earth is built by heavy nuclei, $B - L$ charge of the Earth is approximated as 50% of the number of baryons in the Earth, i.e. $Y_{nE} \simeq 0.5$. Then, assuming that λ is much larger than the radius of the Earth, $\lambda \gg R_\oplus = 6.4 \times 10^8$ cm, and taking the upper bound $\tilde{\alpha} < 1.7 \times 10^{-11}$ [34], we get an upper limit:

$$V_{nE} = -\tilde{\alpha}\, Y_{nE}\, V^{\text{gr}}_{nE} < 5.6 \times 10^{-12}\ \text{eV}. \quad (13)$$

For λ smaller than the Earth diameter, the limits of Ref. [34] allow for larger values of α_{B-L} (see Fig. 1) but the available volume of the source drops roughly as $(\lambda/R_\oplus)^3$ and thus the limit on V_n sharply decreases. The upper limit on V_{nE} as a function of λ corresponding to the limits on α_{B-L} of

[5] Notice that one deals with the gravitational potential which behaves as $\propto 1/r$, while the torsion-balance experiments test gravitational force which falls as $1/r^2$. Thus the hierarchy between the gravitational forces induced by the Earth, sun and the galaxy, becomes reordered in an opposite way with respect to the hierarchy of corresponding potentials. This is the reason why for λ exceeding the Earth diameter, the experimental limits of Ref. [34] become independent of λ.

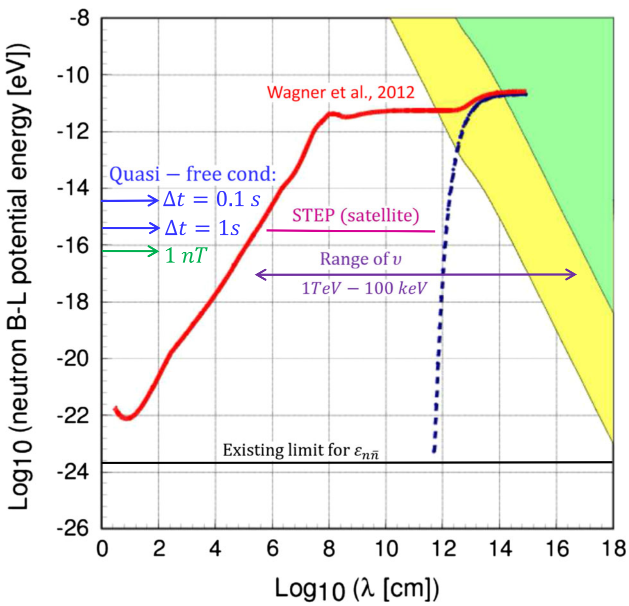

Fig. 2 Potential energy V_n of the neutron in $B-L$ field of Sun and Earth (see text for explanations). The region of potentials V_n above the *red curve* is excluded by torsion-balance experiment limits [34]. The claimed sensitivity limit of the canceled space-mission project STEP [38] is also shown for comparison

Ref. [34], obtained by the integration over the Earth density profile, is shown on Fig. 2.

The sun dominantly consists of hydrogen with vanishing $B-L$, and thus its fifth force is essentially determined by the mass fraction of heavier nuclei (helium, etc.), yielding $Y_{nS} \simeq 0.13$ as one can estimate from the known chemical composition [37]. Hence, taking the upper bound $\tilde{\alpha} < 1.7 \times 10^{-11}$, we obtain[6]

$$V_{nS} = -\tilde{\alpha}\, V_{nS}\, V^{\rm gr}_{nS} \exp(-{\rm AU}/\lambda) < 2 \times 10^{-11}\ \text{eV}, \tag{14}$$

which upper limit can be approached provided that $\lambda \gg 1$ AU, where AU$= 1.5 \times 10^{13}$ cm is the sun–Earth distance (astronomical units). The sun's contribution as a function of λ is shown by a red dashed curve on Fig. 2.

The Yukawa radius λ is related to the baryophoton coupling constant g and the $B-L$ symmetry breaking scale v via Eq. (10). A priori, this scale could be arbitrarily small, and thus for given g the Yukawa radius could be very large, eventually covering the whole galaxy and even the whole Universe. However, in the presence of n–$\tilde{n}$ mixing, the scale $v > v_\chi$ is limited by the nuclear stability bound (7). In fact, Eq. (10) can be conveniently rewritten as

[6] We neglect annual modulation of $V^{\rm gr}_S$ due to a small variation of the sun–Earth distance, as well as potentials induced by other planets and the Moon.

$$\left(\frac{\alpha_{B-L}}{10^{-49}}\right)^{1/2} \frac{\lambda}{1\ \text{AU}} = \frac{0.6\ \text{MeV}}{v} \leq \frac{6 \cdot 10^{-26}\ \text{eV}}{\varepsilon_{n\tilde{n}}}. \tag{15}$$

This dependence is shown on Fig. 1 for various values of v from 1 keV to 1 PeV. Thus, discovery of n–$\tilde{n}$ oscillation with free neutrons at the present experimental limit $\varepsilon_{n\tilde{n}} = 2.5 \times 10^{-24}$ eV would imply $v > 25$ MeV. (Let us recall that this limit is set by nuclear stability, while the direct limit on the n–$\tilde{n}$ oscillation time [12] is about three times weaker, $\varepsilon_{n\tilde{n}} < 7.8 \times 10^{-24}$ eV.) For $\alpha_{B-L} = 10^{-49}$ this implies $\lambda < 0.024$ AU, so that the sun's contribution (14) is exponentially suppressed and becomes negligible with respect to the Earth-induced potential (13). Even for $\varepsilon_{n\tilde{n}} > 10^{-25}$ eV, within the reach of the ESS experiments, we have $v > 1$ MeV and the sun-induced $B-L$ potential is still less than the Earth's contribution (13), so that $V_n = V_{nE} + V_{nS}$ in total cannot be higher than 10^{-11} eV or so.

In a hypothetical experiment with very cold neutrons vertically falling in 1 km length mine with magnetic field reduced below 1 nT one could approach the sensitivity down to $\varepsilon_{n\tilde{n}} \sim 10^{-26}$ eV, but better limits can hardly be achieved even in the far future. Then nuclear stability limits would allow for $v \sim 100$ keV or so, in which case the sun-induced potential (14) could reach its maximal value and hence the total potential $V_n = V_{nE} + V_{nS}$ could amount to up to 2.5×10^{-11} eV.

In any case, provided that n–$\tilde{n}$ mixing exists at the experimentally testable level $\varepsilon_{n\tilde{n}} > 10^{-26}$ eV or so, the limit (15) excludes the Yukawa range λ to be much larger than the size of the solar system. By this reasoning, the contributions of the galaxy and more distant cosmological objects become insignificant. Due to the limited interaction length λ also the contribution of relic neutrinos would be irrelevant even if the latter could store a large leptonic asymmetry. By the same reasoning, the contribution of dark matter is irrelevant even though some dark matter candidates, as e.g. mirror matter discussed in Sect. 4, could have common $B-L$ interactions with ordinary matter.

3 $B-L$ fifth force and n–$\tilde{n}$ oscillation

CPT invariance implies that the neutron and antineutron should have exactly equal masses and thus also their gravitational potentials should be equal, while their magnetic moments should have the opposite signs. Therefore, omitting universal contributions, the Hamiltonian that describes n–$\tilde{n}$ oscillation in a magnetic field B in the presence of $B-L$ potentials $V_n = -V_{\tilde{n}}$ can be presented as a 4×4 matrix acting on the state vector $(n_+, n_-, \tilde{n}_+, \tilde{n}_-)$ containing the neutron and antineutron states of two spin polarizations:

$$H = \begin{pmatrix} V_n + \mu_n B\sigma_3 & \varepsilon_{n\tilde{n}} \\ \varepsilon_{n\tilde{n}} & V_{\tilde{n}} - \mu_n B\sigma_3 \end{pmatrix}, \tag{16}$$

where $\mu_n = -6 \times 10^{-17}$ eV/nT is the magnetic moment of the neutron and σ_3 is the third Pauli matrix as far as the spin quantization axis can be chosen as the direction of the magnetic field. In this basis one has no spin precession and the Hamiltonian (16) is diagonal:

$$H = \begin{pmatrix} V_n - \Omega_B & 0 & \varepsilon_{n\tilde{n}} & 0 \\ 0 & V_n + \Omega_B & 0 & \varepsilon_{n\tilde{n}} \\ \varepsilon_{n\tilde{n}} & 0 & -V_n + \Omega_B & 0 \\ 0 & \varepsilon_{n\tilde{n}} & 0 & -V_n - \Omega_B \end{pmatrix} \quad (17)$$

where $\Omega_B = |\mu_n B| = (B/\mathrm{nT}) \times 6 \cdot 10^{-17}$ eV is the Zeeman energy shift induced by the magnetic field, and $V_n = -V_{\tilde{n}} > 0$ is the $B - L$ (repulsive) potential for the neutron.

Let us consider first the case when the fifth force is absent, $V_n = 0$. In this case the probability of n–$\tilde{n}$ oscillation for a free-flight time t reads

$$P_{n\tilde{n}}(t) = \frac{{\varepsilon_{n\tilde{n}}}^2}{\varepsilon_{n\tilde{n}}^2 + \Omega_B^2} \sin^2\left(t\sqrt{\varepsilon_{n\tilde{n}}^2 + \Omega_B^2}\right) \quad (18)$$

This probability depends on the value $\Omega_B t = 0.1(B/\mathrm{nT})(t/\mathrm{s})$. If $\Omega_B t < 1$, then the argument of sine wave is small and probability (18) becomes independent on Ω_B:

$$P_{n\tilde{n}}(t) \approx (\varepsilon_{n\tilde{n}} t)^2 = (t/\tau_{n\tilde{n}})^2 \quad (19)$$

The latter condition, $\Omega_B < 1/t$, is known as the "quasi-free" condition since the oscillation probability becomes in practice the same as for vanishing magnetic field. On the other hand, when $\Omega_B t \gg 1$, the probability (18) should be averaged in time, and so we get $P_{n\tilde{n}} = \varepsilon_{n\tilde{n}}^2/2\Omega_B^2$, by a factor $2(\Omega_B t)^2$ less with respect to Eq. (19).

Hence, for rendering n–$\tilde{n}$ oscillation effective during a flight time t, the magnetic field should be reduced to a certain necessary degree to achieve the quasi-free limit $\Omega_B \leq 1/t$, and its further suppression makes no sense for improving experimental sensitivity. The level of the magnetic field needed depends on the neutron free flight time under experimental conditions. In realistic experiments t cannot be very large, e.g. it was ~0.1 s in the experiment [12], it can be up to ~1 s in the experimental setup for cold neutrons at the ESS [8], and in principle it could reach ~10 s in the far future experiments with the neutrons vertically falling down in a deep mine. For $t = 1$ s, the condition $\Omega_B t < 1$ implies $\Omega_B < 10^{-16}$ eV, and thus one needs $B < 10$ nT. Thus, reduction of magnetic background to the level of a few nT would be sufficient for realistic experimental times order 1 s for discovering the n–$\tilde{n}$ oscillation at the level $\varepsilon_{n\tilde{n}} \sim 10^{-25}$ eV.

Let us consider now the case with non-zero potential energy V_n induced by the Earth and perhaps also the sun for a neutron experiment at the Earth surface. The red curve in Fig. 2 shows the upper limit on V_n as a function of the Yukawa radius λ, which is calculated via integration of Eq. (12) over the Earth density profile for a given value of λ and by taking the experimental upper limit on $\tilde{\alpha}(\lambda)$ as a function of λ from Fig. 6 of Ref. [34] (shown also directly for α_{B-L} on our Fig. 1.) For λ larger than the Earth diameter the limit on α_{B-L} becomes independent of λ (see Fig. 1) and it corresponds to $\alpha_{B-L} \leq 10^{-49}$, or $\tilde{\alpha} \leq 1.7 \times 10^{-11}$, in which case the Earth-induced potential reaches the maximal value (13), $V_{nE} = 5.6 \times 10^{-12}$ eV shown by the red curve in Fig. 2 as a plateau starting at $\lambda \geq 10^9$ cm. At the smaller values of λ the limit shown in the red curve is a result of the interplay of weaker limits $\tilde{\alpha}(\lambda)$ provided by Ref. [34] and fast reduction of the available mass due to the Yukawa term in the integration of the Earth potential for λ smaller than the Earth diameter. For larger values of λ, approaching the sun–Earth distance, the sun-induced potential (14) shown by dashed line starts to contribute building up the total potential up to 2.5×10^{-11} eV corresponding to a plateau at $\lambda > 10^{13}$ cm or so. The parameter areas of Fig. 1 excluded by condition (15) for $\varepsilon_{n\tilde{n}} = 2.5 \times 10^{-24}$ eV and $\varepsilon_{n\tilde{n}} = 10^{-26}$ eV, respectively, are shown by shaded regions, while the limit for $\varepsilon_{n\tilde{n}} = 10^{-25}$ eV corresponding to the sensitivity of the ESS long-baseline experiment passes in the middle.

The presence of the $B - L$ potential $V_n \sim 10^{-11}$ eV is equivalent to a Zeeman energy Ω_B for the magnetic field $B \sim 10^5$ nT. Thus, as far as $B - L$ forces cannot be screened, n–$\tilde{n}$ oscillation can remain strongly suppressed even if the magnetic field is reduced to the level $\Omega_B t < 1$. In this case the oscillations still should be averaged and one would have $P_{n\tilde{n}} = \varepsilon_{n\tilde{n}}^2/2V_n^2$, by a factor of $2(V_n t)^2$ less with respect to the value (19) expected in the absence of the $B - L$ potential, if $V_n t = 1.5 \times (V_n/10^{-15}\,\mathrm{eV})(t/\mathrm{s}) \gg 1$.

Therefore, if n–$\tilde{n}$ oscillation will be discovered in free neutron experiments with the values of the magnetic field satisfying the quasi-free condition $\Omega_B t < 1$, this would mean that the quasi-free condition is satisfied also by $B - L$ potential, i.e. $V_n t < 1$. Levels of potential energy V_n corresponding to quasi-free conditions for $n \to \tilde{n}$ observation time $t = 0.1$ s and $t = 1$ s are shown in Fig. 2 in comparison with Ω_B corresponding to the magnetic field 1 nT. We see that V_n can exceed the limit of the quasi-free condition, in which case $n \to \tilde{n}$ oscillation will be suppressed, in the range of λ between $10^6 - 10^{15}$ cm. Thus, the discovery of n–$\tilde{n}$ oscillation at the experiment with $t = 1$ s would exclude the $B - L$ potentials $V_n > 10^{-15}$ eV or so. This, in turn, would lead to the upper limit on the $U(1)_{B-L}$ gauge coupling constant being much stronger than the limit of Ref. [34] which is shown by a blue curve on Fig. 1. As we see, for the range of λ between 10^6–10^{15} cm these limits are stronger by up to four orders of magnitude than the limits of the torsion-balance experiments [34].[7]

On the other hand, if a considerably large $B-L$ potential is indeed present, with $V_n > 10^{-15}$ eV or so, it would lead to a ,strong suppression of n–$\tilde{n}$ oscillation even if the magnetic field value meets the quasi-free condition while n–$\tilde{n}$ mixing indeed exists at the level $\varepsilon_{n\tilde{n}} \sim 10^{-24}$ eV. One can envisage e.g. a situation where $n \to \tilde{n}$ will be discovered in large underground experiments via intranuclear transformations $n \to \tilde{n}$ for which tiny fifth forces have no effect, although it will not be observed in experiments with free neutrons at the corresponding level in magnetic quasi-free conditions, e.g. at the ESS. This can be an indication that some extra potential different between the neutron and antineutron is in play, which could be the one mediated by $B-L$ photons.

Nevertheless, even in this case one can continue experiments with free neutrons and discover free n–$\tilde{n}$ oscillation if, instead of suppressing the magnetic field, one would vary it in the range up to 1 G or so, to adjust to the value of Zeeman energy Ω_B to the value of the fifth potential V_n. In fact, the Hamiltonian (17) with V_n and Ω_B both non-zero leads to n–$\tilde{n}$ oscillation probabilities different between the + and − polarization states:

$$P^{\pm}_{n\tilde{n}}(t) = \frac{\varepsilon^2_{n\tilde{n}}}{\varepsilon^2_{n\tilde{n}} + \Delta^2_{\pm}} \sin^2\left(t\sqrt{\varepsilon^2_{n\tilde{n}} + \Delta^2_{\pm}}\right), \qquad (20)$$

where $\Delta_{\pm} = |V_n \mp \Omega_B|$. Therefore, for achieving the quasi-free condition for n–$\tilde{n}$ oscillation, $\Delta_{\pm} t < 1$, which is different for the + and − polarization states, the value of the magnetic field should be tuned with a precision of 10–100 nT to a resonance value, for achieving $\Omega_B = V_n$ with the precision of 10^{-15} eV or so. Let us notice that, since the oscillation probabilities of + and − polarization states are different, see Eq. (20), resonance can occur only for one polarization, namely for the + polarization since $V_n > 0$. This situation can be checked by applying the magnetic field with the programmed magnitude and direction in the whole neutron flight path and by varying it to find the resonance value for which Ω_B would compensate the effect of $B-L$ potential V_n. An example of such a variation of the magnetic field is shown in Fig. 3.

The following remark is in order. If n–$\tilde{n}$ oscillation will be discovered under these conditions, this could be interpreted as an indication for the CPT violation due to different masses of the neutron and antineutron, $m_n - m_{\tilde{n}} \neq 0$ [39]. In this case, however, one has to renounce at least one of the basic principles of quantum field theories: locality, unitarity and Lorentz invariance. In our interpretation, CPT violation is simulated by environmental effects without renouncing consistency of the field-theoretical description, as a fifth force mediated by $B-L$ gauge field: in fact, a non-zero difference of the $B-L$ potentials for the neutron and antineutron, $V_n - V_{\tilde{n}} = 2V_n$, imitates their mass difference $m_n - m_{\tilde{n}} \neq 0$. However, the fundamental CPT violation would imply $m_n - m_{\tilde{n}}$ to be constant everywhere whereas the value of V_n is environment dependent and it must vary in space, depending on the distance to the gravitating bodies and their masses. Hence, for compensating the effect of V_n in the experiments proposed above to achieve the resonance conditions $\Omega_B = V_n$, one has to adjust the magnetic field B at different values depending on whether experiments are done at the surface of the Earth, inside the Earth or on the Moon.

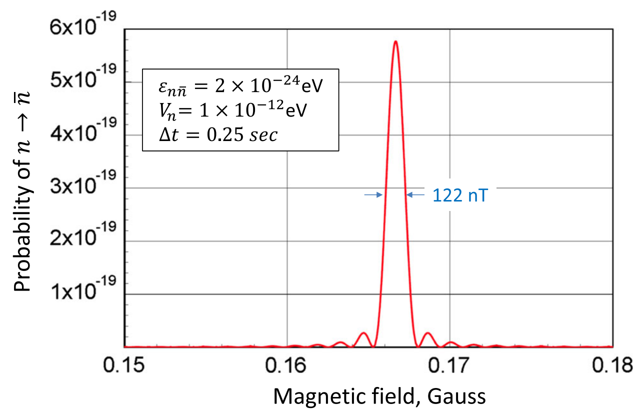

Fig. 3 Variation of the external magnetic field that would compensate for the effect of the $B-L$ potential $V_n = 10^{-12}$ eV for flight time 0.25 s

4 $B-L$ forces and neutron–mirror neutron oscillation

For a demonstration, let us discuss the example of a consistent renormalizable model for n–$\tilde{n}$ mixing via the operators (3) in which $B-L$ is spontaneously broken at rather low scales, so that the baryophoton couplings to the neutron can have an effect on the laboratory search of the n–$\tilde{n}$ oscillation.

One can consider a simple seesaw-like scenario along the lines suggested in Ref. [4,17]. We introduce a couple of gauge singlet right-handed Weyl fermions, N with $Q=-1$ and N' with $Q=1$. These two together form a heavy Dirac particle with a large mass M_D, while both of N and N' can be coupled to the scalar χ ($Q=2$) and get the Majorana mass terms from its VEV $\sim \langle\chi\rangle = v_\chi$. Let us introduce also a color-triplet scalar S with $Q=-2/3$, having the mass M_S and precisely the same gauge quantum numbers as the right down-quark $d_{(R)}$. Consider now the Lagrangian terms

[7] A potential energy $V_n - V_{\tilde{n}}$ smaller than 10^{-16} eV, i.e. below the quasi-free condition limit, practically will not be sensed by $n \to \tilde{n}$ oscillation and therefore cannot be excluded in this way. The STEP experiment for satellite proposed some years ago claimed a sensitivity of testing the equivalence principle to the level 10^{-18} [38]. The STEP mission was not pursued. The corresponding levels of magnitude of the $B-L$ potential energy that could be excluded by the STEP test are also shown in Fig. 2.

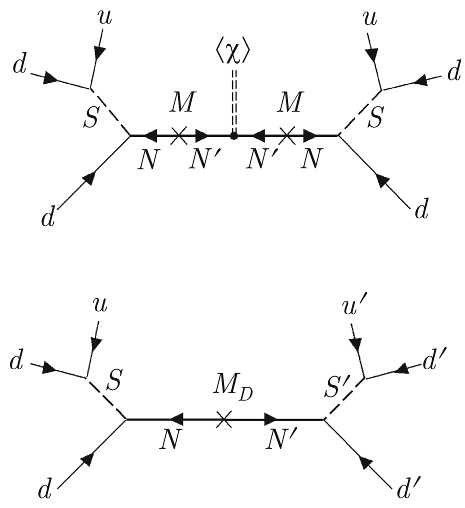

Fig. 4 The *upper* diagram with M_D being the Dirac mass of heavy fermions generates n–$\tilde{n}$ mixing after inserting the $B-L$ symmetry breaking VEV $\langle\chi\rangle$. The *lower* diagram with twin mirror quarks u', d', connected to N', induces (without insertion of χ field) n–$\tilde{n}'$ mixing which conserves the combination of charges $\bar{Q}=(B-L)-(B'-L')$

$$Sud + Sqq + S^{\dagger}d\,N + M_D NN' + \chi^{\dagger}N^2 + \chi N'^2 + \text{h.c.} \tag{21}$$

(Yukawa constants are omitted). After integrating out the heavy fermions N, N', this Lagrangian induces the $D=10$ operators (3) with $M^6 \sim M_D^2 M_S^4$, as shown in the upper diagram of Fig. 4.

Such a scenario of low-scale baryon violation was suggested in Ref. [4], in a model which was mainly designed for inducing neutron–mirror neutron oscillation. This model treats N and N' states symmetrically: they get equal Majorana masses from Yukawa couplings with χ in (21), while in addition to the couplings (21), there are analogous terms that couple N' to q', u', d' and S' states, partners of u, d and S from the hidden mirror sector which has a particle content exactly identical to that of ordinary one (for a review, see, e.g. [40–43]). One can prescribe to mirror particles the respective quantum numbers $Q'=B'-L'$. The mass term $M_D NN'$ violates $Q=B-L$ and $Q'=B'-L'$ separately but conserves a combined quantum number $\bar{Q}=Q-Q'$. Hence, in this context one can promote $U(1)_{\bar{Q}}$ as a local gauge symmetry rather than $U(1)_{B-L}$ and $U(1)_{B'-L'}$ separately, so that both ordinary and mirror particles couple to the gauge field b_μ of $U(1)_{\bar{Q}}$ proportionally to their quantum number $\bar{Q}=Q-Q'$ (mirror particles have $\bar{Q}$ of the opposite sign with respect to their ordinary partners). Then the lower diagram of Fig. 4 induces the $D=9$ operators

$$\frac{1}{\mathscr{M}^5}\left(uddu'd'd' + qqdq'q'd' + uddq'q'd' + qqdu'd'd'\right) + \text{h.c.} \tag{22}$$

with $\mathscr{M}^5 = M_D M_S^4$, which gives rise to n–$\tilde{n}'$ mixing:

$$\varepsilon_{n\tilde{n}'} \sim \frac{\Lambda_{\text{QCD}}^6}{\mathscr{M}^5} \sim \left(\frac{10\ \text{TeV}}{\mathscr{M}}\right)^5 \times 10^{-15}\ \text{eV} \tag{23}$$

For definiteness, one can call the mirror neutron a state n' consisting of three valence mirror quarks $u'd'd'$, so that the ordinary neutron n should oscillate into a mirror antineutron state $\tilde{n}'$ [44]. As far as n–$\tilde{n}'$ mixing conserves $\bar{Q}=(B-L)-(B'-L')$, the baryophoton associated with $U(1)_{\bar{Q}}$ should interact symmetrically with the ordinary neutron n and mirror antineutron $\tilde{n}'$, and thus the corresponding environmental potential should have no effect on n–$\tilde{n}'$ (or $\tilde{n}$–n') conversion.[8]

Hence, in this case n–$\tilde{n}'$ mixing can be a dominant effect since it can be induced without breaking of $U(1)_{\bar{Q}}$ symmetry, while n–$\tilde{n}$ (and n'–$\tilde{n}'$) mixings which break $\bar{Q}$ are suppressed by the small VEV v_χ. Assuming that both n–$\tilde{n}'$ and n–$\tilde{n}$ mixings are induced by the diagrams of Fig. 4 and setting, respectively, $\mathscr{M}^5 = M_D M_S^4$ in Eq. (23) and $M^6 \sim M_D^2 M_S^4$ in Eq. (6), we obtain

$$\frac{\varepsilon_{n\tilde{n}'}}{M_D} \sim \frac{\varepsilon_{n\tilde{n}}}{v_\chi}. \tag{24}$$

Therefore, assuming that $\varepsilon_{n\tilde{n}} > 10^{-25}$ eV and taking into account the nuclear stability limit $v_\chi > 1$ MeV (7), we find that e.g. for $M_D \sim 100$ TeV the n–n' mixing mass can be as large as $\varepsilon_{n\tilde{n}'} \sim 10^{-17}$ eV, corresponding to an n–$\tilde{n}'$ oscillation time $\tau_{n\tilde{n}'} = \varepsilon_{n\tilde{n}'}^{-1}$ of the order of 100 s.

As a matter of fact, n–$\tilde{n}'$ mixing can indeed be much larger than n–$\tilde{n}$. Existing experimental limits [47–52] allow the n–$\tilde{n}$ oscillation time to be less than the neutron lifetime, with interesting implications for astrophysics and particle phenomenology [4,44,53–56].

Concluding this section, let us remark about the possibility of the kinetic mixing of $B-L$ photons with the regular QED photons. Such mixing could make the equivalence principle tests potentially different for electrically neutral and charged objects, e.g. neutrons and also neutrinos having non-zero $B-L$ could acquire also tiny electric charges. As a matter of fact, the considered $B-L$ potentials can have no effect on the oscillations between the three neutrinos ν_e, ν_μ and ν_τ since their $B-L$ charges are equal, but they can be relevant for the active–sterile neutrino (e.g. mirror neutrino) oscillations and can suppress them in certain situations. Also, the $B-L$ charge of the Earth would create a $B-L$ magnetic field due to the rotation of the Earth, which could have observable effects.

[8] One has to note, however, that n–$\tilde{n}'$ oscillation can also be affected by environmental effects if ordinary and mirror worlds are associated with different gravities, in the context of bigravity models [45,46].

In the case that a mirror world exists, the photon associated with $U(1)_{\tilde{Q}}$ should have equal kinetic mixings with both ordinary and mirror photons, along with the direct kinetic mixing between the latter. The interplay of kinetic mixings between three vector fields have interesting cosmological implications e.g. for the direct search of dark matter [57,58] or for generating large-scale magnetic fields due to the relative flow of ordinary and mirror matter components in the early Universe [59].

5 Conclusions

$U(1)_{B-L}$ can be a local gauge symmetry, and the relative gauge boson – $B-L$ photon or baryophoton – can have a tiny gauge coupling constant and long-range interaction radius which are limited by experimental tests of the weak equivalence principle [34]. Discovery or the neutron–antineutron oscillation would mean that $B-L$ is spontaneously broken, and thus this baryophoton should be massive. Namely, for a $B-L$ symmetry breaking scale v varying between 100 keV and 1 TeV, the limits of Ref. [34] imply that the baryophoton mass, respectively, should be larger than 10^{-19}–10^{-11} eV. Nevertheless, these limits still allow for $B-L$ potentials induced by the Earth and sun to be considerably large, up to $\sim 10^{-11}$ eV, which can suppress $n \to \tilde{n}$ oscillation even if n–$\tilde{n}$ mixing exists at the level $\varepsilon_{n\tilde{n}} > 10^{-25}$ eV which is experimentally accessible in the future long-baseline experiments at the ESS.

If these experiments will discover n–$\tilde{n}$ oscillation in quasi-free conditions, with properly suppressed external magnetic fields, this would exclude the neutron $B-L$ potential at the Earth surface down to the values of about 10^{-16} eV, independently on the interaction radius λ of the baryophotons. This in turn will imply limits on the $B-L$ photon coupling constant considerably stronger than the present limits from the tests of the equivalence principle [34].

If instead n–$\tilde{n}$ conversion will be discovered via nuclear instability, but not in free neutron oscillations at the corresponding level, this would indicate that free n–$\tilde{n}$ oscillation is suppressed by such an environmental potential distinguishing between the neutron and antineutron states. Then this suppression in principle can be removed by the tuning of the external magnetic field in the experiment, in which case free n–$\tilde{n}$ oscillation can be discovered but for only one polarization state, whereas for the other polarization it will remain suppressed. Namely, a potential mediated by $B-L$ baryophotons would imply that n–$\tilde{n}$ oscillation will be observed only for neutrons polarized towards the magnetic field.

Concluding, if the neutron–antineutron oscillation will be discovered in free neutron oscillation experiments, this will imply much stronger limits on fifth forces mediated by the $B-L$ photon than the present limits from the tests of the equivalence principle. If instead n–$\tilde{n}$ oscillation will be discovered via nuclear instability but not in free neutron experiments at the corresponding level, this would indicate the presence of a fifth force mediated by such baryophotons; these could be revealed in experiments with controlled magnetic fields.

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