Neutron–Antineutron Oscillation as a Signal of CP Violation

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Abstract

Assuming the Lorentz and CPT invariances we show that neutron-antineutron oscillation implies breaking of CP along with baryon number violation – i.e. two of Sakharov conditions for baryogenesis. The oscillation is produced by the unique operator in the effective Hamiltonian. This operator mixing neutron and antineutron preserves charge conjugation C and breaks P and T. External magnetic field always leads to suppression of oscillations. Its presence does not lead to any new operator mixing neutron and antineutron.
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Assuming the Lorentz and CPT invariances we show that neutron-antineutron oscillation implies breaking of CP along with baryon number violation – i.e. two of Sakharov conditions for baryogenesis. The oscillation is produced by the unique operator in the effective Hamiltonian. This operator mixing neutron and antineutron preserves charge conjugation \( C \) and breaks \( P \) and \( T \). External magnetic field always leads to suppression of oscillations. Its presence does not lead to any new operator mixing neutron and antineutron.

1. Experimental search for neutron-antineutron oscillation \([1]\) is under active discussion nowadays (see the recent review \([2]\)). Its discovery would be a clear evidence of baryon charge nonconservation, \(|\Delta B| = 2\). In this note we would like to emphasize that neutron-antineutron oscillation also breaks CP invariance. This conclusion is based on the Lorentz invariance and CPT.

To demonstrate our assertion let us start with the Dirac Lagrangian

\[
\mathcal{L} = i\bar{n}\gamma^\mu \partial_\mu n - m \bar{n}n
\]

(1)

with four-component spinor \( n \) and the mass parameter \( m \) which is real and positive. The Lagrangian gives the Lorentz-invariant description of free neutron and antineutron states and preserves the baryon charge, \( B = 1 \) for \( n \) and \( B = -1 \) for \( \bar{n} \). This charge corresponds to the continuous symmetry

\[
n \to e^{i\alpha}n, \quad \bar{n} \to e^{-i\alpha}\bar{n}
\]

(2)

of Lagrangian \([1]\). At each spatial momentum there are four degenerate states, two spin doublets which differ by the baryon charge \( B \).

Another bilinear mass term,

\[
\Delta \mathcal{L}_m = -im\bar{n}\gamma_5 n,
\]

(3)

consistent with the baryon charge conservation, can be rotated away by chiral transformation \( n \to e^{i\beta\gamma_5}n \) if there is no terms breaking the baryon charge. As we will see it is not the case when the baryon charge is broken.

How the baryon number non-conservation shows up at the level of free one-particle states? In Lagrangian description it could be only modification of the bilinear mass term. We show below that the most generic Lorentz invariant modification of Eq. \([1]\) reduces to one possibility for the baryon charge breaking by two units,

\[
\Delta \mathcal{L}_B = -\frac{1}{2}\epsilon \left[ n^T C n + \bar{n}C\bar{n}^T \right].
\]

(4)

Here \( C = i\gamma^2\gamma^0 \) is the charge conjugation matrix in the standard representation of gamma matrices, and \( \epsilon \) is a real positive parameter. The reality of \( \epsilon \) as a coefficient for \( n^T C n \) can be always achieved by the phase rotation \([2]\) of \( n \) field.

One could add also \(|\Delta B| = 2\) term of the form \( n^T C\gamma_5 n \). However, it can be rotated away by the chiral rotation \( n \to e^{i\beta\gamma_5}n \). The price for this is, as we mentioned above, an appearance of the \( \gamma_5 \) mass term \([3]\). Also mixed kinetic terms \( \propto i\bar{n}\gamma^\mu C \partial_\mu \bar{n} + h.c. \) can be turned away with redefinition of the fermion field.

Hence, a generic Lagrangian containing the fermion bilinears can always be brought to a form containing only the terms \([1], [3] \) and \([4]\).

What is the status of discrete \( C, P \) and \( T \) symmetries in this situation? It is simple to verify that the Lagrangian terms \([1], [3] \) and \([4]\) are all invariant under the charge conjugation \( C \),

\[
n \to n^c = Cn^T.
\]

(5)

In fact, the expression \([4]\) can be rewritten in the form \( -\frac{1}{2}(1/2)\epsilon \left[ n^T n - n^c n^c \right] \), which makes its \( C \) invariance explicit.

The parity transformation \( P \) involves (besides reflection of the space coordinates) the substitution

\[
n \to \gamma^0 n.
\]

(6)

This substitution changes \( \Delta \mathcal{L}_B \) to \(-\Delta \mathcal{L}_B \) because \( \gamma^0 C\gamma^0 = -C \). The breaking of parity in neutron-antineutron transition reflects the well-known feature of the opposite parity of fermion and antifermion. The term \( \Delta \mathcal{L}_m \) also breaks \( P \) parity, it is evidently pseudoscalar. Clearly, the parity violation comes together with breaking of \( T \) invariance since CPT invariance is guaranteed by a local, Lorentz invariant form of the Lagrangian.

Thus, we demonstrated that observation of neutron-antineutron oscillation signals breaking of CP invariance together with breaking of baryon charge.

2. To show that the above consideration covers indeed a generic case it is convenient to introduce two left-handed Weyl spinors \([3]\), forming a flavor doublet

\[
\psi^i\alpha, \quad i = 1, 2, \quad \alpha = 1, 2,
\]

(7)
with their complex conjugates, representing the right-handed spinors,
\[ \overline{\psi}_{i}^{\dot{\alpha}} = (\psi_{i}^{\dot{\alpha}})^{\ast}, \quad i = 1, 2, \quad \dot{\alpha} = 1, 2. \] (8)

One can raise and lower space \( \alpha, \dot{\alpha} \) and flavor \( i \) indices using \( \epsilon_{\alpha \dot{\beta}}, \epsilon_{\dot{\alpha} \beta} \) and \( \epsilon_{i k} \). In terms of Dirac spinor \( n \) these two left-handed Weyl spinors are \( n_{L} \) and \( (n_{R})^{*} \). The most generic Lagrangian is
\[ \mathcal{L} = \psi_{i}^{\dot{\alpha}} i \partial_{\mu} \overline{\psi}_{i}^{\dot{\alpha}} + \frac{1}{2} \left[ m_{i k} \psi_{i}^{\dot{\alpha}} \psi_{k}^{\dot{\alpha}} + \overline{m}_{i k} \overline{\psi}_{i}^{\dot{\alpha}} \overline{\psi}_{k}^{\dot{\alpha}} \right], \] (9)

where \( \partial_{\mu} = \sigma_{\alpha \dot{\beta}} \partial_{\mu}, \sigma^{\alpha} = \{ 1, \sigma \} \), \( m_{i k} \) is the symmetric mass matrix, \( m_{i k} = m_{k i} \) and \( \overline{m}_{k i} = (m_{i k})^{*} \) is its conjugate.

The kinetic term in (11) is U(2) symmetric: besides SU(2) rotations of the doublets, there is U(1) associated with the overall phase rotation of the doublet (17). The mass terms break both, U(1) and SU(2) flavor symmetries, so, generically, no continuous symmetry remains.

To see how the symmetry (2) associated with the baryon charge could arise note that one can interpret U(2) transformations as acting on the external mass matrix \( m_{i k} \). This matrix is charged under U(1), the overall phase rotation, so this U(1) symmetry is always broken. In respect to SU(2) transformations the symmetric tensor \( m_{i k} \) is the adjoint representation, i.e., can be viewed as an isovector \( \mu^{a} \), \( a = 1, 2, 3 \),
\[ m_{i k}^{a} = \epsilon^{i j k} m_{j k} = \mu^{a}(\tau^{a})_{i k}, \quad a = 1, 2, 3. \] (10)

Because \( \mu^{a} \) is complex, we are actually dealing with two real isovectors, \( \text{Re} \mu^{a} \) and \( \text{Im} \mu^{a} \). The SU(2) transformations are equivalent to simultaneous rotation of real and imaginary vectors, while U(1) changes phases of all \( \mu^{a} \) simultaneously, which is equivalent to SO(2) rotation inside each couple \( \text{Re} \mu^{a}, \text{Im} \mu^{a} \). Only in case when these vectors are parallel we have an invariance of the mass matrix which is just a rotation around this common direction. (In this case, all \( \text{Im} \mu^{a} \) can be absorbed in \( \text{Re} \mu^{a} \) by U(1) transformation.) This symmetry is the one identified with the baryonic U(1) in Eq. (4).

Let us show now that in the absence of the common direction we get two spin 1/2 Majorana fermions with different masses. From equations of motion
\[ i \partial_{\alpha} \psi^{i \alpha} + \overline{m}_{i k} \psi_{k}^{i \alpha} = 0, \]
\[ i \partial_{\alpha} \psi_{i}^{\dot{\alpha}} + m_{i k} \psi_{k}^{\dot{\alpha}} = 0 \] (11)

to exclude \( \overline{\psi}_{i}^{\dot{\alpha}} \) we come to the eigenvalue problem for \( M^{2} = p \mu p^{\dagger} \),
\[ M^{2} \psi^{i \alpha} - \overline{m}_{i k} m_{k \alpha} \psi^{i \alpha} = 0. \] (12)

Using definition (10) of \( \mu^{a} \) the squared mass matrix can be presented as a combination of isoscalar and isovector pieces:
\[ \overline{m} \cdot m = \mu^{a} \overline{\mu}^{a} \delta_{n}^{k} + i \epsilon^{a b c} \mu^{a} \overline{\mu}^{b} (\tau^{c})_{n}^{k}. \] (13)

Correspondingly, there are two invariants defining \( M^{2} \). The isoscalar part gives the sum of eigenvalues,
\[ \frac{M^{2} + M_{3}}{2} = \mu^{a} \overline{\mu}^{a} = (\text{Re} \mu^{a})^{2} + (\text{Im} \mu^{a})^{2} \] (14)

while the length of the isovector part defines the splitting of the eigenvalues,
\[ \frac{M^{2} - M_{3}}{2} = 2 \sqrt{\epsilon^{a b c} \text{Re} \mu^{a} \text{Im} \mu^{b}}. \] (15)

Thus, we see the splitting associated with the breaking of the baryon charge.

To follow the discrete symmetries we can orient the mass matrix \( m_{i k} \) in a convenient way. In terms of \( \mu^{a} \) the matrix has the form
\[ m_{i k} = \begin{pmatrix} -\mu^{1} & -i \mu^{2} & \mu^{3} \\ -i \mu^{2} & \mu^{3} & \mu^{1} \\ \mu^{3} & \mu^{1} & \mu^{2} \end{pmatrix} \] (16)

Without lost of generality we can put both, \( \text{Re} \mu^{a} \) and \( \text{Im} \mu^{a} \), onto the 23 plane, i.e., put \( \mu^{1} = 0 \). Moreover, we can orient \( \text{Re} \mu^{a} \) along the the third axis, i.e., put \( \mu^{2} = 0 \). Then, only 3 nonvanishing parameters, \( \text{Re} \mu^{3}, \text{Im} \mu^{3}, \mu^{2} \), remain and the mass matrix takes the form,
\[ m_{i k} = \begin{pmatrix} \text{Re} \mu^{2} & \text{Im} \mu^{3} + i \text{Im} \mu^{3} \\ \text{Re} \mu^{3} & \text{Im} \mu^{2} \end{pmatrix}. \] (17)

In the Weyl description the charge conjugation \( C \) is just an interchange of \( \psi^{1 \alpha} \) and \( \psi^{2 \alpha} \), the symmetry which implies that \( m_{12} = m_{21} \) and \( m_{11} = m_{22} \). The matrix (17) clearly satisfies these conditions. The \( \mathbf{P} \) reflection involves the interchange \( \psi^{i \alpha} \leftrightarrow \overline{\psi}^{i \alpha} = \epsilon^{i k} \overline{\psi}^{k \alpha} \). This symmetry is broken by nonvanishing \( \text{Im} \mu^{a} \).

Now it is simple to establish a correspondence with parameters introduced earlier in four-component spinor notations. Namely,
\[ \text{Re} \mu^{3} = m, \quad \text{Im} \mu^{3} = m', \quad \text{Im} \mu^{2} = \epsilon. \] (18)

So while \( C \) parity is preserved, we have \( \mathbf{P} \) even, Eq. (11), and \( \mathbf{P} \) odd, Eqs. (3), (4), mass terms. Thus, we proved for generic case the association of baryon charge breaking with \( \mathbf{CP} \) violation.

Note that in terms of remaining 3 parameters the masses of \( C \) even and \( C \) odd Majorana fermions are
\[ M_{1}^{2} = (m + \epsilon)^{2} + (m')^{2}, \quad M_{2}^{2} = (m - \epsilon)^{2} + (m')^{2}, \] (19)

what different from standard expressions when \( m' \) is nonvanishing. In particular, it implies that the oscillation time \( \tau_{n \overline{n}} \) in free neutron transition probability,
\[ P_{n \overline{n}}(t) = \sin^{2}(t/\tau_{n \overline{n}}) \] instead of \( 1/\epsilon \).

The \( \mathbf{CP} \) odd nature of the operator (4) was noted recently in Ref. [3]. However, the authors of this paper discussed also the \( \mathbf{CP} \) even operator \( n^{T} \gamma_{5} C n \) which, as we showed, can be rotated away by field redefinition. These authors also analyzed modifications induced by
external magnetic field claiming an existence of a new $n - \bar{n}$ transition magnetic moment and also an existence of the usual suppression of $n - \bar{n}$ oscillation in presence of magnetic field. We will show below that both claims are invalid.

3. Our consideration above refers to the neutron-antineutron oscillation in vacuum. Now we show that even in the presence of magnetic field no new $|\Delta \mathcal{B}| = 2$ operator appears. Similar consideration was done in Ref. [3] in application to magnetic moment of neutrinos.

In the Weyl formalism the field strengths tensor $F_{\mu\nu}$ is substituted by the symmetric tensor $F_{\alpha\beta}$ and its complex conjugate $\bar{F}_{\alpha\beta}$. They correspond to $\tilde{E} \pm i\tilde{B}$ combinations of electric and magnetic fields. Then Lorentz invariance allows only two structures involving electromagnetic fields,

$$F_{\alpha\beta} \psi^i \gamma^\alpha \gamma^\beta \psi^j \epsilon_{ik} \quad \bar{F}_{\alpha\beta} \bar{\psi}^i \gamma^\alpha \gamma^\beta \bar{\psi}^j \epsilon^{ik}$$  \hspace{1cm} (20)

Antisymmetry in flavor indices implies that spinors with the opposite baryon charge enter. So both operators preserve the baryon charge, they describe interactions with the magnetic and electric dipole moments of the neutron.

The authors of [3] realize that the operator $n^T \sigma^{\mu\nu} C \bar{n} F_{\mu\nu}$ is vanishing due to Fermi statistics. They believe, however, that a composite nature of neutron changes the situation and a new type of magnetic moment in $\Delta \mathcal{B} = \pm 2$ transitions may present. In other words they think that the effective Lagrangian description is broken for composite particles.

To show that is not the case let us consider the process

$$n(p_1) + n(p_2) \rightarrow \gamma^*(k)$$  \hspace{1cm} (21)

in the crossing channel to $n - \bar{n}$ transition magnetic moment. The number of invariant amplitudes for the process (21) which is $1/2^+ + 1/2^+ \rightarrow 1^-$ transition is equal to one. Only orbital momentum $L = 1$ and total spin $S = 1$ in two neutron system are allowed by angular momentum conservation and Fermi statistics. The gauge-invariant form of the amplitude is

$$u^T(p_1) C \gamma^\mu \gamma^\nu u(p_2) k^\nu F_{\mu\nu}, \quad F_{\mu\nu} = k_\mu \epsilon_\nu - k_\nu \epsilon_\mu,$$  \hspace{1cm} (22)

where $u_{1,2}$ are Dirac spinors describing neutrons and $\epsilon_\mu$ refers to the gauge potential. In space representation we deal with $\partial^\mu F_{\mu\nu}$ which vanishes outside of the source of the electromagnetic field, and, in particular, for the distributed magnetic field. It proves that there is no place for magnetic moment of $n - \bar{n}$ transition, and effective Lagrangian description does work.

Even in the absence of new $n - \bar{n}$ magnetic moment the authors of [3] claim that suppression of $n - \bar{n}$ oscillations by external magnetic field can be overcome by applying the magnetic field transversal to quantization axis.

In their first example where the transversal field is time-independent (after switching) they obtained four different energy eigenvalues (Eq. (26) in [4]) which depend on direction of magnetic field. This clearly breaks rotational invariance. The source of this breaking is the wrong sign of the $\mathcal{H}_{34}$ and $\mathcal{H}_{43}$ in the Hamiltonian matrix $\mathcal{H}$ in Eq. (20) The existing sign implies that $\Delta \mathcal{B} = 2$ amplitude is of different sign for spins up and down. Changing sign of $\mathcal{H}_{34}$ and $\mathcal{H}_{43}$ restores rotational invariance. In their second example, where the transversal field is rotating, the result of [4] is also incorrect – after a change of variables indicated in [4] the consideration is similar to the first example with time-independent field.

As a consequence the magnetic field suppression does present indeed, and the suggestion in [3] that $n - \bar{n}$ oscillations can be measured without minimizing magnetic field does not work.

4. Our use of the effective Lagrangian for the proof means that the Lorentz invariance and CPT are crucial inputs. Once constraints of Lorentz invariance are lifted new $|\Delta \mathcal{B}| = 2$ operators could show up.

Such operators were analyzed in Ref. [2] for putting limits on the Lorentz invariance breaking. In particular, the authors suggested the operator $n^T C \gamma^\mu \gamma^\nu n$ as an example which involves spin flip and, correspondingly, less dependent on magnetic field surrounding.

Note, however, that besides breaking of Lorentz invariance this operator breaks also 3d rotational invariance, i.e., isotropy of space. Such anisotropy could be studied by measuring spin effects in neutron-antineutron transitions.

5. In the Standard Model (SM) conservations of baryon $B$ and lepton $L$ numbers are related to accidental global symmetries of the SM Lagrangian. (Nonperturbative breaking of $B$ and $L$, preserving $B - L$, is extremely small.) The violation of $B$ by two units can be originated only from new physics beyond SM which could induce the effective six-quark operators

$$\mathcal{O} = \frac{1}{M^5} u d u d u$$  \hspace{1cm} (23)

involving $u$ and $d$ quarks of different families in different color and Lorentz invariant combinations (all possible

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1. Let us also remark that $n - \bar{n}$ transition with a virtual photon connected to the proton, as well as $n n \rightarrow \gamma \gamma$ annihilation, would destabilise the nuclei even in the absence of $n - \bar{n}$ mass mixing.

2. The situation is different if one considers oscillation $n - n'$ where $n'$ is a mirror neutron, twin of the neutron from hidden mirror sector [4]. In this case, operators $\bar{m} \sigma^{\mu\nu} F_{\mu\nu}$ and/or $n^T \sigma^{\mu\nu} n' F_{\mu\nu}$ are allowed. Hence, $n - n'$ and/or $n - n'$ transition probabilities may not depend on the value of magnetic field provided that it is large enough, with possible implications for the experimental search of neutron–mirror neutron oscillations.
convolutions of spinor indices are omitted). The smallness of baryon violation is related to the large mass scale \(M\) related to new physics.

In fact, the \(\mathcal{B}\) breaking mass term (4) emerges by taking matrix element between \(n\) and \(\bar{n}\) states of the operator structures (23), see diagram in Fig. 1:

\[
\begin{align*}
-\frac{1}{2} \epsilon \langle \bar{n} | n^T C u | n \rangle &= \langle \bar{n} | O | n \rangle.
\end{align*}
\] (24)

It gives an estimate of order \(\Lambda_{QCD}^6/M^5\) for the parameter \(\epsilon\) which describes the oscillation time.

Our consideration shows that only operators which are \(C\) even and \(P\) odd contribute to the above matrix element (up to small corrections due to electroweak interactions where the discrete symmetries are broken). In general, operators coming from physics beyond SM do not respect any of discrete symmetries \(C, P\) and \(\mathcal{C}\). If, however, a new physics model produces \(\mathcal{B}\) violating operators which do not satisfy the selection rules of \(n - \bar{n}\) transition, their effect will show up in instability of nuclei but not in free neutron-antineutron oscillations. Indeed, such operators would induce processes of annihilation of two nucleons like \(N + N \rightarrow \pi + \pi\) inside nucleus, as shown on Fig. 2.

The operators of the type of (23) involving strange quark, \(udsuds\), could induce \(\Lambda - \bar{\Lambda}\) mixing. However, such operators would also lead to nuclear instability via nucleon annihilation into kaons \(N + N \rightarrow K + K\), see the diagram in Fig. 2 where in upper lines \(d\) quark is substituted by \(s\) quark (and \(\pi^+\) by \(K^+\)). In fact, nuclear instability bounds on \(\Lambda - \bar{\Lambda}\) mixing are only mildly, within an order of magnitude, weaker than with respect to \(n - \bar{n}\) mixing which makes hopeless the possibility to detect \(\Lambda - \bar{\Lambda}\) oscillation in the hyperon beam. (Instead, it can be of interest to search for the nuclear decays into kaons in the large volume detectors.) The nuclear instability limits on \(\Lambda - \bar{\Lambda}\) mixing are about 15 orders of magnitude stronger than the sensitivity \(\delta\Lambda\Lambda \sim 10^{-6}\) eV which can be achieved in the laboratory conditions \(\mathcal{B}\). The nuclear stability limits make hopeless also the laboratory search of \(b\bar{u}\)-like baryon oscillation due to operator \(wsb\) suggested in Ref. 9.

6. The construction we used for neutron-antineutron transition could be applied to mixing of massive neutrinos. As an example, let us take the system of left-handed \(\nu_e\) and \(\nu_\mu\) and their conjugated partners, right-handed \(\bar{\nu}_e\) and \(\bar{\nu}_\mu\). One can ascribe them a flavor charge \(\mathcal{F} = \mathcal{L}_e - \mathcal{L}_\mu\) (analog of \(\mathcal{B}\)), to be (+1) for \(\nu_e\) and (-1) for \(\nu_\mu\). Then, \(\mathcal{C}\) conjugation is interchange of \(\nu_e\) and \(\nu_\mu\). Again, \(\mathcal{F}\) breaking mass term would be \(\mathcal{C}\) even and \(\mathcal{P}\) odd.

A similar scenario can be played in case of Dirac massive neutrino.

7. In summary, we show that the Lorentz and \(\mathcal{CPT}\) invariance lead to the unique \(\text{[\Delta]\mathcal{B}} = 2\) operator in the neutron-antineutron mixing. This operator is \(\mathcal{CPT}\) odd. Switching on external magnetic field influences the level splitting which suppresses \(n - \bar{n}\) oscillations but does not add any new \(\text{[\Delta]\mathcal{B}} = 2\) operator in contradistinction with recent claims in literature.

Interesting to note that observation of neutron-antineutron transition would show that two of three Sakharov conditions for baryogenesis are satisfied, violations of \(\mathcal{B} - \mathcal{L}\) and \(\mathcal{CP}\). However, it would be honest to say that primordial baryogenesis in the Early Universe should be related to underlying physics that induces operators (23) rather than to neutron-antineutron oscillation phenomenon itself. On the other hand, for new physics involving contact operators (23) (or heavy particles mediating these operators) the third, out-of-equilibrium condition is also automatically satisfied when the universe temperature drops below the relevant mass scales. Thus, discovery of neutron-antineutron oscillation would make it manifest that these operators contain \(\mathcal{CP}\) violating terms which could be at the origin of the baryon asymmetry of the Universe.

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