$n - \bar{n}$ OSCILLATIONS - THEORY AND EXPERIMENT

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The violation of baryon (and total lepton) number is a natural prediction of grand unified theories (GUT's), since in such theories quarks can undergo transitions into (anti)leptons or (anti)-quarks, mediated by gauge or Higgs bosons. Of course, such a violation has always been a phenomenological possibility, but with the development of GUT's, its status has advanced to that of a prediction. The manifestations of B and L violation include proton and bound neutron decay, $n - \bar{n}$ oscillations,[1-6] and Majorana neutrino masses. In this report we shall consider the second of these effects.

In the simplest GUT based on the gauge group SU(5) with a $\underline{5}_R = \psi^\alpha_R$ and $\underline{10}_L = \psi^{\alpha\beta}_L$ of fermions in each generation, and a $\underline{5}$, $\underline{24}$, and optionally a $\underline{45}$ of Higgs,[7] although B and L are individually violated, the combination B-L is exactly conserved, both at the Lagrangian level and in the presence of nonzero vacuum expectation values of the $\underline{24}$ and $\underline{5}$, and $\underline{45}$ of Higgs.[8] Hence, processes such as $n \to \bar{n}$, and $nn \to \pi$'s, as well as $p \to e^-\pi^+\pi^+$, are strictly forbidden. However, it is easy to add Higgs to the theory which break B-L at the Lagrangian level, even before spontaneous breaking of the SU(5) gauge invariance; for example, a $\underline{10}$ or $\underline{15}$ of Higgs will eliminate the B-L symmetry.[1,2,6,8] One might naively view such an enlargement of the Higgs sector as being unaesthetic. But this view may be misguided. If SU((5) is considered as embedded in SO(10), then the $\underline{126}$ of Higgs, which is probably necessary, contains not just a $\underline{5}$, $\underline{45}$, and $\underline{50}$ of Higgs, but also the (B-L)-violating $\underline{10}$ and $\underline{15}$.

Let us then consider the operators responsible for $n \to \bar{n}$ transitions and the issues involved in experimentally searching for such transitions. The effective Lagrangian is of the form

$$\mathcal{L}_{eff}(n \to \bar{n}) = \sum c_i \mathcal{O}_i + \text{h.c.} \,, \qquad (1)$$

where the $\mathcal{O}_i$ are six-quark operators, and the c_i are c-number coefficients, which thus have dimension $(\text{mass})^{-5}$. Such an effective Lagrangian is obtained in a straightforward manner from the actual Feynman diagrams. For example, the diagram of Fig. 1(a) yields the effective local operator

$$\frac{h_5^2 h_{15} M}{m_5^4 m_{15}^2} [\psi^{T\alpha}_{R_r} C \psi^\gamma_{R_s}] [\bar{\psi}_{L\,\alpha\beta_t} \psi^\beta_{R_x}] [\bar{\psi}_{L\,\gamma\delta_y} \psi^\delta_{R_z}] \quad (2)$$

in terms of SU(5) fields, where α, β, etc. are SU(5) indices; r, s, t, etc. are generation indices; the h's denote (dimensionless) Yukawa couplings, M denotes the 5-5-15 triple Higgs coupling; and m_5 and m_{15} represent the masses of the respective components of Higgs fields. The part of this operator which contributes to $n \to \bar{n}$ transitions corresponds to graph 1(b) and is

$$\frac{h_5^2 h_{15} M}{m_5^4 m_{15}^2} [d^{T\alpha}_R C d^\gamma_R] [u^{T\lambda}_R C d^\mu_R] \times$$
$$\times [u^{T\nu}_R C d^\eta_R] (\epsilon_{\alpha\lambda\mu} \epsilon_{\gamma\nu\eta} + \epsilon_{\gamma\lambda\mu} \epsilon_{\alpha\nu\eta}) \qquad (3)$$

If it happens that all of the masses contributing to the c_i are $\sim M_{GU} \sim 10^{14}$ GeV, then it is easy to see that $n \to \bar{n}$ transitions are negligible: the time $\tau_{n\bar{n}}$ which characterizes such transitions would be

$$\tau_{n\bar{n}} = |\langle\bar{n}| - \int \mathcal{L}_{eff} |n\rangle|^{-1} \sim \frac{M_{GU}^5}{M_N^6} \,, \qquad (4)$$

and hence

$$\tau_{n\bar{n}} \sim \left(\frac{M_{GU}}{M_N}\right) \tau_p >> \tau_p \,, \qquad (5)$$

where τ_p denotes the proton lifetime. Thus, $n-\bar{n}$ oscillations are a very sensitive probe of the so-called desert hypothesis; if they are observed, one will have established not only B-violation and (B-L) violation, but also the existence of a new mass scale much less than 10^{14} GeV.

In complete generality, without recourse to specific diagrams or Higgs, one can show that there are six one-generation SU(5) operators which contribute to $n \to \bar{n}$ transitions.[2,9] They are

$$R_1 = [\bar{\psi}_{L\,\alpha\beta} C \bar{\psi}^T_{L\gamma\delta}] [\psi^{T\alpha}_R C \psi^\gamma_R] [\psi^{T\beta}_R C \psi^\delta_R] \qquad (6)$$

$$R_2 = [\bar{\psi}_{L\,\alpha\beta} \psi^\gamma_R] [\bar{\psi}_{L\,\gamma\delta} \psi^\alpha_R] [\psi^{T\beta}_R C \psi^\delta_R] \qquad (7)$$

$$R_3 = [\bar{\psi}_{L\,\alpha\beta} \psi^\beta_R] [\bar{\psi}_{L\gamma\delta} \psi^\delta_R] [\psi^{T\alpha}_R C \psi^\gamma_R] \qquad (8)$$

$$R_4 = \epsilon_{\alpha\beta\gamma\delta\lambda} [\psi^{T\alpha\beta}_L C \psi^{\gamma\delta}_L] [\bar{\psi}_{L\mu\nu} \psi^\nu_R] [\psi^{T\mu}_R C \psi^\lambda_R] \qquad (9)$$

$$R_5 = \epsilon_{\rho\alpha\beta\gamma\delta} \epsilon_{\sigma\lambda\mu\nu\eta} [\psi^{T\alpha\beta}_L C \psi^{\gamma\delta}_L] [\psi^{T\lambda\mu}_L C \psi^{\nu\eta}_L] [\psi^{T\rho}_R C \psi^\sigma_R] \qquad (10)$$

$$R_6 = \epsilon_{\rho\alpha\beta\lambda\mu} \epsilon_{\sigma\gamma\delta\nu\eta} [\psi^{T\alpha\beta}_L C \psi^{\gamma\delta}_L] [\psi^{T\lambda\mu}_L C \psi^{\nu\eta}_L] [\psi^{T\rho}_R C \psi^\sigma_R] \qquad (11)$$

These yield the following six SU(3) × SU(2) × U(1) operators:[4,9]

$$P_1 = [u_R^{T\alpha} C u_R^{\beta}] [d_R^{T\gamma} C d_R^{\delta}] [d_R^{T\rho} C d_R^{\sigma}] (T_s)_{\alpha\beta\gamma\delta\rho\sigma} \quad (12)$$

$$P_2 = [u_R^{T\alpha} C d_R^{\beta}] [u_R^{T\gamma} C d_R^{\delta}] [d_R^{T\rho} C d_R^{\sigma}] (T_s)_{\alpha\beta\gamma\delta\rho\sigma} \quad (13)$$

$$P_3 = [u_R^{T\alpha} C d_R^{\beta}] [u_R^{T\gamma} C d_R^{\delta}] [d_R^{T\rho} C d_R^{\sigma}] (T_a)_{\alpha\beta\gamma\delta\rho\sigma} \quad (14)$$

$$\begin{aligned} P_4 &= [q_L^{Ti\alpha} C q_L^{j\beta}] [u_R^{T\gamma} C d_R^{\delta}] [d_R^{T\rho} C d_R^{\sigma}] \epsilon_{ij} (T_a)_{\alpha\beta\gamma\delta\rho\sigma} \\ &= 2 [u_L^{T\alpha} C d_L^{\beta}] [u_R^{T\gamma} C d_R^{\delta}] [d_R^{T\rho} C d_R^{\sigma}] (T_a)_{\alpha\beta\gamma\delta\rho\sigma} \end{aligned} \quad (15)$$

$$\begin{aligned} P_5 &= [q_L^{Ti\alpha} C q_L^{j\beta}] [q_L^{Tk\gamma} C q_L^{\ell\delta}] [d_R^{T\rho} C d_R^{\sigma}] \epsilon_{ij}\epsilon_{k\ell}(T_a)_{\alpha\beta\gamma\delta\rho\sigma} \\ &= 4[u_L^{T\alpha} C d_L^{\beta}] [u_L^{T\gamma} C d_L^{\delta}] [d_R^{T\rho} C d_R^{\sigma}] (T_a)_{\alpha\beta\gamma\delta\rho\sigma} \end{aligned} \quad (16)$$

and

$$\begin{aligned} P_6 &= [q_L^{Ti\alpha} C q_L^{j\beta}] [q_L^{Tk\gamma} C q_L^{\ell\delta}] [d_R^{T\rho} C d_R^{\sigma}] (\epsilon_{ik}\epsilon_{j\ell} + \epsilon_{i\ell}\epsilon_{jk}) (T_s)_{\alpha\beta\gamma\delta\rho\sigma} \\ &= 4 \{[u_L^{T\alpha} C u_L^{\beta}] [d_L^{T\gamma} C d_L^{\delta}] [d_R^{T\rho} C d_R^{\sigma}] \\ &\quad -[u_L^{T\alpha} C d_L^{\beta}] [u_L^{T\gamma} C d_L^{\delta}] [d_L^{T\rho} C d_R^{\sigma}]\} (T_s)_{\alpha\beta\gamma\delta\rho\sigma} \end{aligned} \quad (17)$$

where

$$(T_a)_{\alpha\beta\gamma\delta\rho\sigma} = \epsilon_{\rho\alpha\beta}\, \epsilon_{\sigma\gamma\delta} + \epsilon_{\sigma\alpha\beta}\, \epsilon_{\rho\gamma\delta} \quad (18)$$

and

$$(T_s)_{\alpha\beta\gamma\delta\rho\sigma} = \epsilon_{\rho\alpha\gamma}\, \epsilon_{\sigma\beta\delta} + \epsilon_{\sigma\alpha\gamma}\, \epsilon_{\rho\beta\gamma} + \epsilon_{\rho\beta\gamma}\, \epsilon_{\sigma\alpha\delta} + \epsilon_{\sigma\beta\gamma}\, \epsilon_{\rho\alpha\delta} \quad (19)$$

and here the indices α, β, etc. represent SU(3) indices while i, j, k, etc. represent SU(2) indices. The precise correspondence is[9]: $R_m \to P_m$, m = 1-5 and $R_6 \to -P_5 + P_6$. One may also enumerate the total number of SU(3) × U(1)-invariant operators which contribute to n-$\bar{n}$ oscillations: there are eighteen of these.[9] Although the additional twelve non SU(2)-invariant operators may be suppressed somewhat relative to the fully SU(3) × SU(2) × U(1)-invariant operators, they do not necessarily give negligible contributions. The matrix elements $\langle \bar{n} | \int \mathcal{O}_1 | n \rangle$ of these operators can be reasonably reliably calculated in the framework of the standard MIT bag model of hadron structure.[9]

The analysis of n ↔ $\bar{n}$ transitions proceeds by considering the 2 × 2 mass matrix

$$m \equiv \begin{pmatrix} \langle n|H|n\rangle & \langle n|H|\bar{n}\rangle \\ \langle \bar{n}|H|n\rangle & \langle \bar{n}|H|\bar{n}\rangle \end{pmatrix} \quad (20)$$

This analysis will be independent of any particular grand unified theory. In field-free vacuum the mass matrix takes the form

$$m = \begin{pmatrix} m_n - i\frac{\lambda}{2} & \delta m \\ \delta m & m_n - \frac{i\lambda}{2} \end{pmatrix} \quad (21)$$

where $\delta m = \langle \bar{n}| - \int \mathcal{L}_{eff} | n \rangle$ and $\lambda = 1/\tau_n$ is the rate for the regular weak decay of the neutron. The resulting mass eigenstates are

$$|n_\pm\rangle = \frac{1}{\sqrt{2}} [|n\rangle \pm |\bar{n}\rangle] \quad (22)$$

with masses

$$m_\pm = m_n \pm \delta_m \quad (23)$$

Thus, if one starts with a pure beam of neutrons at proper time $\tau = 0$, after propagation for a time τ, there is a nonzero probability to find a component of antineutrons in the beam given by

$$\begin{aligned} P(n(\tau) = \bar{n}) &\equiv |\langle \bar{n}|n(\tau)\rangle|^2 = \sin^2\{(\delta m)\tau\}\, e^{-\lambda\tau} \\ &\equiv \sin^2(\tau/\tau_{n\bar{n}})\, e^{-\lambda\tau} \end{aligned} \quad (24)$$

This is evidently a situation of maximum mixing of n and $\bar{n}$ to form mass eigenstates.

The n ↔ $\bar{n}$ transition amplitude gives rise to matter instability. Indeed, an important lower bound on $\tau_{n\bar{n}}$ can be derived from the existing lower bound on $\tau_{matter\ (instability)} \gtrsim 3 \times 10^{30}$ yr. To obtain this limit, consider bound neutrons in matter. The crucial point is that an n and an $\bar{n}$ are subject to different nuclear potentials, V_n and $V_{\bar{n}}$. The mass matrix for this case is

$$m = \begin{pmatrix} (m_n)_{eff} & \delta m \\ \delta m & (m_{\bar{n}})_{eff} \end{pmatrix} \quad (25)$$

where

$$(m_{(\bar{n})})_{eff} \equiv m_n + V_{(\bar{n})} \quad .$$

Numerically,[10] $V_n = \mathrm{Re}\,(V_n) = -50$ MeV. There is, however, considerable uncertainty in the value of $V_{\bar{n}}$. One of the main determinations of this quantity comes from $\bar{p}$-atom de-excitation spectra; a recent analysis[11] gives $V_{\bar{n}} = -(130 + 200i)$ MeV. Another source of information on Im($V_{\bar{n}}$) is low-energy $\bar{p}$-p scattering[12]; this yields the larger values Im($V_{\bar{n}}$) = -850 MeV for I (isospin) = 1 and -659 MeV for I = 0. Moreover, a recent theoretical estimate[13] is $V_{\bar{n}} = -(1 + 0.7i) \times 10^3$ MeV, much higher in magnitude than the values inferred from $\bar{p}$-atom data. We shall comment on the implications of this uncertainty further below.

The diagonalization of the matrix m in Eq. (25) yields results quite different from those of Eqs. (22) and (23); the eigenstates are

$$|n_+\rangle = \cos\theta|n\rangle + \sin\theta|\bar{n}\rangle$$

and

$$|n_-\rangle = -\sin\theta|n\rangle + \cos\theta|\bar{n}\rangle \qquad (26)$$

where

$$\tan\theta = \frac{\delta m}{\frac{1}{2}\left[(m_n)_{eff} - (m_{\bar{n}})_{eff} + \{[(m_n)_{eff} - (m_{\bar{n}})_{eff}]^2 + 4(\delta m)^2\}^{1/2}\right]} \qquad (27)$$

with eigenvalues

$$m_\pm = \tfrac{1}{2}[\{(m_n)_{eff} + (m_{\bar{n}})_{eff}\} \pm \{[(m_n)_{eff} - (m_{\bar{n}})_{eff}]^2 + 4(\delta m)^2\}^{1/2}] \qquad (28a)$$

i.e.,

$$m_+ \simeq m_n + V_n + \frac{(\delta m)^2[(V_n - \mathrm{Re}(V_{\bar{n}})) + i\,\mathrm{Im}(V_{\bar{n}})]}{[\{V_n - \mathrm{Re}(V_{\bar{n}})\}^2 + \{\mathrm{Im}(V_{\bar{n}})\}^2]}. \qquad (28b)$$

and

$$m_- \simeq m_n + \mathrm{Re}(V_{\bar{n}}) + i\,\mathrm{Im}(V_{\bar{n}}) + \mathcal{O}\left(\frac{(\delta m)^2}{|V_{\bar{n}}|}\right) \qquad (28c)$$

As will be shown, $(\delta m)_{max} \simeq 10^{-28}$ MeV $<< |V_{n,\bar{n}}|$, so $|\theta| << 1$. That is, the $n \to \bar{n}$ mixing is extremely strongly suppressed in matter. The imaginary term in m_+ represents the decay of the $|n_+\rangle$ state via annihilation of the $\bar{n}$ with other nucleons. The mean rate for this decay is

$$\Gamma_{|n_+\rangle\ decay} = \frac{2(\delta m)^2|\mathrm{Im}\ (V_{\bar{n}})|}{[\{V_n - \mathrm{Re}\ (V_{\bar{n}})\}^2 + \{\mathrm{Im}\ (V_{\bar{n}})\}^2]} \quad . \qquad (29)$$

Then

$$\Gamma_{\substack{matter\\decay}} \simeq \Gamma_{\substack{nucleon\\decay}} + \Gamma_{|n_+\rangle\ decay} \quad . \qquad (30)$$

To obtain an upper limit on the $|n_+\rangle$ decay rate, assume that it saturates the right hand side of Eq. (30); then

$$\tau_{matter} \simeq \frac{[\{V_n - \mathrm{Re}\ (V_{\bar{n}})\}^2 + \{\mathrm{Im}\ (V_{\bar{n}})\}^2]\tau^2_{n\bar{n}}}{2|\mathrm{Im}\ (V_{\bar{n}})|} \qquad (31)$$

As is evident from Eq. (31), if one can experimentally increase the lower bound on $\tau_{n\bar{n}}$ by a certain amount, one thereby increases the lower bound on τ_{matter} instability by the square of this amount. The great sensitivity of $n - \bar{n}$ mixing to B-violation arises from the fact that it is an interference effect and depends on an amplitude rather than the square of an amplitude. Numerically, the lower limit[14] $\tau_{matter} \gtrsim 3 \times 10^{30}$ yr and the value[11] $V_{\bar{n}} = -(130 + 200i)$ imply that $\tau_{n\bar{n}} \gtrsim 2 \times 10^7$ sec. It is striking that such a large value of τ_{matter} allows such a small value of $\tau_{n\bar{n}}$. In Table I, we list some corresponding values of lower limits of τ_{matter} and $\tau_{n\bar{n}}$. We estimate that there is a theoretical uncertainty of order ~ 3-5 in the determination of $(\tau_{n\bar{n}})_{min}$ from $(\tau_{matter})_{min}$. This table will serve later as the basis for our comparison of the relative sensitivities and advantages of reactor and nucleon decay experiments as probes of $n \to \bar{n}$ transitions. If one used values of $V_{\bar{n}}$ which have larger magnitudes, such as those given in Refs. 12 and 13, one would obtain commensurately lower values of $(\tau_{n\bar{n}})_{min}$ for a given $(\tau_{matter})_{min}$; this, in turn would increase the inferred sensitivity of $n - \bar{n}$ propagation experiments relative to nucleon decay experiments.

We next consider the techniques and status of current searches for $n \to \bar{n}$ transitions, and proposals for future experiments.[15] There are two classes of experiments to search for such transitions: (1) those which try to detect antineutrons in a beam which is initially comprised purely of neutrons, after it has been propagated for a certain distance, and (2) those which try to detect the decay of the $|n_+\rangle$ state of the mixed n-$\bar{n}$ system in matter. We shall deal with these in turn.

The propagation of a neutron or antineutron in a terrestrial experiment is determined mainly by its interaction with the earth's magnetic field, $\vec{B}$. For this case, the n-$\bar{n}$ mass matrix is

$$m = \begin{pmatrix} m_n - \vec{\mu}_n\cdot\vec{B} & \delta m \\ \delta m & m_n + \vec{\mu}_n\cdot\vec{B} \end{pmatrix} \qquad (32)$$

The resulting mass eigenstates are given by Eq. (26), with

$$\tan\theta = \frac{\delta m}{[-\vec{\mu}_n\cdot\vec{B} + \{(\vec{\mu}\cdot\vec{B}_n)^2 + (\delta m)^2\}^{1/2}} \qquad (33)$$

and

$$m_\pm = m_n \pm \{(\vec{\mu}_n\cdot\vec{B})^2 + (\delta m)^2\}^{1/2} \qquad (34)$$

In practical experiments, it is possible to shield the earth's magnetic field down to the level $|B| \sim 2 \times 10^{-4}$ G whence $\mu_n|\vec{B}| \sim 10^{-21}$ MeV $>> (\delta m)_{max} \sim 10^{-28}$ MeV. Therefore, $|\theta| << 1$, as was the case with bound neutrons in matter. The oscillation probability is

$$P(n(t) = \bar{n}) = \sin^2 2\theta \sin^2 \left(\frac{(\Delta m)\tau}{2}\right) e^{-\lambda\tau}, \quad (35)$$

where $\Delta m \equiv m_+ - m_-$. At first sight, such an experiment seems hopeless, since $|\theta| << 1$ and Δm is very insensitive to δm. However, if one arranges that $1/2|\Delta m|\tau \sim \mu_n|\vec{B}|\tau << 1$, then there is a cancellation of the $\vec{\mu}_n \cdot \vec{B}$ term

$$P(n(t) = \bar{n}) \approx ((\delta m)\tau)^2 = \left(\frac{\tau}{\tau_{n\bar{n}}}\right)^2 \quad (36)$$

In a propagation experiment, one thus uses slow neutrons from a reactor (to maximize the flux) and lets them propagate for a certain length L, carefully shielding this region from the earth's magnetic field. Since under the assumptions of Eq. (36), the $\bar{n}$ yield is proportional to τ^2, one designs the experiment so as to maximize τ, subject to the crucial condition that $\mu_n|\vec{B}|\tau << 1$. The beam then impinges on a target, and one searches for the $n\bar{n}$ annihilation products, which are pions (primarily) with mean multiplicity 4-5, total energy $\sim 2m_n$, total charge zero, and total momentum $\sim$ zero, up to a correction due to the Fermi momentum of the bound nucleon. There are several checks in such an experiment, including running with the reactor on and off and with the magnetic field shielded and unshielded.

Neutron-antineutron propagation experiments which are running or proposed include, first, an experiment by a CERN-ILL-Padova-REHL-Sussex collaboration at the Grenoble reactor.[16] The first phase of this experiment is now complete and the group has reported the lower bound $\tau_{n\bar{n}} > 1 \times 10^5$ sec (90% CL).[16] The experiment uses ultracold neutrons transported via completely reflecting guides which are magnetically shielded down to the level of a few $\times 10^{-4}$ G. It employs a plastic calorimeter with lead layers in order to detect the photons from π^0's produced by the $\bar{n}n$ annihilation. The group reports that the photons from the reactor constitute a negligible background. Cosmic rays are vetoed with 99.95% efficiency but still represent the main source of background. The general sensitivity is determined by the parameters of the experiment: neutron current $I \simeq 10^9$ n/sec, total running time $T \simeq 10^7$ sec, and neutron travel time $t = \tau \simeq 10^{-1} - 10^{-2}$ sec. A second generation experiment by the same collaboration at Grenoble is in progress at the present time. It improves on the first search by having a better calorimeter made up of a layered streamer chamber for good track reconstruction. Further, this chamber covers a greater solid angle: 90% of 4π sr. In the most advanced stage planned for the experiment, the neutron flux will be increased to $\sim 10^{12}$ n/sec.

A second propagation experiment is being conducted by a Pavia-Rome collaboration at the Pavia reactor.[17] The group uses thermal neutrons with a flux of 2×10^{11} n/sec at the target. The detector is a lead flash chamber calorimeter and scintillator hodoscope with 65% solid angle coverage. This experiment claims an estimated sensitivity of $(\tau_{n\bar{n}})_{min} \simeq 10^6$ sec (90% CL).

In the U.S., a Harvard-Oak Ridge-University of Tennessee collaboration is planning to perform a propagation experiment at the Oak Ridge Research (ORR) reactor using thermal neutrons with a 20 m flight path and a flux on target of $\gtrsim 2 \times 10^{13}$ n/sec.[18] The central detector consists of a lead-glass Cherenkov counter and position-sensitive detectors covering >80% of the full solid angle. The estimated sensitivity of this experiment is $(\tau_{n\bar{n}})_{min} \simeq 2 \times 10^8$ sec (90% CL) for a run of 120 days. This sensitivity could be increased in the future with the addition of a cold source at the reactor port and a longer neutron flight path. There are, finally, two proposals for similar experiments at LASL.[12,16,17] The first[19] would use LAMPF as a source of thermal neutrons with neutrons with a flux of 2-4 $\times 10^{12}$ n/sec and a detector consisting of time-of-flight and tracking chambers, but no calorimeter. The second[20] would use the Omega West reactor for a thermal neutron flux of 3×10^{11} n/sec and a detector comprised of a liquid scintillator which covers 95% of 4π sr and yields time of flight and calorimetry information. The estimated sensitivities of these two LASL proposals are $(\tau_{n\bar{n}})_{min} \approx 4 \times 10^7$ sec and 2×10^7 sec (90% CL), respectively. Thus, existing and proposed propagation experiments to test for $n \to \bar{n}$ transitions hope to achieve a lower limit of about $(\tau_{n\bar{n}}) \simeq 10^8$ sec (90% CL) or detect such transitions if they occur at or below this level.

We next consider the second class of n-$\bar{n}$ search experiments, namely those which use nucleon decay detectors.[21] The signature for the decay of the $|n_+\rangle$ state is the same as in the propagation experiments, since in the latter case the momentum of the incident neutron beam is negligibly small. The great suppression of n-$\bar{n}$ mixing for bound neutrons is compensated for by the huge mass of the detector and the long time of observation. Since such nucleon decay detectors are positioned in deep mines (or auto tunnels, in Europe), they are well shielded from cosmic rays. Water Cherenkov experiments such as that of the Irvine-Michigan-Brookhaven collaboration are not optimized to search for $n\bar{n}$ annihilation, since, although they can detect the photons from the π^0's, they are relatively insensitive to the charged π's, and hence cannot test the constraints of $2m_n$ total energy, zero charge, or zero momentum to within Fermi motion corrections. In contrast, multikiloton fine-grained detectors such as those proposed for the next generation of proton decay experiments[21] would be able to detect charged as well as neutral pions with high efficiency and apply the above-mentioned constraints. Indeed, in view of this fact, we feel that such "nucleon decay" experiments might more aptly be called "matter instability" experiments, since they can search for both nucleon decay and $n\bar{n}$ annihilation due to $n \to \bar{n}$ transitions with very good sensitivity.

The values of $\tau_{n\bar{n}}$ and τ_{matter} calculated in Table 1 provide a comparison of the relative sensitivities of n-$\bar{n}$ propagation experiments and matter instability experiments. Since existing or proposed propagation searches can reach the level of $\tau_{n\bar{n}} \simeq 10^8$ which corresponds approximately to $\tau_{matter} \simeq 10^{32}$ they are a valuable complement to nucleon decay experiments in this region of matter instability lifetimes. It should, of course, be kept in mind that the nucleon decay detectors can observe both (B-L)-conserving nucleon decays and (B-L)-violating decays such as $p \to e^-\pi^+\pi^+$ and $n \to \bar{n}$ transitions (via $|n_+\rangle$ decay), while the propagation experiments can only test for the third of these processes. If matter instability lifetimes turn out to be $\gtrsim 10^{32}$ yr, then baryon number as well as (B-L) violation would be best searched for with massive fine-grained nucleon decay detectors.

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Footnotes and References

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Table 1

$(t_{\text{matter instability}})_{min}$ (yr)	$(\tau_{n\bar{n}})_{min}$ (sec)
3×10^{30}	2×10^{7}
1×10^{31}	3×10^{7}
1×10^{32}	1×10^{8}
	→
1×10^{33}	3×10^{8}
1×10^{34}	1×10^{9}

Corresponding values of $(\tau_{\text{matter instability}})_{min}$ and $(\tau_{n\bar{n}})_{min}$. The arrow indicates the anticipated ultimate lower bound on $\tau_{n\bar{n}}$ which can be achieved from n-$\bar{n}$ propagation experiments. See text for further details.

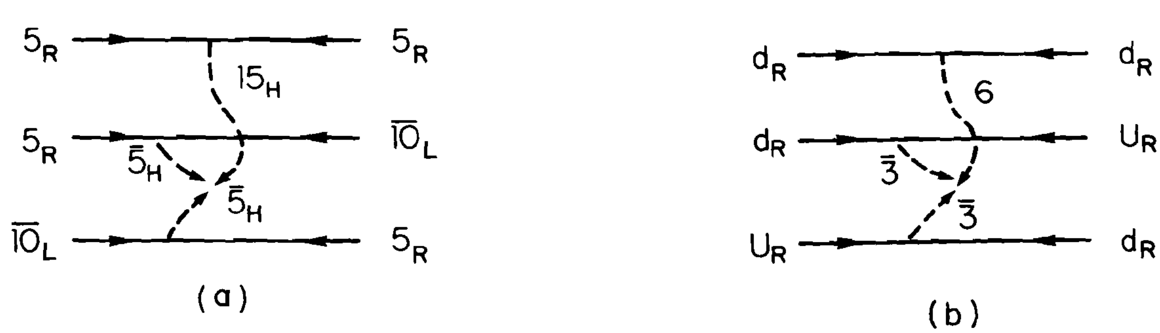

Figure 1

A graph contributing to n → $\bar{n}$ transitions, in terms of SU(5) fields.

The part of graph (a) which contributes to n → $\bar{n}$ transitions, in terms of quarks. The number on the Higgs lines refer to SU(3) representation.