Dark particle interpretation of the neutron decay anomaly

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Abstract. There is a long-standing discrepancy between the neutron lifetime measured in beam and bottle experiments. We propose to explain this anomaly by a dark decay channel for the neutron, involving a dark sector particle in the final state. If this particle is stable, it can be the dark matter. Its mass is close to the neutron mass, suggesting a connection between dark and baryonic matter. In the most interesting scenario a monochromatic photon with energy in the range $0.782 \text{ MeV} - 1.664 \text{ MeV}$ and branching fraction $1\%$ is expected in the final state. We construct representative particle physics models consistent with all experimental constraints.

1. Introduction

In the previous talk, A. Holley presented the status of measurements of neutron lifetime. We will not repeat this here. The purpose of this talk is to contemplate the possibility that the discrepancy between neutron lifetime measurements in bottle and beam experiments is due to a new phenomenon, rather than a systematic error or a statistical fluctuation [1]. The measurements differ by about $1\%$; in terms of the decay widths, \( \Delta \Gamma_n^\text{exp} = \Gamma_n^\text{bottle} - \Gamma_n^\text{beam} \approx 7.1 \times 10^{-30} \text{ GeV} \approx (9.5 \times 10^{-3}) \Gamma_n^\text{bottle} \).

Basic principles demand that the particle lifetime does not depend on the measurement technique or decay mode. For a new physics explanation of the discrepancy we posit that one of the experiments is not reporting a lifetime. The bottle experiment fits to a decaying in time exponential, measuring the lifetime by definition. An important observation is that the beam experiment, by contrast, measures the instantaneous \( \beta \)-decay rate:

\[
\tau_n^\text{beam} = \frac{\tau_n}{\text{Br}(n \rightarrow p + \text{anything})}.
\]

This gives the decay lifetime only if the neutron \( \beta \)-decays $100\%$ of the time. Therefore, the discrepancy could be explained if the neutron decays about $1\%$ of the time into a protonless final state. A very direct test of this hypothesis could be performed, but has not been done, in both bottle and beam experiments. In a bottle experiment one could measure the number of decays that result in final state protons: the hypothesis asserts that these would account only for $99\%$ of the neutrons lost. In a beam experiment one could fit the number of observed protons to an exponential decaying along the beam: here the fit to the exponential would result in a $1\%$ smaller lifetime than the one determined from the instantaneous rate at one point along the beam.
The question immediately arises as to what is the new decay channel of the neutron, a protonless state with 1% branching fraction. It is incumbent upon us to present a scenario that is consistent with our current understanding of elementary particle interactions.

2. Model independent analysis

The final state of neutron decay must have unit fermion number and be overall electrically neutral. The simplest possibility for a protonless decay channel consists of a chargeless, “dark” fermion, $\chi$, and a neutral boson. In principle the dark fermion could be a neutrino and the neutral boson could be a photon, but soon we will see that the sum of their masses must be within 1.572 MeV of the neutron mass, so at least one of the two must be a new elementary particle. We consider two possible final states: (i) a dark fermion, $\chi$, and a photon, and (ii) a dark fermion, $\chi$, and a new spinless elementary dark boson, $\phi$.

In either case, the mass $M_f$ of the final state $f$ must be sufficiently large that proton stability is not vitiated: to avoid $p \to n^+ e^- \nu$ followed by $n^+ \to f$, we must have $M_f > m_p - m_e$. A slightly stronger bound follows from stability of $^9$Be against decay to $^2$H + $f$, yielding [2]:

$$937.992 \text{ MeV} < M_f < 939.565 \text{ MeV}.$$  \hfill (2)

For the case $f = \chi + \gamma$ this immediately gives $937.992 \text{ MeV} < m_\chi < 939.565 \text{ MeV}$. For $m_\chi < m_p + m_e \approx 938.783 \text{ MeV}$, $\beta$ decay of $\chi$ is not allowed giving the interesting possibility that $\chi$ is stable (or very long lived) and could be a candidate for a dark matter component of the universe. Similarly, for the case $f = \chi + \phi$ both $\chi$ and $\phi$ are dark matter candidates if $|m_\chi - m_\phi| < m_p + m_e$.

2.1. Effective theory for $\chi \gamma$ final state

To describe the decay $n \to \chi \gamma$ in a quantitative way, we consider theories with a mass mixing term $\chi n$, and an induced interaction $\chi n \gamma$. An example of such a theory is given by the effective Lagrangian

$$L^\text{eff}_1 = \bar{n} \left( i \partial - m_n + \frac{g_n e}{2 m_n} \sigma^{\mu \nu} F_{\mu \nu} \right) n + \bar{\chi} \left( i \partial - m_\chi \right) \chi + \varepsilon (\bar{n} \chi + \bar{\chi} n),$$  \hfill (3)

where $g_n \approx -3.826$ is the neutron $g$-factor and $\varepsilon$ is the mixing parameter with dimension of mass. The term corresponding to $n \to \chi \gamma$ is obtained by transforming Eq. (3) to the mass eigenstate basis and, for $\varepsilon \ll m_n - m_\chi$, yields

$$L^\text{eff}_{n \to \chi \gamma} = \frac{g_n e}{2 m_n} \varepsilon \bar{\chi} \gamma^\mu F_{\mu \nu} n.$$  \hfill (4)

Therefore, the neutron dark decay rate is

$$\Delta \Gamma_{n \to \chi \gamma} = \frac{g_n^2 e^2}{8 \pi} \left( 1 - \frac{m_\chi^2}{m_n^2} \right)^3 \frac{m_n \varepsilon^2}{(m_n - m_\chi)^2} \approx \Delta \Gamma_{n \to \chi \gamma}^\text{exp} \left( \frac{1+x}{2} \right)^3 \left( \frac{1-x}{1.8 \times 10^{-3}} \right) \left( \varepsilon [\text{GeV}] \right)^2,$$

where $x = m_\chi/m_n$. The rate is maximized when $m_\chi$ saturates the lower bound, $m_\chi = 937.992 \text{ MeV}$. Below we will give a microscopic particle physics realization of this case.

The testable prediction of this class of models is a monochromatic photon with an energy in the range $0 < E_\gamma < 1.572 \text{ MeV}$ and a branching fraction

$$\frac{\Delta \Gamma_{n \to \chi \gamma}}{\Gamma_n} \approx 1\%.$$  

If the dark fermion $\chi$ is to be sufficiently light that it may be stable, and hence a DM candidate, then $0.782 \text{ MeV} < E_\gamma < 1.572 \text{ MeV}$. A null search for monochromatic photons rules out this
hypothesis in the range 0.782 to 1.664 MeV at the 97% confidence level [3]. A signature involving
an $e^+e^-$ pair with total energy $E_{e^+e^-} < 1.572$ MeV is also expected, but with a suppressed
branching fraction of at most $1.1 \times 10^{-6}$. A null search for $n \rightarrow \chi e^+e^-$ sets a bound of $\sim 10^{-4}$
on the branching fraction with 99.9997% significance for 937.900 MeV $< m_\chi < 938.543$ MeV [4].

2.2. Effective theory for $\chi \phi$ final state
Consider now a scenario in which the neutron decays to an intermediate off-shell dark fermion,
$\tilde{\chi}$, which subsequently decays to a final state dark fermion and a scalar, $\chi$ and $\phi$. The masses
of the final state particles are constrained by Eq. (2), with $M_f = m_\chi + m_\phi$, and both $\chi$ and $\phi$ are stable if $|m_\chi - m_\phi| < 938.783$ MeV. It is also possible to have an on-shell decay to $\tilde{\chi}\gamma$ if
937.992 MeV $< m_\tilde{\chi} < 939.565$ MeV, but we are mostly interested in the case $m_\tilde{\chi} > 939.565$ MeV
for which the neutron decay is dark.

An example of such a theory is

$$\mathcal{L}_2^{\text{eff}} = \bar{n} \left( i \frac{\partial}{\partial t} - m_n + \frac{g_{\nu
u} \sigma^{\mu\nu}}{2m_n} F_{\mu\nu} \right) n + \tilde{\chi} \left( i \frac{\partial}{\partial t} - m_{\tilde{\chi}} \right) \tilde{\chi} + \tilde{\chi} \left( i \frac{\partial}{\partial t} - m_\chi \right) \chi + \partial_\mu \phi \partial^\mu \phi - \frac{m_\phi^2}{2} |\phi|^2 + \varepsilon (n\tilde{\chi} + \tilde{\chi}n) + (\lambda_\phi \bar{n} \chi \phi + \text{h.c.}). \quad (5)$$

Going to the mass eigenstate basis, the term corresponding to $n \rightarrow \chi \phi$ is

$$\mathcal{L}_{n \rightarrow \chi \phi}^{\text{eff}} = \frac{\lambda_\phi \varepsilon}{m_n - m_{\tilde{\chi}}} \bar{n} \phi^* \chi \phi, \quad (6)$$

which yields the neutron dark decay rate

$$\Delta \Gamma_{n \rightarrow \chi \phi} = \frac{|\lambda_\phi|^2}{16\pi} \sqrt{f(x,y)} \frac{m_n \varepsilon^2}{(m_n - m_{\tilde{\chi}})^2}, \quad (7)$$

where $f(x,y) = [(1-x)^2 - y^2]/[(1+x)^2 - y^2]^3$ with $x = m_{\tilde{\chi}}/m_n$ and $y = m_\phi/m_n$. A microscopic,
particle physics realization of this scenario is provided below.

For $m_{\tilde{\chi}} > m_n$ the missing energy signature has a branching fraction $\approx 1\%$. There will also be
a very suppressed radiative process involving a photon in the final state with a branching fraction $3.5 \times 10^{-10}$ or smaller.

As discussed earlier, in the case 937.992 MeV $< m_{\tilde{\chi}} < m_n$ both the visible and invisible
neutron dark decay channels are present. The ratio of their branching fractions is

$$\frac{\Delta \Gamma_{n \rightarrow \chi \gamma}}{\Delta \Gamma_{n \rightarrow \chi \phi}} = \frac{2g_\gamma^2 \varepsilon^2 (1 - \tilde{x}^2)^3}{|\lambda_\phi|^2 \sqrt{f(x,y)}}, \quad (8)$$

where $\tilde{x} = m_{\tilde{\chi}}/m_n$, while their sum accounts for the neutron decay anomaly, i.e.

$$\frac{\Delta \Gamma_{n \rightarrow \chi \gamma} + \Delta \Gamma_{n \rightarrow \chi \phi}}{\Gamma_n} \approx 1\% \quad (9)$$

The branching fraction for the process involving a photon in the final state ranges from about
0 to 1%, depending on the masses and couplings. A suppressed decay channel involving $e^+e^-$
is also present.

3. Microscopic Models
While the effective theory analysis presented above is useful in estimating the dark decay
parameters that a fundamental microscopic model must produce, the existence of a viable model
with renormalizable interactions that satisfies all experimental constraints is not obvious. We will
present one example of a microscopic model for each of the scenarios above. Other possibilities may exist, but we are satisfied for now by presenting these working examples.

The minimal model for the $n \rightarrow \chi \gamma$ requires only two particles beyond the SM: a scalar $\Phi = (3,1)_{-1/3}$ (color triplet, weak singlet, hypercharge $-1/3$), and a Dirac fermion $\chi$ (SM singlet). The interaction Lagrangian of the model includes

$$\mathcal{L}_1 = \lambda_\Phi \epsilon^{ijk} u_{Li}^{\dagger} d_{Rj} \Phi_k + \lambda_\chi \Phi_{\ast} \bar{\chi} d_{Ri} + \text{h.c.},$$

(10)

where $u_i^L$ is the charge conjugate of $u_R$. Assigning baryon numbers $B_\chi = 1$, $B_\Phi = -2/3$ we see that proton decay is forbidden [5–7]. The rate for $n \rightarrow \chi \gamma$ is given by Eq. (5) with

$$\varepsilon = \frac{\beta \lambda_\Phi \lambda_\chi}{M_\Phi^2},$$

(11)

and $\beta$ defined by $\langle 0| \epsilon^{ijk} (u_{Li}^{\dagger} d_{Rj}) d_{Rk}^\sigma |n \rangle = \beta \frac{1}{2} (1 + \gamma_5) u^\sigma$, where $u$ is the neutron spinor, $\sigma$ is the spinor index and the parenthesis denote spinor contraction. Lattice QCD calculations give $\beta = 0.0144(3)\, 21$ GeV$^3$ [8], where the errors are statistical and systematic, respectively. Assuming $m_\chi = 937.992$ MeV to maximize the rate, in order to explain the anomaly the parameters must satisfy

$$\frac{|\lambda_\Phi \lambda_\chi|}{M_\Phi^2} \approx 6.7 \times 10^{-6} \, \text{TeV}^{-2}.$$  

(12)

In addition to the monochromatic photon with energy $E_\gamma < 1.572$ MeV and the $e^+e^-$ signal, one may search directly also for $\Phi$. It can be singly produced through $pp \rightarrow \Phi$ or pair produced via gluon fusion $gg \rightarrow \Phi \Phi$. This results in a dijet or four-jet signal from $\Phi \rightarrow d \chi$. Given Eq. (12), $\Phi$ is not excluded by recent LHC analyses [9–14] provided $M_\Phi \gtrsim 1$ TeV.

The parameter choice in Eq. (12) is excluded if $\chi$ is a Majorana particle, as in the model proposed in [15], by the neutron-antineutron oscillation and dinucleon decay constraints [16,17]. Neutron decays considered in [18] are too suppressed to account for the neutron decay anomaly.

A representative model for the case $n \rightarrow \chi \phi$ involves four new particles: the scalar $\Phi = (3,1)_{-1/3}$, two Dirac fermions $\bar{\chi}$, $\chi$, and a complex scalar $\phi$, the last three being SM singlets. The interaction Lagrangian is an extension of the one in (10)

$$\mathcal{L}_2 = \mathcal{L}_1 (\chi \rightarrow \bar{\chi}) + (\lambda_\phi \bar{\chi} \chi \phi + \text{h.c.}) .$$

(13)

We have imposed an additional $U(1)$ symmetry under which $\chi$ and $\phi$ have opposite charges. For $m_\chi > m_\phi$ the annihilation channel $\chi \bar{\chi} \rightarrow \phi \bar{\phi}$ via a $t$-channel $\chi$ exchange is open. The observed DM relic density, assuming $m_\chi = 937.992$ MeV and $m_\phi \approx 0$, is obtained for $\lambda_\phi \simeq 0.037$. Alternatively, the DM can be non-thermally produced.

The rate for $n \rightarrow \chi \phi$ is described by Eq. (7) with $\varepsilon = \beta \lambda_\Phi \lambda_\chi / M_\Phi^2$. For $m_\chi = m_\phi$, the anomaly is explained with

$$\frac{|\lambda_\Phi \lambda_\chi|}{M_\Phi^2} \approx 4.9 \times 10^{-7} \, \text{TeV}^{-2}.$$  

(14)

For $\lambda_\phi \approx 0.04$ this is consistent with LHC searches, provided again that $M_\Phi \gtrsim 1$ TeV. Direct DM detection searches present no constraints. For similar reasons as before, $\chi$ and $\bar{\chi}$ cannot be Majorana particles.

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1 A similar model with a scalar $\Phi = (3,1)_{2/3}$ and a Dirac fermion $\chi$ would also work. Since the scalar $(3,1)_{2/3}$ cannot couple to two first generation quarks, the rate in Eq. (12) would be suppressed by the strange quark content of the neutron and would require a larger value of $|\lambda_\Phi \lambda_\chi| / M_\Phi^2$. Another viable option for $\Phi$ is the vector $(3,2)_{1/6}$.
As discussed above, in this model the branching fractions for the visible (including a photon) and invisible final states can be comparable, and their relative size is described by Eq. (8). A final state containing an $e^+e^-$ pair is also possible. The same LHC signatures are expected as in model 1.

4. Further Developments

4.1. Unstable Nucleus Decay

While the lower bound in (2) with $M_f = m_\chi$ forbids the decay $(A, Z) \rightarrow (A - 1, Z) + \chi$ for any stable nucleus, it is still permitted for some unstable nuclei. There are several unstable isotopes with a neutron binding energy $S(n) < 1.572$ MeV and a sufficiently long lifetime to probe the dark decay channel when the dark particle mass $m_\chi < m_n - S(n)$. One example presented in [1] is $^{11}\text{Li}$, for which $S(n) = 0.396$ MeV. $^{11}\text{Li} \beta$ decays with a lifetime 8.75 ms. However, in the presence of a dark particle $\chi$ the decay chain $^{11}\text{Li} \rightarrow ^{10}\text{Li} + \chi \rightarrow ^9\text{Li} + n + \chi$ becomes available. Ref. [2] proposes to search for $^{11}\text{Be} \rightarrow ^{10}\text{Be} + \chi$. This has $S(n) = 501.6(3)$ keV and half-life 13.76(7) s. Since this is a halo nucleus, the calculation of the decay rate $^{11}\text{Be} \rightarrow ^{10}\text{Be} + \chi + \gamma$ can be well estimated from that of $n \rightarrow \chi + \gamma$ by accounting for the phase-space factor. This leads to a branching fraction of $10^{-6} - 10^{-4}$ for 938 MeV $\lesssim m_\chi \lesssim 939$ MeV. However, the dark decay of $^{11}\text{Be}$ does not necessarily involve a photon in the final state. For example, if the model (3) is supplemented by a $\pi$-nucleon interaction, $\pi n\bar{n}$, or a two-body nucleon interaction, $\bar{n}n\overline{n}n$, then in terms of the mass eigenstates one obtains interactions with the dark fermion, $\pi\bar{\chi}n$ and $\bar{\chi}n\bar{n}n$ that can directly mediate $(A, Z) \rightarrow (A - 1, Z) + \chi$. Ref. [19] estimates the decay width for $^{11}\text{Be} \rightarrow ^{10}\text{Be} + \chi$ to be an order of magnitude larger than the measured total width of $^{11}\text{Be}$ for $m_\chi \lesssim 939$ MeV, excluding this model. However, we note that the result of Ref. [19] exhibits unphysical behavior: the decay amplitude, expressed as a function of $x = m_n - m_\chi - S(n)$, has singular behavior as $x \rightarrow 0$. No independent verification of the result is available. Yet, the computation uses the mixing term $\bar{\chi}n$, rather than the induced interactions, $\pi\bar{\chi}n$ or $\bar{\chi}n\bar{n}n$, to mediate the transition, and it is likely this artifact is responsible for the unphysical behavior.

The calculated branching fraction for dark $^{11}\text{Be}$ decay in Ref. [2] applies just as well to the purely dark decay of the model in (5). The branching fraction for $^{11}\text{Be}$ decay into $^{10}\text{Be}$ has been measured, and the result is some 400 times larger than that expected from $\beta$-delayed proton emission [20]. This could be explained if a yet unobserved very low lying and very narrow $^{11}\text{B}$ resonance existed. The authors of Ref. [2] propose an experiment measuring protons in $^{11}\text{Be}$ decay. This may discover the putative resonance, or, in the absence of the expected proton signal, give supporting evidence for the dark decay hypothesis.

4.2. Neutron Stars

The impact of neutron dark decays on neutron stars was considered in Refs. [21–23]. The resulting production of dark particles changes the energy density and pressure inside a neutron star, modifying its equation of state. This in turn changes the predictions for the maximum allowed neutron star masses, since they are derived from integrating the Tolman-Oppenheimer-Volkoff equation that explicitly depends on the equation of state.

It was shown that the observed neutron star masses ($2M_\odot$ for the heaviest neutron stars discovered) are allowed if strong repulsive self-interactions are present in the dark sector of our models. Such interactions are easily introduced in the representative Models 1 and 2 discussed in Sec. 3 by simply adding a dark vector boson coupled strongly to the dark particle $\chi$.

Interestingly, a strongly self-interacting dark sector lies along the lines of the self-interacting dark matter paradigm, which was introduced two decades ago [24] to solve the core-cusp and missing satellite problem of the $\Lambda$CDM model.
A repulsive interaction between neutrons and dark particles can also be effective in allowing observed, $\lesssim 2M_\odot$, neutron star masses [25]. For appropriate strength and range of the interaction there is an energy cost associated with creating dark particles and therefore pure neutron matter is preferred.

4.3. Models with a self-interacting dark sector

A model of this type was constructed in Ref. [26], where a neutron dark decay involving a dark fermion and a dark photon in the final state was considered, i.e., $n \rightarrow \chi A'$. The effective Lagrangian is

$$\mathcal{L}^{\text{eff}} = \bar{n} \left( i D - m_n + \frac{g e}{2m_n} \sigma^{\mu\nu} F_{\mu\nu} \right) n + \bar{\chi} \left( i \not{D} - m_\chi \right) \chi + \varepsilon (\bar{n} \chi + \bar{\chi} n) - \frac{1}{2} F'_{\mu\nu} F'^{\mu\nu} - \frac{g}{2} F_{\mu\nu} F^{\mu\nu} - \frac{1}{2} m^2_{A'} A'_\mu A'^{\mu} ,$$  \hspace{1cm} (15)

where the covariant derivative $D_\mu = \partial_\mu - igA'_\mu$. It was shown that the strength of the dark photon coupling to the dark particle $\chi$, governed by the parameter $g'$ and resulting in repulsive interactions between the $\chi$ particles, can be chosen such that the neutron lifetime discrepancy is explained and, at the same time, all astrophysical bounds are satisfied, including constraints from neutron stars, galaxy clusters, cosmic microwave background, Big Bang nucleosynthesis and supernovae. If the dark particle $\chi$ in this model is stable, it can contribute to the dark matter in the universe, but cannot account for all of the dark matter.

Many of the astrophysical constraints are alleviated if one assumes non-thermal dark matter production. This was shown in Ref. [27], where a model for the neutron dark decay $n \rightarrow \chi \phi$ was constructed, based on our Model 2, but with a dark boson introduced to mediate large self-interactions of $\chi$. The Lagrangian for the dark sector is

$$\mathcal{L}_D = g \bar{\chi} Z_D \chi + (\lambda_\phi \bar{\chi} \chi \phi + \text{h.c.}) - i g Z^\mu_D \left( \phi^* \partial_\mu \phi - \phi \partial_\mu \phi^* \right).$$  \hspace{1cm} (16)

There exists a choice of parameters for which this model satisfies neutron star constraints, remains consistent with all other astrophysical bounds and $\chi$ makes up all of the dark matter in the universe. In addition, due to the self-interactions of $\chi$, the model is shown to solve the small-scale structure problems of the $\Lambda$CDM model.

4.4. Hadron dark decays

The idea of dark decays can be applied also to other neutral hadrons. In Ref. [28] it was argued that the mesons $K^0_L$ and $B^0$ can decay to dark sector particles at measurable rates. An explicit model was constructed with a dark sector consisting of several families of dark fermions. An analogous mechanism that prevents neutron beta decays in neutron stars, i.e., Pauli blocking, also forbids neutron dark decays inside a neutron star in this model.

4.5. Baryogenesis

It has recently been shown that the model addressing the neutron lifetime puzzle based on the Lagrangian in Eq. (15) provides a successful framework for low-scale baryogenesis [29]. In addition, a model very similar to our Model 2, with couplings of $\bar{\chi}$ to other quark flavors and a Majorana (instead of Dirac) fermion $\chi$, has been proposed in the context of low-scale baryogenesis as well [30].

4.6. Related solutions

Taking into consideration only the experimental data for $g_A$ from experiments performed after the year 2002, the bottle neutron lifetime is favored [31]. Based on this observation, explanations
of the neutron lifetime discrepancy have been put forward in which it is the bottle lifetime that is equivalent to the Standard Model prediction for $\tau_n$. The difference in outcomes of the bottle and beam measurements is explained via neutron-mirror neutron oscillations resonantly enhanced in large magnetic fields thus affecting only beam measurements \cite{32}, or by invoking a sizable Fierz interference term canceling the dark decay contribution to the neutron decay rate \cite{33}.

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