PHYSICAL REVIEW D **84,** 085008 (2011)

Neutrino physics with dark matter experiments and the signature of new baryonic neutral currents

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(Received 15 May 2011; published 11 October 2011)

New neutrino states ν_b, sterile under the standard model interactions, can be coupled to baryons via the isoscalar vector currents that are much stronger than the standard model weak interactions. If some fraction of solar neutrinos oscillate into ν_b on their way to Earth, the coherently enhanced elastic ν_b-nucleus scattering can generate a strong signal in the dark matter detectors. For the interaction strength a few hundred times stronger than the weak force, the elastic ν_b-nucleus scattering via new baryonic currents may account for the existing anomalies in the direct detection dark matter experiments at low recoil. We point out that for solar-neutrino energies, the baryon-current-induced inelastic scattering is suppressed, so that the possible enhancement of a new force is not in conflict with signals at dedicated neutrino detectors. We check this explicitly by calculating the ν_b-induced deuteron breakup, and the excitation of a 4.4 MeV γ line in ^{12}C. A stronger-than-weak force coupled to the baryonic current implies the existence of a new Abelian gauge group $\mathrm{U}(1)_B$ with a relatively light gauge boson.

DOI: 10.1103/PhysRevD.84.085008 PACS numbers: 95.35.+d, 12.60.Cn

I. INTRODUCTION

The standard model (SM) of particles and fields must be augmented to include neutrino mass physics and perhaps extended even further to account for the ‘‘missing mass’’ of the Universe, or cold dark matter (DM). During the last decade the underground experiments [1–8] aimed at direct detection of DM in the form of weakly interacting massive particles (WIMPs) [9] have made significant inroads into the WIMP-nucleon cross section vs WIMP mass parameter space. Since no DM-induced nuclear recoil signal was found with the exception of two hints to be discussed below, they constrained many models of dark matter and ruled out some portion of the parameter space in the best motivated cases such as e.g. supersymmetric neutralino DM [10], Higgs-portal singlet DM [11], etc.

It has been argued by some authors that, although primarily designed to search for WIMPs, these experiments are in fact multipurpose devices that can also be used for alternative signatures of other effects beyond the standard model. In particular, using the same instruments one can look for the absorption of keV-scale bosonic super-WIMPs [12], search for the axion emission from the Sun [13], and also investigate some additional signatures of WIMP-atom scattering that exist in ‘‘nonminimal’’ WIMP models [14]. This paper extends this point further and opens a new direction: We show that neutrino physics beyond the SM can also be probed with the dark matter experiments.

Elastic scattering of neutrinos on nuclei [15] is enhanced by the coherence factor N^2, where N is the number of neutrons. Straightforward analyses [16,17] of the SM solar-neutrino elastic scattering rates on nuclei used in DM experiments reveal several basic points:

(i) Despite the coherent enhancement, the scattering rates are way too small, leading to counting rates not exceeding 10^{-3} kg^{-1} day^{-1} keV^{-1}. Such low rates do not introduce any ν-background to WIMP searches at the current levels of sensitivity.

(ii) The nuclear recoil spectrum is usually very soft, $E_r \sim (E_\nu)^2/M_{\mathrm{Nucl}} \sim$ few KeV or less.

(iii) Solar boron (^{8}B) neutrinos are the best candidates for producing an observable signal, because of the compromise between the relatively large flux and the energy spectrum extending to 15 MeV.

Of course, at this point the DM experiments typically target much harder recoil and are far away from low counting rates induced by solar neutrinos. On the other hand, the last generation of dedicated solar-neutrino experiments [18–20] has been extremely successful in detecting solar neutrinos via charged current reactions (CC), elastic scattering on electrons (ES), and Z-exchange mediated (NC) deuteron breakup [19]. It is the combination of all three of these signals that leads to a very credible resolution of the long-standing solar-neutrino deficit problem via the neutrino oscillation and the Mikheev-Smirnov-Wolfenstein mechanism [21,22].

However, it is easy to imagine that three active SM species ν_e, ν_μ, ν_τ with their (almost completely) established mass/mixing parameters may not be the last word in the neutrino story. In this paper we consider a model of the ‘‘quasisterile’’ neutrino ν_b that has no charged currents with normal matter and no ES on electrons or other leptons, but that has much enhanced NC with baryons (NCB). We shall consider the strength of the new NCB interaction to be much larger than the Fermi constant, $G_B \sim (10^2–10^3) \times G_F$. Such interactions can be mediated by new vector bosons of the gauged baryon number $\mathrm{U}(1)_B$, and for that reason we call this new hypothetical neutrino state the "baryonic" neutrino ν_b. If any of the solar SM neutrino flavors oscillate into ν_b within 1 AU, then the current DM experiments will, in principle, be able to pick it up via the coherently enhanced NCB signal. Whether such strong NCB would lead to a measurable energy deposition in the standard neutrino experiments requires a special investigation and is addressed in this paper. We find that although the inelastic NCB scattering is enhanced by a huge factor G_B^2/G_F^2, it is also suppressed by a tiny factor $E_\nu^4 R_N^4$, where R_N is a nuclear radius-related parameter. The resulting rate for NCB processes in neutrino detectors can then be made comparable to or smaller than the regular neutrino counting rates.

It is somewhat tempting to relate the proposed ν_b model with the recently reported anomalies/signals in the direct DM detection. For a long time, of course, the DAMA experiment and its successor, DAMA/LIBRA, have been claiming [3] the annual modulation of the energy deposition in NaI crystals with the maximum in early June and minimum in December, which is consistent with the expected seasonal modulation of the WIMP-nucleus scattering rate. Then, last year the CoGeNT Collaboration [4] reported an unexpected (in the null hypothesis) rise of their signal at recoil energies below $E_r = 1$ keVee. Given the mass of Ge nuclei, and typical quenching factors in germanium, it is plausible that the rise of the CoGeNT signal at low keVee can be produced by ν_b resulting from oscillations of boron neutrinos, $\nu_{\rm SM} \to \nu_b$, and a hypothesized enhancement of NCB can compensate for small neutrino scattering rates. It is also clear that mimicking the DAMA signal with ν_b is possible in a more restricted sense as well. Of course, the usual seasonal modulation of the neutrino flux due to the eccentricity of the Earth's orbit will have a minimum in early July and a maximum in early January. However, the neutrino oscillation phenomenon is not monotonic in distance [23], and if the oscillation length for $\nu_{\rm SM} \to \nu_b$ is comparable to the Sun-Earth distance, the annual modulation phase of the ν_b scattering signal can be reversed by π. We investigate this opportunity, and conclude that both the CoGeNT and DAMA signals can be described with ν_b-type models (provided that DAMA data can tolerate a ~1 month phase shift). We further argue that if indeed this is the case, the model is very predictive, and there will be further ample opportunities for probing ν_b both at DM and neutrino detectors, as well as at more conventional particle physics experiments.

This paper is organized as follows: In the next section we introduce the class of ν_b models and specify a parameter range that is the most perspective for the DM experiments. Section III addresses NCB elastic and inelastic scattering, including calculations of ν_b-induced activation of carbon and the deuteron breakup reactions. Section IV studies the possibility of phase reversal in the annual modulation signal. Our conclusions are reached in Sec. V.

II. BARYONIC NEUTRINO AND BARYONIC NEUTRAL CURRENTS

The basic features of the model are as follows: We introduce a new gauge group $\mathrm{U}(1)_B$, and give all quark fields of the SM, Q, U, D, the same charge under $\mathrm{U}(1)_B$, which we call $g_b/3$. We also introduce a new left-handed neutrino species ν_b that has a charge g_l under this new group and no charge under any of the SM gauge groups. In the interest of anomaly cancellations it is also desirable to introduce a right-handed partner of ν_b with the same charge. Then the new group couples to the "vectorlike" matter multiplets, and although SM + $\mathrm{U}(1)_B$ will, in general, be anomalous, the anomalies can be canceled at some heavy scales. Variants of this model may include some partial gauging of the SM lepton species under $\mathrm{U}(1)_B$. Neither the right-handed ν_b nor the details of the anomaly cancellation will be important for this paper. Furthermore, we assume that $\mathrm{U}(1)_B$ is spontaneously broken by the $\mathrm{U}(1)_B$-Higgs vacuum expectation value $\langle \phi_b \rangle$, and exactly how this happens will not be of direct consequence for us either. The relevant gauge part of the Lagrangian is then given by

$$\mathcal{L} = -\frac{1}{4}V_{\mu\nu}^2 + \frac{1}{2}m_V^2 V_\mu^2 + \bar{\nu}_b\gamma_\mu(i\partial_\mu + g_l V_\mu)\nu_b + \sum_q \bar{q}\Big(i\not{D}_{\rm SM}/ + \frac{1}{3}g_b\gamma_\mu V_\mu\Big)q + \mathcal{L}_m. \tag{1}$$

The first two terms in (1) are the standard Maxwell-Proca terms for V_μ; the sum extends over all quark types and flavors, and $D_{\rm SM}$ is the SM covariant derivative that includes gauge interactions appropriate for each quark species q. The mass part of the Lagrangian $\mathcal{L}_m$, besides the usual SM mass terms, should also account for neutrino masses and generate mixing to a new state ν_b. In this paper we are not going to consider vector exchange with virtualities beyond $O(10\ \mathrm{MeV})$, and therefore it is convenient to switch from quarks to nucleons,

$$\frac{1}{3}V_\mu g_b \sum_q \bar{q}\,\gamma_\mu q \to g_b V_\mu(\bar{p}\gamma_\mu p + \bar{n}\gamma_\mu n) + \ldots. \tag{2}$$

The ellipsis stands for $O(m_N^{-1})$ terms associated with the $V_{\mu\nu}\bar{N}\sigma_{\mu\nu}N$ part of the form factor, which will be small for any process we consider in this paper. The coupling of V_μ to the isoscalar vector current of nucleons $J_\mu^{(0)} = \bar{p}\gamma_\mu p + \bar{n}\gamma_\mu n$ will have important implications for both the elastic and inelastic scattering of ν_b on nuclei. The exchange by the $\mathrm{U}(1)_B$ gauge boson creates the NCB Lagrangian

$$\mathcal{L}_{\rm NCB} = \bar{\nu}_b\gamma_\mu\nu_b\frac{g_l g_b}{m_V^2 + \Box}J_\mu^{(0)}, \tag{3}$$

which in the limit of $m_V^2 \gg Q^2$ is just a new contact, dimension-6 operator with the effective coupling constant G_B:

$$\mathcal{L}_{\rm NCB} = G_B \times \bar{\nu}_b \gamma_\mu \nu_b J_\mu^{(0)};$$
$$G_B = \frac{g_l g_b}{m_V^2} \equiv \mathcal{N} \times \frac{10^{-5}}{\rm GeV^2}. \tag{4}$$

Here we have introduced an "enhancement" parameter $\mathcal{N}$ that quantifies how much stronger G_B is compared to the weak-scale value of 10^{-5} GeV^{-2}. We note that stronger-than-weak interactions among four neutrino species were considered earlier in e.g. Ref. [24]. The use of the baryonic force as a mediator between the SM and dark matter was considered recently in [25].

One may wonder if $\mathcal{N}$ as large as 100 or 1000 can be consistent with low-energy data on meson decays, such as $K \to \pi \bar{\nu}_b \nu_b$. It turns out that, due to the conservation of the baryon current, the loop amplitude for the underlying $s \to d\bar{\nu}_b \nu_b$ decay is additionally suppressed by $G_F Q^2$, which compensates for all possible enhancements due to $\mathcal{N}$. [In contrast, the quark axial-vector analogue of (1) will be strongly constrained to have $\mathcal{N} \lesssim 1$.] Thus, from the quark flavor perspective, the baryonic portal (1) is one of the two "safe" portals (the other being the kinetic mixing with hypercharge [26]) that allow attaching stronger-than-weak interactions to the quark currents. We shall not pursue the meson decay constraints on the model any further in this paper, and we turn our attention to the neutrino mass sector.

The most natural way of having a UV-complete theory of neutrino masses is via the introduction of right-handed neutrino states N_R. We can use the same singlet right-handed neutrinos coupled to the Higgs-lepton bilinears LH and Higgs$_b$-neutrino ν_b bilinears $\nu_b \phi_b$ in a gauge-invariant way,

$$\mathcal{L}_m = LH\mathbf{Y}N + \nu_{bL}\phi \mathbf{b} N + (\text{H.c.}) + \tfrac{1}{2} N^T \mathbf{M}_R N. \tag{5}$$

Here $\mathbf{M}_R$ and $\mathbf{Y}$ are the familiar 3×3 right-handed neutrino mass matrix and Yukawa matrix, while $\mathbf{b}$ is the new Yukawa vector parametrizing the couplings of the left-handed part of ν_b to N. Integrating out N states results in the low-energy 4×4 neutrino mass and mixing matrices, M_{ij}, where i, j run over e, μ, τ, b flavors. While of course a full four-state analysis can be done, we shall simplify our discussion by the following assumptions:

(1) The entries of the 3×3 submatrix $M_{\text{active,active}}$ will, in general, be somewhat larger than the $M_{\text{active},b}$ and $M_{b,b}$ components so that the mixing pattern can be addressed sequentially: first the mixing of the SM neutrinos and then the admixture of the ν_b.

(2) A tribimaximal ansatz will be taken for the 3×3 mixing of the SM neutrino species for simplification, although having $\theta_{13} = 0$ is not crucial.

(3) We shall assume a preferential mixing of ν_b to ν_2, with the relevant parameters that we call Δm_b^2 and θ_b, so that the true mass eigenstates are $\nu_I = \cos\theta_b \nu_2 + \sin\theta_b \nu_b$, $\nu_{II} = -\sin\theta_b \nu_2 + \cos\theta_b \nu_b$.

(4) The sign of G_B will be chosen to ensure that the matter effects for ν_b will not lead to the matter-induced $\nu_{\text{active}} \to \nu_b$ transitions.

The combination of these assumptions forms the following (simplified) picture of neutrino oscillations: Inside the Sun the neutrino oscillations occur largely between ν_e and $\nu_+ \equiv (\nu_\mu + \nu_\tau)/\sqrt{2}$,

$$\begin{pmatrix} \nu_e \\ \nu_+ \end{pmatrix} \simeq \begin{pmatrix} \sqrt{\frac{2}{3}} & \sqrt{\frac{1}{3}} \\ -\sqrt{\frac{1}{3}} & \sqrt{\frac{2}{3}} \end{pmatrix} \begin{pmatrix} \nu_1 \\ \nu_2 \end{pmatrix}, \tag{6}$$

while the "$-$" combination and ν_b stay unexcited. We would need only the higher end of the Boron neutrino spectrum, where the Mikheev-Smirnov-Wolfenstein effect dominates. Upon the neutrino exit from the dense region of the Sun, it represents an almost pure ν_2 state, $\nu_2 = \sqrt{\frac{1}{3}}\nu_e + \sqrt{\frac{2}{3}}\nu_+$, with individual flavor probabilities

$$P_e(\text{Sun}) \simeq \tfrac{1}{3}; \qquad P_+(\text{Sun}) \simeq \tfrac{2}{3}; \qquad P_b(\text{Sun}) = 0. \tag{7}$$

Then vacuum oscillations start building a nonzero probability for ν_b due to ν_I and ν_{II} being the true mass eigenstates in vacuum. Upon the arrival at the Earth, the following energy-dependent probabilities will approximate the neutrino flavor composition:

$$\begin{aligned} P_b(\text{Earth}) &\simeq \sin^2(2\theta_b)\sin^2\left[\frac{\Delta m_b^2 L(t)}{4E}\right], \\ P_e(\text{Earth}) &\simeq \frac{1}{3}\left(1 - \sin^2(2\theta_b)\sin^2\left[\frac{\Delta m_b^2 L(t)}{4E}\right]\right), \\ P_+(\text{Earth}) &\simeq \frac{2}{3}\left(1 - \sin^2(2\theta_b)\sin^2\left[\frac{\Delta m_b^2 L(t)}{4E}\right]\right), \end{aligned} \tag{8}$$

where $L(t)$ is the Earth-Sun distance with a slight eccentricity modulation,

$$L(t) \simeq L_0\left(1 - \epsilon \cos\left[\frac{2\pi(t - t_0)}{T}\right]\right); \quad L_0 = 1.5 \times 10^8\ \text{km};$$
$$\epsilon \simeq 0.0167; \quad t_0 \simeq 3\ \text{Jan}. \tag{9}$$

The most interesting range of Δm_b^2 to consider is

$$10^{-10}\ \text{eV}^2 \lesssim \Delta m_b^2 \ll \Delta m^2_{\text{Solar,atm}}. \tag{10}$$

A scale of $O(10^{-10})$ eV^2 is the so-called "just so" mass splitting that may introduce significant changes to the otherwise very predictable $\propto L^{-2}$ seasonal variations of the ν_b flux at the Earth's location. With Δm_b^2 being much smaller than 10^{-5} eV^2, there is no danger of distorting KamLAND results [23] even for a relatively large angle θ_b, although the matter effects for $\bar{\nu}_b$ could be significant. We also find it convenient to introduce the energy parameter E_0 directly related to the mass splitting,

$$E_0 = \frac{\Delta m_b^2 L_0}{4\pi} = 6.05\ \text{MeV} \times \frac{\Delta m_b^2}{10^{-10}\ \text{eV}^2}, \tag{11}$$

which defines the last zero of P_b as a function of energy, $P_b(L = L_0, E = E_0) = 0$. Since in all NCB rates P_b will enter in the combination with G_B^2, it is also convenient to define

$$\mathcal{N}_{\rm eff}^2 = \mathcal{N}^2 \times \tfrac{1}{2} \times \sin^2(2\theta_b), \tag{12}$$

so that in the limit of large E_0 the oscillations average out and $P_b G_B^2 \to \mathcal{N}_{\rm eff}^2 \times 10^{-10}\ \mathrm{GeV}^{-4}$.

The pattern of masses and mixing considered here is not the most natural: We assume a pair of very degenerate ν_I and ν_{II} mass eigenstates replacing ν_2 and ν_b. Given that the mass of ν_2, regardless of the hierarchy pattern, is always in between ν_1 and ν_3, this will require some specific adjustments of the full 4×4 mass matrix. The search for more natural realizations of $\nu_{\rm SM} \to \nu_b$ oscillations with long oscillation lengths, including matter effects for a different sign of G_B, goes beyond the scope of this paper. The goal of the next two sections will be to find the sensitivity to $\mathcal{N}_{\rm eff}$ in various processes involving the elastic and inelastic scattering of ν_b.

We would like to close this section with some model-building comments. A very intriguing question to ask is whether SM neutrinos would tolerate new large NCB. A conventional answer is "no," as the so-called nonstandard neutrino interactions (NSI) with quarks and charged leptons were addressed in a number of papers [27], and almost no room at the $O(1)G_F$ level was found, let alone the much enhanced NCBs hypothesized in this paper. However, NSI studies [27], with rare exceptions [28], assume that the scale of the mediation is comparable to the weak scale, and ignore the possibility of light vector bosons communicating between neutrinos and baryons. As a counterexample, one could consider a model with two new gauge groups, $\mathrm{U}(1)_B$ and a quantized lepton flavor, e.g. L_μ or L_τ. The connection between two new vector sectors is given by the kinetic mixing term $\eta V_{\mu\nu}^{(1)} V_{\mu\nu}^{(2)}$. Then, an additional effective interaction of a SM neutrino with the baryonic current is given by

$$\mathcal{L}_{\rm eff} \propto \bar{\nu}_{\rm SM} \gamma_\alpha \nu_{\rm SM} \frac{\eta g_l g_b \Box}{(m_{V1}^2 + \Box)(m_{V2}^2 + \Box)} J_\alpha^{(0)}. \tag{13}$$

Such an interaction gives no contribution to the forward scattering amplitude and thus does not affect neutrino oscillation, and it is $1/Q^2$ suppressed in the large Q^2 regime, avoiding strong constraints from deep-inelastic neutrino scattering. It is then clear that the choice of m_{V1}, m_{V2} in the MeV range may allow having (13) at $Q^2 \sim (1\text{–}10)\ \mathrm{MeV}^2$ to be considerably stronger than the SM weak force.[1] The interactions of type (13) can lead to the detectable recoil signal from elastic scattering of solar SM neutrinos, along the same lines as the ν_b-scattering idea advocated in this paper. The possibility of modified SM neutrino interactions such as (13) can be very effectively tested using the proposed neutrino-nucleus elastic scattering detectors placed near the intense source of stopped pions [29].

[1]The author would like to acknowledge very stimulating discussions with B. Batell and I. Yavin on the possibility of the NSI enhancement.

III. ELASTIC AND INELASTIC SCATTERING OF ν_b

Elastic scattering of ν_b on nuclei will create a recoil signal regulated by the strength of NCB and the probability of oscillation (8). It can be picked up by the direct dark matter detection experiments with low recoil thresholds. Also, ν_b neutrinos can deposit a significant amount of energy on the order of a few MeV by activating excited nuclear states or via extra neutrons created by nuclear breakup. The main finding of this section can be summarized as follows: The ratio of the elastic to inelastic cross sections in the interesting neutrino energy range $E_\nu \lesssim 15$ MeV is governed by the following relation:

$$\frac{\sigma_{\nu_b-\mathrm{Nucl}}(\text{elastic})}{\sigma_{\nu_b-\mathrm{Nucl}}(\text{inelastic})} \sim \frac{A^2}{E_\nu^4 R_N^4} \sim 10^8, \tag{14}$$

where we took $A \sim 100$, $R_N^{-1} \sim 100$ MeV, and $E_\nu \sim 10$ MeV. It is this huge ratio that makes small-scale experiments such as in [4] competitive in sensitivity to ν_b with the large-scale neutrino detectors.

A. Elastic scattering

The differential cross section for the NCB elastic scattering of left-handed ν_b on a nucleus of mass M_N with A nucleons is given by

$$\frac{d\sigma_{\rm el}}{d(\cos\theta)} = \frac{E^2 A^2 g_b^2 g_l^2 (1+\cos\theta)}{4\pi (M_V^2 + \mathbf{q}^2)^2} \simeq \frac{1}{4\pi} \times G_B^2 E^2 A^2 (1 + \cos\theta), \tag{15}$$

where the elastic scattering momentum transfer is $q = (\mathbf{q}^2)^{-1/2} = 2E\sin(\theta/2)$ and it cannot exceed twice the neutrino energy E. In the second relation we took $M_V \gg E$, which allows us to shrink the vector propagator. Using relations between the neutrino scattering angle θ, nuclear recoil energy E_r, and the minimum neutrino energy required to produce the E_r-recoiling nucleus,

$$E_r = \frac{E^2}{m_N} \times (1 - \cos\theta); \qquad E^{\rm min} = \sqrt{\frac{E_r M_N}{2}}, \tag{16}$$

we rewrite the elastic cross section (15) in the following form:

$$\frac{d\sigma_{\rm el}}{dE_r} = \frac{1}{2\pi} \times G_B^2 A^2 m_N \left(1 - \frac{(E^{\rm min})^2}{E^2}\right). \tag{17}$$

One can readily see that the NCB cross section (17) is related to the SM elastic neutrino-nucleus cross section

[15] by $G_B^2 A^2 \to G_F^2 (N/2)^2$ substitution, where N is the number of neutrons (with small corrections in the $1 - 4\sin^2\theta_W$ parameter). For the momentum transfers and nuclei considered in this paper, the form factor corrections are <5% percent and are ignored.

Using cross section (17), the flux and the energy distribution of ^{8}B neutrinos [30] (*hep* solar neutrinos provide a small correction), we derive the counting rates as a function of interaction strength and the oscillation probability. For the moment, we neglect small seasonal modulations and take the limit of $\epsilon \to 0$. For the medium composed only of atoms with atomic number A, we approximate these rates by

$$\frac{dR}{dE_r} \simeq \frac{A^2 m_N}{2\pi} \times \frac{1}{2}\sin^2(2\theta_b) G_B^2 \Phi_{^8\mathrm{B}} \times I(E_r, E_0)$$
$$\to 85 \frac{\text{recoils}}{\text{day} \times \text{kg} \times \text{KeV}} \times \left(\frac{A}{70}\right)^2 \times \frac{\mathcal{N}_{\text{eff}}^2}{10^4} \times I(E_r, E_0). \tag{18}$$

The input (total flavor) flux of ^{8}B neutrinos is taken to be $\Phi_{^8\mathrm{B}} = 5.7 \times 10^6\ \mathrm{cm}^{-2}\,\mathrm{s}^{-1}$ and $m_N \propto A m_p$.

The recoil integral $I(E_r, E_0)$ in Eq. (18) is given by the convolution of the ^{8}B energy distribution, the energy-dependent part of the oscillation probability, and the kinematic factor in the cross section reflecting neutrino helicity conservation:

$$I(E_r, E_0) = \int_{E^{\min}(E_r)}^{\infty} dE \left(1 - \frac{(E^{\min})^2}{E^2}\right) \times f_{^8\mathrm{B}}(E) \times 2\sin^2\left[\frac{\pi E_0}{E}\right]. \tag{19}$$

Here the distribution function is normalized as $\int_{\text{all } E} f_{^8\mathrm{B}}(E) dE = 1$. For the limit of large E_0 (fast oscillations), the last multiplier in (19) becomes 1. If a detector threshold corresponds to recoil energies that require $E^{\min}$ to be above the endpoint of the ^{8}B neutrino spectrum, $I \equiv 0$ (apart from small corrections from *hep* and diffuse supernova neutrinos). This is the case for most of the existing WIMP detectors, but not for all of them.

In real detectors registering ionization such as in [3,4], it is the electron equivalent of the energy release rather than the recoil energy that is detected. We take the relation between the two by following recent DM-related analyses [31,32],

$$\text{Ge: } E_r(\text{keVee}) \simeq 0.2 \times (E_r(\text{keV}))^{1.12},$$
$$\text{Na in NaI: } E_r(\text{keVee}) \approx 0.33 \times E_r(\text{keV}). \tag{20}$$

The second relation is far less precise than the first one [32].

The counting rates in germanium resulting from scattering of ν_b created by the oscillations of ^{8}B and *hep* solar neutrinos are presented in Fig. 1. We have taken three cases of mass splitting: $E_0 \gg E_\nu^{\max}$, and $E_0 = 12$, 14 MeV.

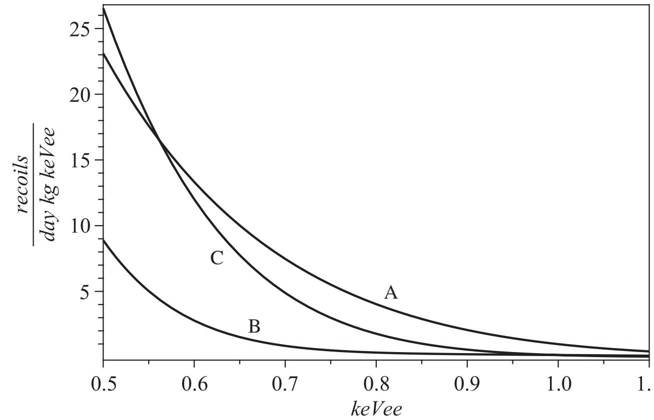

FIG. 1. Expected recoil event rate in Germanium in units of recoils/day/kg/keVee as a function of E_r in keVee. The NCB enhancement factor $\mathcal{N}_{\text{eff}} = 100$. The lines A, B, and C correspond to $E_0 = \infty$, $E_0 = 12$ MeV, and $E_0 = 14$ MeV.

The NCB rates are plotted for the value of $\mathcal{N}_{\text{eff}}^2 = 10^4$. For this enhancement factor, the resulting counting rates are clearly within reach of the current generation of low-threshold germanium detectors (i.e. CoGeNT).

Inspection of Fig. 1 shows that the choice of different mass splittings that make the oscillation length comparable to 1 A.U. influences the shape of the recoil spectrum. This is because the most important part of the spectrum for the recoil in excess of 0.5 keVee is above neutrino energies of 10 MeV, where the ^{8}B neutrino spectrum is already sharply falling. The neutrino oscillations with E_0 close to 12 MeV will lead to the suppression of higher E_r and to the steep rise of the signal at lower E_r. The sharp end of the neutrino spectrum prevents other Ge experiments with higher thresholds, like CDMS [1], to probe the NCB scattering in the regime of large recoil where CDMS [1] has strong sensitivity. The signal from the recoil due to ν_b neutrinos is very similar in morphology to that of sub-10 GeV-scale WIMPs. This is because a typical momentum transfer in a heavy nucleus–light WIMP collision is $q \sim m_{\text{wimp}} v \sim 10$ MeV, which is about the same for ^{8}B neutrino scattering. There is one kinematic difference though: The backscattering of WIMPs that produces the hardest recoil is kinematically allowed, while for neutrinos it is forbidden by helicity conservation. This additionally limits the capabilities of high-threshold experiments to detect ν_b neutrinos in comparison with light WIMPs.

Is it possible to use ν_b as another speculative explanation of CoGeNT results [4]? The overall event rate can indeed be reproduced well with $\mathcal{N}_{\text{eff}} \sim 10^2 - \text{few} \times 10^2$, depending on E_0. For the large E_0 parameter, the enhancement factor of $\mathcal{N}_{\text{eff}} = 10^2$ seems sufficient: It gives 7 recoils/day/keVee at $E_r = 0.7$ keVee, which is about the same as the experimental data suggest after accounting for the efficiency [4,32]. The shape of the predicted signal is also similar to the counting rate profiles observed by

CDMS at the Stanford underground facility [33]. Fitting the exact spectral shape of excess events at CoGeNT falls outside the scope of our current investigation. We should also note that the expected total counting rate for the material used in the CRESST experiment [5] due to neutrino-oxygen scattering is given by

$$R_{\mathrm{O\,in\,CaWO_4}}(E_r > 10\,\mathrm{keV}) \simeq 0.2 \times \frac{\mathrm{recoils}}{\mathrm{day}\times\mathrm{kg}} \times \frac{\mathcal{N}_{\mathrm{eff}}^2}{10^4}, \quad (21)$$

which is well within their detection capabilities for $\mathcal{N}_{\mathrm{eff}} \sim 100$. Other methods in development that use liquid helium as a detecting medium with a potentially very low-energy threshold [34] also look promising for detecting ν_b-induced recoil. It is also important that the choice of a very low-mass target such as ^{4}He will allow discriminating between $\gtrsim 5$ GeV WIMPs and ν_b's: The effective recoil energy decreases at $M_N < M_{\mathrm{WIMP}}$, while it becomes larger for ν_b scattering.

B. Inelastic scattering

Unlike light WIMPs that can carry significant momentum but very little energy, ν_b can easily lead to an MeV-scale energy deposition. Here we turn our attention to the NCB inelastic processes and will address the following issues: the NCB deuteron breakup, and the NCB excitation of the first 2^+ resonance in ^{12}C resulting in a 4.4 MeV γ line:

$$d + \nu_b \to \nu_b + n + p, \quad (22)$$

$$^{12}\mathrm{C} + \nu_b \to \nu_b + {}^{12}\mathrm{C}^*(4.44\ \mathrm{MeV}) \to \nu_b + {}^{12}\mathrm{C} + \gamma. \quad (23)$$

The main scientific question to answer is whether the enhanced values of $G_B^2 P_b$ can be consistent with the constraints provided by SNO on "extra neutrons" from (22) and by Borexino and other liquid scintillator detectors on "extra gammas" from (23). There are of course other processes that one has to consider in a more comprehensive study, including the excitation of oxygen, the breakup of ^{13}C to ^{12}C + n, etc., but they will all follow the scaling in Eq. (14). The earlier studies of the nuclear excitations due to the different types of neutrino couplings can be found in [35]. There are also elastic channels of energy deposition via $\nu_b + p \to \nu_b + p$ [36], but the proton recoil from the scattering of ^{8}B neutrinos would fall below the detector thresholds.

To understand the origin of the ratio (14), one does not have to perform any sophisticated calculations. We consider the scattering of $\sim$10 MeV energy neutrinos, so that their wavelengths are much larger than the characteristic nuclear size of a few fm. Therefore, one can safely expand the nuclear matrix elements in series in q, or in neutrino energy E, as q is bounded by E. Here is how the inelastic matrix element of the $\mu = 0$ component of the isoscalar vector current $J_\mu^{(0)}$ between the deuteron bound state and np continuum will look in this expansion:

$$\begin{aligned}\langle d|\exp(i\mathbf{q}\mathbf{r}^{(n)}) + \exp(i\mathbf{q}\mathbf{r}^{(p)})|np\rangle &= 2\langle d|np\rangle + i\mathbf{q}\cdot\langle d|\mathbf{r}^{(n)} + \mathbf{r}^{(p)}|np\rangle \\ &\quad - \frac{q_k q_l}{2}\langle d|r_k^{(n)} r_l^{(n)} + r_k^{(p)} r_l^{(p)}|np\rangle \\ &= -\frac{q_k q_l}{4}\langle d|r_k r_l|np\rangle, \end{aligned} \quad (24)$$

where $\mathbf{r}^{(n)}$, $\mathbf{r}^{(p)}$ are the position operators for the neutron and the proton. The zeroth and first order terms in $q = |\mathbf{q}|$ are trivially zero due to the orthogonality of the wave functions ($\mathbf{r}^{(n)} + \mathbf{r}^{(p)}$ is the center-of-mass operator and cannot mediate inelastic transitions). In the last line we have introduced the relative position vector $\mathbf{r} = \mathbf{r}^{(n)} - \mathbf{r}^{(p)}$, and the quadratic in r operator can be further separated into the isotropic "charge-radius" and quadrupole components. For the $0^+ \to 2^+$ transition in ^{12}C, only the quadrupole part will matter. It is of course instructive to revisit the SM deuteron breakup [37] and observe that the isoscalar vector component of the standard weak current gives a very minor contribution to the total cross section at low E_ν due to this q^2 suppression of the amplitude. The SM rate is of course dominated by the isovector axial-vector current that corresponds to the difference of nucleon spins $\mathbf{s}^{(n)} - \mathbf{s}^{(p)}$, an operator that has nonzero inelastic matrix elements even in the $O(q^0)$ order.

We perform the calculation of the NCB-induced deuteron breakup using the "zero-radius" approximation of the initial and final state wave functions,

$$\psi_{\mathrm{in}}(\mathbf{r}) = \frac{\sqrt{2\kappa}}{\sqrt{4\pi}r}\exp(-\kappa r), \qquad \psi_f(\mathbf{r}) = \exp(i\mathbf{p}\mathbf{r}), \quad (25)$$

$$\kappa = \sqrt{2E_b\mu} = 45\ \mathrm{MeV}, \qquad p^2 = \mathbf{p}^2 = 2\mu(E - E_f - E_b). \quad (26)$$

Here E and E_f are the initial and final energy of ν_b, $E_b = 2.2$ MeV is the absolute value of the deuteron binding energy, and $\mu \simeq (m_n + m_p)/4$ is the reduced mass of the proton-neutron two-body system. Relative momentum p of the final state is fully determined from the neutrino energies, as the recoil of the deuteron center-of-mass is negligible. Parameter κ is the familiar "bound state momentum," related to the inverse size of the deuteron, $\kappa \sim R_d^{-1}$, and its relative smallness reflects the large spatial extent of the deuteron. In a language that is slightly excessive for the problem at hand, our calculations correspond to the leading order of the pionless effective field theory [37,38]. They can be systematically improved if needed, or treated with the more elaborate nuclear physics tools (see e.g. [39]). Of course, none of this will change the order of q in which the effect first occurs.

Straightforward calculations give the differential over the final neutrino energy cross section:

$$\frac{d\sigma_{d\to np}}{dE_f} = \frac{G_B^2 E_f^2 m_p}{8\pi^2} \frac{\kappa^5 p}{(p^2+\kappa^2)^6}\left[E^2E_f^2 + \frac{12p^4}{5\kappa^4} \times \left(E^4 - \frac{2}{3}E^3E_f + \frac{10}{9}E^2E_f^2 - \frac{2}{3}EE_f^3 + E_f^4\right)\right]. \quad (27)$$

The result for $d\sigma_{d\to np}$ shows an $O(E^4\kappa^{-4})$ suppression in agreement with $O(q^2)$ of the deuteron matrix element and in agreement with (14). Judging by the size of the subleading corrections in the SM calculations [37], we expect this answer to hold within $\sim 20\%$ accuracy.

The final integral over E_f in the interval from 0 to $E - E_b$ gives the total NCB deuteron breakup cross section. In Fig. 2, upper panel, we plot $\sigma_{d\to np}$ for the standard model neutrino and for ν_b with the choice of enhancement factor $\mathcal{N}^2 = 10^4$. As expected, the NCB cross section has a faster rise with neutrino energy due to the quadrupolar nature of the NCB interaction. In Fig. 2, lower panel, we also show the convolution of the cross section with the energy distribution of boron neutrinos times the energy-dependent part of the oscillation probability, $2\sin^2(\pi E_0/E)$. As in the previous subsection, the NCB rates are considered for large E_0 and for $E_0 = 12$, 14 MeV, while the enhancement factor is kept at $\mathcal{N}_{\rm eff} = 100$. The areas under curves give the total effective cross sections and, when multiplied by $\Phi_{^8\rm B}$, correspond to the breakup rate per deuteron. For the three NCB cases considered here, we have the following comparison to the SM rate:

$$\frac{\sigma_{\rm NCB}}{\sigma_{\rm SM\,NC}} \simeq \frac{\mathcal{N}_{\rm eff}^2}{10^4} \times (0.14, 0.06, 0.13) \quad \text{at } E_0 = \infty, \quad 12\text{ MeV}, \quad 14\text{ MeV}. \quad (28)$$

A 15% increase in the neutron production rate at SNO can be tolerated, and if one chooses a sizable θ_b so that the active neutrino flux is slightly less than that predicted by SM + SSM, the total neutral current rate may not even change. We conclude that SNO NC events leave enough room for the possible $\mathcal{N}_{\rm eff}^2 \sim O(10^4)$ (and slightly higher) enhancement of the NCB rate.

We now address the $^{12}\rm C \to {}^{12}C^*$ activation due to ν_b. To avoid the complications arising from nuclear physics, we shall assume that both the ground state and the first excited state of ^{12}C are given by 3α configurations. This is a very well-justified assumption, which leads to a relation between the matrix elements of the baryonic current and electric current,

$$\langle 0^+|J_i^{(0)}|2^+\rangle = 2\langle 0^+|J_i^{\rm em}|2^+\rangle = 2\frac{(E_{2^+} - E_{0^+})q_j}{6}\langle 0^+|Q_{ij}^{\rm em}|2^+\rangle. \quad (29)$$

Only the lowest order in q terms are retained here, and the $\mu = 0$ component can be restored from gauge invariance. The factor of 2 in (29) comes from the fact that the baryonic charge of α particles is twice as large as their electric charge. The information on the value of the transitional quadrupole moment $\langle Q_{ij}\rangle$ can be extracted from the $^{12}\rm C^*$ decay width $\Gamma = 1.08 \times 10^{-2}$ eV:

$$\overline{|\langle Q_{ij}^{\rm em}\rangle|^2} = \frac{90\Gamma}{\alpha(\Delta E)^5} = (3.3\text{ fm})^4, \quad (30)$$

where the value of the quadrupole moment squared is averaged over the arbitrary projection of the 2^+ angular momentum, and $\Delta E = E_{2^+} - E_{0^+} = E - E_f = 4.439$ MeV. Note that we define electric quadrupole and electric current *without* the "e," and account for α explicitly in (30). The total inelastic cross section for the ν_b-induced $^{12}\rm C \to {}^{12}C^*$ transition is given by

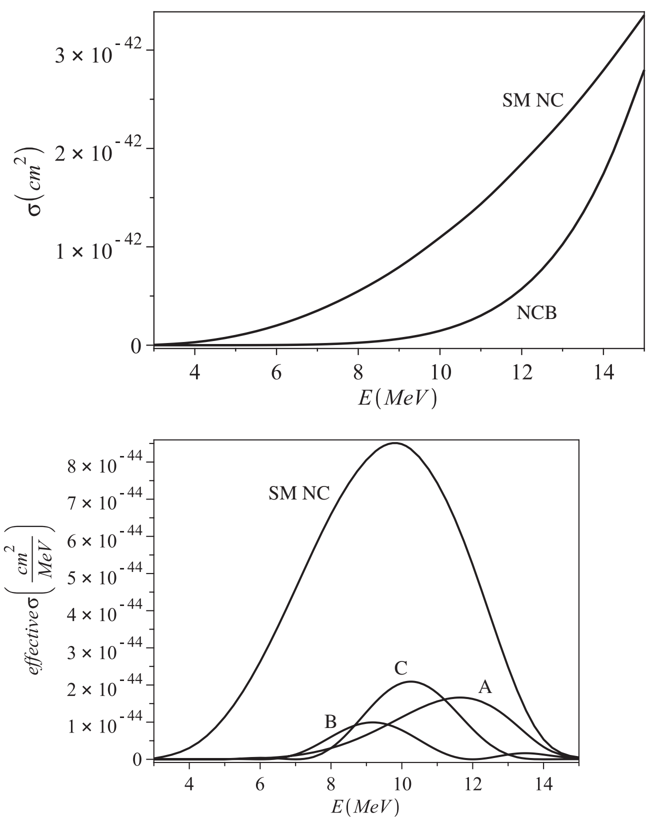

FIG. 2. Top panel: Deuteron breakup cross section for the SM NC processes (top curve) and for the NCB ν_b neutrinos (bottom curve). The NCB cross section is plotted for the $\mathcal{N}^2$ enhancement factor of 10^4. Bottom panel: The same cross sections convoluted with a ^{8}B energy distribution and the energy-dependent part of the oscillation probability. The top curve is the SM NC distribution of the effective cross section, and the lines A, B, and C are the same for the NCB with $E_0 = \infty$, 12, 14 MeV and $\mathcal{N}_{\rm eff}^2 = 10^4$. The areas under the curves give the proportion of neutrons produced via SM NC and NCB processes.

$$\sigma_{^{12}\mathrm{C}\to^{12}\mathrm{C}^*} = \frac{8G_B^2E^5(E-\Delta E)\overline{|\langle Q_{ij}^{\mathrm{em}}\rangle|^2}}{81\pi} \times \left[1 - 3x + \frac{39}{8}x^2 - \frac{19}{4}x^3 + \frac{39}{16}x^4 - \frac{9}{16}x^5\right], \tag{31}$$

where $x = \Delta E/E$. The benchmark value for this cross section at $E = 8$ MeV and $\mathcal{N} = 1$ is 2.5×10^{-48} cm^2.

With this cross section the effective rate of injection of 4.4 MeV gamma quanta in pseudocumene (scintillating material used by the Borexino experiment) is estimated to be

$$R(4.4\ \mathrm{MeV}) \sim (0.05\text{–}0.15) \times \frac{\gamma\ \mathrm{injections}}{100\ \mathrm{tons} \times \mathrm{day}} \times \frac{\mathcal{N}_{\mathrm{eff}}^2}{10^4}. \tag{32}$$

This is not a large rate by any measure, but it is nevertheless comparable to the counting rates in the 3–5 MeV window from ^{8}B ES processes and from the ^{208}Tl background events [40]. The actual counting rate should be obtained by applying to (32) the efficiency factor that the collaboration can extract from their calibration data and simulations. At this point we can only conclude that there must be some sensitivity to NCB at the $\mathcal{N}_{\mathrm{eff}} \sim 10^2$ level at the large-scale neutrino detectors that use carbon-based scintillators. More definitive statements and perhaps stronger sensitivity to ν_b can be derived from dedicated analyses. Moreover, the search for the extra γ lines in a different energy range was already performed by the Borexino Collaboration in connection with hypothetical Pauli-forbidden ^{12}C decays [41]. The search for NCB would represent a far less exotic physics goal in our opinion. One could also conduct similar searches of a ν_b-induced excitation of ^{16}O nuclei using SNO and SuperK data.

IV. ANNUAL MODULATION OF ν_b RATES

In this section we would like to address the question of seasonal modulation of the NCB rate. The seasonal modulation of the solar-neutrino rate was observed by the SNO and SuperK collaborations [18,42]. It exhibits full agreement with the expected $\propto L^{-2}(t)$, 3.3% modulation of the neutrino flux, with an appropriate minimum in the summer (northern hemisphere). The hypothetical NCB elastic scattering rate will have the same modulation pattern as long as Δm_b^2 is large or, in other words, at $E_0 \gg E_{\nu\,\mathrm{solar}}$. In the opposite limit of $E_0 \ll E_{\nu\,\mathrm{solar}}$ the modulation effects are suppressed because $\Phi_\nu \propto L_0^2L^{-2}(t)\sin^2[\pi E_0 L(t)(L_0E)^{-1}] \to (\pi E_0/E)^2$, which is time independent.

However, it is easy to imagine that the flux of ν_b neutrinos can have a more intricate seasonal modulation pattern. For example, if E_0 is between the maximum of the ^{8}B neutrino spectrum and its endpoint, the high-energy fraction of the distribution will have a higher flux in the summer. This is best illustrated in Fig. 3, where the expected flux of ν_b resulting from oscillations of boron neutrinos is convoluted with the time-dependent part of P_b, $L_0^2L^{-2}(t)\sin^2[\pi E_0L(t)(L_0E)^{-1}]$ at $E_0 = 12$ MeV. The two curves correspond to $t = t_{\mathrm{perihelion}}$ (~ 3 Jan) and t_{aphelion} (~ 4 July). Although, on average, there are more ν_b neutrinos arriving at the Earth in January, in the most relevant range of energies, $E > 10$ MeV, the flux in July is larger. Therefore, for this fraction of neutrinos there is a phase reversal, and the elastic scattering rates will reflect that.

Next we calculate the expected seasonal modulation in the counting rate,

$$\frac{dR_{\mathrm{mod}}}{dE_r} = \frac{1}{2}\left(\frac{dR}{dE_r}\bigg|_{\mathrm{Jul}} - \frac{dR}{dE_r}\bigg|_{\mathrm{Jan}}\right), \tag{33}$$

for NaI detectors using the quenching factor from Eq. (20). We would like to remark in passing that for some ranges of neutrino energies there can be a significant departure from a simple time-sinusoidal function, but to observe such effects one would probably require very high statistics and very good energy resolution. Modulation rates, dR_{mod}/dE_r, as defined in Eq. (33) are plotted in Fig. 4 for the same three choices of E_0 and $\mathcal{N}_{\mathrm{eff}} = 100$ as before. One can see that, indeed, modulation of both signs is possible, and that the rate of the modulated NCB signal at this $\mathcal{N}_{\mathrm{eff}}$ is indeed probed by the DAMA/LIBRA experiment [3], which is sensitive to modulation amplitudes $O(10^{-2})$ cpd/kg/keVee.

Is it possible that ν_b-Na elastic scattering is behind the DAMA/LIBRA seasonal modulation anomaly? The magnitude of the predicted modulation can be in very good agreement with DAMA results [3]. Moreover, as we saw in the previous section, $\mathcal{N}_{\mathrm{eff}}^2 \sim 10^4$ is thus far consistent with

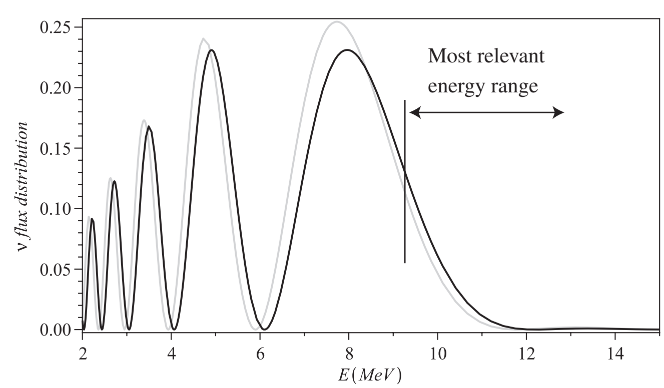

FIG. 3. Boron ν_b neutrino flux modified by the time-dependent part of the oscillation probability $2L_0^2L^{-2}(t)\sin^2[\pi E_0L(t) \times (L_0E)^{-1}]$ with $E_0 = 12$ MeV. The black curve is for July, and the gray curve is for January. Although the total integral under the gray curve is bigger than under the black one, it is the high end of the spectrum that would determine rates at the existing DM detectors, where July rates are larger.

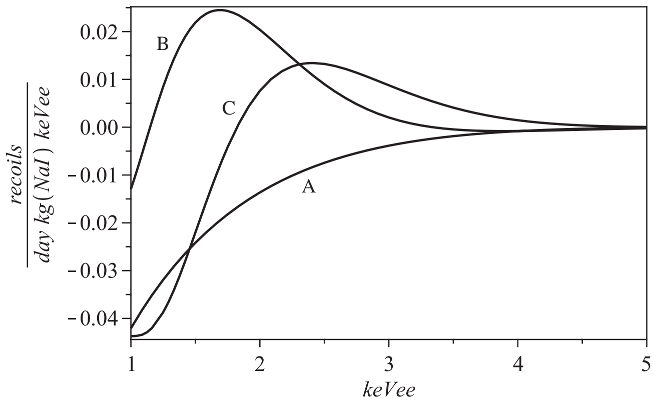

FIG. 4. Modulation of the counting rate in recoils/kg(NaI)/keVee for ν_b scattering on Na. As before, curve A is for large E_0, B is for $E_0 = 12$ MeV, and C is for $E_0 = 14$ MeV, while $\mathcal{N}_{\rm eff}^2 = 10^4$. Both signs of modulation are possible.

other observations and constraints (and with the simultaneous explanation of the CoGeNT low-E_r anomaly). Of course, the phase of the DAMA results will require E_0 to be in the right range. But even then, would the early July maximum be consistent with DAMA/LIBRA claims of the oscillation phase? The best fit point for the maximum is about 4–5 weeks different from the $t_{\rm aphelion}$ [3]. It would be interesting to find out whether the early July maximum is actually excluded by DAMA data, and if the criticism expressed in Ref. [43] about the errors on the phase being too tight is properly addressed by the collaboration.

Moving away from the ν_b idea, one can also notice that many other exotic physics explanations of the DAMA signal can be invoked (if it allows us to tolerate a phase shift of $\sim$1 month). For example, the emission of solar axions with their subsequent absorption in DM detection experiments can be a cause of low-energy ionization signals [13]. On this picture, one can superimpose the oscillations of axions into some ‘‘sterile axions’’ with an oscillation length similar to L_0 in order to break the monotonic $L(t)$ dependence, and flip the phase of the modulation, achieving results similar to those of Figs. 3 and 4.

V. DISCUSSION AND CONCLUSIONS

The oscillation of SM neutrinos to some new neutrino state on the way from the Sun to the Earth is a realistic possibility. We have shown that there exists a whole class of models where neutrino physics beyond the SM can be probed by the low-threshold DM detectors, which become equally sensitive or even more sensitive to this type of neutrino than the large-scale neutrino detectors. Such models require that new neutrino states ν_b [or modified SM neutrino interactions in the spirit of Eq. (13)] couple almost exclusively to the baryon current. The isoscalar vector properties of this current lead to a very strong enhancement of the elastic over inelastic scattering, $\sigma_{\rm elastic}/\sigma_{\rm inelastic} \sim 10^8$, providing an unexpected competitiveness factor to small-scale experiments such as CoGeNT. We have shown that the effective strength of the NCB can be much larger than the weak-scale value without being in conflict with any of the observational data.

We have also shown that the recent anomalies in direct DM detection, such as DAMA and CoGeNT, can be explained by the $\nu_{\rm SM} \rightarrow \nu_b$ oscillation of ^{8}B neutrinos with subsequent scattering of ν_b on Ge and Na nuclei. (This statement relies on the assumption of $t_{\rm aphelion}$ being consistent with the DAMA/LIBRA modulation phase.) This may look counterintuitive at first, but we have shown that the phase flip of seasonal modulations is possible for the high-energy end of the ^{8}B spectrum. This is a very speculative explanation (and perhaps as equally speculative as the WIMP recoil explanation), but it is interesting enough to motivate further studies. In particular, we believe that the Borexino Collaboration can perform the search of the ν_b-activated 4.4 MeV line in ^{12}C, and probably surpass the sensitivity to the NCB enhancement factor $\mathcal{N}_{\rm eff}$ of 100. At the same time, it seems apparent that further technological developments of low-threshold WIMP/ν_b detectors are required. Should the current low-energy anomalies in DM detectors firm up to constitute a definitive signal of new-physics-induced recoil, some significant efforts and different mass targets would be required to observationally distinguish between the low-mass WIMP and ν_b signals.

Below, we would like to discuss further implications of the models involving new neutrino states with enhanced baryonic currents.

(i) *Collider implications.* If the $G_B \gtrsim 100 G_F$, and if the new interaction is truly contact, the proton-antiproton collisions will lead to strong new sources of missing energy signals in the $\nu_b \bar{\nu}_b$ channel and will most likely be excluded by the Tevatron experiments. This will not happen, however, in models of relatively weakly coupled mediators with sub-GeV mass. Therefore, the collider searches should be able to place an *upper* bound on m_V.

(ii) *Fixed target implications.* A GeV-scale U(1)$_B$ baryonic vector, the carrier of the NCB interaction, can be produced in the collisions of energetic proton beams with a target. Immediate decays of these vectors will generate a flux of the ν_b state that can be searched for at near detectors via their NCB interactions. This is very similar to the ideas of the ‘‘MeV-scale DM beams’’ discussed previously in [44]. There can also be implications for terrestrial antineutrino physics, as matter effects induced by V exchange can be large for ν_b antineutrinos [28]. Enhanced neutral currents of ν_b neutrinos may help explain the long-standing puzzle of the LSND anomaly [45], perhaps borrowing some elements of the recent suggestion [46]. It also has to be said that, over the last two years, much effort has been invested in systematically searching for the "kinetic mixing" (or hypercharge) portal (see e.g. [47] and references therein). The baryonic portal is another example of a perfectly safe way, from the model-building perspective, of introducing stronger-than-weak forces at low energy, and therefore it should be systematically searched for using proton-on-target facilities. But perhaps the most NCB-search effective type of experiments to perform with proton beams is the proposal [29] of a neutrino-nucleus elastic scattering detector.

(iii) *Cosmological implications.* A new light neutrino state (and two neutrino states if the right-handed copy of ν_b is also light) can be at the borderline of what is allowed by early cosmology and observations of light elemental abundances (the most recent analysis can be found in [48]). Does the model with such a strong enhancement of baryonic currents have a chance to be consistent with these constraints? Actually, despite the interaction strength of G_F of 100 or 1000, ν_b will decouple from thermal plasma *earlier* than the SM neutrinos. That is because its thermalization rate will be proportional to the baryon-to-photon ratio, which is a small number $O(10^{-10})$. Therefore, the actual decoupling of ν_b may happen with the decays and annihilations of abundant hadronic species at temperatures of $\sim$100 MeV, and therefore big bang nucleosynthesis bounds from overpopulation of radiative degrees of freedom can be easily evaded.

(iv) *Astrophysical implications.* Another interesting aspect of ν_b models is their NCB production in stars. In the SN, ν_b will not provide new effective energy sinks because they would not freely escape the explosion zone. However, one should expect that a number of ν_b neutrinos comparable to the SM is created, so that one could detect them using the same DM/ν_b detectors. Should a nearby SN explosion happen, the existing neutrino scintillator detectors can pick up the ν_b-NCB signal that would appear as a much enhanced ν_μ, ν_τ-NC signal considered in [49] (modulo the uncertainty in the effective temperature for ν_b).

(v) *Rare decay implications.* Relatively large NCB currents should open new channels for the missing energy decays of B and K mesons. As argued in this paper, the conservation of the baryonic current makes it a relatively safe portal compared to e.g. scalar or axial-vector portals. Nevertheless, if the $K \to \pi V$ decay is kinematically allowed, it may lead to the underlying two-body signature of $K \to \pi$ plus missing energy decays, making it an appealing target for the next generation of precision kaon physics experiments.

ACKNOWLEDGMENTS

I would like to thank B. Batell, D. McKeen, J. Pradler, and I. Yavin for helpful discussions. I am also grateful to A. de Gouvea, R. Harnik, and J. Kopp for pointing out inconsistencies of the neutrino oscillation picture in version 1 of this work. The work is supported in part by the Government of Canada through NSERC and by the Province of Ontario through MEDT.

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