

Probing a light scalar boson with a few-MeV proton beam deep underground

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We propose to investigate the production of a light scalar boson ϕ in low-energy proton-nucleus interactions using the 3.5 MV accelerator of the Bellotti Ion Beam Facility, located in the underground Gran Sasso National Laboratory. Nuclear reactions induced by a few-MeV proton beam on suitable target materials can act as a controlled source of ϕ particles. Owing to the deep-underground location, the facility benefits from substantial cosmic-ray shielding, enabling searches for rare processes with minimal background. The produced ϕ particles will be sought with large-volume, low-background detectors already operating or currently under construction at the Gran Sasso National Laboratory. This approach combines a tunable accelerator-based production mechanism with the exceptional sensitivity of underground rare-event searches, offering a novel avenue to probe light scalar bosons beyond the Standard Model.

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I. INTRODUCTION

Light bosonic particles beyond the Standard Model (SM) are theoretically well motivated and arise in a variety of extensions of the SM, including models with additional scalar fields and dark-sector frameworks. Owing to their feeble couplings, such particles are difficult to probe at high-energy colliders, but may be produced and detected in low-energy laboratory experiments or through rare processes in astrophysics and cosmology. Among these candidates, light scalar bosons with MeV-scale masses provide particularly intriguing targets: they can play a role in dark matter dynamics, impact stellar evolution, and modify precision observables. Dedicated laboratory searches are therefore crucial to test their parameter space in a controlled and systematic way.

A promising strategy is to exploit nuclear reactions at low energies as a source of new light states, employing either underground accelerators or radioactive sources, as suggested, e.g., in Refs. [1,2]. Proton-nucleus interactions in the few-MeV regime can efficiently produce scalar particles via nuclear transitions, while the resulting flux is relatively tunable through the choice of beam energy and target

material. When combined with large-volume, ultra-low-background detectors, such as those employed in underground rare-event searches, this approach provides a unique opportunity to investigate light scalars with unprecedented sensitivity.

As a concrete realization of this idea, we consider the 3.5 MV accelerator of the Bellotti Ion Beam Facility (Bellotti-IBF) [3,4], located in the Gran Sasso National Laboratory (LNGS). The underground location ensures substantial shielding from cosmic rays, strongly reducing backgrounds and enabling the study of rare processes. In this setup, proton beams impinging on selected target nuclei act as a controlled source of scalar bosons ϕ , which can then be searched for with the existing or forthcoming generation of LNGS detectors, such as XENONnT [5] and DarkSide-20k [6].

The structure of this paper is as follows. In Sec. II we describe the production mechanisms for a light scalar boson in nuclear reactions and provide an estimate of the expected flux. In Sec. III we discuss detection strategies, focusing on the signatures in large-volume underground detectors. We present the results and outline future prospects in Sec. IV, before concluding in Sec. V. Appendix is devoted to current constraints on MeV-scale scalar bosons from astrophysics, flavor physics, and laboratory probes.

II. PRODUCTION OF A LIGHT SCALAR BOSON

The 3.5 MV accelerator at the Bellotti-IBF [3,4] can deliver proton, helium, and carbon beams [3,4]. In this

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work we focus on the proton beam, which can reach currents up to 1 mA. Proton fusion with a target nucleus (T) produces a new nucleus either in its ground state (N) or in an excited state (N_*). The excited state subsequently deexcites via the standard channels, i.e., gamma or particle emission, or, if the quantum numbers allow, through the emission of a new scalar boson, ϕ .

A. Nuclear production mechanisms

We can schematically define two classes of photon/scalar production mechanisms:

- (1) Direct nuclear production $p + T \rightarrow N + X$; and
- (2) Nuclear reaction $p + T \rightarrow N_* + \dots$ with $N_* \rightarrow N + X$.

where $X = \gamma/\phi$. For simplicity, we will consider only the last production mechanism, which is the easiest to evaluate, thus producing conservative bounds. We then need the decay widths for

$$\Gamma(N_i \rightarrow N_f X), \quad (2.1)$$

where $N_{i,f}$ are two energy levels of the same nucleus. Nuclear states possess definite spin and parity quantum numbers, respectively $J_{i,f}$ and $\pi_{i,f}$. Labeling the total (orbital) angular momentum of the emitted boson as J (L) and its parity as π , conservation laws impose,

$$|J_i - J_f| \leq J \leq J_i + J_f \quad \text{and} \quad \pi_i = \pi_f \pi (-1)^L. \quad (2.2)$$

If X is a photon, $J \geq 1$ and we identify two types of transition:

- (i) electric type transition EJ if $\pi_i \pi_f = (-1)^J$; and
- (ii) magnetic type transition MJ if $\pi_i \pi_f = (-1)^{J+1}$.

If X is a scalar, one has $J = L \geq 0$ and $\pi_i \pi_f = (-1)^L$, corresponding to an electric-type transition, with the exception that a photon cannot be emitted for $J = 0$.¹

1. Multipole expansion

We now expand Eq. (2.1) in multipoles, with k denoting the boson momentum,

$$\Gamma_\phi = \frac{2k}{2J_i + 1} \left\{ \sum_{J \geq 0} |\langle J_f | \mathcal{G}_J | J_i \rangle|^2 \right\}, \quad (2.3)$$

$$\Gamma_\gamma = \frac{2k}{2J_i + 1} \left\{ \sum_{J \geq 1} [|\langle J_f | \mathcal{T}_J^{\text{el}} | J_i \rangle|^2 + |\langle J_f | \mathcal{T}_J^{\text{mag}} | J_i \rangle|^2] \right\}, \quad (2.4)$$

where the spherical operators have been introduced (see, e.g., Ref. [7] for details and references)

¹If X is an axion one finds $J = L \geq 0$ and $\pi_i \pi_f = (-1)^{L+1}$, as in a magnetic-type transition, again with the exception that a photon cannot be emitted for $J = 0$.

$$\mathcal{G}_{JM} = \int d^3\vec{r} j_J(kr) Y_{JM}(\hat{r}) \mathcal{S}(\vec{r}), \quad (2.5)$$

$$\mathcal{T}_{JM}^{\text{el}} = \frac{1}{k} \int d^3\vec{r} \vec{\nabla} \times [j_J(kr) \mathbf{Y}_{JM}(\hat{r})] \cdot \vec{\mathcal{J}}_\gamma(\vec{r}), \quad (2.6)$$

$$\mathcal{T}_{JM}^{\text{mag}} = \int d^3\vec{r} [j_J(kr) \mathbf{Y}_{JM}(\hat{r})] \cdot \vec{\mathcal{J}}_\gamma(\vec{r}). \quad (2.7)$$

We then consider the nucleon-photon/scalar effective interaction Lagrangians, defined via

$$\mathcal{L}_{\gamma NN} = A_\mu(x) \mathcal{J}_\gamma^\mu(x) = A^\mu \left[e Q_p \bar{p} \gamma_\mu p + e Q_n \bar{n} \gamma_\mu n + \frac{e \kappa_p^\gamma}{2m_p} \partial^\nu (\bar{p} \sigma_{\mu\nu} p) + \frac{e \kappa_n^\gamma}{2m_n} \partial^\nu (\bar{n} \sigma_{\mu\nu} n) \right], \quad (2.8)$$

$$\mathcal{L}_{\phi NN} = \phi(x) \mathcal{S}(x) = \phi [g_p \bar{p} p + g_n \bar{n} n], \quad (2.9)$$

where $\kappa_p^\gamma = \mu_p - Q_p = +1.792847351(28)$, $\kappa_n^\gamma = \mu_n - Q_n = -1.9130427(5)$ and $Q_{p,n}$ indicates the electric charge of the nucleon in units of the absolute electron charge. The same methodology can be applied to internal e^+e^- pair creation, see Appendix B of Ref. [7] and references thereby.

Consider the case of proton impinging on a ^{19}F target. The nuclear fusion will produce ^{20}Ne , which will mainly decay via α emission to ^{16}O . Among the possible ^{16}O states populated by ^{20}Ne α decay, we focus on the $^{16}\text{O}(6.05) \rightarrow ^{16}\text{O}(g.s.)$, where both the initial and final states are $J^\pi = 0^+$ nuclei with isospin $I = 0$. In this case, due to angular momentum conservation, the single- γ emission is forbidden and the leading SM decay channel proceeds via internal e^+e^- pair creation, $^{16}\text{O}(6.05) \rightarrow ^{16}\text{O}(g.s.) + e^+e^-$. Thus, the relative branching to new physics can be greatly enhanced [1].² Employing isospin symmetry, the scalar emission rate can be related to the internal e^+e^- pair creation rate, and is given by [8]³

$$\mathcal{B}(^{16}\text{O}(6.05) \rightarrow ^{16}\text{O} + \phi) \approx \frac{15}{8} \left[1 - \left(\frac{m_\phi}{6.05 \text{ MeV}} \right)^2 \right]^{\frac{5}{2}} \left(\frac{g_p + g_n}{\alpha} \right)^2. \quad (2.10)$$

²In contrast, in the axion case, the SM background is dominated by the large single- γ decay, resulting in a strong suppression.

³The calculations of Refs. [7,9] reproduce the result of Ref. [8] in the full nonrelativistic limit of the nuclear current, and further include the next to leading order (NLO) term in such expansion. A naive estimate suggests that the NLO contribution could enhance the predicted branching ratio for scalar emission by a factor of $\sim 10^3$. However, this term depends on the nuclear matrix element of the kinetic operator, which has not been computed in the literature. For this reason, we conservatively neglect this enhancement factor. We also note that Eq. (2.10) numerically agrees with the result in Ref. [2].

B. Estimate of the scalar flux

The scalar flux from fusion reactions at the Bellotti-IBF receives contributions from all the production mechanisms discussed above. For simplicity, we neglect the direct nuclear contribution, which is difficult to evaluate, and focus instead on scalars produced through resonant reactions followed by subsequent decays. The resulting flux therefore represents a conservative estimate of the total yield. Accordingly, the number of scalars produced in fusion reactions can be approximated as

$$N_\phi = N_{\text{POT}} \times \sum_{N_*} \left\{ f_{N_*}(T) \times \sum_{N_f} \mathcal{B}(N_* \rightarrow N_f + \phi) \right\}, \quad (2.11)$$

where $f_{N_*}(T)$ is the multiplicity of the N_* states produced for each proton on target T , while N_{POT} is the number of protons on target (POT) delivered at the Bellotti-IBF. From this viewpoint, the scalar flux is a superposition of monochromatic components

$$N_\phi = \sum_{\omega_\phi} N_\phi(\omega_\phi), \quad (2.12)$$

summing over all possible transition energies. The multiplicity of the excited state can be evaluated as [1]

$$f_{N_*}(T) = n_T \times \int_0^{E_p} dE \frac{\sigma(p + T \rightarrow N_* + \dots)}{|dE/dx|}, \quad (2.13)$$

where n_T denotes the target density and the stopping power $|dE/dx|$ depends on the target material. The ellipsis in the cross section indicates that we perform an inclusive sum over all possible byproducts (γ , α , etc.) accompanying the excited state of interest.

We recall that the final nucleus N_* is produced approximately at rest, so the angular distribution of the scalar flux is entirely given by the nuclear transition and it can be shown that the angular distribution of the emitted boson is isotropic.

On a practical level, we consider a 3 MeV proton beam with an intensity of 1 mA, equivalent to 6.2×10^{15} POT per second, impinging on a TaF₃ target of about 3 mm thickness. The proton beam will stop inside the target, allowing the cross section to be integrated over the full beam energy. In this scenario, the expected multiplicity of the $^{16}\text{O}(6.05)$ state is found to be,

$$f_{^{16}\text{O}(6.05)}(\text{TaF}_3) = 2.3 \times 10^{-6}. \quad (2.14)$$

The N_* production cross section was taken from the comprehensive R-matrix analysis of Ref. [10]. To be conservative we assumed a vanishing contribution from the 2^+ (13095 keV) resonant state in ^{20}Ne . The effective

stopping power has been calculated using the program SRIM [11]. The profile of the target was assumed to be boxlike, with about 10^{18} atoms/cm². It must be noted that compared to calculation reported in [1], we used an updated cross section, based on experimental data rather than a simple model, and a realistic thick solid target.

III. DETECTION OF A LIGHT SCALAR BOSON

The scalar flux emerging from the target can then be intercepted by a suitable detector. We consider two scenarios: detection with the existing XENONnT experiment [5], and with the DarkSide-20k setup [6] currently under construction.

A. Experimental signal signatures

The experimental signatures of the signal include:

- (i) a photon from scattering, $\phi + e^- \rightarrow e^- + \gamma$;
- (ii) two photons from the decay, $\phi \rightarrow \gamma\gamma$; and
- (iii) an electron-positron pair from the decay, $\phi \rightarrow e^+e^-$.

We consider here also a lepton-scalar coupling through the Lagrangian in Eq. (4.1). In our scenario the scalar can only decay into an e^+e^- pair or into two photons, thus

$$\Gamma_\phi = \Gamma(\phi \rightarrow \gamma\gamma) + \Gamma(\phi \rightarrow e^+e^-). \quad (3.1)$$

In the following, we report the number of expected events for each channel.

1. $\phi + e^- \rightarrow e^- + \gamma$

The number of photons produced in the signal volume, with the interaction point set at $\vec{x}_{\text{IP}} = 0$, is given by

$$N_\gamma = \sum_{\omega_\phi} N_\phi(\omega_\phi) \int \frac{d\Omega_\phi}{4\pi} \int_0^\infty dt \exp\left\{-\frac{t}{\tau_\phi \gamma_\phi}\right\} + \int_0^t dt' n_e(\vec{\beta}_\phi t') \sigma_{\phi \rightarrow \gamma} \beta_\phi \left\{ n_e(\vec{\beta}_\phi t) \sigma_{\phi \rightarrow \gamma}^{\text{cut}} \beta_\phi, \right. \quad (3.2)$$

where $\vec{\beta}_\phi$ is the scalar velocity, Ω_ϕ the scalar emission solid angle, $\beta_\phi = |\vec{\beta}_\phi| = \sqrt{\omega_\phi^2 - m_\phi^2}/\omega_\phi$, $\gamma_\phi = 1/\sqrt{1 - \beta_\phi^2}$, $\tau_\phi = \Gamma_\phi^{-1}$ the scalar lifetime at rest, and $n_e(\vec{x})$ the electron number density of the detector. The cross section reads

$$\sigma_{\phi \rightarrow \gamma}^{(\text{cut})} = \left(\frac{g_e^2 \alpha}{2m_e \omega_\phi^2 \beta_\phi^2} \right) \times \int_{(\text{cut})} d\omega_\gamma F(m_e^2 + m_\phi^2 + 2m_e \omega_\phi, m_e^2 - 2m_e \omega_\gamma), \quad (3.3)$$

with

$$\begin{aligned}
F(s, t) &= \frac{1}{2} \sum_{\text{pol}} |\mathcal{M}(\phi + e^- \rightarrow e^- + \gamma)|^2 \\
&= -\frac{2}{(m_e^2 - s)^2 (m_e^2 - t)^2} \\
&\quad \times \{20m_e^8 - 4m_e^6(3m_\phi^2 + 4(s+t)) \\
&\quad + m_e^4(2m_\phi^4 + 5s^2 - 2st + 5t^2 + 10m_\phi^2(s+t)) \\
&\quad - m_e^2(4m_\phi^2 st + 2m_\phi^4(s+t) + (s-t)^2(s+t)) \\
&\quad + st(2m_\phi^4 - 2m_\phi^2(s+t) + (s+t)^2)\}. \quad (3.4)
\end{aligned}$$

The minimal and maximal value of the photon energy are

$$\begin{aligned}
(\omega_\gamma)_{\min} &= \frac{(m_e + \omega_\phi(1 - \beta_\phi))(m_\phi^2 + 2m_e\omega_\phi)}{2(m_e^2 + m_\phi^2 + 2m_e\omega_\phi)}, \\
(\omega_\gamma)_{\max} &= \frac{(m_e + \omega_\phi(1 + \beta_\phi))(m_\phi^2 + 2m_e\omega_\phi)}{2(m_e^2 + m_\phi^2 + 2m_e\omega_\phi)}. \quad (3.5)
\end{aligned}$$

However, to identify the signal, the detector may require veto conditions, thereby modifying the integral into $\sigma_{\phi \rightarrow \gamma}^{\text{cut}}$ to account for this effect.

In the limit where the scalars arrive at the detector approximately collinearly, i.e., when the detector size is much smaller than the distance L between the production target and the detector, so that the solid angle of the captured scalar flux is negligible, the integration can be simplified and the expression reduces to

$$N_\gamma = \frac{N_\phi}{4\pi L^2} \times V_{\text{eff}} \times n_e \times \sigma_{\phi \rightarrow \gamma}^{\text{cut}} \times \exp\left\{-\frac{L}{\tau_\phi \gamma_\phi \beta_\phi}\right\}, \quad (3.6)$$

which is similar to what has been used in Ref. [1].

2. $\phi \rightarrow \gamma\gamma$

Scalar couplings to charged fermions induce the diphoton decay of the scalar at one-loop, with a decay rate given by

$$\Gamma(\phi \rightarrow \gamma\gamma) = \frac{\alpha^2 m_\phi^3}{512\pi^3} \left| \sum_f q_f^2 \frac{g_f}{m_f} A_{1/2}\left(\frac{m_\phi^2}{4m_f^2}\right) \right|^2, \quad (3.7)$$

with

$$\begin{aligned}
A_{1/2}(x) &= \frac{2[x + (x-1)f(x)]}{x^2}, \\
f(x) &= \begin{cases} \arcsin^2 \sqrt{x} & \text{for } x \leq 1, \\ -\frac{1}{4} \left(\ln \frac{1+\sqrt{1-x^{-1}}}{1-\sqrt{1-x^{-1}}} - i\pi \right)^2 & \text{for } x > 1, \end{cases} \quad (3.8)
\end{aligned}$$

where the sum in Eq. (3.7) is performed over all the fermion fields of the low-energy theory with electric charge q_f . The scalar couplings to nucleons originate at the fundamental

level from scalar interactions with quarks and gluons (cf. discussion in Sec. IV). To estimate the contribution of the light-quark couplings, we then rely on the calculation of Ref. [12], which incorporates the one-loop effects of pseudoscalar mesons within chiral perturbation theory.

The number of diphoton decays then produced in the detector, assuming that the interaction target is set at the origin, i.e., $\vec{x}_{\text{IP}} = 0$, is given by

$$\begin{aligned}
N_{\gamma\gamma} &= \sum_{\omega_\phi} N_\phi(\omega_\phi) \int \frac{d\Omega_\phi}{4\pi} \int_0^\infty dt \exp\left\{-\frac{t}{\tau_\phi \gamma_\phi}\right\} \\
&\quad + \int_0^t dt' n_e(\vec{\beta}_\phi t') \sigma_{\phi \rightarrow \gamma\gamma} \beta_\phi \left\{ \frac{\mathcal{B}(\phi \rightarrow \gamma\gamma)}{\tau_\phi \gamma_\phi} \Theta(\vec{\beta}_\phi t) \right\}, \quad (3.9)
\end{aligned}$$

where

$$\Theta(\vec{x}) = \begin{cases} 1 & \text{if } \vec{x} \text{ is inside the detector,} \\ 0 & \text{if } \vec{x} \text{ is not inside the detector.} \end{cases} \quad (3.10)$$

In the limit of L much larger than the detector size, we get

$$N_{\gamma\gamma} = \frac{N_\phi}{4\pi L^2} \times V_{\text{eff}} \times \frac{\mathcal{B}(\phi \rightarrow \gamma\gamma)}{\tau_\phi \gamma_\phi \beta_\phi} \times \exp\left\{-\frac{L}{\tau_\phi \gamma_\phi \beta_\phi}\right\}. \quad (3.11)$$

3. $\phi \rightarrow e^+ e^-$

For $m_\phi > 2m_e$, Eq. (4.1) gives

$$\Gamma(\phi \rightarrow e^+ e^-) = g_e^2 \frac{m_\phi}{8\pi} \left(1 - \frac{4m_e^2}{m_\phi^2}\right)^{3/2}. \quad (3.12)$$

The number of $e^+ e^-$ pair decays produced in the detector, assuming that the target is set at the origin, i.e., $\vec{x}_{\text{IP}} = 0$, is then given by

$$\begin{aligned}
N_{ee} &= \sum_{\omega_\phi} N_\phi(\omega_\phi) \int \frac{d\Omega_\phi}{4\pi} \int_0^\infty dt \exp\left\{-\frac{t}{\tau_\phi \gamma_\phi}\right\} \\
&\quad + \int_0^t dt' n_e(\vec{\beta}_\phi t') \sigma_{\phi \rightarrow e^+ e^-} \beta_\phi \left\{ \frac{\mathcal{B}(\phi \rightarrow e^+ e^-)}{\tau_\phi \gamma_\phi} \Theta(\vec{\beta}_\phi t) \right\}. \quad (3.13)
\end{aligned}$$

In the limit of L much larger than the detector size, we get

$$N_{ee} = \frac{N_\phi}{4\pi L^2} \times V_{\text{eff}} \times \frac{\mathcal{B}(\phi \rightarrow e^+ e^-)}{\tau_\phi \gamma_\phi \beta_\phi} \times \exp\left\{-\frac{L}{\tau_\phi \gamma_\phi \beta_\phi}\right\}. \quad (3.14)$$

TABLE I. Parameters used in our analysis: distance from the Bellotti-IBF (L), active volume (V_{eff}), and electron number density (n_e) for XENONnT and DarkSide-20k.

	L	V_{eff}	n_e
XENONnT [5]	60 m	2.0 m ³	$7.3 \times 10^{29} \text{ m}^{-3}$
DarkSide-20k [6]	90 m	500 m ³	$3.8 \times 10^{29} \text{ m}^{-3}$

B. Sensitivity of XENONnT and DarkSide-20k

We emphasize that in the previous calculations we assumed the scalar interactions to occur and be detected inside the central detector of XENONnT, namely the liquid xenon time projection chamber (TPC) with an active target mass of 5.9 tons [5]. We did not include possible interactions in the two larger water Cherenkov vetoes surrounding the cryostat, since their limited energy resolution and high background levels are unlikely to yield a competitive signal-to-background ratio.

In contrast, for the DarkSide-20k setup (under construction) described in [6], we account for scalar interactions both in the central liquid argon (LAr) TPC and in the surrounding LAr vetoes, corresponding to a total LAr mass of about 700 tons. Unlike water Cherenkov vetoes, the LAr veto operates via scintillation light detection, potentially offering significantly better energy resolution.

In our analysis we use the detector parameters summarized in Table I, and assume that both XENONnT and DarkSide-20k operate with unit detection efficiency. For the sensitivity projections we require at least ten signal events per year, assuming an equal number of background events. This benchmark corresponds to the observation of a 2σ excess: with $N_S + N_B = 20$ total events and $N_B = 10$ background events per year, the statistical uncertainty is $\sqrt{N_S + N_B} \simeq 5$, so that the signal stands at about 2σ above the background expectation.

To estimate the background rate, we consider the background levels achieved by the Borexino experiment at the LNGS with a 280-ton active-volume liquid-scintillator detector [13] (the detector, which required ultralow backgrounds as in DarkSide-20k, has now been dismantled). In a 145-ton fiducial volume, over 11 years, and after applying cosmogenic temporal and spatial vetoes, fewer than ten background events were collected in a 40 keV energy bin around 6 MeV. A similar background rate will likely be achieved by DarkSide-20k, at least in the central core of the detector.

We finally point out that one of the most important backgrounds in accelerator experiments, i.e., the

beam-induced backgrounds that produce γ rays and neutrons, is absent in our case. First, the beam energy is only 3 MeV, and the accelerator-target room has 80-cm-thick concrete walls. In addition, the accelerator-target complex is located in Gran Sasso Hall-B, whereas the DarkSide-20k detector is in Hall-C, with more than 50 m of rock in between.

IV. RESULTS AND FUTURE PROSPECTS

In the Appendix we summarize the existing bounds on an MeV-scale scalar coupled to nucleons and electrons. These constraints are essential when comparing with the future sensitivities of the experimental setup proposed in this work.

As a validation of our framework, we have verified that our results reproduce those of Ref. [2] when adopting the same simplified model. Our assumptions, however, differ in important respects. The simplified model of Ref. [2] was originally motivated by the proton radius anomaly [14] and the muon $g - 2$ discrepancy [15,16], both of which are now resolved. In particular, explaining the proton radius anomaly required $g_n \ll g_p$, while the muon $g - 2$ motivated scalar couplings to heavy leptons, which also affected the $\phi \rightarrow \gamma\gamma$ rate.

In this work we adopt a simpler setup, assuming couplings only to electrons and light quarks at the fundamental level,

$$\mathcal{L} \supset \phi \sum_{f=e,u,d,s} g_f \bar{\psi}_f \psi_f. \quad (4.1)$$

This choice is further motivated by the strong constraints from $K^+ \rightarrow \pi^+ \phi$ decay (cf. Appendix A 2), which requires the flavor-alignment condition [17]

$$\frac{g_u}{m_u} = \frac{g_d}{m_d} = \frac{g_s}{m_s} \equiv \frac{1}{f_\phi}, \quad (4.2)$$

that we will assume in the following.

The above condition has direct implications for nucleon couplings. The matching between scalar couplings to quarks and nucleons, which is tied to nucleon mass generation, reads (see, e.g., [18])

$$g_N = m_N \sum_{q=u,d,s} f_{Tq}^{(N)} \frac{g_q}{m_q}, \quad (4.3)$$

with $N = p, n$, and neglecting couplings to gluons. The nucleon mass fractions are [19,20]

$$f_{Tu}^{(p)} = 0.020 \pm 0.004, \quad f_{Td}^{(p)} = 0.026 \pm 0.005, \quad f_{Ts}^{(p)} = 0.118 \pm 0.062, \quad (4.4)$$

$$f_{Tu}^{(n)} = 0.014 \pm 0.003, \quad f_{Td}^{(n)} = 0.036 \pm 0.008, \quad f_{Ts}^{(n)} = 0.118 \pm 0.062. \quad (4.5)$$

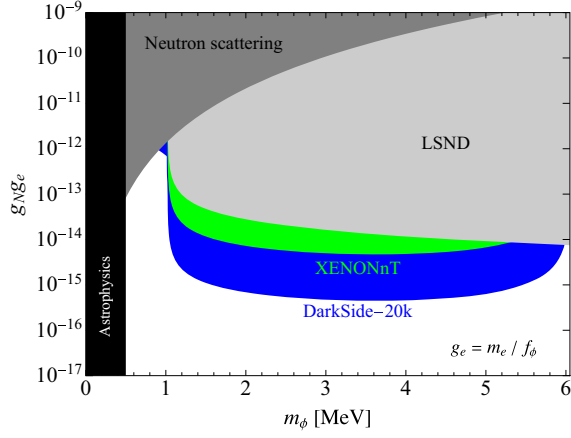


FIG. 1. Projected sensitivity of XENONnT (green) and DarkSide-20k (blue), combined with scalar production at the Bellotti-IBF. The projections assume $g_e = m_e/f_\phi$ and correspond to 1 year of data taking with a 3 MeV, 1 mA proton beam on a 3 mm thick TaF₃ target. Constraints from LSND (grey), neutron scattering (dark grey), and astrophysics (black) are also shown, while those from the electron $g-2$, SNO and Borexino are subdominant (cf. Appendix).

Importantly, in the flavor-aligned scenario, the scalar couplings to protons and neutrons cannot be tuned independently. Numerically one finds

$$g_p \approx 1.54 \times \left(\frac{100 \text{ MeV}}{f_\phi} \right), \quad g_n \approx 1.58 \times \left(\frac{100 \text{ MeV}}{f_\phi} \right), \quad (4.6)$$

so that $g_p \simeq g_n$, up to $\mathcal{O}(3\%)$ isospin-breaking corrections. The scalar emission rate is then estimated from Eq. (2.10) as

$$\mathcal{B}(^{16}\text{O}(6.05) \rightarrow ^{16}\text{O} + \phi) \approx 1.4 \times 10^{-9} \left(\frac{g_N}{10^{-7}} \right)^2, \quad (4.7)$$

where we neglected the mass term in the phase space factor and we normalized the scalar-nucleon coupling to a typical benchmark probed by the proposed setup. In contrast to the analysis of Refs. [1,2], which assumed $g_n = 0$, our framework therefore predicts an enhancement of a factor 4 in the ϕ production rate [cf. Eq. (2.10)] and a factor 2 in the liquid scintillator neutrino detector (LSND)

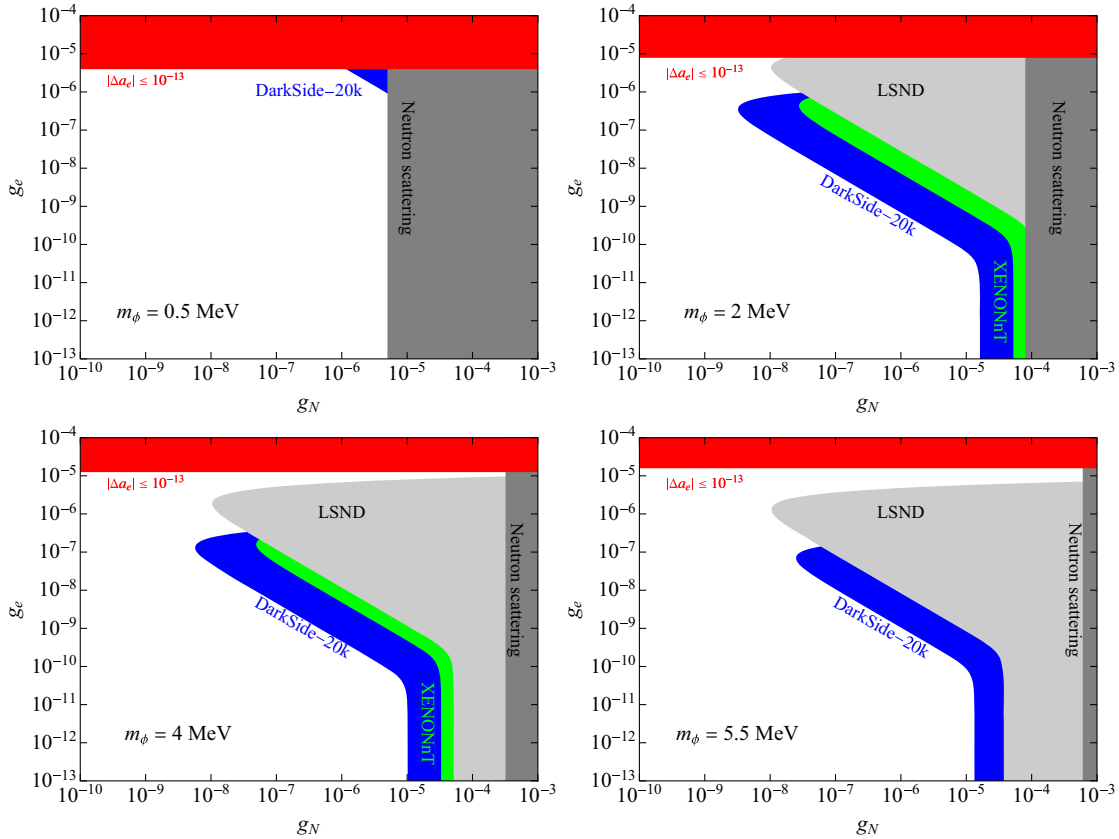


FIG. 2. Projected sensitivity of XENONnT (green) and DarkSide-20k (blue) combined with scalar production at the Bellotti-IBF, for fixed values of the scalar mass. The projections assume a data acquisition time of 1 year with the same beam and target as in Fig. 1. Constraints from LSND (grey), neutron scattering (dark grey), and the electron $g-2$ (red) are also shown, while SNO and Borexino remain subdominant.

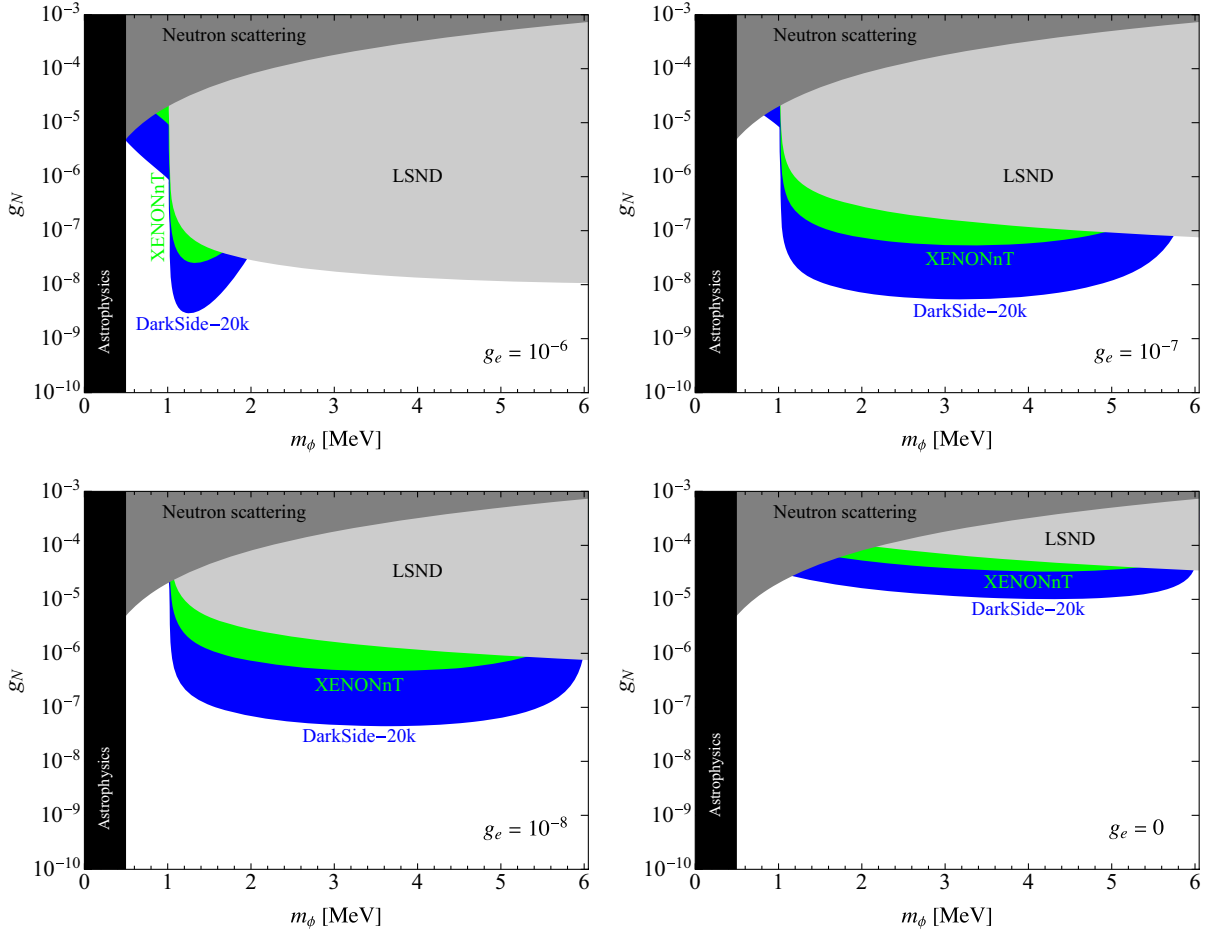


FIG. 3. Projected sensitivity of XENONnT (green) and DarkSide-20k (blue) combined with scalar production at the Bellotti-IBF, for fixed values of the electron coupling. The projections assume a data acquisition time of 1 year with the same beam and target as in Fig. 1. Constraints from LSND (grey), neutron scattering (dark grey), and astrophysics (black) are also shown, while SNO and Borexino are subdominant.

constraint (cf. Appendix A 4). Furthermore, the solar reaction $p + {}^2\text{H} \rightarrow {}^3\text{He} + \phi$, which produces scalars with energy $E_\phi = 5.49$ MeV, is proportional to the isovector combination $g_p - g_n$ (cf. Eq. (A8)). As a result, constraints from scalar flux detection in the Sudbury Neutrino Observatory (SNO) experiment (via g_N) and Borexino (via g_e) are strongly suppressed in this limit.

The projected reach of the setup discussed in this work is shown in Fig. 1, where we display the sensitivity of XENONnT (green) and DarkSide-20k (blue) to the product coupling $g_N g_e$ as a function of the scalar mass m_ϕ . For comparison, we also show existing bounds from astrophysics (cf. Appendix A 1), neutron scattering (cf. Appendix A 3), and the beam-dump experiment LSND (cf. Appendix A 4). To generate this plot we fixed the electron coupling to $g_e = m_e/f_\phi$.

Further insight can be obtained by fixing m_ϕ and presenting the reach in the (g_N, g_e) plane. This is illustrated in Fig. 2, where we show benchmarks for $m_\phi = (0.5, 2, 4, 5.5)$ MeV. In this representation, the anomalous

magnetic moment of the electron (cf. Appendix A 6) provides important constraints on g_e . The sensitivity curves terminate at large g_e values, since in this regime the scalar lifetime becomes too short for ϕ to reach the detector. A pronounced change also occurs across the e^+e^- threshold at $m_\phi \simeq 1$ MeV, where the opening of the efficient $\phi \rightarrow e^+e^-$ detection channel significantly enhances the sensitivity. Conversely, as m_ϕ approaches the kinematic threshold for scalar production, the LSND bounds become dominant over the projected signal reach.

Finally, Fig. 3 presents the complementary case of fixing g_e to benchmark values, starting from $g_e \lesssim 10^{-6}$, as allowed by the electron $g - 2$. This allows one to visualize the reach in terms of the nucleon coupling g_N , as a function of m_ϕ , under different assumptions for the electron coupling.

V. CONCLUSIONS

In this work we have proposed and analyzed a novel strategy to search for MeV-scale scalar bosons. The central idea is to exploit nuclear reactions induced by low-energy

accelerators as controlled sources of new light particles, and to search for their flux with large-volume, low-background detectors typically designed for rare-event physics. As a concrete realization, we have considered the Gran Sasso National Laboratory, where the 3.5 MV accelerator at the Bellotti-IBF can serve as a tunable source of nuclear reactions. The resulting scalar flux could then be searched for with existing/forthcoming detectors such as XENONnT and DarkSide-20k.

We have discussed the nuclear production mechanisms, focusing on scalar emission from excited nuclear states, and provided estimates of the resulting fluxes. We then studied the detection channels—electron scattering, diphoton decays, and e^+e^- decays—and presented the projected sensitivities of XENONnT and DarkSide-20k. These sensitivities were compared with existing bounds from astrophysics, neutron scattering, and LSND.

Our analysis shows that the proposed setup can probe previously unexplored regions of parameter space, complementary to astrophysical and laboratory constraints. The approach is parasitic in nature, relying on detectors already in operation or construction, and can thus provide a cost-effective and innovative probe of light scalar bosons.

Future work may include a more refined treatment of nuclear production mechanisms, as well as dedicated efficiency, energy resolution, and background studies for the specific detectors. Moreover, this proposal could serve as a test bed for probing other light bosonic particles beyond scalars.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available because they are owned by a third party and the terms of use prevent public distribution. The data are available from the authors upon reasonable request.

APPENDIX: CONSTRAINTS ON MeV-SCALE SCALAR BOSONS

In this appendix, we collect present bounds on light scalar bosons coupled to electrons and nucleons, specifically for $m_\phi \lesssim 6$ MeV, that are relevant for the new search discussed in this work.

1. Astrophysical constraints

Thermal production of light scalar bosons may lead to energy loss in stars, thus providing strong constraints on the scalar coupling to electrons (g_e) and nucleons (g_N), see, e.g., [21–24]. These bounds are exponentially suppressed if the mass of the scalar boson is larger than the typical temperature of the astrophysical object. Therefore, we focus on the region $m_\phi \gtrsim 0.5$ MeV, where the constraints on g_e from horizontal-branch stars, red giants, and white-dwarf cooling become negligible, and the leading astrophysical limits on g_N instead originate from neutron-star cooling [24] and SN 1987A [23]. However, we do not display explicit astrophysical bounds in our plots, as existing studies typically consider one coupling at a time, while the simultaneous presence of multiple interactions could change the picture. In the case of supernovae, constraints can be very strong in the free-streaming regime, but the trapping regime is affected by sizable uncertainties. In particular, the reabsorption of beyond the SM particles inside the protoneutron star can influence multidimensional processes such as convection and accretion, whose impact on the neutrino signal is not yet reliably quantified. For these reasons, we simply refer the reader to the discussion in Ref. [23], emphasizing these caveats.

2. $K^+ \rightarrow \pi^+ \phi$

Flavor changing Kaon decays involving a light particle with a missing energy signature are well constrained by collider searches at NA62, namely $\mathcal{B}(K^+ \rightarrow \pi^+ \phi) \lesssim \mathcal{O}(10^{-11})$ [25–27]. Even if the scalar couplings are all diagonal in flavor space, such decay can be induced by a W loop. For a recent calculation, including next-to-leading order corrections, see Ref. [17]. Using these results, one finds that if the light scalar couples only to light quarks (i.e., not to gluons) through the Lagrangian term in Eq. (4.1), and if these couplings are proportional to the quark masses, $g_u/m_u = g_d/m_d = g_s/m_s \equiv 1/f_\phi$, then the leading-order expression for $K^+ \rightarrow \pi^+ \phi$ in the chiral expansion vanishes. In this flavor-aligned scenario, the constraint on g_N from $K^+ \rightarrow \pi^+ \phi$ searches is significantly weakened⁴ and becomes subleading compared to other bounds discussed below.

3. Neutron scattering experiments

As discussed in Ref. [14], neutron scattering experiments provide stringent bounds on new light mediators that couple to neutrons. In particular, a mediator with mass in the MeV range induces corrections to the neutron-nucleus scattering cross section, which can interfere with the strong interaction amplitude. The resulting angular distortions in the differential cross section allow for bounds

⁴For comparison, in the case of a Higgs-mixed light scalar, the bound from $K^+ \rightarrow \pi^+ \phi$ corresponds to $g_N \lesssim 10^{-7}$ [28].

on the neutron coupling strength, originally derived in Ref. [29] (see also [30,31]). For a mediator with equal couplings to protons and neutrons, in the relevant mass range considered in this work, this leads to

$$g_N \lesssim 2 \times 10^{-5} \left(\frac{m_\phi}{\text{MeV}} \right)^2. \quad (\text{A1})$$

In our case, since we do not require the mediator to address the muon $g-2$ or muonic hydrogen anomalies (as, e.g., in [14]), couplings of comparable strength to protons and neutrons remain consistent with current bounds, while still being testable in the parameter space accessible to the setup proposed in this work.

4. LSND

The LSND measurements of the elastic electron-neutrino cross section [32,33] can be reinterpreted as constraints on light scalar particles. A full recast would require a dedicated analysis; here, as a first approximation, we follow the approach of Ref. [2].

Assuming that the decay of the Δ resonance saturates the pion production inside the target, one can estimate the scalar production from the decay $\Delta \rightarrow N + \pi + \phi$, with the scalar emitted from the nucleon leg, as

$$N_\phi^{\text{LSND}} \approx N_\pi^{\text{LSND}} \times \sum_{N=p,n} \frac{\Gamma(\Delta \rightarrow N + \pi + \phi)}{\Gamma(\Delta \rightarrow N + \pi)} \times P_{\text{signal}}. \quad (\text{A2})$$

The last term in the above expression accounts for the probability that the emitted scalar produces a signal inside the detector, given by

$$P_{\text{signal}} = \left[\exp \left\{ -\frac{L_{\text{LSND}} + \frac{d_{\text{LSND}}}{2}}{\tau_\phi \gamma_\phi \beta_\phi} \right\} - \exp \left\{ -\frac{L_{\text{LSND}} - \frac{d_{\text{LSND}}}{2}}{\tau_\phi \gamma_\phi \beta_\phi} \right\} \right] \left(\frac{A_{\text{LSND}}}{4\pi L_{\text{LSND}}^2} \right), \quad (\text{A3})$$

where $L_{\text{LSND}} = 30$ m, $d_{\text{LSND}} = 8.3$ m and $A_{\text{LSND}} \approx 25$ m². Relying again on Ref. [2], we consider the approximated result

$$\frac{\Gamma(\Delta \rightarrow N + \pi + \phi)}{\Gamma(\Delta \rightarrow N + \pi)} \approx 0.04 \times g_N^2, \quad (\text{A4})$$

with the average energy of the scalar estimated as 300 MeV. Finally, a conservative estimate of the number of pions produced in the experiment is $N_\pi^{\text{LSND}} \sim 10^{22}$, while the number of signal events is taken to be less than 20 as in Ref. [2].

5. Solar ϕ flux from $p + {}^2\text{H}$ fusion

Light scalars can be produced via nuclear reactions occurring in the Sun, and later be detected on Earth. We focus here on the reaction $p + {}^2\text{H} \rightarrow {}^3\text{He} + \phi$, which generates a flux of scalars with energy of 5.49 MeV.

In the Sun, deuterium (${}^2\text{H}$) is produced in 99.6% of cases through the proton-proton chain, $p + p \rightarrow {}^2\text{H} + e^+ + \nu$, whose neutrino flux $\Phi_{pp\nu} \approx 6 \times 10^{10} \text{ cm}^{-2} \text{ s}^{-1}$ has been measured by Borexino [34], while the remaining 0.4% is due to $p + p + e^- \rightarrow {}^2\text{H} + \nu$. The deuterium is immediately converted into Helium-3 via $p + {}^2\text{H} \rightarrow {}^3\text{He} + \gamma$, or, in our scenario, also via scalar boson emission (in place of a photon), though with a suppressed rate. The scalar flux can be thus estimated by appropriately rescaling the $\Phi_{pp\nu}$ flux (see also [2])

$$\Phi_\phi^{\text{Sun}} \approx \frac{\sigma(p + {}^2\text{H} \rightarrow {}^3\text{He} + \phi)}{\sigma(p + {}^2\text{H} \rightarrow {}^3\text{He} + \gamma)} \times \Phi_{pp\nu} \times P_{\text{exit}} \times P_{\text{survive}}, \quad (\text{A5})$$

where we have also included terms accounting for the probability of ϕ escaping the star and reaching the Earth. The former is given by

$$P_{\text{exit}} = \exp \left\{ - \int^{R_\odot} dr n_\odot(r) \times \sigma_{\phi \rightarrow \gamma} \right\}, \quad (\text{A6})$$

with $R_\odot = 6.96 \times 10^{10}$ cm denoting the solar radius, and $n_\odot$ the mean solar electron density, which is approximately described by the exponential function [35]

$$\frac{n_\odot(r)}{N_A} \approx 245 \exp \left\{ -10.54 \frac{r}{R_\odot} \right\} \text{ cm}^{-3}, \quad (\text{A7})$$

where $N_A = 6.022 \times 10^{23}$ is the Avogadro's number. The survival probability until the Earth is instead given by

$$P_{\text{survive}} = \exp \left\{ -\frac{L_\odot}{\tau_\phi \gamma_\phi \beta_\phi} \right\}, \quad (\text{A8})$$

where $L_\odot = 1.5 \times 10^{11}$ m is the Sun-Earth distance.

The cross section of the γ emission is a sum of M1 (s-wave) and E1 (p-wave) multipole contributions with a ratio $\sigma_{\text{M1}}/\sigma_{\text{E1}} \approx 1.3/2.9$ [36], with both of them predominantly isovectorial [37]. The scalar emission cross section can then be obtained by rescaling the E1 contribution of the electromagnetic one yielding [2,9]

$$\begin{aligned} & \frac{\sigma(p + {}^2\text{H} \rightarrow {}^3\text{He} + \phi)}{\sigma(p + {}^2\text{H} \rightarrow {}^3\text{He} + \gamma)} \\ &= \frac{2.9}{1.3 + 2.9} \times \frac{1}{2} \left[1 - \left(\frac{m_\phi}{5.49 \text{ MeV}} \right)^2 \right]^{\frac{3}{2}} \left(\frac{g_p - g_n}{e} \right)^2. \end{aligned} \quad (\text{A9})$$

a. SNO (detection via g_N)

A significant constraint on light scalar interactions with nucleons can be derived by recasting the SNO experiment [38] bound originally set for axionlike particles (ALPs) [39], as discussed in Ref. [28]. This bound relies on the assumption that the new particles produced in solar nuclear transitions are sufficiently long-lived to reach the Earth-based SNO detector, where they could be observed via deuterium dissociation events.

In analogy to the ALP case, we reinterpret the SNO bound in the case of a light scalar boson ϕ that couples to nucleons. The scalar is assumed to be emitted from nuclear transitions in the Sun and then detected via its interaction with deuterium in the SNO detector, provided it survives the propagation from the solar core to Earth. The scalar flux from the Sun is taken to be the one in Eq. (A5). The detection cross section is given by Eq. (4.3) in Ref. [28], which describes deuterium dissociation induced by scalar absorption.

This bound becomes ineffective in regions of parameter space where the scalar decays before reaching the Earth. Nevertheless, in the regime where the scalar is sufficiently long-lived, the recast SNO data provide a relevant and competitive constraint on light scalar couplings to nucleons.

b. Borexino (detection via g_e)

Similarly to SNO, the scalar flux on Earth is detectable at Borexino if the particle is long-lived, see Ref. [40]. The expected number of $\phi + e^- \rightarrow e^- + \gamma$ events is constrained as

$$\Phi_\phi^{\text{Sun}} \times \sigma_{\phi \rightarrow \gamma} \times n_e^{\text{Bor.}} \times T^{\text{Bor.}} \times \epsilon_{\text{eff}}^{\text{Bor.}} < 6.9 \text{ at 90\% CL}, \quad (\text{A10})$$

where $n_e^{\text{Bor.}} = 9.17 \times 10^{31}$ is the number of electrons, $T^{\text{Bor.}} = 4.63 \times 10^7$ s is the exposure time, and $\epsilon_{\text{eff}}^{\text{Bor.}} = 0.358$ is the efficiency. Furthermore the expected number of $\phi \rightarrow \gamma\gamma$ decay inside the Borexino detector is constrained as

$$\Phi_\phi^{\text{Sun}} \times V_{\text{eff}}^{\text{Bor.}} \times \frac{\mathcal{B}(\phi \rightarrow \gamma\gamma)}{\tau_\phi \gamma_\phi \beta_\phi} < 8.4 \text{ at 90\% CL}, \quad (\text{A11})$$

where $V_{\text{eff}}^{\text{Bor.}}$ is the Borexino active volume described as a sphere of radius $R = 3.02$ m.

6. Anomalous magnetic moment of the electron

The anomalous magnetic moment of the electron, $a_e \equiv (g_e - 2)/2$, has been commonly used to extract the

value of the fine-structure constant, α . However, recent improvements in atomic-physics experiments using Cesium (Cs) and Rubidium (Rb) interferometry have led to the following results for α :

$$\alpha(\text{Cs}) = 1/137.035999046(27), \quad (\text{A12})$$

$$\alpha(\text{Rb}) = 1/137.035999206(11), \quad (\text{A13})$$

showing a disagreement of 5.5σ [41,42]. Using the above determinations of α to predict the SM value a_e^{SM} and comparing it with the latest experimental measurement of $a_e^{\text{exp}} = (115965218059 \pm 13) \times 10^{-14}$ [43], yields the following values of $\Delta a_e \equiv a_e^{\text{exp}} - a_e^{\text{SM}}$:

$$(\Delta a_e)_{\text{Cs}} = (-102 \pm 26) \times 10^{-14}, \quad (\text{A14})$$

$$(\Delta a_e)_{\text{Rb}} = (34 \pm 16) \times 10^{-14}. \quad (\text{A15})$$

The contribution of the scalar ϕ , stemming from the electron coupling defined in Eq. (4.1), reads [44]

$$\begin{aligned} \Delta a_e^\phi &= \frac{g_e^2 m_e^2}{4\pi^2 m_\phi^2} \frac{1}{2} \int_0^1 dx \frac{x^2(2-x)}{1-x + \frac{m_e^2}{m_\phi^2} x^2} \\ &\approx \frac{g_e^2 m_e^2}{4\pi^2 m_\phi^2} \left(\ln \frac{m_\phi}{m_e} - \frac{7}{12} \right), \end{aligned} \quad (\text{A16})$$

where the approximation in the last step is valid for $m_\phi \gg m_e$. Using the latter expression, we obtain the following numerical estimate

$$\Delta a_e^\phi \approx 1.3 \times 10^{-13} \left(\frac{g_e}{10^{-5}} \right)^2 \left(\frac{2 \text{ MeV}}{m_\phi} \right)^2, \quad (\text{A17})$$

which, for $m_\phi = 2$ MeV, approximates the exact result at the 20% level. In our numerical analysis, however, we employ the full expression given in the first line of Eq. (A16).

Here, one could assume three benchmark scenarios: (i) $|\Delta a_e| \leq 10^{-12}$, where we inflated the current experimental errors on $a_e^{\alpha(\text{Cs})}$ and $a_e^{\alpha(\text{Rb})}$ to make Eqs. (A14) and (A15) consistent, (ii) $|\Delta a_e| \leq 10^{-13}$, assuming a resolution of the current discrepancy in the measurements of α with a precision of $\mathcal{O}(10^{-13})$, and (iii) $|\Delta a_e| \leq 10^{-14}$, which is the ultimate expected uncertainty on Δa_e if both the errors on a_e^{exp} and a_e^α will improve by roughly one order of magnitude [45]. In practice, we adopt scenario (ii) for our analysis, noting that the results for the other cases can be obtained by a straightforward rescaling.

- [1] E. Izaguirre, G. Krnjaic, and M. Pospelov, Probing new physics with underground accelerators and radioactive sources, *Phys. Lett. B* **740**, 61 (2015).
- [2] M. Pospelov and Y.-D. Tsai, Light scalars and dark photons in Borexino and LSND experiments, *Phys. Lett. B* **785**, 288 (2018).
- [3] A. Sen, G. Domínguez-Cañizares, N. Podaru, D. Mous, M. Junker, G. Imbriani, and V. Rigato, A high intensity, high stability 3.5 MV singletronTM accelerator, *Nucl. Instrum. Methods Phys. Res., Sect. B* **450**, 390 (2019).
- [4] <https://www.lngs.infn.it/en/pagine/bellotti-facility-en>
- [5] E. Aprile *et al.* (XENON Collaboration), The XENONnT dark matter experiment, *Eur. Phys. J. C* **84**, 784 (2024).
- [6] F. Acerbi *et al.* (DarkSide-20k Collaboration), DarkSide-20k sensitivity to light dark matter particles, *Commun. Phys.* **7**, 422 (2024).
- [7] D. Barducci and C. Toni, An updated view on the ATOMKI nuclear anomalies, *J. High Energy Phys.* **02** (2023) 154; **07** (2023) 168(E).
- [8] L. Resnick, M. K. Sundaresan, and P. J. S. Watson, Is there a light scalar boson?, *Phys. Rev. D* **8**, 172 (1973).
- [9] D. Barducci, D. Germani, M. Nardecchia, S. Scacco, and C. Toni, On the Atomki nuclear anomaly after the MEG-II result, *J. High Energy Phys.* **04** (2025) 035.
- [10] I. Lombardo, D. Dell'Aquila, J.-J. He, G. Spadaccini, and M. Vigilante, New analysis of $p + {}^{19}\text{F}$ reactions at low energies and the spectroscopy of natural-parity states in ${}^{20}\text{Ne}$, *Phys. Rev. C* **100**, 044307 (2019).
- [11] J. F. Ziegler, SRIM-2003, *Nucl. Instrum. Methods Phys. Res., Sect. B* **219**, 1027 (2004).
- [12] F. Feruglio and G. Levati (to be published).
- [13] D. Basilico *et al.* (BOREXINO Collaboration), Search for high energy 5.5 MeV solar axions with the complete Borexino dataset, *Eur. Phys. J. C* **85**, 1182 (2025).
- [14] D. Tucker-Smith and I. Yavin, Muonic hydrogen and MeV forces, *Phys. Rev. D* **83**, 101702 (2011).
- [15] C.-Y. Chen, H. Davoudiasl, W. J. Marciano, and C. Zhang, Implications of a light “dark Higgs” solution to the $g_\mu - 2$ discrepancy, *Phys. Rev. D* **93**, 035006 (2016).
- [16] B. Batell, N. Lange, D. McKeen, M. Pospelov, and A. Ritz, Muon anomalous magnetic moment through the leptonic Higgs portal, *Phys. Rev. D* **95**, 075003 (2017).
- [17] C. Delaunay, T. Kitahara, Y. Soreq, and J. Zupan, Light scalar beyond the Higgs mixing limit, [arXiv:2501.16477](https://arxiv.org/abs/2501.16477).
- [18] J. Fan, M. Reece, and L.-T. Wang, Non-relativistic effective theory of dark matter direct detection, *J. Cosmol. Astropart. Phys.* **11** (2010) 042.
- [19] J. R. Ellis, A. Ferstl, and K. A. Olive, Reevaluation of the elastic scattering of supersymmetric dark matter, *Phys. Lett. B* **481**, 304 (2000).
- [20] H.-Y. Cheng and C.-W. Chiang, Revisiting scalar and pseudoscalar couplings with nucleons, *J. High Energy Phys.* **07** (2012) 009.
- [21] E. Hardy and R. Lasenby, Stellar cooling bounds on new light particles: Plasma mixing effects, *J. High Energy Phys.* **02** (2017) 033.
- [22] S. Bottaro, A. Caputo, G. Raffelt, and E. Vitagliano, Stellar limits on scalars from electron-nucleus bremsstrahlung, *J. Cosmol. Astropart. Phys.* **07** (2023) 071.
- [23] E. Hardy, A. Sokolov, and H. Stubbs, Supernova bounds on new scalars from resonant and soft emission, *J. High Energy Phys.* **04** (2025) 013.
- [24] D. F. G. Fiorillo, A. Lella, C. A. J. O’Hare, and E. Vitagliano, Leading bounds on micro- to picometer fifth forces from neutron star cooling, [arXiv:2506.19906](https://arxiv.org/abs/2506.19906).
- [25] E. Cortina Gil *et al.* (NA62 Collaboration), Observation of the $K^+ \rightarrow \pi^+ \nu \bar{\nu}$ decay and measurement of its branching ratio, *J. High Energy Phys.* **02** (2025) 191.
- [26] D. Guadagnoli, A. Iohner, C. Lazzeroni, D. Martinez Santos, J. C. Swallow, and C. Toni, New bound on the vectorial axion-down-strange coupling from $K^+ \rightarrow \pi^+ \nu \bar{\nu}$ data, [arXiv:2503.05865](https://arxiv.org/abs/2503.05865).
- [27] E. Cortina Gil *et al.* (NA62 Collaboration), Searches for hidden sectors using $K^+ \rightarrow \pi^+ X$ decays, [arXiv:2507.17286](https://arxiv.org/abs/2507.17286).
- [28] C. Baruch, P. J. Fitzpatrick, T. Menzo, Y. Soreq, S. Trifinopoulos, and J. Zupan, Searching for exotic scalars at fusion reactors, [arXiv:2502.12314](https://arxiv.org/abs/2502.12314).
- [29] R. Barbieri and T. E. O. Ericson, Evidence against the existence of a low mass scalar boson from neutron-nucleus scattering, *Phys. Lett.* **57B**, 270 (1975).
- [30] J. Schmiedmayer, P. Riehs, J. A. Harvey, and N. W. Hill, Measurement of the electric polarizability of the neutron, *Phys. Rev. Lett.* **66**, 1015 (1991).
- [31] H. Leeb and J. Schmiedmayer, Constraint on hypothetical light interacting bosons from low-energy neutron experiments, *Phys. Rev. Lett.* **68**, 1472 (1992).
- [32] A. Aguilar *et al.* (LSND Collaboration), Evidence for neutrino oscillations from the observation of $\bar{\nu}_e$ appearance in a $\bar{\nu}_\mu$ beam, *Phys. Rev. D* **64**, 112007 (2001).
- [33] L. B. Auerbach *et al.* (LSND Collaboration), Measurement of electron—neutrino—electron elastic scattering, *Phys. Rev. D* **63**, 112001 (2001).
- [34] M. Agostini *et al.* (BOREXINO Collaboration), Comprehensive measurement of pp -chain solar neutrinos, *Nature (London)* **562**, 505 (2018).
- [35] J. N. Bahcall, M. H. Pinsonneault, and S. Basu, Solar models: Current epoch and time dependences, neutrinos, and helioseismological properties, *Astrophys. J.* **555**, 990 (2001).
- [36] G. Raffelt and L. Stodolsky, New particles from nuclear reactions in the sun, *Phys. Lett. B* **119**, 323 (1982).
- [37] U. Eichmann, Zum photoeffekt an den kernen der masse 3, *Z. Phys.* **175**, 115 (1963).
- [38] B. Aharmim *et al.* (SNO Collaboration), Combined analysis of all three phases of solar neutrino data from the sudbury neutrino observatory, *Phys. Rev. C* **88**, 025501 (2013).
- [39] A. Bhusal, N. Houston, and T. Li, Searching for solar axions using data from the sudbury neutrino observatory, *Phys. Rev. Lett.* **126**, 091601 (2021).
- [40] G. Bellini *et al.* (Borexino Collaboration), Search for solar axions produced in $p(d, {}^3\text{He})A$ reaction with Borexino detector, *Phys. Rev. D* **85**, 092003 (2012).
- [41] R. H. Parker, C. Yu, W. Zhong, B. Estey, and H. Müller, Measurement of the fine-structure constant as a test of the standard model, *Science* **360**, 191 (2018).

- [42] L. Morel, Z. Yao, P. Cladé, and S. Guellati-Khélifa, Determination of the fine-structure constant with an accuracy of 81 parts per trillion, *Nature (London)* **588**, 61 (2020).
- [43] X. Fan, T. G. Myers, B. A. D. Sukra, and G. Gabrielse, Measurement of the Electron magnetic moment, *Phys. Rev. Lett.* **130**, 071801 (2023).
- [44] F. Jegerlehner, *The Anomalous Magnetic Moment of the Muon*, Vol. 274 (Springer, Cham, 2017).
- [45] L. Di Luzio, A. Keshavarzi, A. Masiero, and P. Paradisi, Model-independent tests of the hadronic vacuum polarization contribution to the Muon $g-2$, *Phys. Rev. Lett.* **134**, 011902 (2025).