

New physics searches via beam normal spin asymmetry in Bhabha scattering

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We examine the sensitivity of the beam normal spin asymmetry in Bhabha scattering to beyond the Standard Model mediators, in the context of the JLab polarized positron program. A key property of this observable is that the Standard Model contribution exhibits a zero crossing at a fixed scattering angle, providing a clean, effectively background-free point for these searches. We consider scalar, vector, and axial vector mediators and present projected bounds, finding that scalar and vector scenarios allow a significant extension of the search ranges beyond existing constraints.

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I. INTRODUCTION

The search for light dark sector particles in the MeV to GeV mass range has intensified over the past decade [1–6], with a particular focus on axionlike particles [7–21] and dark photons [22–34]. The coverage across this mass range is primarily provided by e^+e^- colliders, beam-dump, and fixed-target experiments, mostly relying on “bump hunt” and missing energy searches. These have proven powerful but are not exhaustive. In particular, the few MeV to few-hundred MeV mediator-mass range remains relatively underexplored [15,17,20]. The detection of low-energy final states is challenged by large QED backgrounds and detector-threshold effects, which complicate reconstruction [11,18,21], motivating complementary low-energy precision programs such as PADME [35], the MAGIX@MESA facility [36], and the Jefferson Lab (JLab) polarized positron program [37].

The latter is the main subject of this work. While existing studies have placed constraints on BSM effects using unpolarized Bhabha scattering [38–40], the uniqueness of the JLab positron project stems from its high-intensity positron beam with a high degree of polarization. The interaction of this beam with unpolarized atomic electrons in a fixed target enables access to polarization-sensitive observables in Bhabha scattering $e^+e^- \rightarrow e^+e^-$. The beam normal spin asymmetry, which is proportional to an absorptive (imaginary) part of the scattering amplitude, is of particular interest.

The 11 GeV beam at JLab will allow access to the Bhabha process up to center-of-momentum (c.m.) energies $\sqrt{s} \simeq 106$ MeV. Being below any hadron production threshold this process is theoretically very clean as it cannot get contributions from any hadronic absorptive part. Moreover, the leading QED contribution to this asymmetry vanishes at a fixed scattering angle, yielding a kinematic point with strongly reduced SM background for BSM searches. Since QED interactions at energies much larger than the electron mass are helicity-conserving, particularly strong bounds can be derived for dark sector mediators which induce a lepton helicity flip, such as scalars. Furthermore, our findings show that the projected dark photon search range can also be extended significantly beyond the existing limits, demonstrating the complementarity and broad potential of e^+e^- polarization observables in light messenger particle searches.

The paper is organized as follows. In Sec. II, we outline the beam normal spin asymmetry in Bhabha scattering. Section III details the QED background to this observable, while in Sec. IV we investigate the BSM contributions for generic scalar, vector, and axial vector mediators. Section V presents the sensitivity projections for JLab kinematics. Finally, Sec. VI provides a summary of our work.

II. BEAM NORMAL SPIN ASYMMETRY B_n

In this work, we consider the Bhabha scattering process

$$e^-(p_1, s_1) + e^+(p_2, s_2) \rightarrow e^-(p_3, s_3) + e^+(p_4, s_4), \quad (1)$$

where p_1 and p_2 (p_3 and p_4) denote the initial (final) electron and positron four-momenta, and the s_i denote the corresponding helicities.

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Imposing parity (P), time-reversal (T), and charge-conjugation (C) invariance leads to five independent helicity amplitudes for this process [41]. For the purpose of exploring the contribution of P, T, and C conserving dark messenger particles to this process it is convenient to parametrize the amplitude $\mathcal{M}$ through the following covariants:

$$\begin{aligned}\mathcal{M} = & S\bar{u}_3u_1\bar{v}_2v_4 + P\bar{u}_3\gamma^5u_1\bar{v}_2\gamma^5v_4 \\ & + V\bar{u}_3\gamma_\mu u_1\bar{v}_2\gamma^\mu v_4 + A\bar{u}_3\gamma_\mu\gamma^5u_1\bar{v}_2\gamma^\mu\gamma^5v_4 \\ & + T\bar{u}_3\sigma_{\mu\nu}u_1\bar{v}_2\sigma^{\mu\nu}v_4,\end{aligned}\quad (2)$$

where for compactness we adopt the notation $u_i \equiv u(p_i, s_i)$ and $v_i \equiv v(p_i, s_i)$ for the Dirac spinors, and use $\sigma^{\mu\nu} = i/2[\gamma^\mu, \gamma^\nu]$ in terms of the Dirac matrices. In Eq. (2), the invariant amplitudes S , P , V , A , and T —corresponding to the scalar, pseudoscalar, vector, axial vector, and tensor structures, respectively—are functions of the Mandelstam variables, defined in the standard way

$$s = (p_1 + p_2)^2, \quad t = (p_1 - p_3)^2, \quad u = (p_1 - p_4)^2. \quad (3)$$

The parametrization (2) is arguably the most natural choice—particularly in the context of BSM searches—although it is not unique. We also note that while the decomposition above is written in terms of scattering-channel structures¹ $\bar{u}_3\Gamma_i u_1\bar{v}_2\Gamma_i v_4$, the annihilation channel structures $\bar{v}_2\Gamma_i u_1\bar{u}_3\Gamma_i v_4$ can be straightforwardly projected onto these using Fierz identities [42].

The single-spin asymmetry is defined through the beam-polarized cross sections $d\sigma_\uparrow$ and $d\sigma_\downarrow$, differential in the c.m. scattering angle θ , as

$$B_n = \frac{d\sigma_\uparrow - d\sigma_\downarrow}{d\sigma_\uparrow + d\sigma_\downarrow}. \quad (4)$$

It is expressed in terms of the invariants of Eq. (2) as

$$B_n = (S \cdot \hat{n}) \frac{m_e}{4\pi s} \frac{\sqrt{stu}}{\sigma_0} \text{Im}[(S - A)(V^* + 2T^*)], \quad (5)$$

with m_e the lepton mass, σ_0 the unpolarized cross section, S^μ the beam particle polarization four-vector, and

$$\hat{n}_\mu \equiv \frac{2}{\sqrt{stu}} \varepsilon_{\mu\alpha\beta\gamma} p_1^\alpha p_2^\beta p_3^\gamma, \quad (6)$$

where we adopt the choice $\varepsilon_{0123} = +1$ for the antisymmetric Levi-Civita tensor. In the c.m. frame the four-vector n^μ can be expressed as $\hat{n}^\mu = (0, \hat{\mathbf{n}})$, with $\hat{\mathbf{n}}$ the unit

¹The t -channel is chosen since, in the kinematic regime relevant to the JLab measurement, its contribution dominates across most scattering angles.

three-vector normal to the scattering plane. One thus sees that the factor $(S \cdot \hat{n})$ in Eq. (5) leads to a nonzero single spin asymmetry for a polarization normal to the scattering plane. In the following the asymmetry is shown using the sign convention

$$S \cdot \hat{n} = +1. \quad (7)$$

The normal beam spin asymmetry Eq. (5) requires the absorptive (imaginary) part of the amplitude to be nonzero. Furthermore, one notices that an amplitude containing a pseudoscalar exchange—the P term in Eq. (2)—does not contribute to B_n , as all resulting Dirac traces involving γ^5 identically vanish in the cross section difference $d\sigma_\uparrow - d\sigma_\downarrow$, given the three unpolarized leptons and the spin projection operator.

III. QED CONTRIBUTIONS TO B_n

We now discuss subsequently the QED contribution to B_n as well as the s -channel resonance excitation of a BSM dark sector messenger particle of different quantum numbers. The tree-level QED diagrams to the Bhabha process are shown in Figs. 1(a) and 1(b). Projected on the basis of Eq. (2) they yield the amplitudes,

$$\begin{aligned}V &= -e^2 \left(\frac{1}{t} + \frac{1}{2s} \right), \\ S &= -P = -2A = \frac{e^2}{s}, \quad T = 0,\end{aligned}\quad (8)$$

where e is the electric charge. As the tree-level QED diagrams do not contribute to the imaginary part of the Bhabha scattering amplitude, B_n is zero at tree-level in QED. The leading nonvanishing contribution to the observable arises from the term $\propto 2 \text{Im}[\mathcal{M}_{\text{tree}} \mathcal{M}_{1\text{-loop}}^*]$, involving the tree-level amplitude and 1-loop corrections. The relevant 1-loop diagrams generating these amplitudes are

- (i) s -channel vacuum polarization with intermediate e^+e^- loop, Fig. 1(c),
- (ii) s -channel vertex corrections on both the initial- and final-state fermion lines combined with the t -channel two-photon box diagrams, for which only the direct graph yields a nonzero imaginary part, Figs. 1(d)–1(f),
- (iii) s -channel direct and crossed two-photon box diagrams, Figs. 1(g) and 1(h).

All other 1-loop diagrams have vanishing discontinuities and therefore do not contribute to B_n . A number of cancellations were found among the individual contributions. The vertex correction diagrams and t -channel box diagrams of Figs. 1(d)–1(f) require an infrared regulator, which cancels out when calculating their contribution to B_n , rendering it an infrared-safe observable. Furthermore, any logarithmic dependence on the electron mass cancels out in the final result.

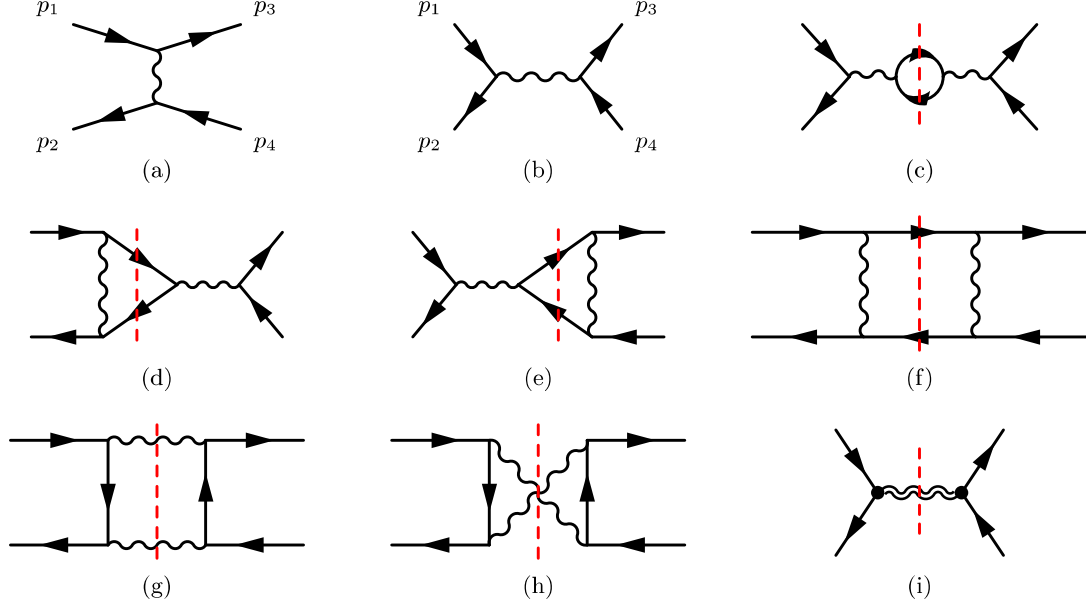


FIG. 1. Feynman diagrams contributing to B_n at one-loop order (a)–(h), and the BSM contribution (i). The red vertical line shows the physical-region cuts.

We calculated the relevant 1-loop QED diagrams and projected them onto the Dirac basis in Eq. (2). Our result for B_n obtained from Eq. (5) verifies the expression given in the literature [43]. In the ultrarelativistic limit the resulting QED expression at 1-loop level can be written in the compact form,

$$B_n^{\text{QED,1-loop}} = \alpha \frac{m_e}{\sqrt{s}} \frac{x}{2\sqrt{1-x^2}(1-x^2+x^4)^2} \times \left[2x^2(1-x^2)(-5+x^4) - 3x^2(1+x^2)\ln(x^2) + 3(1-x^2)(1-2x^2)\ln(1-x^2) + \mathcal{O}\left(\frac{m_e^2}{s}\right) \right], \quad (9)$$

with $\alpha \equiv e^2/(4\pi) \simeq 1/137$, and $x \equiv \sin(\theta/2)$.

The QED result for B_n , shown on Fig. 2, exhibits a zero crossing at $\theta \approx 120.4^\circ$ due to cancellations between different discontinuities. In the ultrarelativistic regime, where the only kinematic dependence appears in the overall prefactor, the position of this zero is independent of $\sqrt{s}$, thus providing a kinematic point that is effectively background free.

While two-loop effects or finite-mass corrections could in principle shift this zero, these contributions are expected to lie well below the anticipated experimental precision for B_n , estimated at ~ 1 ppm [44]. Furthermore, a real experiment will measure over a finite angular bin $\Delta\theta$, expected to be 10° or smaller. Integrating the smoothly varying SM distribution over a symmetric bin centered on the zero-crossing cancels the SM background to first order. Other potential backgrounds from positron-nucleus elastic

scattering $e^+N \rightarrow e^+N$ can be effectively eliminated by utilizing the distinct kinematics of the Bhabha process, specifically by requiring a coincidence detection of the final state e^+e^- pair.

IV. BSM CONTRIBUTIONS TO B_n

We next discuss the sensitivity of B_n to BSM messengers. As the exchange of a pseudoscalar particle provides no contribution to B_n , seen from Eq. (5), three BSM mediators with minimal couplings to SM leptons are considered below: a scalar, a vector, and an axial vector. Their interaction terms are

$$\begin{aligned} \mathcal{L}_S &= -g\tilde{S}\bar{\ell}\ell, \\ \mathcal{L}_V &= -e\bar{\ell}\gamma^\mu\ell V_\mu, \\ \mathcal{L}_A &= -e\bar{\ell}\gamma^\mu\gamma^5\ell A_\mu, \end{aligned} \quad (10)$$

where $\tilde{S}$ is a scalar field coupled to SM leptons with strength g , and V^μ (A^μ) denote the vector (axial vector) mediator fields respectively, with e due to kinetic mixing. The considered BSM contribution to B_n arises from the term $\propto 2\text{Im}[\mathcal{M}_{\text{tree}}\mathcal{M}_{\text{BSM}}^*]$, involving the tree-level QED amplitude [Figs. 1(a) and 1(b)] and the s -channel tree-level BSM amplitude [Fig. 1(i)], with the imaginary part resulting from the propagator

$$\frac{1}{s - m_X^2 + im_X\Gamma_X}, \quad (11)$$

where m_X and Γ_X denote the mediator mass and total decay width. In the kinematic region of interest, the only

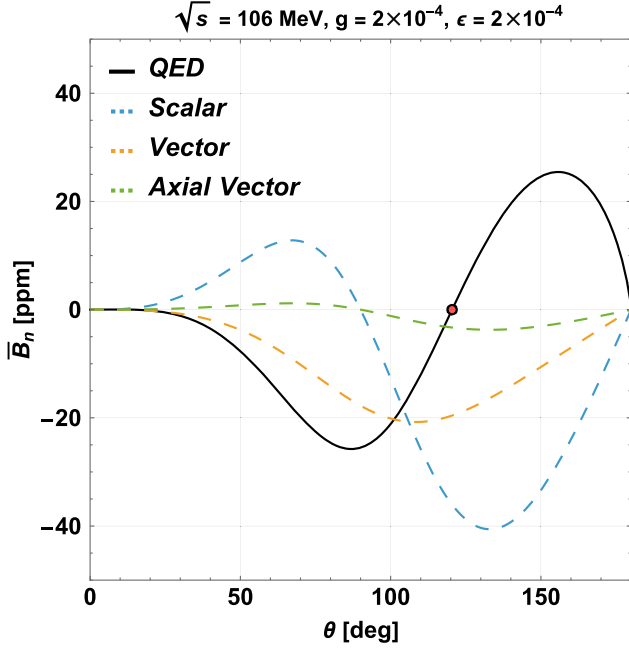


FIG. 2. Comparison of the angular dependence of the QED and BSM contributions to the Bhabha beam normal spin asymmetry $\bar{B}_n$ for indicated values of BSM coupling strengths, for an 11 GeV e^+ beam at JLab. The numerator and denominator in $\bar{B}_n$ are integrated over a bin in s of size $2m_e\delta E_+$, with $\delta E_+ = 0.5$ MeV. The QED zero-crossing point is marked with a red dot.

accessible final states are e^+e^- and $\gamma\gamma$. While the Landau-Yang theorem forbids massive spin-1 particles from decaying into two photons, this channel remains open for scalars. However, the corresponding decay width scales as m_S^3 , in contrast to the leptonic channel, which scales linearly with m_S , as discussed in [21]. In the relatively low-mass range studied here, the diphoton mode is therefore strongly suppressed. Consequently, in the minimal model considered here the total width consists of a single contribution from the mediator's coupling to electrons.

Existing constraints imply that Γ_X is a small quantity—well below the experimental resolution of the invariant mass of the decay products—so the narrow width approximation for the imaginary part amounts to the replacement

$$\frac{-im_X\Gamma_X}{(s-m_X^2)^2+m_X^2\Gamma_X^2} \rightarrow -i\pi\delta(s-m_X^2). \quad (12)$$

To determine the sensitivity reach of the dark messenger coupling as a function of its mass m_X , the differential cross sections in the numerator and denominator of B_n must be integrated over a bin in the e^+e^- invariant mass around a given central value s_0 . For a smoothly varying function, such as the QED cross section, this amounts to the factor

$$\int_{s_0-\frac{\delta s}{2}}^{s_0+\frac{\delta s}{2}} f(s)ds = 2m_e\delta E_+ f(s_0), \quad (13)$$

with δE_+ the bin width in the positron (*lab*) beam energy. To remain conservative, we adopt a bin width of $\delta E_+ = 0.5$ MeV. For the QED contribution, this factor cancels out between numerator and denominator, but it becomes relevant for the BSM-induced result. Denoting the asymmetry in which numerator and denominator are integrated over a bin in s of size $2m_e\delta E_+$ as $\bar{B}_n$, we find that this results in

$$\bar{B}_n^{\text{BSM}} = \frac{\sqrt{s}}{\delta E_+} \frac{x\sqrt{1-x^2}}{8} \frac{N_i}{(1-x^2+x^4)^2}, \quad (14)$$

where the numerator function N_i for scalar (N_S), vector (N_V), and axial vector (N_A) exchange is given by,

$$N_S = \frac{g^2}{\alpha} x^2(1-2x^2), \quad (15)$$

$$N_V = -4\pi\epsilon^2 3x^2, \quad (16)$$

$$N_A = 4\pi\epsilon^2 x^2(1-2x^2). \quad (17)$$

Note that in the ultrarelativistic regime the scalar and axial vector contributions have the same angular (x) dependence up to the coupling constant. In contrast to the QED contribution, which decreases with energy as $1/\sqrt{s}$, the BSM-induced spin asymmetry grows proportionally to $\sqrt{s}$, improving the signal-to-background ratio at higher beam energies. Furthermore, one notices from Eqs. (15)–(17) that the observable B_n depends quadratically on the BSM couplings, as it results from the product of a large tree-level QED amplitude and the absorptive part of a BSM amplitude. This is in contrast to BSM searches in unpolarized observables, for which the BSM signal goes with the fourth power of the small coupling constant. One thus sees that the tree-level QED process serves as an amplifier of a possible BSM signal.

The result for the angular dependence of B_n is shown in Fig. 2 for an 11 GeV positron beam at JLab. The coupling strengths $g = 2 \times 10^{-4}$ and $\epsilon = 2 \times 10^{-4}$ are chosen for illustration. Even for such small values, one sees that the BSM effects in the vicinity of the QED zero crossing can be sizable compared to the expected ppm-level experimental precision, enabling enhanced sensitivity or, in the absence of a signal, competitive limits to be set.

V. PROJECTED BOUNDS FOR JLAB KINEMATICS

We also made an assessment of the parameter space that can be probed at JLab, where the scan over the range of m_X can be performed with the initial state radiation technique. The resulting limits, obtained by requiring the BSM contribution to remain within the projected experimental uncertainty of 1 ppm, are shown in Fig. 3, together with existing bounds from $(g-2)_e$ measurements [45,46] and

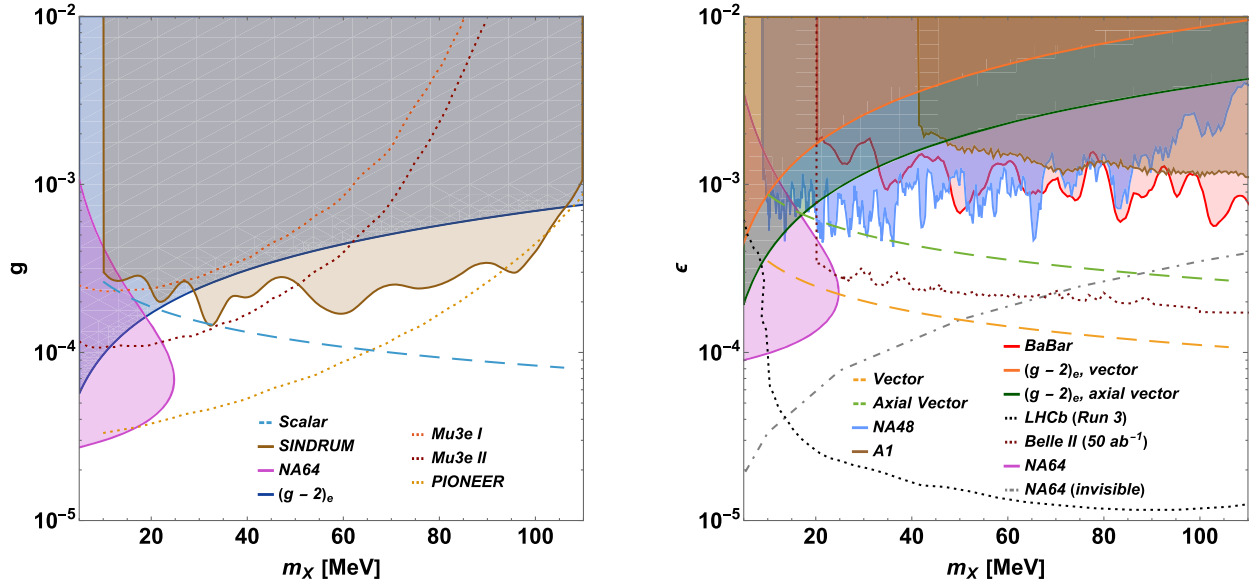


FIG. 3. Constraints on the allowed BSM parameter space from $(g-2)_e$ and from existing experiments searching for visible e^+e^- decays (shaded exclusion regions), shown for the scalar (left) and the vector/axial vector scenarios (right). The projected sensitivity of a future JLab polarized-positron experiment, assuming a 1 ppm precision on B_n , is indicated by dashed lines, while other near-future experiments are shown as dotted lines and are described in the text. For comparison, the gray dot-dashed curve shows limits from NA64 searches for invisibly decaying mediators.

visible decay constraints on dark photons from A1@MAMI [23], BABAR [47], NA48/2 studies of π^0 decays [24], and the NA64 analyses [48]. We also include the projected sensitivity from the Belle II analysis [49], as well as the LHCb Run 3 projection based on studies of dark photon radiation from charm mesons [50].

The NA64 bound for visible decaying mediators was also translated to the scalar case. Additionally, we incorporate projections for the Mu3e experiment based on the $\mu^+ \rightarrow e^+ \bar{\nu}_\mu \nu_e \tilde{S}$ decay, assuming 2.5×10^{15} and 5.5×10^{16} produced μ^+ , respectively (labeled I and II in Fig. 3), adapted from [51]. We further reproduced limits from the SINDRUM experiment [52], which searched for the $\pi^+ \rightarrow e^+ \nu \tilde{S}$ process, alongside the projected sensitivity of PIONEER, which aims to improve the branching ratio limit on this decay to $\mathcal{O}(10^{-11})$. In both aforementioned processes, the scalar $\tilde{S}$ is emitted from the final-state positron and subsequently decays as $\tilde{S} \rightarrow e^+e^-$.

Our results indicate that up to an order of magnitude extension in reach, as compared to current bounds, can be achieved for the scalar and vector scenarios for masses around 100 MeV without additional model assumptions. The ultimate reach, however, will depend on the precision achieved at JLab.

We also emphasize that while recent missing-energy searches at NA64 [34] have placed severe constraints on the parameter space of vector mediators decaying invisibly into dark sector states, these bounds are inherently model-dependent. The NA64 invisible limit in Fig. 3 (gray dot-dashed line) assumes a dark coupling $\alpha_D = 0.1$ and

dark fermion masses $m_\chi = m'_A/3$. These constraints do not apply to the minimal scenario considered in this work, where the mediator decays with a 100% branching ratio into visible e^+e^- pairs; however, they show the complementarity of the two search strategies.

VI. CONCLUSION

In conclusion, we have shown that the Bhabha beam normal spin asymmetry, which can be accessed using a polarized positron beam at JLab, provides a uniquely sensitive probe of BSM mediators. The uniqueness is threefold: first the JLab energy region, which corresponds with a maximum c.m. energy of 106 MeV of the Bhabha process prevents any hadronic discontinuity to contribute. Second, the leading QED result has a zero in the angular distribution due to cancellations between different discontinuities, allowing for a kinematic point that is nearly background-free. Third, the enhanced sensitivity of the observable B_n to a BSM contribution results from the quadratic dependence on the coupling, in contrast to BSM searches in unpolarized processes, for which the BSM signal goes with the fourth power of the small coupling constant.

We have shown that up to an order of magnitude extension in reach can be obtained, in particular for the scalar and vector scenarios. Our estimates were done in a conservative way, and additional improvements in energy resolution or experimental precision could strengthen the final results. While the analysis presented in this work focused on the most natural scenarios involving scalar, vector, and axial vector mediators, the method can be

straightforwardly extended to a broader class of models, including tensor-coupled interactions [53].

In this work we assumed a minimal model with a 100% branching ratio for the messenger decay into electron-positron pairs. If invisible decays into dark sector states are considered possible, the proposed framework remains fully applicable. However, since a larger total width Γ_X suppresses the visible signal, the resulting constraints must be reinterpreted as bounds on the effective coupling scaled by the branching ratio, rather than the electron coupling alone, provided the total width Γ_X remains smaller than the experimental invariant mass resolution.

Finally, also exploring the potential for BSM searches in double-polarization observables of the Bhabha process may be of interest for a future study.

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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.

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- [1] J. Beacham *et al.*, *J. Phys. G* **47**, 010501 (2020).
 - [2] G. Lanfranchi, M. Pospelov, and P. Schuster, *Annu. Rev. Nucl. Part. Sci.* **71**, 279 (2021).
 - [3] P. Agrawal *et al.*, *Eur. Phys. J. C* **81**, 1015 (2021).
 - [4] B. Batell, N. Blinov, C. Hearty, and R. McGehee, in *Snowmass 2021* (2022).
 - [5] C. Antel *et al.*, *Eur. Phys. J. C* **83**, 1122 (2023).
 - [6] S. Balan *et al.*, *J. Cosmol. Astropart. Phys.* **01** (2025) 053.
 - [7] J. Jaeckel and M. Spannowsky, *Phys. Lett. B* **753**, 482 (2016).
 - [8] E. Izaguirre, T. Lin, and B. Shuve, *Phys. Rev. Lett.* **118**, 111802 (2017).
 - [9] S. Knapen, T. Lin, H. K. Lou, and T. Melia, *Phys. Rev. Lett.* **118**, 171801 (2017).
 - [10] M. Bauer, M. Neubert, and A. Thamm, *J. High Energy Phys.* **12** (2017) 044.
 - [11] M. J. Dolan, T. Ferber, C. Hearty, F. Kahlhoefer, and K. Schmidt-Hoberg, *J. High Energy Phys.* **12** (2017) 094; **03** (2021) 190(E).
 - [12] L. Merlo, F. Pöbe, S. Rigolin, and O. Sumensari, *J. High Energy Phys.* **06** (2019) 091.
 - [13] D. Banerjee *et al.* (NA64 Collaboration), *Phys. Rev. Lett.* **125**, 081801 (2020).
 - [14] G. Aad *et al.* (ATLAS Collaboration), *J. High Energy Phys.* **03** (2021) 243; **11** (2021) 50.
 - [15] F. Abudinén *et al.* (Belle-II Collaboration), *Phys. Rev. Lett.* **125**, 161806 (2020).
 - [16] J. P. Lees *et al.* (BABAR Collaboration), *Phys. Rev. Lett.* **128**, 131802 (2022).
 - [17] M. Ablikim *et al.* (BESIII Collaboration), *Phys. Lett. B* **838**, 137698 (2023).
 - [18] A. Pustyntsev and M. Vanderhaeghen, *Eur. Phys. J. C* **84**, 546 (2024).
 - [19] P. R. Teles, D. d’Enterria, V. P. Gonçalves, and D. E. Martins, *Phys. Rev. D* **109**, 055003 (2024).
 - [20] M. Ablikim *et al.* (BESIII Collaboration), *Phys. Rev. D* **110**, L031101 (2024).
 - [21] A. Pustyntsev and M. Vanderhaeghen, *Phys. Rev. D* **110**, 075027 (2024).
 - [22] S. Abrahamyan *et al.* (APEX Collaboration), *Phys. Rev. Lett.* **107**, 191804 (2011).
 - [23] H. Merkel *et al.*, *Phys. Rev. Lett.* **112**, 221802 (2014).
 - [24] J. R. Batley *et al.* (NA48/2 Collaboration), *Phys. Lett. B* **746**, 178 (2015).
 - [25] A. Anastasi *et al.* (KLOE-2 Collaboration), *Phys. Lett. B* **757**, 356 (2016).
 - [26] J. H. Chang, R. Essig, and S. D. McDermott, *J. High Energy Phys.* **01** (2017) 107.
 - [27] J. L. Feng and J. Smolinsky, *Phys. Rev. D* **96**, 095022 (2017).
 - [28] P. H. Adrian *et al.* (HPS Collaboration), *Phys. Rev. D* **98**, 091101 (2018).
 - [29] R. Aaij *et al.* (LHCb Collaboration), *Phys. Rev. Lett.* **124**, 041801 (2020).
 - [30] P. H. Adrian *et al.*, *Phys. Rev. D* **108**, 012015 (2023).
 - [31] M. Ablikim *et al.* (BESIII Collaboration), *Phys. Lett. B* **839**, 137785 (2023).
 - [32] E. Cortina Gil *et al.* (NA62 Collaboration), *J. High Energy Phys.* **09** (2023) 035.
 - [33] H. Abreu *et al.* (FASER Collaboration), *Phys. Lett. B* **848**, 138378 (2024).
 - [34] Y. M. Andreev *et al.* (NA64 Collaboration), *Phys. Rev. Lett.* **131**, 161801 (2023).

- [35] F. Bossi *et al.* (PADME Collaboration), *J. High Energy Phys.* **11** (2025) 007.
- [36] S. Schlimme *et al.*, *EPJ Web Conf.* **303**, 06002 (2024).
- [37] A. Accardi *et al.*, *Eur. Phys. J. A* **57**, 261 (2021).
- [38] D. Bourilkov, *J. High Energy Phys.* **08** (1999) 006.
- [39] D. Bourilkov, *Phys. Rev. D* **62**, 076005 (2000).
- [40] M. Acciarri *et al.* (L3 Collaboration), *Phys. Lett. B* **489**, 81 (2000).
- [41] M. Goldberger, Y. Nambu, and R. Oehme, *Ann. Phys. (N.Y.)* **2**, 226 (1957).
- [42] M. E. Peskin and D. V. Schroeder, *An Introduction to Quantum Field Theory* (Addison-Wesley, Reading, MA, 1995).
- [43] C. Fronsdal and B. Jakšić, *Phys. Rev.* **121**, 916 (1961).
- [44] D. Mack, Searches for BSM interference terms in unpolarized and polarized Bhabha scattering, Talk at the “*Hadron Physics 2030*” workshop (2024), <https://indico.ijclab.in2p3.fr/event/10641/>.
- [45] X. Fan, T. G. Myers, B. A. D. Sukra, and G. Gabrielse, *Phys. Rev. Lett.* **130**, 071801 (2023).
- [46] A. Pustyntsev and M. Vanderhaeghen, *Phys. Rev. D* **112**, 095001 (2025).
- [47] J. P. Lees *et al.* (BABAR Collaboration), *Phys. Rev. Lett.* **113**, 201801 (2014).
- [48] D. Banerjee *et al.* (NA64 Collaboration), *Phys. Rev. D* **101**, 071101 (2020).
- [49] T. Ferber, *Acta Phys. Pol. B* **46**, 2285 (2015).
- [50] P. Ilten, J. Thaler, M. Williams, and W. Xue, *Phys. Rev. D* **92**, 115017 (2015).
- [51] L. Di Luzio, P. Paradisi, and N. Selimovic, *Nucl. Phys. B* **1021**, 117177 (2025).
- [52] S. Egli *et al.* (SINDRUM Collaboration), *Phys. Lett. B* **222**, 533 (1989).
- [53] V. Kozhuharov and M. Naydenov, *Nucl. Phys. B* **986**, 116044 (2023).