Single arm interferometry to probe the scalar field dark matter

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We analyze the interaction of photons with a scalar dark matter field ϕ, and we propose to use a single arm interferometer to reveal this interaction and constrain the parameters of the scalar dark matter model. By considering a beam of coherent light and two spatially separated squeezing operations, we show that the interaction of photons with scalar dark matter leads to an observable deviation in the outgoing light state, with respect to free evolution. Therefore the single-arm interferometer may yield a novel revelation method for scalar dark matter.

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I. INTRODUCTION

The nature of the dark sector of the universe is one of the main puzzles in modern cosmology. Dark energy is held responsible for the observed accelerated expansion of the universe, while dark matter is required to explain the gravitational binding of galaxies and galaxy clusters, as well as the rotation curve of spiral galaxies [1–8]. Prominent dark matter candidates are axions and axion-like particles [9–33], supersymmetric particles [34–50], ultralight scalar particles, known as scalar dark matter (SDM) [51–61], and fuzzy cold dark matter [62]. Condensates from quantum field theory have also been suggested as possible contributions to the dark matter [63–69]. In this work, we focus our attention on SDM. Several techniques have been proposed to test the SDM hypothesis. These include the use of highly sensitive optical cavities [70], gravitational wave detectors [71–73], and analyses of astrophysical phenomena that could provide indirect evidence for its existence [74,75]. Constraints on SDM have been obtained in various experiments such as DAMA [76,77] and precision atomic and molecular spectroscopy [78–82].

Here, we present a novel approach to test the SDM model by means of single-arm laser interferometry [83]. The method is aimed at probing the phase shift and the intensity variation experienced by a laser beam due to the interaction with the SDM. By considering laser light passing through two spatially separated squeezing operations, we show that the interaction between photons and scalar dark matter causes a detectable alteration in the outgoing laser state compared to its free evolution. This interaction cannot be revealed in standard interferometry experiments, since it affects the evolution of photons in the two arms equally and, therefore, no effect can be detected by phase difference analyses. In the following work, we provide a discussion of the interaction between SDM and photons, and then we introduce the interferometric setup and compute the phase shift and the intensity variation due to SDM-photon interaction. By using single-arm interferometers with an effective long baseline, similar to the one of the gravitational wave detector arm cavities, it is possible to verify the existence of the coupling between SDM and photons and, consequently, test the presence of scalar dark matter.

In Sec. II, we introduce the interaction between scalar dark matter and photons and examine its impact on the evolution of laser light. In Sec. III, we demonstrate how this interaction can be probed by means of a coherent light beam and two spatially separated squeezing operations within a single-arm interferometric setup. Numerical results are presented in Sec. IV, followed by a discussion of the implications and our conclusions in Sec. V.

II. SCALAR DARK MATTER AND PHOTON INTERACTION

The evolution of the scalar field ϕ is described by the Klein-Gordon equation $(\Box + m_\phi^2)\phi = 0$, where m_ϕ

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represents the mass of the field. A key aspect of SDM's phenomenology is its possible interaction with the photon field. This interaction can be described by an effective coupling term in the Lagrangian density[1] $\mathcal{L}_{\rm int} = \frac{1}{4}\frac{\phi}{\Lambda_\gamma} F_{\mu\nu}F^{\mu\nu}$ [84,85], where $F_{\mu\nu}$ is the electromagnetic field tensor, and Λ_γ parametrizes the interaction strength.[2] The interaction term affects the propagation of photons through regions permeated by SDM. The field ϕ is assumed to have an extremely small mass ($\ll 1$ eV) [51]; then, its Compton wavelength, $\lambda_c = \frac{h}{mc}$, is very large, and the field ϕ can be approximated as a coherently oscillating classical field: $\phi(t) \approx \phi_0 \cos(\omega t + \tilde{\theta})$. Here, $\tilde{\theta}$ is a random phase factor, $\phi_0 = \frac{1}{m_\phi}\sqrt{2\rho_{\rm local}}$ is the amplitude of the field, with $\rho_{\rm local}$ being local dark matter density, m_ϕ is the mass of the scalar particle, and ω is the oscillation frequency, which is of the order of the mass of the SDM, $\omega \simeq m_\phi$ [72,74,86]. We consider natural units.

We impose the radiation gauge, where the electromagnetic potential $A^\mu = (A^0, \vec{A})$ satisfies the condition $\overrightarrow{\nabla}\cdot\vec{A} = 0$. Then, the Hamiltonian corresponding to the above Lagrangian density $\mathcal{L}_{\rm int}$ is[3]

$$H_{\rm int}(t) = -\frac{\phi(t)}{2\Lambda_\gamma}\int d^3x \left[(\dot{\vec{A}})^2 + (\overrightarrow{\nabla}\times\vec{A})^2\right],$$

Upon quantization, the classical fields are replaced by the corresponding field operators, and the interaction Hamiltonian becomes an operator written in terms of the photon creation and annihilation operators. Normal ordering is then applied to eliminate the vacuum expectation value of the field energy. The resulting interaction Hamiltonian operator reads

$$\hat{H}_{\rm int}(t) = -\frac{\phi(t)}{\Lambda_\gamma}\int d^3k \sum_\lambda |\vec{k}|\hat{N}_{\vec{k}\lambda}, \tag{1}$$

where $\hat{N}_{\vec{k}\lambda} = \hat{a}^\dagger_{\vec{k}\lambda}\hat{a}_{\vec{k}\lambda}$ is the number operator of photons with momentum $\vec{k}$ and polarization λ, and $\hat{a}_{\vec{k}\lambda}$ is the annihilator. States with a definite number of photons are eigenstates of the Hamiltonian (1). Consequently, if the initial photon state is $|n_{\vec{k}\lambda}\rangle = \frac{1}{\sqrt{n!}}(\hat{a}^\dagger_{\vec{k}\lambda})^n|0\rangle$, the evolved state at time t, $|n_{\vec{k}\lambda}(t)\rangle$, in the interaction picture is

$$|n_{\vec{k}\lambda}(t)\rangle = e^{-i\int_0^t \hat{H}_{\rm int}(t')dt'}|n_{\vec{k}\lambda}\rangle = e^{i\hat{N}_{\vec{k}\lambda}\delta_k(t)}|n_{\vec{k}\lambda}\rangle, \tag{2}$$

with phase shift $\delta_k(t)$ given by

$$\delta_k(t) = \frac{|\vec{k}|}{\Lambda_\gamma}\int_0^t \phi(t')dt' = \frac{|\vec{k}|\phi_0}{\Lambda_\gamma\omega}\sin(\omega t). \tag{3}$$

Note that the phase shift is independent on the polarization λ. In the following section, we shall consider polarized (fixed λ) and monochromatic (fixed k) light, and we will omit the indices k, λ. To study the evolution of laser light, we analyze a coherent state $|\alpha\rangle = e^{-\frac{|\alpha|^2}{2}}\sum_{n=0}^\infty \frac{\alpha^n}{\sqrt{n!}}|n\rangle$. For coherent states, one has $e^{i\hat{N}\delta}|\alpha\rangle = e^{-\frac{|\alpha|^2}{2}}\sum_{n=0}^\infty \frac{\alpha^n e^{in\delta}}{\sqrt{n!}}|n\rangle = |\alpha e^{i\delta}\rangle$, i.e., $e^{i\hat{N}\delta}$ rotates the complex amplitude α by an angle δ in the complex plane.

III. PHASE MEASUREMENT SCHEME

The phase shift in Eq. (3) cannot be observed using a double-arm interferometer like the Michelson-Morley type since the interaction Lagrangian $\mathcal{L}_{\rm int}$ influences both arms in the same way, leading to the cancellation of the phase shifts [83]. However, a single-arm interferometer like the one described in Refs. [87–89], can be utilized for our purposes. The single arm interferometer can be schematized as follows. First, a squeezing operation $\hat{S}$ is performed on the laser beam. The beam then evolves along a path of length L, in a vacuum, subject to the interaction with SDM. Later, an antisqueezing operation $\hat{S}^{-1}$ is applied, and the laser light is detected, see Fig. 1. If the beam does not undergo an interaction with the scalar fields, the initial state of the photons coincides, up to a global phase, with the final state since the beam is subjected to the action of $\hat{S}$ and $\hat{S}^{-1}$. If, however, in the path L the photons interact with the scalar field, the final state will differ from the initial one. Thus, an eventual revelation of the difference between the

[1]The total interaction Lagrangian when the electron field is also present is given by $\mathcal{L}_{\rm int} = \frac{1}{4}\frac{\phi}{\Lambda_\gamma}F_{\mu\nu}F^{\mu\nu} - \frac{\phi}{\Lambda_e}\bar{\psi}_e\psi_e$. Since we consider the propagation of the photon in a vacuum, the last term is irrelevant, at least to the lowest order.

[2]We point out that in classical mechanics on electromagnetic field configurations for which $F^2 = 0$, neither the photon nor the scalar field are affected by this interaction. On the contrary, in quantum mechanics both the Hamiltonian and F^2 are operators that do not have a definite value, except when evaluated on specific states. This, in turn, means that the relation holding classically for plane wave $F^2 = 0$ does not hold as an operator identity $\hat{F}^2 \neq 0$, as can be immediately verified by inserting the photon field expansion within the photon Lagrangian. It follows that quantum mechanically both the photon and the scalar fields are affected by the interaction term $\mathcal{L}_{\rm int} = \frac{1}{4}\frac{\phi}{\Lambda_\gamma}F_{\mu\nu}F^{\mu\nu}$.

[3]Notice that the interaction Lagrangian, $\mathcal{L}_{\rm int}$, can be written in terms of the free electromagnetic Lagrangian, $\mathcal{L}_0 = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} = \frac{1}{2}(\vec{E}^2 - \vec{B}^2)$, as $\mathcal{L}_{\rm int} = -\frac{\phi(t)}{\Lambda_\gamma}\mathcal{L}_0$. As a consequence, $\mathcal{H}_{\rm int} = -\frac{\phi(t)}{\Lambda_\gamma}\mathcal{H}_0$, where $\mathcal{H}_0$ is the free electromagnetic Hamiltonian. Since $\mathcal{H}_0 = \frac{1}{2}(\vec{E}^2 + \vec{B}^2)$, then $\mathcal{H}_{\rm int} \neq 0$ even when $F^2 = 0$. This consideration is independent of the classical or quantum mechanical nature of the electromagnetic field. However, in our computation we consider the quantized Hamiltonian which, *a fortiori*, is different from zero (see also footnote 2). Since in the following, we will analyze the quantum dynamics of the photons, it is necessary to employ the quantum mechanical Hamiltonian operator [see Eq. (1)].

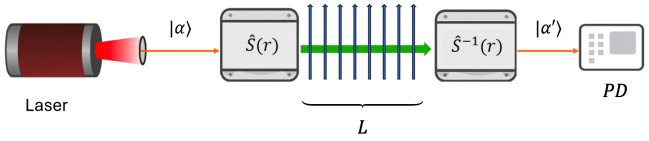


FIG. 1. Schematic depiction of the single-arm interferometer. A coherent laser beam, represented by the quantum state $|\alpha\rangle$, is emitted from the laser source. The beam first undergoes a squeezing operation $\hat{S}(r)$. This is followed by an inverse squeezing (antisqueezing) operation $\hat{S}^{-1}(r)$, which gives the state $|\alpha'\rangle$, which is then probed by a photodetector (PD).

initial and final states could signal the presence of SDM interacting with photons. After the squeezing operation, the squeezed state that we employ are bright, that is, carrying a large average number of photons, and then distinct from the squeezed vacuum. For such states, standard read/out techniques, radio frequency (RF) modulation/demodulation or DC readout techniques can be used to stabilize the phase, as reported in [90–92]. We point out that the coherent control loop will not be affected by SDM. In particular, we consider an initial coherent beam of light $|\alpha\rangle$; it satisfies the condition $\hat{a}|\alpha\rangle = \alpha|\alpha\rangle$, where $\alpha = |\alpha|e^{i\theta}$, and the average photon number is given by $\langle\alpha|\hat{a}^\dagger\hat{a}|\alpha\rangle = |\alpha|^2 = N$. According to Eq. (2), the state $|\alpha\rangle$ at time t is $e^{i\hat{N}\delta(t)}|\alpha\rangle$. Then, at any time t, the average photon number is constant and indeed $\langle\alpha(t)|\hat{a}^\dagger\hat{a}|\alpha(t)\rangle = \langle\alpha|e^{-i\hat{N}\delta(t)}\hat{a}^\dagger\hat{a}e^{i\hat{N}\delta(t)}|\alpha\rangle = |\alpha|^2$. Therefore, the presence of the interaction in Eq. (1) cannot be investigated by analyzing the simple time evolution of the number operator. In the system of Fig. 1, the squeezing operator $\hat{S}$ is $\hat{S}(r) = \exp[\frac{r}{2}(\hat{a}^{\dagger 2} - \hat{a}^2)]$, with r being the squeezing parameter. The squeezed state, obtained by the application of $\hat{S}$, is then $|\Phi(r)\rangle = \hat{S}(r)|\alpha\rangle$. With reference to the device sketched in Fig. 1, the emergent state $|\alpha'\rangle$ in absence of an interaction would be $|\alpha'\rangle_{\delta=0} = \hat{S}^{-1}(r)\hat{S}(r)|\alpha\rangle = |\alpha\rangle$. On the contrary, in the presence of the interaction (1), $|\alpha'\rangle$ experiences a phase shift: $|\alpha'\rangle = \hat{S}^{-1}(r)e^{i\delta\hat{a}^\dagger\hat{a}}\hat{S}(r)|\alpha\rangle$, where δ is the phase of Eq. (6). From the squeezing transformations[4]

$$\hat{S}^{-1}(r)\hat{a}\,\hat{S}(r) = \hat{a}\cosh r + \hat{a}^\dagger \sinh r,$$
$$\hat{S}^{-1}(r)\hat{a}^\dagger\hat{S}(r) = \hat{a}^\dagger\cosh r + \hat{a}\sinh r. \tag{4}$$

Here, $\cosh r$ and $\sinh r$ are Bogoliubov's coeffients, and the number of photons emerging after the application of the squeezing operators $\hat{S}$ and $\hat{S}^{-1}$ will be

$$\begin{aligned} N_{\rm out}(t) &= \langle\alpha'|\hat{a}^\dagger\hat{a}|\alpha'\rangle \\ &= |\alpha|^2\big[1 - \sinh(2r)\sin(2\theta)\sin(2\delta(t)) \\ &\quad + \sinh(4r)\cos(2\theta)\sin^2\delta(t) \\ &\quad + 2\sinh^2(2r)\sin^2\delta(t)\big] \\ &\quad + \sinh^2(2r)\sin^2\delta(t), \end{aligned} \tag{5}$$

where, in the following, we set $\theta = 0$. Notice that for $\delta = 0$, one has $N_{\rm out} = N_{\rm in} = |\alpha|^2$. Therefore, $N_{\rm out} \neq N_{\rm in}$, provides an indirect evidence of a nonzero phase shift induced by the SDM-photon interaction.

Introducing the relations $\phi_0 = \frac{1}{m_\phi}\sqrt{2\rho_{\rm local}}$ and $\omega = m_\phi$ in Eq. (3), and being that (in natural units) $t = L$, the phase shift (3) induced by the scalar field becomes

$$\delta(L) = \frac{g_\gamma\sqrt{2\rho_{\rm local}}}{m_\phi^2}|\vec{k}|\sin(m_\phi L), \tag{6}$$

where $g_\gamma = 1/\Lambda_\gamma$, and $|\vec{k}| = 2\pi/\lambda$ is the photon wavenumber (with wavelength λ). Equation (6) shows that the phase shift depends on the local dark matter density $\rho_{\rm local}$, the photon momentum $|\vec{k}|$, the scalar field mass m_ϕ, and the coupling constant g_γ.

We emphasize that to observe the interaction between SDM and photons it is necessary to have a single-arm interferometer with a long effective optical path, as can be achieved in experimental setups, such as LIGO, VIRGO and PVLAS. LIGO is a two-arm Fabry-Perot-Michelson interferometer that precisely measures the differential arm length. For our purposes, the current configuration cannot be used. One possible use of devices like LIGO and VIRGO would consist of using only a single arm of these interferometers, applying squeezing and antisqueezing operators at the beginning and end of the effective optical path, and using a photodetector to count the output photons. In the case of PLVAS, however, the device is represented by a Fabry-Perot cavity, realizing a long effective optical path, and the photodetectror is already implemented in the setup. Therefore, it would only be necessary to apply the squeezer and antisqueezer operators, removing the magnetic field used in PVLAS and not necessary for our studies. In the following section, we consider values of the parameters compatible with the LIGO interferometer.

IV. NUMERICAL ANALYSIS

In our analysis, we consider $\rho_{\rm local} = 4\times10^5$ GeV$/m^3$ and the following range of values of g_γ and mass m_ϕ [93,94]: $g_\gamma \in [10^{-35}, 10^{-28}]$ eV^{-1}, $m_\phi \in [10^{-22}, 10^{-12}]$ eV.

In Fig. 2, we plot $\delta(L)$ as a function of m_ϕ and g_γ, respectively. We consider $L = 950$ km, which is compatible with the effective length of the LIGO

[4]It is easy to check that if there is no mixing between annihilation and creation operators, the interaction in Eq. (1) does not produce any observable effects. A squeezing-like operation, such as the Bogoliubov transformation in Eq. (4), is necessary to detect the effects of the interaction in Eq. (1) within a single-arm interferometer.

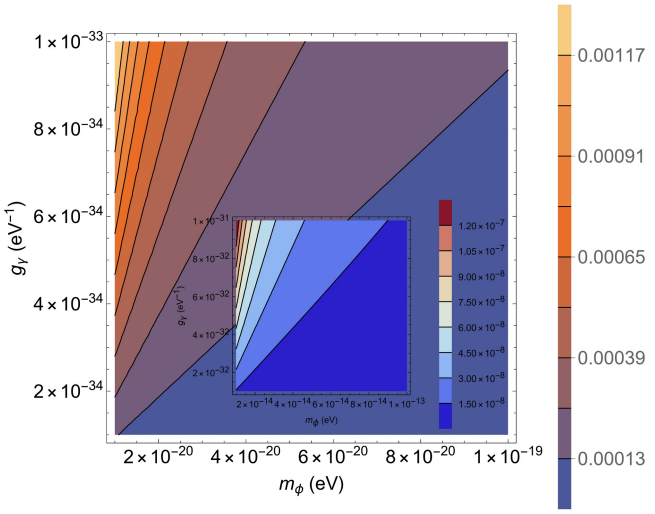


FIG. 2. Main graph: plot of $\delta(L)$ as as function of $g_\gamma \in [10^{-34}, 10^{-33}]$ eV^{-1} and $m_\phi \in [10^{-20}, 10^{-19}]$ eV. In the inset is plotted $\delta(L)$ as a function of $g_\gamma \in [10^{-31}, 10^{-30}]$ eV^{-1} and $m_\phi \in [10^{-13}, 10^{-12}]$ eV. We consider the arm length $L = 950$ km and the photon wavelength $\lambda = 1064 \times 10^{-9}$ m.

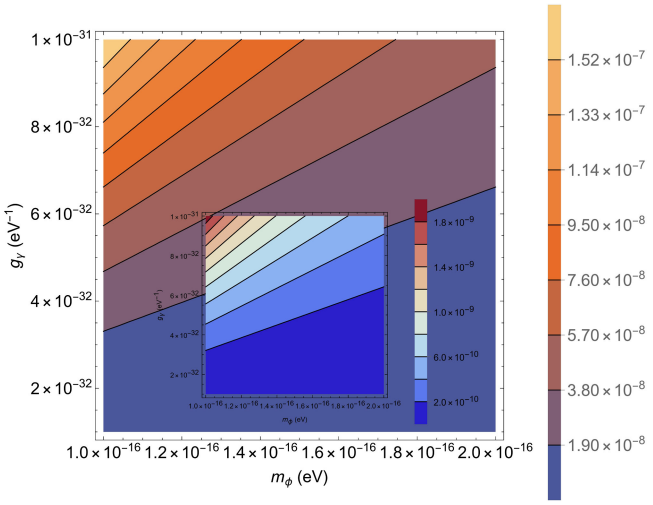


FIG. 3. Plot of $\frac{\Delta N}{N_{\rm in}}$ as a function of $g_\gamma \in [10^{-32}, 10^{-31}]$ eV^{-1} and $m_\phi \in [10^{-16}, 2\times10^{-16}]$ eV. Main graph: we consider the squeezing parameter $r = 1.7$, which corresponds to 15 dB squeezed states of light. Inset: we use $r = 0.6$, which is compatible with the value of the squeezing parameter of the LIGO experiment. We utilized the same values of L and λ considered in Fig. 2.

interferometer[5] [95]. This length is also used in the other figures. A bigger effective length can be achieved by using devices like PVLAS, in which an effective length of 1.47×10^9 m can be reached [96]. In principle, this device can be used for our purposes simply by introducing squeezer-anti squeezer operators at the beginning and at the end of the effective optical path. (We point out that in the case of PVLAS the presence of the magnetic field is not necessary in our configuration.) In Figs. 3–5, we plot the quantity $\Delta N(L)/N_{\rm in}$, where $\Delta N(L) = N_{\rm out}(L) - N_{\rm in}$ and

$$\frac{\Delta N(L)}{N_{\rm in}} = \sin^2\delta(L)\left[\left(\frac{1}{|\alpha|^2}+2\right)\sinh^2 2r + \sinh 4r\right], \quad (7)$$

as functions of m_ϕ and g_γ for fixed values of the length L and of the squeezing parameter r. The signal we are interested in is the $\Delta N(L)$, which quantifies the impact of the scalar interaction on the output photon number. Assuming a Poissonian photon counting statistics, the corrisponding noise is $\sqrt{N_{\rm out}}$, and then the signal-to-noise ratio (SNR) is given by

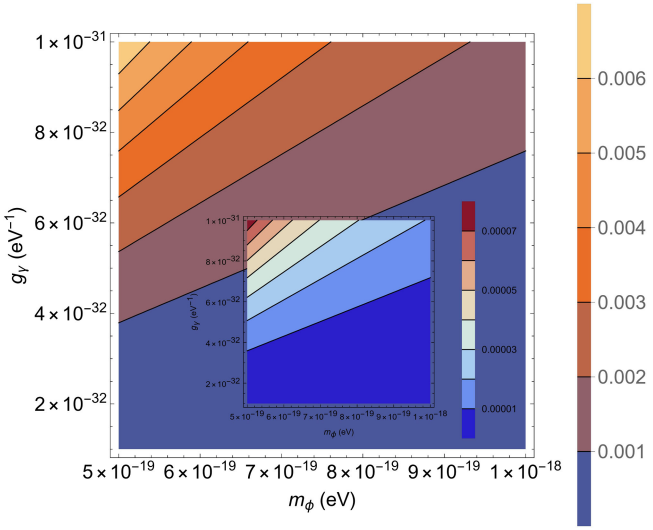


FIG. 4. Plot of $\frac{\Delta N}{N_{\rm in}}$ as a function of $g_\gamma \in [10^{-32}, 10^{-31}]$ eV^{-1} and $m_\phi \in [0.5\times10^{-18}, 10^{-18}]$ eV. Main graph: we consider $r = 1.7$. Inset: $r = 0.6$. The same values of L and λ considered in Figs. 2 and 3 are used.

$$\mathrm{SNR} = \frac{\Delta N(L)}{\sqrt{N_{\rm out}}} = \frac{|\alpha|^2\sin^2\delta(L)\left[\left(\frac{1}{|\alpha|^2}+2\right)\sinh^2 2r + \sinh 4r\right]}{\sqrt{N_{\rm out}}}. \quad (8)$$

We point out that $\Delta N(L)$ is zero for δ, and r is equal to zero, since in this case $N_{\rm out}(L) = N_{\rm in}$. We consider $r = 0.6$, which is compatible with the value of the squeezing parameter of the LIGO experiment, and $r = 1.7$, which

[5]The effective length is approximately computed as $L = \frac{2\mathcal{F}}{\pi}4$ km, where $\mathcal{F}$ is the cavity finesse. This approximation breaks down when the frequency of SDM exceeds $1/\tau$, where τ is the cavity storage time. Consider that storage time is of the order of $10ms$, and the approximation breaks down for $\omega_\phi \geq 50$ Hz. All the masses considered in this work lie below 10^{-13} eV, corresponding to ≈24 Hz. Therefore the effective length employed here is compatible with the mass range considered.

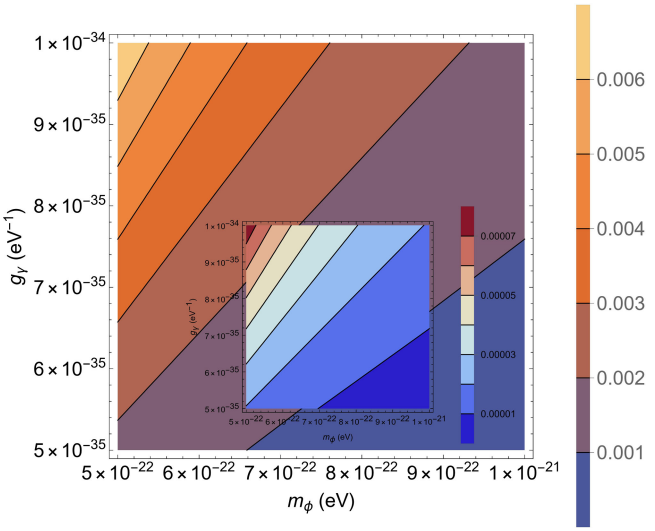


FIG. 5. Plot of $\frac{\Delta N}{N_{\rm in}}$ as a function of $g_\gamma \in [5\times10^{-35}, 10^{-34}]$ eV^{-1} and $m_\phi \in [5\times10^{-22}, 10^{-21}]$ eV. These values of the masses are typical of fuzzy cold dark matter [62]. Main graph: we consider $r = 1.7$. Inset: $r = 0.6$. The same values of L and λ considered in Figs. 2 and 3 are used.

corresponds to 15 dB squeezed states of light [97]. In Fig. 6, $\Delta N(L)/N_{\rm in}$ is plotted as a function of L for fixed values of m_ϕ, g_γ, and r. In the inset of Fig. 6, $\Delta N(L)/N_{\rm in}$ is shown as a function of r for fixed values of m_ϕ, g_γ, and L. In the LIGO experiments, the initial number of photons is $N_{\rm in} \sim 10^{23}$ [95], and the noise is approximately $\sqrt{N_{\rm out}}$. In the ranges of the parameters where g_γ and m_ϕ analyzed, $N_{\rm out}$ is of the same order of $N_{\rm in}$ and $\Delta N > \sqrt{N_{\rm out}}$. Another source of uncertainty is represented by the optical losses, which increases the quantum noise in detectors like LIGO

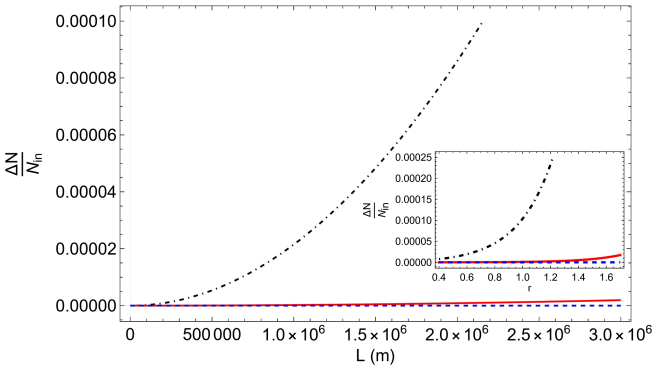


FIG. 6. Main picture: plot of $\frac{\Delta N}{N_{\rm in}}$ as a function of the arm length L for $g_\gamma = 10^{-33}$ eV^{-1} and mass $m_\phi = 10^{-20}$ eV (the black dotted line), for $g_\gamma = 10^{-32}$ eV^{-1} and mass $m_\phi = 10^{-18}$ eV (the red line), and for $g_\gamma = 10^{-30}$ eV^{-1} and mass $m_\phi = 10^{-12}$ eV (the blue dashed line), and for the value of squeezing parameters $r = 0.6$. In the inset is plotted $\frac{\Delta N}{N_{\rm in}}$ as a function of r for $L =$ 300 km and for the values of g_γ and m_ϕ used in the corresponding line of the same colors.

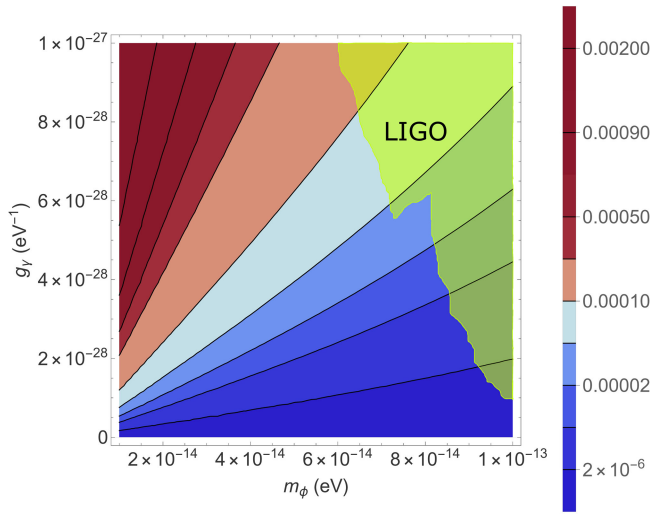


FIG. 7. Plot of $\frac{\Delta N}{N_{\rm in}}$ and comparison with the results of [85] (green shaded region) as a function of $g_\gamma \in [10^{-33}, 10^{-27}]$ eV^{-1} and $m_\phi \in [10^{-14}, 10^{-13}]$ eV.

and PVLAS. How to reduce these optical losses has been analyzed in [98], where the noise associated with such reduction is estimated in 42 and 87 parts per million in any case less than $\Delta N/N_{\rm in}$. Other forms of noise such as thermal noise and seismic noise are expected to be subleading with respect to the Poissonian noise $\sqrt{N_{\rm out}}$. The instrument sensitivity varies significantly depending on the device considered (LIGO, VIRGO, PVLAS, etc.). Therefore, our plots show that with gravitational wave detectors and PVLAS-like detectors, the interaction between SDM (including the fuzzy cold dark matter [62]) and photons can be revealed, eventually constraining the parameters of the SDM model.

In Fig. 7, we compare the plots obtained with our method with the one derived in the recent studies regarding the upper limits on SDM, as derived from LIGO [85], and compare our bindings to their prediction. As it is evident from Fig. 7, the approach we propose can probe a region of parameter space encompassing the range considered in [85] and beyond.

V. DISCUSSIONS AND CONCLUSIONS

The single-arm interferometric scheme, here introduced, provides a novel strategy to probe scalar dark matter (SDM) through its coupling to photons. Limitations of the present analysis deserve careful consideration.

First, the scalar dark matter field has been modeled as a coherently oscillating, monochromatic, and spatially homogeneous classical field [see Eq. (3)]. This approximation is valid only within the coherence time and coherence length associated with the galactic dark matter velocity distribution. In realistic scenarios, the finite velocity dispersion of SDM induces a frequency spread and phase decoherence, which may reduce the effective phase

accumulation over the interferometer arm length L. Since the predicted phase shift in Eq. (6) scales linearly with the interaction time (or equivalently with L), such effects could modify the signal profiles shown in Figs. 2–6, especially for the longest effective optical paths considered.

Second, the analysis of the interferometric signal presented relies on an idealized description of the experimental setup. In particular, the expressions for the output photon number [Eq. (5)] and for the relative variation $\Delta N/N_{\rm in}$ [Eq. (7)] neglect several realistic noise sources, including optical losses, residual phase noise, technical noise associated with squeezing generation, and fluctuations in the squeezing angle. While we argue that Poissonian photon-counting noise is subdominant in the parameter regions explored in Figs. 3–6, a complete noise budget analysis is required to quantitatively assess the experimental sensitivity.

Third, the proposed detection strategy critically depends on the implementation of two spatially separated squeezing and antisqueezing operations, as illustrated schematically in Fig. 1. In realistic interferometers, finite squeezing levels, imperfect mode matching, and optical losses can degrade the effective signal encoded in ΔN. Since the enhancement of the signal scales strongly with the squeezing parameter r [see Eqs. (5)–(8) and Fig. 6], deviations from the ideal Bogoliubov transformations in Eq. (4) may lead to a reduction of the observable effect.

Moreover, the interaction described by the effective coupling term $\phi F^{\mu\nu}F_{\mu\nu}/\Lambda_\gamma$, is generic to scalar fields coupled to electromagnetism and is not unique to cosmological dark matter. As a consequence, the signal characterized by the phase shift in Eq. (6) and by the photon-number variation shown in Figs. 3–5 may, in principle, also arise from other scalar backgrounds or from effective variations of electromagnetic parameters. Additional discriminants—such as characteristic temporal modulations linked to Earth's motion through the galactic halo—will therefore be necessary to unambiguously attribute a detected signal to scalar dark matter.

Despite these limitations, the results presented demonstrate that single-arm interferometry combined with nonclassical states of light can probe regions of the SDM parameter space that are complementary to those explored by existing interferometric searches. Future work will focus on extending the present framework by incorporating realistic noise models, accounting for the stochastic nature of the SDM field, and developing a concrete experimental implementation within existing or planned interferometric facilities. In this perspective, the approach proposed here may represent a promising and complementary avenue for probing ultralight scalar dark matter beyond current experimental techniques.

In conclusion, we presented a novel approach based on single-arm interferometry to prove the existence of scalar dark matter. Specifically, we investigated the possible coupling between scalar dark matter and photons. We have shown that should this interaction be present, a single-arm interferometer would detect a phase shift and an intensity variation directly resulting from the scalar-photon coupling. The numerical analysis of the phase shift and of the intensity variation indicates that the proposed approach can probe SDM for a wide range of parameters. Our setup can be developed by using gravitational wave detectors and two spatially separated squeezing operations, and it can, in principle, be employed to test any interaction proportional to the photon number operator.

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DATA AVAILABILITY

The data that support the findings of this article are available within this article.

APPENDIX: INTERACTION HAMILTONIAN

In this Appendix, we derive the interaction Hamiltonian density in Eq. (1), starting from the Lagrangian density $\mathcal{L}_{\rm int} = \frac{1}{4}\frac{\phi}{\Lambda_\gamma}F_{\mu\nu}F^{\mu\nu}$. We consider the radiation gauge [99], and we set $\eta_{\mu\nu} = \mathrm{diag}(1,-1,-1,-1)$.

Notice that $\mathcal{L}_{\rm int} = -\frac{\phi(t)}{\Lambda_\gamma}\mathcal{L}_0$, where $\mathcal{L}_0 = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} = \frac{1}{2}(\vec{E}^2 - \vec{B}^2)$ is the free electromagnetic Lagrangian.

As a consequence, $\mathcal{H}_{\rm int} = -\frac{\phi(t)}{\Lambda_\gamma}\mathcal{H}_0$, where $\mathcal{H}_0 = \frac{1}{2}(\vec{E}^2 + \vec{B}^2)$ is the free electromagnetic Hamiltonian.

We derive explicitly this result. We consider the radiation gauge $\vec{\nabla}\cdot\vec{A} = 0$ and $A^0 = 0$, where $\vec{A}$ is the vector potential defined by $\vec{B} = \vec{\nabla}\times\vec{A}$, and then the electric field is given by $\vec{E} = -\dot{\vec{A}}$. The canonical momentum conjugate to the generalized coordinate $\vec{A}$ is

$$\Pi^\mu = \frac{\partial\mathcal{L}}{\partial\dot{A}_\mu} = -\left(1 - \frac{\phi(t)}{\Lambda_\gamma}\right)F^{0\mu}, \tag{A1}$$

or

$$\Pi^0 = 0 \quad \text{and} \quad \Pi^i = \left(1 - \frac{\phi(t)}{\Lambda_\gamma}\right)E^i, \tag{A2}$$

where $i = 1, 2, 3$ denote the spatial component. The Hamiltonian density is $\mathcal{H} = \Pi_\mu\dot{A}^\mu - \mathcal{L}$. Being that $\dot{\vec{A}} = -\vec{E}$ and using Eq. (A1), one has

$$\mathcal{H} = \frac{1}{2}\left(1 - \frac{\phi(t)}{\Lambda_\gamma}\right)(\vec{E}^2 + \vec{B}^2). \tag{A3}$$

The quantum field operators $\hat{\vec{E}}$ and $\hat{\vec{B}}$ are given, respectively, by [99]

$$\hat{\vec{E}}(\vec{r},t) = \int \frac{d^3k}{\sqrt{2|\vec{k}|(2\pi)^3}} \sum_\lambda i|\vec{k}| \left[\hat{a}_{\vec{k}\lambda}\vec{\epsilon}_{\vec{k}\lambda}e^{-ikx} - \hat{a}^\dagger_{\vec{k}\lambda}\vec{\epsilon}_{\vec{k}\lambda}e^{ikx}\right],$$
$$\hat{\vec{B}}(\vec{r},t) = \int \frac{d^3k}{\sqrt{2|\vec{k}|(2\pi)^3}} \sum_\lambda i|\vec{k}| \left[\hat{a}_{\vec{k}\lambda}(\hat{k}\times\vec{\epsilon}_{\vec{k}\lambda})e^{-ikx} - \hat{a}^\dagger_{\vec{k}\lambda}(\hat{k}\times\vec{\epsilon}_{\vec{k}\lambda})e^{ikx}\right], \quad \text{(A4)}$$

where $\vec{\epsilon}_{\vec{k}\lambda}$ are the polarization vectors, and $\hat{a}_{\vec{k}\lambda}$ are the annihilator satisfying the commutation relations $[\hat{a}_{\vec{k}\lambda}, \hat{a}^\dagger_{\vec{k}'\lambda'}] = \delta^{(3)}(\vec{k}-\vec{k}')\delta_{\lambda\lambda'}$. Introducing Eqs. (A4) in Eq. (A3) and integrating on the volume, we obtain

$$\hat{H}(t) = \hat{H}_0 + \hat{H}_{\rm int} = \left(1 - \frac{\phi(t)}{\Lambda_\gamma}\right) \int d^3k \sum_\lambda |\vec{k}| \hat{a}^\dagger_{\vec{k}\lambda}\hat{a}_{\vec{k}\lambda}, \quad \text{(A5)}$$

where the second term in Eq. (A5) is the interaction Hamiltonian in Eq. (1).

We point out that in quantum mechanics if no mixing occurs between annihilation and creation operators, as in the Bogoliubov transformation [Eq. (4)], the scalar-photon interaction, Eq. (1), does not produce any observable effect. Indeed, when the squeezing-antisqueezing operation is removed, the expected number of output photons coincides with the number of input photons [see Eq. (5)], in which $N_{\rm out} = N_{\rm in}$, for $r = 0$ and $\delta = 0$, which correspond to the absence of Bogoliubov transformations. While in principle also classic techniques may be employed, they must necessarily involve a Bogoliubov transformation and a form of squeezing as in Eq. (4). The efficiency of such techniques in highlighting the effect of the SDM is, in each case, dependent on an effective squeezing parameter r. We note that our SNR, Eq. (8), has no strict classical correspondence, since in our case, SNR $= 0$, for $r = 0$ and $\delta = 0$. In conclusion, our approach relies on quantum rather than classical aspects.

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