

Resonant enhancement of axion dark matter decay

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(Received 4 August 2025; revised 31 October 2025; accepted 23 February 2026; published 27 March 2026)

The axion is a particularly well-motivated candidate for the dark matter comprising most of the mass of our visible Universe, leading to worldwide experimental and observational efforts toward its discovery. As is well known, resonant cavities are a primary tool in this search, where they are used to enhance the conversion rate of axions into photons in a background electromagnetic field. What is perhaps less well appreciated is that resonant cavities can also enhance the rate at which axions decay to two photons, a manifestation of the Purcell effect. We examine this possibility and show that it offers a previously unexplored method to search the axion parameter space in a way that is competitive and complementary to other approaches, with the capability to probe a well-motivated region for QCD axion dark matter. Furthermore, we establish that this method can be implemented via preexisting axion heterodyne detection schemes, enabling such experiments to perform these searches with minimal modification.

DOI: [10.1103/9c76-zcym](https://doi.org/10.1103/9c76-zcym)

I. INTRODUCTION

Overwhelming evidence exists that most of the matter in our Universe is “dark,” in that it interacts very feebly or not at all with the standard model (SM) [1–6]. In the standard Λ cold dark matter (CDM) cosmology, this dark matter (DM) is weakly interacting, nonrelativistic, and cosmologically stable. Beyond this, not much is known about the underlying nature of DM or its interactions with SM particles, which has led to a profusion of DM models and candidates.

A particularly well-motivated candidate for this DM is the axion, and by extension axionlike particles [7–9]. As they arise straightforwardly from a minimal extension of the standard model, the Peccei-Quinn solution to the strong CP problem [10–12], axions occupy a rare focal point in theoretical physics, in that they are also a generic prediction of the exotic physics of string and M-theory compactifications [13,14]. Despite the profound differences between these two contexts, the resulting axion properties are also largely universal, presenting an easily characterizable theoretical target. Furthermore, as a light DM candidate coupled to the SM, axions also offer discovery potential via small-scale “tabletop” experiments.

This being the case, many experiments are ongoing worldwide, seeking to discover the axion by a variety of means [15]. Among these experiments, a primary technique is the cavity

haloscope, which relies upon a powerful magnetic field and a resonant cavity to enhance the rate of Primakoff conversion of axions into photons [16]. At present, some of the most sensitive experiments searching for light DM are of this type [17–31].

In this paper, we explore an alternative channel to axion discovery via the same haloscope technology. As we will demonstrate, in addition to enhancing the rate of Primakoff conversion of axions to single photons, resonant cavities can also be used to enhance the rate of axion decay into two photons. This enhancement can be understood as an example of a more general phenomenon, more typically seen in the context of quantum optics, known as the Purcell effect.

We will in the following outline and explore the capability of a resonant cavity to enhance the decay rate of DM axions, before examining the resultant discovery potential of experiments making use of this principle. Conclusions and discussion are presented in closing.

II. AXION DECAY INSIDE A RESONANT CAVITY

Following the basic principles of quantum mechanics, it is generally believed that particle decays occur spontaneously, independently of outside influence. The Purcell effect, originally explored in Ref. [32], is a rather notable counterexample to this, where the enhancement or suppression of decay processes occurs due to the environment in which these decays take place.

We can see this effect most straightforwardly in Fermi’s golden rule for transition rates,

$$\Gamma = 2\pi | \langle f | H_I | i \rangle |^2 \rho, \quad \rho = dN/dE, \quad (1)$$

where H_I is an interaction Hamiltonian and ρ is the density of final states. Altering the interaction strength in H_I will affect Γ ; however, in principle so will modifications to ρ . While we

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expect a continuum of states in free space, inside a resonant cavity there will be a spectrum of discrete modes due to the cavity boundary conditions, which will enhance ρ at certain frequencies and suppress it at others, modifying the transition rate. This typically leads to an enhancement of the form

$$F_{\text{Purcell}} \equiv \frac{\tau_{\text{free}}}{\tau_{\text{cavity}}} \simeq \frac{3\lambda^3}{4\pi^2} \left(\frac{Q}{V} \right), \quad (2)$$

where τ indicates the lifetime of a particular state, λ the corresponding wavelength, Q is the quality factor, and V the corresponding mode volume [32].

Turning to axion physics, from the interaction term

$$\mathcal{L} = g_{a\gamma} a \vec{E} \cdot \vec{B}, \quad (3)$$

where $g_{a\gamma}$ is the axion-photon coupling, axions are able to convert to single photons in a background electromagnetic field, and also to decay to two photons. In the latter case, $\tau_{\text{free}} = 64\pi^3 / g_{a\gamma}^2 m_a^3$, where m_a is the axion mass, leading to a lifetime for typical axion DM parameters that far exceeds the age of our Universe. To enhance this decay rate to an observable level with a cavity experiment, there are a few factors we must account for.

From Eq. (3), axion decay must produce pairs of photons with $\vec{E} \cdot \vec{B} \neq 0$. Therefore, we first require a cavity with two resonant modes and a nonzero form factor

$$\eta \equiv \frac{|\int_V \vec{E} \cdot \vec{B}|}{\sqrt{\int_V |\vec{E}|^2} \sqrt{\int_V |\vec{B}|^2}} \leq 1. \quad (4)$$

Denoting the mode frequencies by $\omega_{s/p}$ (since we will subsequently identify one mode as “pump” and the other as “signal”), energy conservation dictates that $\omega_s + \omega_p = \omega_a \simeq m_a(1 + v_{\text{vir}}^2/2)$, where v_{vir} is the virial velocity, and so in general both modes do not need to be degenerate. Of course, since m_a is in general unknown, we will also need to be able to scan the axion parameter space by tuning the resonant frequency of one or both modes.

Conveniently, unlike ordinary single-mode haloscope experiments, this does not require a strong magnetic field, which can be a significant source of experimental complexity. In turn, this permits the use of ultrahigh Q superconducting radio-frequency (SRF) cavities, which also allow for enhanced sensitivity [23].

Since the Purcell enhancement discussed above applies to both spontaneous and stimulated decays, we have the option to operate with an empty cavity, or with one of the modes already populated with photons. In the latter case, photons from the so-called pump mode induce axion decays, in a manner akin to the stimulated axion decay scenario previously explored in various astrophysical settings [33–42]. One of the resulting decay photons remains in the pump mode, while the other appears in the signal mode, to be subsequently detected.

Pumping one cavity mode does introduce further experimental complexity, namely, additional cooling requirements due to the power deposited with the cavity and the necessity of strongly isolating pump and signal to avoid signal contamination. However, as we will see it is nonetheless this case that offers better sensitivity. There is also a further advantage, which we now explore.

III. AXION UP-CONVERSION

Ordinary axion haloscopes suffer from an $\sim (m_a L)^2$ parametric suppression for small m_a , where the axion Compton wavelength significantly exceeds the characteristic size of the experiment L . This being the case, various authors have explored the potential of “heterodyne” or “up-conversion” detection schemes, where one cavity mode is pumped to generate an oscillating magnetic field, which acts as a background for Primakoff-type conversion but enables the aforementioned suppression to be evaded. Low-frequency axions satisfying $m_a \simeq \omega_s - \omega_p$ can then be absorbed, resulting in a pump photon up-converting to a signal photon of higher energy [43–46], building upon the idea originally presented in Ref. [47].

Similar considerations for photon regeneration experiments have also been explored in Refs. [48–51], and Refs. [52,53] also employed a dual-mode strategy to search for axion-induced frequency shifts. Cavity design requirements for such searches are analyzed in Ref. [54].

The key point for our purposes is that axion up-conversion again requires the use of a tunable resonant cavity, with a signal and pump mode satisfying $\eta \neq 0$. Typical implementations also rely upon ultrahigh Q SRF cavities to maximize sensitivity. Therefore, on the basis of the arguments in the previous section, this type of experiment should also be sensitive to the stimulated and spontaneous decay of axions. Indeed, production of photon pairs via axion decay can be understood as a type of parametric down-conversion, which is a complementary process to the up-conversion outlined above.

The primary difference between these two contexts is kinematic: For axion decay (equivalently down-conversion) $m_a \simeq \omega_s + \omega_p$, while for up-conversion $m_a \simeq \omega_s - \omega_p$. Furthermore, because the experimental conditions required for both scenarios are the same, a single experiment of this type could in principle simultaneously search for axions satisfying either condition.

This being the case, we can leverage many results from previous research into up-conversion experiments. However, before doing so, we must calculate the signal power due to axion decay in a resonant cavity, which is unique to the current situation.

IV. SIGNAL POWER

In this section, we will derive the signal power due to axion decay in a resonant cavity, with full details given in the Supplemental Material [55]. Our starting point is the axion photon interaction in Eq. (3). Expanding in a basis of cavity modes with creation (annihilation) operators c_i ($c_i^\dagger$) for a two-mode system ($i = 1, 2$), we find the interaction Hamiltonian

$$H_{\text{int}} = \frac{ig_{a\gamma} a \sqrt{\omega_1 \omega_2}}{2} (\xi_- (c_1 c_2^\dagger - c_1^\dagger c_2) + \xi_+ (c_1 c_2 - c_1^\dagger c_2^\dagger)), \quad (5)$$

where ω_i is the mode energy and $\xi_\pm$ are form factors analogous to Eq. (4). We can describe our cavity system via a total Hamiltonian H following the approach detailed in Ref. [56], so that working in the Heisenberg picture we then have $dc_i(t)/dt = i[H, c_i]$, with a corresponding equation of

motion for the down-conversion case:

$$\frac{dc_i}{dt} = \left(-i\omega_i - \frac{\gamma_i}{2}\right)c_i - \sqrt{\gamma_i}b_{\text{in}} - \frac{g_{ay}\sqrt{\omega_s\omega_p}}{2}\xi_+ac_j^\dagger, \quad (6)$$

where $i \neq j$. Here, $\gamma_i = \omega_i/Q_i$ and b_{in} capture the effects of damping and pumping, respectively.

We now delineate into signal ($i = s$) and pump ($j = p$) modes; Fourier transforming to solve for $c_s(\omega)$, we find

$$c_s(\omega) = \frac{g_{ay}\sqrt{\omega_s\omega_p}\xi_+}{2} \int \frac{d\omega'}{2\pi} \frac{a(\omega + \omega')c_p^\dagger(\omega')}{(i\omega - i\omega_s - \frac{1}{2}\gamma_s)}, \quad (7)$$

where the axion field satisfies

$$\langle a^\dagger(\omega)a(\omega') \rangle = 2\pi\delta(\omega - \omega')F_a(\omega)\rho_a/m_a^2, \quad (8)$$

where $F_a(\omega)$ is the laboratory-frame axion energy distribution, which follows from the underlying Maxwell-Boltzmann velocity distribution, and ρ_a is the DM density.

The expectation value of the number operator $N_s = c_s^\dagger c_s$, averaged over a sufficiently long timescale T , then gives the number of photons in the signal mode,

$$\langle N_s \rangle_T \equiv \frac{1}{T} \int_{-T/2}^{+T/2} dt \langle c_s(t)^\dagger c_s(t) \rangle. \quad (9)$$

Inserting Eq. (7), we find after some manipulations

$$\langle N_s \rangle_T \simeq \frac{g_{ay}^2 Q_s \omega_p |\xi_+|^2 \rho_a F_a(\omega_s + \omega_p)}{4m_a^2} (1 + \langle N_p \rangle). \quad (10)$$

Q_s is the loaded Q factor of the signal mode, where

$$(Q_s)^{-1} = (Q_{\text{int}})^{-1} + (Q_{\text{cpl}})^{-1}, \quad (11)$$

with Q_{int} the intrinsic Q factor and Q_{cpl} determined by the cavity coupling to the readout. The first term in $(1 + \langle N_p \rangle)$ gives the spontaneous decay contribution, while $\langle N_p \rangle$ (the average steady-state number of photons in the pump mode) corresponds to stimulated decays.

Implicit in our derivation here is the assumption that $Q_s/\omega_s \gg Q_a/m_a$, where Q_a is the axion Q factor, and hence that the axion power spectral density (PSD) is broader than our cavity bandwidth. As is evident in Eq. (10), we therefore are only sensitive to a fraction of the total axion PSD, namely, those axions that satisfy $\omega_a \simeq \omega_s + \omega_p$.

As the energy in the pump mode is approximately $\omega_p \langle N_p \rangle$, the pump mode power escaping the cavity is $P_{\text{out}} \simeq \omega_p \langle N_p \rangle \gamma_p$. Assuming a steady state, the pump power into and out of the cavity is balanced, so that

$$\langle N_p \rangle \simeq P_{\text{in}}/(\omega_p \gamma_p) = P_{\text{in}} Q_p / \omega_p^2. \quad (12)$$

In practice, P_{in} is fixed by the requirement that the resultant magnetic field at the cavity walls does not exceed the critical threshold to disrupt superconductivity. We can therefore equate $\max(P_{\text{in}})$ with the dissipated power $P_{\text{diss}} \simeq w_p U_{\text{max}}/Q_{\text{int}}$, where the maximum stored energy U_{max} is determined numerically from the cavity geometry, mode profiles, and material properties.

Similar logic applies to the signal power, although the Q dependence is slightly different as we are interested only in the power reaching the readout rather than simply escaping the cavity. This then yields

$$P_s \simeq \omega_s^2 \langle N_s \rangle_T / Q_{\text{cpl}}. \quad (13)$$

In practice, it will be more convenient to use the corresponding PSD, which we can find by the same procedure,

$$S_s \simeq \frac{g_{ay}^2 \omega_s^3 \omega_p |\xi_+|^2 \rho_a (1 + \langle N_p \rangle) F_a(\omega + \omega_p)}{4Q_{\text{cpl}} m_a^2 ((\omega - \omega_s)^2 + \frac{1}{4}\gamma_s^2)}. \quad (14)$$

As an interesting aside, we can identify the Purcell-derived origin of P_s in the following way. As outlined in Eq. (2), a Purcell-enhanced decay rate characteristically appears with a factor (Q/V) , while the signal power at an experiment of this type depends on this rate and the amount of DM inside the cavity, which is proportional to V . Since these factors cancel the Purcell-enhanced signal, power should be V -independent, and indeed this is precisely reflected in Eqs. (10) and (13). In contrast, in the ordinary haloscope case, the signal power is directly proportional to V .

Of course, there are certain experimental considerations that curtail this “ V independence,” most obviously that the mode frequencies are directly related to the dimensions of the cavity. Furthermore, the necessity of cooling the cavity introduces further V dependence, since cavities with a smaller surface area cannot dissipate power as effectively.

V. NOISE POWER

As explored in Refs. [43–46], experiments of this type receive noise contributions from a variety of sources, including leakage noise, mechanical mode mixing, thermal noise, and amplifier noise. At the higher frequencies relevant to the axion decay scenario, where $m_a \simeq \omega_s + \omega_p \sim \mathcal{O}(\text{GHz})$, the primary contributions are expected to be thermal and amplifier noise. Due to mode pumping, cooling the cavity to $\mathcal{O}(\text{K})$ temperatures is much more difficult than for ordinary haloscopes, likely requiring liquid helium cooling. Further discussion of this point can be found in the references directly above.

Following the experimental configuration given in Ref. [44], we have the noise PSD

$$S_n \simeq \frac{4\beta}{(1 + \beta)^2} \frac{S_T - S_{T_c}}{1 + 4Q_s(\frac{\omega}{\omega_s} - 1)^2} + S_{T_c} + S_a, \quad (15)$$

where $\beta \equiv Q_{\text{int}}/Q_{\text{cpl}}$ is the cavity coupling parameter, $S_T = n_T \omega \equiv (e^{\omega/T} - 1)^{-1} \omega$, and likewise for S_{T_c} , where T_c is the effective temperature of backaction noise from the detector. We assume here that our system is in equilibrium at the cavity temperature, so that $T \simeq T_c \gg \omega$, and hence $S_{T_c} \simeq S_T \simeq T$. S_a is the amplifier noise contribution, which we assume to be subdominant to thermal noise. Improved performance could be achieved via standard quantum limit (SQL)-limited amplification, so that $S_a = \omega$ and $T_c = 0$, but we will not rely upon this here.

VI. PROJECTED SENSITIVITY

In this section, we examine the potential sensitivity of this technique. Since we expect the noise power in different bins to be independent, we can sum the contributions from individual bins in quadrature, giving

$$\text{SNR}^2 \simeq t_{\text{int}} \int_0^\infty \frac{d\omega}{2\pi} \left(\frac{S_s}{S_n} \right)^2, \quad (16)$$

with t_{int} the integration time [57]. Optimal sensitivity comes from the stimulated decay case, as then $\langle N_p \rangle \gg 1$. Our benchmark cavity follows Ref. [44], using the TE₀₁₂ (pump) and TM₀₁₃ (signal) modes of a cylindrical Nb cavity of length/radius $\simeq 2.35$, with the pump mode storing $U_{\text{max}} \simeq 690$ J, $Q_{\text{int}} \simeq 2 \times 10^{11}$, and $V = 60$ L at a frequency $f_c = 1.1$ GHz, with a form factor of 0.19. To map to other values of f_c , we simply rescale the cavity size, where U_{max} scales in proportion to V .

Our choice of this benchmark case is motivated in part by this ease of rescaling to reach other frequencies. As previously mentioned, the signal power in this case is “ V -independent,” modulo certain experimental considerations, which has the potential to evade a significant problem for cavity haloscopes, namely, the diminishing sensitivity at higher frequencies due to increasingly small cavity size. As such, it is important to assess the performance of this approach as a function of frequency. In contrast, the cavity concepts explored in Refs. [43,45,46] may not be as easily mapped to different frequencies.

Introducing the signal mode resonant frequency $f_s = \omega_s/2\pi$, Eq. (16) yields

$$g_{ay}^{95\%} \simeq \frac{5 \times 10^{-15}}{\text{GeV}} \left(\frac{2 \times 10^{11}}{Q_{\text{int}}} \right)^{\frac{1}{4}} \left(\frac{10^6}{Q_a} \right)^{\frac{1}{2}} \left(\frac{f_s}{1.3 \text{ GHz}} \right)^{\frac{9}{4}} \times \left(\frac{0.45 \text{ GeV/cm}^3}{\rho_a} \right)^{\frac{1}{2}} \left(\frac{T}{1.8 \text{ K}} \right)^{\frac{1}{2}} \left(\frac{100 \text{ s}}{t_{\text{int}}} \right)^{\frac{1}{4}}, \quad (17)$$

corresponding to an axion mass $m_a \simeq 2 \times \omega_s \simeq 10.7 \mu\text{eV}$. Here, we assume $\omega_s \simeq \omega_p \simeq m_a/2$ (bearing in mind that ω_s and ω_p should not be exactly equal, to enable better pump mode rejection), and that $(\omega_s + \omega_p)$ is aligned with the peak of the axion energy distribution, so that $F_a(\omega_s + \omega_p) = 60\sqrt{2\pi}/e/(17m_a v_{\text{vir}}^2)$ with $v_{\text{vir}} \simeq 9 \times 10^{-4}$ the virial velocity, where $Q_a \simeq 1/v_{\text{vir}}^2$. β is set to the optimal value of 2/3, and we assume all form factors are $\mathcal{O}(1)$.

Some remarks are in order regarding the tuning range of such an experiment. In practice, tuning of SRF cavities is typically achieved via mechanical adjustment of the overall cavity length, with an $\mathcal{O}(\text{MHz})$ maximum tuning range [92]. Alternatives to this do exist, however, notably the SERAPH experiment aims to develop an SRF cavity tunable from 4 to 7 GHz for dark photon searches [23]. Since we must wait for the cavity to ring up after each tuning adjustment, it is preferable to use the largest possible tuning steps, namely, $\mathcal{O}(m_a/Q_a)$. For each such step, we have a sensitivity bandwidth $\Delta f = f_c/Q_s$. As this is narrower than the axion PSD, we do not capture the full signal power available. However, this can nonetheless result in better performance overall since a single tuning step can be sensitive to a range of axion mass values [23]. Taking the benchmark values above, the radiometer equation yields an instantaneous scan rate for Kim-Shifman-Vainshtein-Zakharov (KSVZ) sensitivity at $m_a \simeq 10.7 \mu\text{eV}$ of 1050 kHz/day.

To assess these results, we can first compare with the analogous expression for an ordinary single-mode Primakoff haloscope. The world-leading constraint at $m_a \simeq 10.7 \mu\text{eV}$ comes from CAPP-PACE [21], which has close to KSVZ axion sensitivity at this frequency. Taking their benchmark

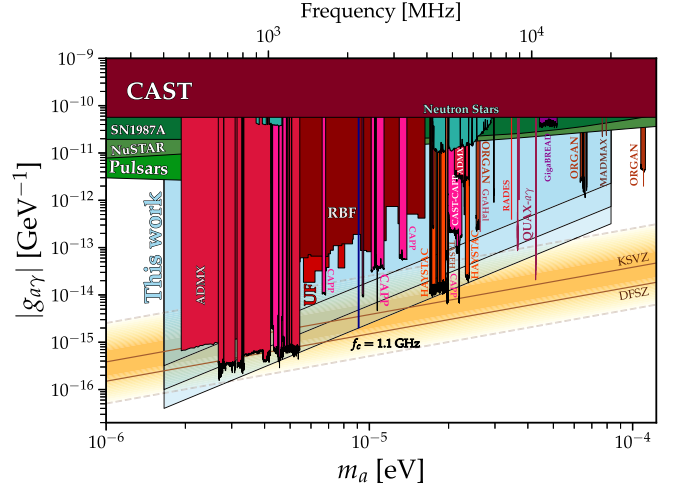


FIG. 1. Sensitivity for resonantly enhanced stimulated axion decay experiments based on our benchmark case, rescaling cavity dimensions to vary f_c from 0.2 to 10 GHz. Due to tuning range limitations multiple experiments would be required to cover this entire range, the dark blue region corresponds to a cavity exactly matching the benchmark case, as described in the text below. Contours shown are for $Q_{\text{int}} \simeq 2 \times 10^9$, $t_{\text{int}} \simeq 100$ s, $T \simeq 2$ K (least sensitive), $Q_{\text{int}} \simeq 2 \times 10^{11}$, $t_{\text{int}} \simeq 100$ s, $T \simeq 2$ K, and $Q_{\text{int}} \simeq 2 \times 10^{11}$, $t_{\text{int}} \simeq 1000$ s, $T \simeq 1$ K (most sensitive). The upper end of this range may be optimistic, since achieving such large Q values at high frequency may not be feasible. Following Refs. [43–46], we also assume sufficient pump mode rejection so as to not impact sensitivity. Figure production uses AxionLimits [58] and Refs. [17,18,20–22,25,29,59–91].

parameters, we find

$$g_{ay}^{95\%} \simeq \frac{10^{-14}}{\text{GeV}} \left(\frac{10^5}{Q} \right)^{\frac{1}{2}} \left(\frac{1.1 \text{ L}}{V} \right)^{\frac{1}{2}} \left(\frac{7.2 \text{ T}}{B_0} \right) \left(\frac{f}{2.6 \text{ GHz}} \right)^{\frac{3}{4}} \times \left(\frac{0.45 \text{ GeV/cm}^3}{\rho_a} \right)^{\frac{1}{2}} \left(\frac{T_{\text{sys}}}{1.2 \text{ K}} \right)^{\frac{1}{2}} \left(\frac{100 \text{ s}}{t_{\text{int}}} \right)^{\frac{1}{4}}, \quad (18)$$

along with a scan rate at KSVZ sensitivity at $m_a \simeq 10.7 \mu\text{eV}$ of 55 kHz/day, with (optimal) $\beta = 2$.

Evidently, this technique holds promise to explore a well-motivated region for QCD axion DM. Going beyond this point of comparison, in Fig. 1 we show the sensitivity reach for various cavity sizes by rescaling our benchmark case to cover the range $f_c = 0.2$ –10 GHz. In Fig. 2, we also reproduce the sensitivity reach from Ref. [44] to schematically show the capability of such experiments to probe axions satisfying $m_a = \omega_s - \omega_p$ and $m_a = \omega_s + \omega_p$.

VII. DISCUSSION AND CONCLUSIONS

The axion is a particularly well-motivated candidate for the DM that makes up most of the matter in our Universe. As a result, experimental efforts are ongoing worldwide toward its discovery. Resonant cavities are a primary tool in this search, where they enhance the rate of axion conversion to photons in a background electromagnetic field.

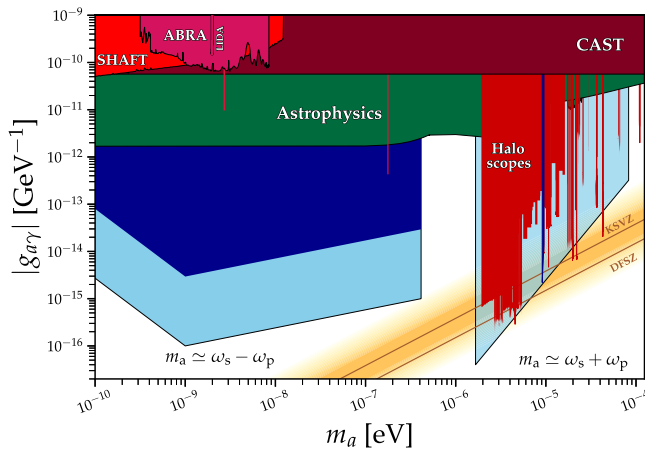


FIG. 2. Sensitivity reach for experiments of this type, where the leftmost light blue region is reproduced from Ref. [44] without modification, and the rightmost light blue region is the most sensitive contour from Fig. 1. Darker blue regions indicate the performance of a single-cavity experiment corresponding exactly to our benchmark case at $f_c = 1.1$ GHz with the “most sensitive” parameter choices from Fig. 1, and assuming 1 h of integration time per e-fold in axion mass values for the left-hand region. The regions shown should be understood as somewhat heuristic, as they depend on specific details about scanning strategy and experimental configuration, which are beyond the current scope of discussion. As we have emphasized, a single experiment can probe parts of both regions shown, potentially simultaneously. Other constraints are from Refs. [93–96].

In this work, we have identified and examined a previously unexplored channel for axion discovery, which also relies upon resonant cavities, by establishing that the rate of spontaneous and stimulated axion decays to two photons inside such a cavity can also be enhanced. As we have discussed, this method allows the parameter space to be explored in a way that is competitive and complementary to other approaches,

with the capability to probe a well-motivated region for QCD axion DM.

Furthermore, we have established that SRF cavity up-conversion experiments of the type already explored in Refs. [43–46], which are designed to search for low-mass axions satisfying $m_a = \omega_s - \omega_p$, already have the capability to search for axions satisfying $m_a = \omega_s + \omega_p$ via this method. As such, an experiment of this type could in principle search two distinct regions of the axion parameter space simultaneously. Because these experiments are already in preparation [46], the scenario we explore here could also soon be directly tested, albeit potentially over a limited frequency range, owing to those experiments being optimized for up- rather than down-conversion.

Going forward, these findings could be improved via further exploration and optimization of potential experimental realizations, especially with regard to enhancing the available tuning range. We also note that the photon pairs produced via axion decay will be polarization entangled, potentially offering further sensitivity enhancements via quantum-enhanced measurement schemes. The applicability of this approach to “light shining through a wall” experiments is also a topic of ongoing investigation.

ACKNOWLEDGMENTS

We thank Yifan Chen, Ariel Arza, and Qiaoli Yang for useful discussions on related topics. This work was supported by the Beijing Natural Science Foundation (Grant No. IS23025) and the National Natural Science Foundation of China (Grant No. 12150410317).

DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

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