

QuGrav: Bringing gravitational waves to light with qumodes

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We propose using qumodes, quantum bosonic modes, for detecting high-frequency gravitational waves via the inverse Gertsenshtein effect, where a gravitational wave resonantly converts into a single photon in a magnetized cavity. For an occupation number n of the photon field in a qumode, the conversion probability is enhanced by a factor of $n + 1$ due to Bose-Einstein statistics. Unlocking this increased sensitivity entails the ability to continuously prepare the qumode and perform nondemolition measurement on the qumode-qubit system within the qumode coherence time. Our results indicate that, at microwave frequencies and with existing technology, the proposed setup can attain sensitivities within 1.7 orders of magnitude of the cosmological bound. With anticipated near-future improvements, it has the potential to surpass this limit and pave the way for the first exploration of high-frequency cosmological gravitational wave backgrounds. At optical frequencies, it can enhance the sensitivity of current detectors by one order of magnitude. That further enhances their potential in reaching the single-graviton level.

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I. INTRODUCTION

At present, precise observations of gravitational waves (GWs) are limited to frequencies below 10 kHz. The detection of higher-frequency GWs (HFGWs) remains an open frontier, currently under active research and development. These HFGWs may arise from cosmological or astrophysical sources; they may also be produced in a laboratory by conversion of high-intensity laser fields. Once detected, HFGWs can reveal a wealth of information about the early Universe, particle physics, and the quantum nature of gravity. Remarkably, electromagnetic (EM) waves and GWs can interconvert in the presence of a static magnetic field—a phenomenon known as the Gertsenshtein effect [1]. Photon regeneration is a leading approach to probe HFGWs, which is based on this mechanism [2]. In such experiments, GWs traverse a photon-shielded cavity permeated by a strong magnetic field, allowing for resonant graviton-to-photon conversion. The current active experiments are Optical

Search for QED vacuum magnetic birefringence, Axions and photon Regeneration (OSQAR) [3] and Any Light Particle Search (ALPS) [4] at optical frequencies, and CERN Axion Solar Telescope (CAST) [5] at x-ray frequencies. In addition, BabyIAXO and International Axion Observatory (IAXO) [6,7] are under development, and Joint Undertaken Research on Axion/ALPs (JURA) [8] is proposed. In Table I, we summarize these experiments. Some of these instruments operate well below the standard quantum limit. Ongoing research and development efforts aim to design next-generation instruments with enhanced power and sensitivity. That includes using stronger magnetic fields, bigger and higher quality cavities, and so on. For a recent review that covers these developments comprehensively, see Ref. [9] and references therein.

In this article, we propose an approach to enhancing the sensitivity of cavity-based gravitational wave searches and reducing the signal-to-noise ratio. Our approach is based on combining the following three ideas developed previously in different fields: (1) photon regeneration, i.e., the resonant conversion of GWs into single photons in a background magnetic field; (2) bosonic stimulation, i.e., an enhancement of a single boson production by a factor of $n + 1$ when the bosonic state has an occupation number n ; and (3) the use of qumodes (states with a fixed number of photons in cavities), which have been developed recently for bosonic quantum computing [10–12].

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The first of these ingredients underlies a number of current and planned experiments, as discussed above. The second one, *bosonic stimulation*, has been established experimentally in the formation of Bose-Einstein condensates [13], and in light scattering in cold atom condensates [14]. The third central pillar of our proposal is n -photon qumodes [15,16]; the measurement of photon number up to ~ 100 has already been demonstrated [17], and is a key ingredient of bosonic quantum computation.

II. THEORETICAL SETUP

The Maxwell-Einstein action describes the EM fields in curved space-time. Consider the space-time metric as

$$g_{\mu\nu} = \eta_{\mu\nu} + h_{\mu\nu}, \quad (1)$$

where $\eta_{\mu\nu} = \text{diag}(-1, 1, 1, 1)$ is the flat space-time metric, $\mu, \nu = 0, 1, 2, 3$, and $h_{\mu\nu}$ denotes small gravitational perturbations, i.e., $|h_{\mu\nu}| \ll 1$. The EM-gravitational interaction Hamiltonian density at first order in $h_{\mu\nu}$ is

$$\mathcal{H}_{\text{int}} = -\frac{1}{2}\eta^{\lambda\sigma}(h^{\mu\nu} - \frac{1}{2}h\eta^{\mu\nu})F_{\mu\lambda}F_{\nu\sigma}, \quad (2)$$

where $F_{\mu\nu}$ is the EM field tensor. Throughout this work, we adopt natural units by setting $c = \hbar = \epsilon_0 = 1$.

Consider a cylindrical cavity of length L and cross-sectional area A . We assume that the cavity is made of a perfect conductor and turn on a constant magnetic field in the z direction inside it: $\mathbf{B}(t, \mathbf{x}) = B_0 \hat{z}$.

At this point, we need to fix the gauge. The frame in which we measure the EM fields in the laboratory is the detector proper frame, i.e., the local inertial frame in which the detector is in free fall. In contrast, theoretical descriptions often employ the transverse-traceless (TT) gauge, where test masses remain at fixed coordinate positions while the metric itself oscillates due to the passing wave. In this work, we adopt the TT gauge, i.e., $h_{00} = h_{0i} = \partial_i h_{ij} = 0$, where $i, j = 1, 2, 3$. We assume that $\mathbf{B}_0$ is static in the TT frame. While, in general, the transformation between frames in the presence of GWs requires careful treatment, it is justified in the high-frequency regime $\omega L \gg v_s$ [18]. Note that the sound velocity inside solids is typically small, i.e., $v_s < 10^{-4}$.

III. PHOTON-GRAVITON TRANSITION AMPLITUDE INSIDE CAVITY

As the initial state, we consider a qumode state made of an n -photon state at a resonant frequency ω :

$$|Q_n\rangle = |n(\omega)\rangle_\gamma, \quad (3)$$

where $|n(\omega)\rangle_\gamma = (b_\omega^\dagger)^n |0\rangle_\gamma$. This state is achievable with current experimental capabilities and is well suited for detecting narrow-band signals. To complement this, we also introduce a more involved qumode configuration to illustrate the theoretical potential and long-term prospects of this approach for broadband detection. The second state is a qumode state made of an n -photon state at each resonant frequency of the cavity, $\omega_l = \frac{\pi l}{L}$, i.e.,

$$|Q_n\rangle = \prod_{l=1}^{\lfloor 2\Delta f L \rfloor} \otimes |n(\omega_l)\rangle_\gamma. \quad (4)$$

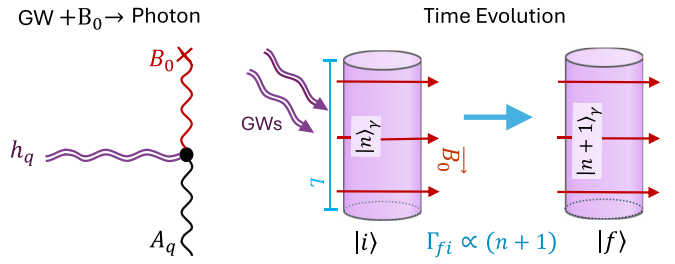


FIG. 1. Left: the inverse Gertsenshtein effect; right: toy model of qumode detector.

where $|n(\omega_l)\rangle_\gamma = (b_{\omega_l}^\dagger)^n |0\rangle_\gamma$ and $\lfloor 2\Delta f L / c \rfloor$ is the number of resonant modes inside ($f = \frac{\omega}{2\pi}$). The gravitational interaction Eq. (2) allows for the conversion between gravitons and photons. The transition amplitude from an initial state $|i\rangle$ to a different final state $|f\rangle$ is given by the S -matrix element

$$\mathcal{A}_{fi} = \langle f | T \exp[-i \int d^4x \mathcal{H}_{\text{int}}] | i \rangle, \quad (5)$$

where T is the time ordering operator. Now, we need to specify the GW signal interacting with the magnetic field in the cavity. In the following, we will consider a gravitational signal passing by, and here we consider three possibilities: (1) a monochromatic GW signal, (2) a stochastic GW field background, and (3) a single-graviton.

(1) *Incoming coherent GWs*: Consider an unpolarized monochromatic GW with frequency ω_g as ($\mathbf{q} = \omega_g \hat{\mathbf{q}}$)

$$h_{ij}(t, \mathbf{x}) = \sum_{\sigma=\pm} e_{ij}^\sigma(\hat{\mathbf{q}}) \hat{h}_{\omega_g} e^{iqx} + \text{H.c.}, \quad (6)$$

where $e_{ij}^\sigma(\hat{\mathbf{q}})$ are the polarization tensors, $\sigma = \pm$ are the circular polarizations of the GWs, and $\hat{h}_{\omega_g}$ is a dimensionless complex-valued operator such that $\hat{h}_{\omega_g} |\alpha\rangle_g = h_\omega |\alpha\rangle_g$, where $|\alpha\rangle_g$ is a coherent GW state, h_ω is the mode function, and $\langle h_{ij}(t, \mathbf{x}) h_{ij}(t, \mathbf{x}) \rangle = 2|h_\omega|^2$. The interaction Hamiltonian density can be written as

$$\mathcal{H}_{\text{int}} = -B_0 \sin \theta_q \sum_{\sigma} \delta B_i(x) e_i^\sigma(\hat{\mathbf{q}}) \hat{h}_{\omega_g} e^{iqx}, \quad (7)$$

where $e_i^\pm(\hat{\mathbf{q}}) = \frac{1}{\sqrt{2}}(\cos \theta_q \cos \phi_q \mp i \sin \phi_q, \cos \theta_q \sin \phi_q \pm i \cos \phi_q, -\sin \theta_q)$ in Cartesian coordinates.

Here, the initial and final photon-graviton states are

$$|i\rangle = |Q_n\rangle_\gamma |\alpha\rangle_g, \quad |f\rangle = b_{\omega_g}^\dagger |Q_n\rangle_\gamma |\alpha\rangle_g, \quad (8)$$

where a photon is created at frequency ω_g , raising the occupation number at this frequency from n to $n+1$. Given that $|\langle n+1 | b^\dagger | n \rangle|^2 \propto (n+1)$, the transfer rate is

$$\Gamma_{fi} = 2(n+1) Q V B_0^2 h_{\omega_g}^2 \sin^2 \theta_q, \quad (9)$$

where Q is the quality factor of the cavity (see Fig. 1). Remarkably, due to Bose-Einstein statistics, once the frequency of the GW matches the frequency of the n -photon mode inside the cavity, the probability of graviton-to-photon conversion is enhanced by a factor of $n+1$. This process is the EM analog to the stimulated absorption of gravitons in the massive quantum acoustic resonators, with the process being most effective at frequencies that match the resonator's natural frequency

[19]. Stimulated conversion in the context of dark matter searches has been demonstrated with $n = 4$ photon states in a microwave cavity [20]. Our work opens a direction by introducing stimulated graviton-photon conversion and systematically exploring different frequency ranges to identify the most promising regime for gravitational wave detection. While a bright coherent drive also enhances the generation rate by the factor $n + 1$, it simultaneously introduces amplitude and phase noise. A quantitative comparison is given in the Supplemental Material [21].

(2) *Cosmic GW background*: A fraction of high-frequency gravitational waves is expected to originate from early Universe phenomena such as inflation or phase transitions. Their total energy density is constrained by cosmological observations, particularly through limits on the effective number of relativistic species, $\Delta N_{\text{eff}} < 0.3$, which gives $h(f) < 2 \times 10^{-30} (\frac{1 \text{ GHz}}{f}) \Delta N_{\text{eff}}^{\frac{1}{2}}$. If the GW background is dominated by a localized feature in frequency, it can be treated as a narrow-band source, with the transition rate given by Eq. (9). Otherwise, it generally exhibits a broadband spectrum, for which we compute the transition rate below.

Consider an unpolarized stochastic background of GWs with cosmological origin. It can be described as

$$h_{ij}(x) = \sum_{\sigma=\pm} \int d \ln \omega e_{ij}^{\sigma}(\hat{q}) \int \frac{d^2 \hat{q}}{2\pi} \hat{h}_{\omega} e^{iqx} + \text{H.c.}, \quad (10)$$

where $\mathbf{q} = \omega \hat{\mathbf{q}}$, h_{ω} is the mode function, and $\langle h_{ij}(t, \mathbf{x}) h_{ij}(t, \mathbf{x}) \rangle = 2 \int d \ln \omega |h_{\omega}|^2$. Let Δf denote the finite frequency range of the detector. For simplicity, assume the GW spectrum is flat across this interval. To improve the sensitivity in this case, we use Eq. (4) and have to sum over all the resonant modes. The total transfer rate is

$$\Gamma_{\text{tot}} = 2(n+1) \Delta f L V B_0^2 h_{\omega}^2 \sin^2 \theta_q. \quad (11)$$

(3) *Single gravitons*: Modern detector architectures can detect single photons, but building a detector sensitive to single gravitons remains an extraordinary challenge [22,23]. More precisely, if gravity is quantized, then the number of gravitons per de Broglie volume is

$$N_{\text{grav}} = n_{\text{grav}} \lambda_{\text{dB}}^3 = \frac{\pi h_{\omega}^2 M_{\text{pl}}^2}{2 f^2}, \quad (12)$$

where $\lambda_{\text{dB}} = f^{-1}$. As an example, a Laser Interferometer Gravitational-Wave Observatory (LIGO) detection signal with $h_{\omega} \sim 10^{-25}$ at $f \sim 10^2$ Hz includes $N_{\text{grav}} \sim 10^{29}$ gravitons. Requiring less than a graviton per de Broglie volume, we find

$$h_{\omega} \leq 2 \times 10^{-31} \left(\frac{f}{10^3 \text{ GHz}} \right), \quad (13)$$

where below that limit, the gravitational field is more accurately described as a highly dilute graviton gas rather than a classical gravitational wave. Notably, establishing the quantum nature of gravitational radiation demands measurements far more challenging than merely detecting individual gravitons [24].

TABLE I. Experimental parameters of the photon-regeneration setups. Here, ε is the single-photon detection efficiency and $\Gamma_D(f)$ is the dark count rate in a frequency bin of width Δf .

	f (Hz)	B_0 (T)	L (m)	A (m ²)	Γ_D (mHz)	ε
OSQAR II	10^{15}	9	14.3	5×10^{-4}	1.14	0.9
ALPS II	10^{15}	5.3	106	10^{-3}	10^{-3}	0.75
JURA	10^{15}	13	960	8×10^{-3}	10^{-6}	1
CAST	10^{18}	9	9.26	2.9×10^{-3}	0.15	0.7
IAXO	10^{18}	2.5	20	3.08	0.1	1
QuGrav	10^{10}	10	1	1	10^{-3}	1

Consider a single graviton propagating in direction $\hat{\mathbf{q}}$ with frequency ω_g , i.e., $|1_{\omega_g}\rangle_g$, as

$$h_{ij}(t, \mathbf{x}) = \sum_{\sigma=\pm} e_{ij}^{\sigma}(\hat{\mathbf{q}}) h_{\omega_g}^{\sigma} \hat{a}_{\omega_g, \sigma} e^{-i\omega_g t + i\mathbf{q} \cdot \mathbf{x}} + \text{H.c.}, \quad (14)$$

where $\hat{a}_{\omega, \sigma}^{\dagger}$ and $\hat{a}_{\omega, \sigma}$ are the graviton creation and annihilation operators, $[\hat{a}_{\omega, \sigma}, \hat{a}_{\omega', \sigma'}^{\dagger}] = \delta_{\sigma\sigma'} \delta(\omega - \omega')$. Here, we set the initial and final states as

$$|i\rangle = |Q_n\rangle |1_{\omega_g}\rangle_g, \quad |f\rangle = \hat{b}_{\omega_g}^{\dagger} |Q_n\rangle |0\rangle_g, \quad (15)$$

and the transition rate is equal to Eq. 9.

IV. SENSITIVITY AND DETECTION PROSPECT

Here, we compute upper limits on the GW dimensionless characteristic strain, $h(f)$, based on the characteristics of each experiment, with the application of the qumode technology. These experiments detect the photon count rate Γ within an energy band $\Delta f = (f_f - f_i)$, with a single-photon detection efficiency ε , and over a cross-sectional area A . The background noise is quantified in terms of the dark count rate in that frequency bin, i.e., $\Gamma_D(f)$. It is the rate at which the detector falsely registers a photon even when no real photon is present. See Table I for the experimental parameters of the photon-regeneration setups. In the following, we refer to the implementation of this technology in each GW detector by appending the prefix “Qu” to its name. For example, the upgraded version of JURA incorporating this technology is denoted as QuJURA. We call these updated experiments collectively as Qu-detectors.

For completeness, we note that n -photon qumodes in the presence of a magnetic field also introduce a subleading noise contribution due to graviton emission, at a rate $\Gamma_{\text{rev}} = \frac{2n}{M_{\text{pl}} \sqrt{V \omega}} Q V B_0^2$. This process introduces a photon-number ambiguity, leading to an additional noise contribution of order $\frac{2\Delta t}{M_{\text{pl}} \sqrt{V \omega}} Q V B_0^2$. As this effect is subleading compared to thermal and other dominant noise sources, we neglect it in the present analysis. While genuine photon-generation backgrounds would also be subject to Bose enhancement, they are suppressed in the cryogenic, shielded regime considered here; the dominant background arises from detector readout noise, which does not create photons in the mode and is therefore not Bose-enhanced (see the Supplemental Material [21]).

V. NARROW-BAND SNR

For resonant detectors with narrow bandwidth sensitivity, a monochromatic gravitational field, or a GW background with a sharply peaked spectrum, Signal-to-noise ratio (SNR) reduces to a contribution from a narrow frequency band $\Delta f = \frac{f}{Q} = \frac{1}{L\mathcal{F}}$, where $\mathcal{F}$ is the finesse of the cavity. Now using Eq. (9) and demanding $\varepsilon\Gamma_{fi} = \Gamma_D$, we have

$$h_{\text{noise}}(f) \approx \frac{1.4 \times 10^{-27}}{\sqrt{n+1}} \left(\frac{1\text{ m}}{L} \right) \left(\frac{10\text{ T}}{B_0} \right) \sqrt{\left(\frac{1\text{ m}^2}{A} \right) \left(\frac{10^5}{\mathcal{F}} \right)} \times \sqrt{\frac{1}{\varepsilon} \left(\frac{\Gamma_D}{1\text{ }\mu\text{Hz}} \right) \left(\frac{1\text{ GHz}}{f} \right)}. \quad (16)$$

VI. BROADBAND SNR

Using the theoretical benchmark qumode state Eq. (4) with Eq. (11) for detectors and signals with broad frequency ranges, we can quantify the SNR by demanding that $\varepsilon\Gamma_{\text{tot}} > \Gamma_D$; we find

$$h_{\text{noise}}(f) \approx \frac{4.5 \times 10^{-25}}{\sqrt{(n+1)}} \left(\frac{1\text{ m}}{L} \right) \left(\frac{10\text{ T}}{B_0} \right) \sqrt{\left(\frac{1\text{ GHz}}{\Delta f} \right)} \times \sqrt{\frac{1}{\varepsilon} \left(\frac{\Gamma_D}{1\text{ }\mu\text{Hz}} \right) \left(\frac{1\text{ m}^2}{A} \right)}. \quad (17)$$

Here, we can compare the qumode implementation of a given detector with its standard single-photon detector:

$$h_{\text{noise}}^{\text{single}} / h_{\text{noise}}^{\text{Qu}} = (n+1)^{-\frac{1}{2}}. \quad (18)$$

Consequently, the use of n -photon states reduces the noise-equivalent strain by a factor of $(n+1)^{-1/2}$. For illustration, we benchmark our estimates with $n = 100$, already demonstrated in the optical domain [17], which improves the sensitivity by about an order of magnitude. While such large- n states are still under development in the microwave regime, steady progress in cavity QED and superconducting systems suggests they represent a realistic long-term target. This highlights the potential of n -photon qumodes to significantly enhance sensitivity across a broad range of frequencies.

The enhanced sensitivity, however, comes at the price of continuously rebuilding the qumode through its coupling to a single qubit, and performing a quantum nondemolition measurement on it within its coherence time [25]. Indeed, it has been demonstrated that an additional single photon can have a big effect on the qubit without ever being absorbed [25]. At present, the demonstrated qumode coherence time already exceeds 10 ms for both microwave [26] and optical cavities [27]. However, at x-ray frequencies, the lifetimes of n -photon states are short compared to the time needed for their preparation, making this technology currently unfeasible for CAST and IAXO.

(1) *QuOSQAR II* and *QuALPS II* (10⁶ GHz): The sensitivities of OSQAR II (active) and ALPS II (in development) probing optical frequencies are, respectively, $h \approx 6 \times 10^{-26}$ and $h \approx 3 \times 10^{-30}$. One of the key advancements in ALPS II is the implementation of a Fabry-Perot cavity with $\mathcal{F} = 4 \times 10^4$ to enhance the photon-regeneration process. The pro-

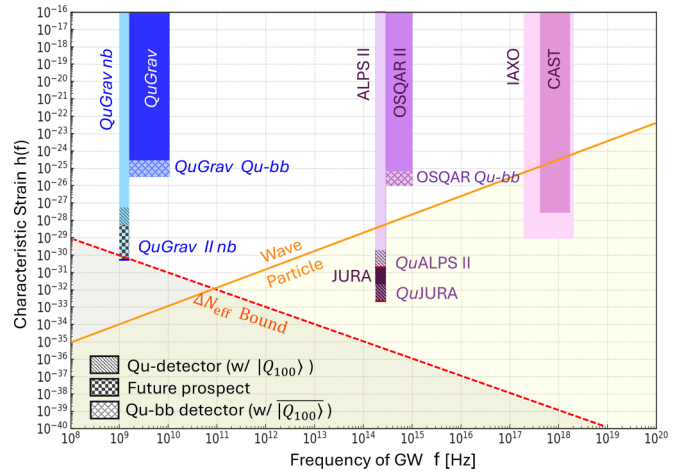


FIG. 2. Strain sensitivity of HFGW detectors based on photon regeneration. The area below the orange line corresponds to gravitational fields that behave like quantum particles (gravitons). The dashed gray curve denotes the ΔN_{eff} cosmological bound. Hatched regions indicate projected improvements in sensitivity from current experiments and proposals employing a 100-photon qumode $|Q_{100}\rangle$ in Eq. (3) within the cavity. The checked (black-blue) represents a sensitivity that relies on technological progress in cavity development anticipated in the near future, i.e., increasing S by 1.7 orders of magnitude [see Eq. (20)]. The crossed regions indicate the enhanced sensitivity to broadband signals achievable with the futuristic broadband 100-photon qumode $|Q_{100}\rangle$ described in Eq. (4).

posed JURA experiment with $\mathcal{F} = 10^5$ aims to reach a strain sensitivity as low as $h \approx 2 \times 10^{-32}$ [2]. As we see in Fig. 2, the ALPS II and JURA are in the quantum particle limit. By employing $n = 100$ qumode states, the sensitivity can be enhanced by an order of magnitude, making the detectors even more responsive to signals of single gravitons.

(2) *QuGrav* (1–10 GHz): Despite the technical challenges at GHz frequencies, rapid progress is underway, driven in part by dark photon and dark matter searches [28,29]. The main challenge is that the photon's energy at this frequency is very low, i.e., around 10^{-6} eV (about 10 mK). Therefore, to dominate over the thermal noise, we need cryogenic temperatures. In addition, even a micron-scale mechanical shift can detune the cavity. Therefore, it requires active feedback stabilization, temperature control, and vibration isolation, which makes it challenging. Note, however, that the cavities coupled to superconducting qubits operate at millikelvin temperatures, and there is a lot of ongoing work on the cryogenic microwave cavities motivated mainly by the needs of bosonic quantum computing (see, e.g., Refs. [15,16]). The minimal frequency of a photon inside the cavity is $f_{\text{min}} \sim 10^{-1} \text{ GHz} \left(\frac{L}{1\text{ m}} \right) \left(\frac{1\text{ m}^2}{A} \right)$ [30]; hence, we will consider frequencies $f \gg f_{\text{min}}$.

Here, we make estimates for a toy detector operating with currently accessible parameters $B_0 = 10\text{ T}$, $\varepsilon = 1$, $\Gamma_D = 1\text{ }\mu\text{Hz}$, $L = 1\text{ m}$, $\mathcal{F} = 5 \times 10^5$, and $A = 1\text{ m}^2$ operating at millikelvin temperature. Setting the resonant frequency at $f = 10f_{\text{min}} = 1\text{--}10\text{ GHz}$, we find the narrow-band sensitivity to h as

$$h_{\text{noise}} \approx 6.3 \times 10^{-29} S^{-1}, \quad (19)$$

where we define

$$S = \left(\frac{L}{1 \text{ m}} \right) \sqrt{\left(\frac{n+1}{10^2} \right) \left(\frac{\mathcal{F}}{5 \times 10^5} \right) \left(\frac{1 \mu\text{Hz}}{\Gamma_D} \right)}. \quad (20)$$

For the currently achievable value of $n = 10^2$, h_{noise} can be as low as 6.3×10^{-29} , which is only 1.7 orders of magnitude above the ΔN_{eff} bound on GWs. While the precise trajectory of future developments remains uncertain, it is natural to expect that continued advancements in qumode technology and improvements in microwave cavity quality at cryogenic temperatures could enhance S by 1.7 orders of magnitudes in the near future. It could then push the sensitivity below the ΔN_{eff} bound, potentially opening a window into the cosmic gravitational wave background, characterized by a sharp peak at this frequency. Enhancing sensitivity across a broadband spectrum is more challenging and necessitates the use of the qumode configuration given in Eq. (4). In Fig. 2, the sensitivity to the narrow band is labeled as *QuGrav nb*, while the sensitivity to the broadband spectrum appears as *QuGrav Qu-bb*.

VII. SUMMARY AND CONCLUSION

In this article, we propose an approach to detecting high-frequency gravitational waves using qumodes, which are used in continuous-variable quantum computation. The detection mechanism relies on the inverse Gertsenshtein effect, wherein a GW is resonantly converted into a single photon within a cavity subjected to a static magnetic field. In this framework, the presence of an n -photon qumode leads to an enhancement in the conversion probability by a factor of $n + 1$, owing to Bose-Einstein statistics. The price to pay for the increased sensitivity is the necessity to rebuild the qumode within its coherence time.

Our analysis based on the qumode state defined in Eq. (3) for narrow-band spectra shows that, at microwave frequencies and with currently achievable technology, the proposed setup can reach sensitivities up to 6.3×10^{-29} [Eq. (16)], i.e., within 1.7 orders of magnitude of the current cosmological bounds. With near-future advancements and improving $L\sqrt{(n+1)\mathcal{F}/\Gamma_D}$, it may even surpass these bounds, opening a window to probe the stochastic GW background of cosmological origin. At optical regimes, the same framework could enhance the sensitivity of existing narrow-band detectors by an order of magnitude [see Eq. (19)], potentially allowing them to approach the single-graviton detection threshold.

Enhancing sensitivity across a broadband spectrum presents greater challenges and calls for further advancements in the generation and stabilization of qumodes. Using the theoretical benchmark qumode state introduced in Eq. (4), we compute the corresponding transition rate in Eq. (11) and sensitivity in Eq. (17), illustrating its long-term potential for HFGW detection in Fig. 2.

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DATA AVAILABILITY

No data were created or analyzed in this study.

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