

Quantum Metrology in the Ultrastrong Coupling Regime of Light-Matter Interactions: Leveraging Virtual Excitations without Extracting Them

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Virtual excitations, inherent to ultrastrongly coupled light-matter systems, induce measurable modifications in system properties, offering a novel resource for quantum technologies. In this Letter, we demonstrate how these virtual excitations and their correlations can be harnessed to enhance precision measurements, without the need to extract them. Building on the paradigmatic Dicke model, which describes the interaction between an ensemble of two-level atoms and a single radiation mode, we propose a method to harness hybridized light-matter modes whose renormalized frequencies encode the effects of virtual excitations for quantum metrology. Remarkably, we find that, for a fixed squeezing parameter ξ , exploiting virtual squeezing through oscillator frequency shifts yields a quadratic enhancement in estimation precision—scaling as $\exp(4\xi)$ —compared to the conventional $\exp(2\xi)$ scaling of real squeezed states. These results show that virtual excitations, though unobservable, can drive metrological performance beyond the standard quantum limit. Our approach establishes a broadly applicable framework for high-precision measurements across a wide class of ultrastrongly coupled quantum systems.

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Introduction—Recent theoretical and experimental advances in understanding and controlling light-matter interactions have brought us closer to exploring the physics of ultrastrong coupling, a regime in which the light-matter coupling strength becomes comparable to the system's transition frequencies [1–3]. While the weak and strong coupling regimes have already driven significant progress in quantum technologies [4,5], the ultrastrong coupling regime remains largely unexplored. This regime gives rise to counterintuitive and intriguing phenomena [6–18] including highly entangled ground states. However, this entanglement is purely virtual, meaning that it does not correspond to directly measurable correlations between atoms and photons. Instead, it manifests indirectly through modifications to observable system properties, such as nonlinear frequency shifts of normal modes, which depend on the system's parameters [19,20]. If such modifications to physical excitations can be precisely controlled, they could potentially enable novel applications in quantum technologies.

In this Letter, we propose a method to leverage virtual excitations in ultrastrongly coupled systems for precision measurements. We consider an ensemble of atoms ultrastrongly coupled to a single mode of radiation, modeled by the paradigmatic Dicke model. Our approach involves coupling the to-be-measured field with a control system to form hybridized modes and performing all measurements on these hybrid modes instead of extracting virtual excitations. The key aspect is the nontrivial dependence of the hybrid mode frequency on the system's parameters as well as on virtual excitations. This Letter introduces a novel direction in quantum metrology, demonstrating how unique properties of ultrastrong coupling can be harnessed for practical applications.

Dicke model in the ultrastrong coupling regime—The Dicke model describes an ensemble of two-level systems collectively coupled to a single mode of electromagnetic radiation [21]. Typically, the coupling strength between light and matter is much smaller than the atomic transition frequency, allowing the use of the rotating wave approximation [22]. This approximation eliminates counterrotating terms, reducing the Dicke model to the Tavis-Cummings model. In the ultrastrong coupling regime, where the coupling strength is comparable to the characteristic frequencies of the system, the rotating wave approximation breaks down, and counterrotating terms must be included [23]. This inclusion fundamentally alters the system's properties. To understand this more clearly, we now consider the Dicke

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model Hamiltonian ($\hbar = 1$)

$$\hat{H} = \omega \hat{a}^\dagger \hat{a} + \Omega \hat{S}_z + \frac{g}{\sqrt{N}} (\hat{a} + \hat{a}^\dagger) \hat{S}_x, \quad (1)$$

where $\hat{a}$ ($\hat{a}^\dagger$) is the annihilation (creation) operator for a photon with frequency ω , $\hat{S}_i = \sum_{n=1}^N \hat{\sigma}_i^n / 2$ —where $\hat{\sigma}_i^n$ is the i th ($i = x, y, z$) Pauli matrix describing the n th atom—are collective spin operators representing N two-level systems with transition frequency Ω , and g is the vacuum Rabi coupling strength. Close to the resonance, $\omega \approx \Omega$, the ultrastrong coupling regime is characterized by $g \lesssim \omega \approx \Omega$, approaching the critical coupling strength $g_c = \sqrt{\omega\Omega}$ for the superradiant phase transition [24].

For large ensembles ($N \gg 1$), the collective spin operators can be approximated by harmonic oscillator operators, which becomes exact in the thermodynamic limit ($N \rightarrow \infty$). Under this approximation, the Dicke Hamiltonian takes the form

$$\hat{H} = \omega \hat{a}^\dagger \hat{a} + \Omega \hat{b}^\dagger \hat{b} + \frac{g}{2} (\hat{a} + \hat{a}^\dagger) (\hat{b} + \hat{b}^\dagger), \quad (2)$$

where $\hat{b}$ and $\hat{b}^\dagger$ represent the collective atomic mode. The ground state of this Hamiltonian is a two-mode squeezed vacuum state, which under the resonance condition ($\omega = \Omega$) can be written as [25]

$$|\xi_-, \xi_+\rangle = \exp\left[\frac{\xi_-}{2} (\hat{c}^2 - \hat{c}^{\dagger 2})\right] \exp\left[\frac{\xi_+}{2} (\hat{d}^2 - \hat{d}^{\dagger 2})\right] |0_a, 0_b\rangle, \quad (3)$$

where $\hat{c} = (\hat{a} - \hat{b})/\sqrt{2}$ and $\hat{d} = (\hat{a} + \hat{b})/\sqrt{2}$ are the light-matter modes that are being squeezed, $\xi_\pm = \frac{1}{4} \log(1 \pm g/g_c)$ are the corresponding squeezing parameters (see Ref. [23] for a general case), and $|0_a\rangle$ and $|0_b\rangle$ represent the vacuum states of mode $\hat{a}$ and $\hat{b}$, respectively.

From the viewpoint of the noninteracting modes $\hat{a}$ and $\hat{b}$, such a ground state contains excitations

$$\langle \hat{a}^\dagger \hat{a} \rangle = \langle \hat{b}^\dagger \hat{b} \rangle = \frac{1}{2} (\sinh^2 \xi_- + \sinh^2 \xi_+) \quad (4)$$

in a correlated form of two-mode squeezing between light ($\hat{a}$) and matter ($\hat{b}$). However, these excitations and the two-mode squeezing are virtual. They exist only on the level of the noninteracting fields $\hat{a}$ and $\hat{b}$ with frequencies ω and Ω . In the ultrastrong coupling limit, however, these modes no longer correspond to any directly measurable quantities; i.e., no correlation between $\hat{a}$ and $\hat{b}$ can be detected directly. The Hamiltonian that describes measurable quantities can be constructed with a standard Bogoliubov transformation

$$\hat{H} = \omega_- \hat{e}^\dagger \hat{e} + \omega_+ \hat{f}^\dagger \hat{f}, \quad (5)$$

where

$$\omega_\pm = \sqrt{\frac{1}{2} \left(\omega^2 + \Omega^2 \pm \sqrt{(\omega^2 - \Omega^2)^2 + 4g^2 \omega \Omega} \right)} \quad (6)$$

are the new frequencies of the physical (normal) modes $\hat{e}$ and $\hat{f}$, typically dubbed as polaritons in the context of light-matter systems [26]. Note that the normal frequencies are nontrivial functions of ω and Ω . In the resonant case, where $\omega = \Omega$, the normal modes are described by

$$\hat{e} = \exp\left[\frac{\xi_-}{2} (\hat{c}^{\dagger 2} - \hat{c}^2)\right] \hat{c} \exp\left[\frac{\xi_-}{2} (\hat{c}^2 - \hat{c}^{\dagger 2})\right] \quad (7)$$

and

$$\hat{f} = \exp\left[\frac{\xi_+}{2} (\hat{d}^{\dagger 2} - \hat{d}^2)\right] \hat{d} \exp\left[\frac{\xi_+}{2} (\hat{d}^2 - \hat{d}^{\dagger 2})\right]. \quad (8)$$

Their corresponding frequencies simplify to

$$\omega_\pm = \omega \sqrt{1 \pm \frac{g}{\omega}} = \omega \sqrt{1 \pm \frac{g}{g_c}} = \omega \exp(2\xi_\pm) \quad (9)$$

and are related to the number of virtual excitations in modes $\hat{c}$ and $\hat{d}$ through the squeezing parameters

$$\langle \hat{c}^\dagger \hat{c} \rangle = \sinh^2 \xi_- \quad \text{and} \quad \langle \hat{d}^\dagger \hat{d} \rangle = \sinh^2 \xi_+. \quad (10)$$

In terms of the normal modes, the ground state contains no excitations, $\langle \hat{e}^\dagger \hat{e} \rangle = \langle \hat{f}^\dagger \hat{f} \rangle = 0$, as it should be for the ground state of the system. At first glance, this virtual two-mode squeezing might seem practically unusable. Converting the virtual excitations into real, physical ones could, in principle, be achieved using techniques such as the dynamical Casimir effect [27,28], for example, by rapidly modulating the light-matter interaction strength g . However, as we demonstrate below, this conversion is not strictly necessary as modifications to system properties induced by these correlated virtual excitations—namely, the frequency shifts—can themselves be leveraged for applications in quantum sensing (see Fig. 1 for a schematic illustration).

Quantum Fisher information perspective—Given the symmetry between ω and Ω , we focus on measuring the frequency of the noninteracting radiation mode ω only. By either directly measuring or converting the virtual excitations into real ones, their squeezing properties can be harnessed for precision measurements. For example, in critical metrology [31–37], the ground state of the system, $|\psi_0\rangle$, near the critical point of a phase transition is used to calculate the quantum Fisher information with respect to ω :

$$\mathcal{I}_\omega = 4(\langle \partial_\omega \psi_0 | \partial_\omega \psi_0 \rangle - |\langle \partial_\omega \psi_0 | \psi_0 \rangle|^2), \quad (11)$$

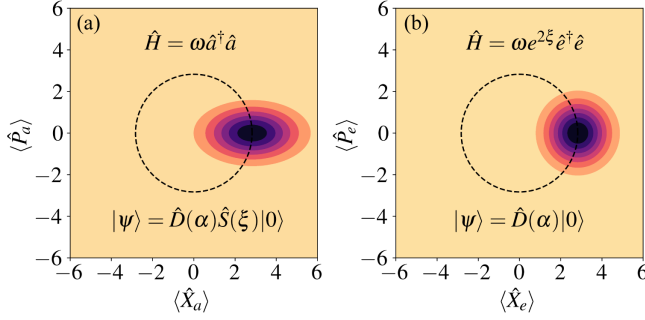


FIG. 1. Schematic illustrating the concept. (a) In conventional quantum metrology, explicit squeezing reduces quantum noise along one direction. (b) Alternatively, virtual squeezing, manifested as a modified system frequency, can be harnessed. In this approach, the signal is amplified through a nontrivial dependence of the effective frequency $\omega e^{2\xi}$ on the unknown parameter ω through the squeezing parameter ξ . For the same level of squeezing, characterized by ξ , and simple quadrature measurements, virtual squeezing offers improved sensitivity [29] and exhibits greater robustness against noise and decoherence [30].

which determines the sensitivity of estimating the parameter ω through the quantum Cramér-Rao bound $\Delta\omega \geq 1/\sqrt{\mathcal{I}_\omega}$. For the Dicke model's ground state, the quantum Fisher information is given by

$$\mathcal{I}_\omega = 2(\partial_\omega \xi_-)^2 + 2(\partial_\omega \xi_+)^2, \quad (12)$$

which reaches its maximum near the critical point of the superradiant phase transition, where

$$\mathcal{I}_\omega \approx 2(\partial_\omega \xi_-)^2 = \frac{g^2}{8\omega^4(1 - g/g_c)^2}. \quad (13)$$

However, leveraging this sensitivity relies on the ability to measure virtual excitations represented by $\hat{a}$ and $\hat{b}$ [2,38].

Here, we demonstrate that transforming virtual particles into measurable ones is not a prerequisite for utilizing them in precision measurements. Instead, measuring the modified properties of physical systems can already yield significant enhancements over weakly and strongly coupled light-matter systems. In order to measure the properties of hybrid modes, the system must be probed. This can be accomplished by using a laser drive tuned near the resonance frequency of a normal mode. When the modes are spectrally well separated, focusing on a single mode becomes feasible. Specifically, we concentrate on the lower polariton mode, as its frequency exhibits a pronounced sensitivity to variations in the bare frequency ω , particularly near the superradiant threshold. The Hamiltonian describing such an isolated normal mode system is given by

$$\hat{H} = \omega_- \hat{e}^\dagger \hat{e}. \quad (14)$$

For simplicity, we now assume that the system is excited into a coherent state $|\alpha\rangle$, satisfying $\hat{e}|\alpha\rangle = \alpha|\alpha\rangle$, where $|\alpha|^2$ is the number of real (measurable) excitations. The subsequent dynamic is governed by the Schrödinger equation using Hamiltonian (14). Under these assumptions, the state of the system as a function of time is given by

$$|\psi(t)\rangle = \exp(-i\omega_- t \hat{e}^\dagger \hat{e})|\alpha\rangle = |\alpha \exp(-i\omega_- t)\rangle. \quad (15)$$

Calculating the quantum Fisher information with respect to the parameter of interest ω yields

$$\mathcal{I}_\omega = 4t^2 \times |\alpha|^2 \times (\partial_\omega \omega_-)^2, \quad (16)$$

which we deliberately separate into three components. The first term, t^2 , highlights that better resolution can be achieved by allowing the system to interact with a perturbation for a longer duration. The second term, $|\alpha|^2$, shows that sensitivity in parameter estimation can be improved by optimizing the initial state. The third term indicates that sensitivity is further enhanced when the frequency is a function of ω with a large derivative. While the first two terms are commonly exploited to enhance sensitivity [39–42], the potential of the third term to increase the quantum Fisher information has been so far overlooked. In the weak and strong coupling regimes, frequency shifts exhibit only weak dependence on the field frequency [22]. As a result, the derivative of the effective frequency with respect to the bare frequency remains close to unity, limiting the quantum Fisher information to the standard quantum limit $\mathcal{I}_\omega = 4t^2 \times |\alpha|^2$. This is not the case in the ultrastrong coupling limit, where the derivative can be extremely large close to the threshold point $g \sim g_c$:

$$\partial_\omega \omega_- = \frac{2 - g/\omega}{2\sqrt{1 - \frac{g}{\omega}}}, \quad (17)$$

and the quantum Fisher information becomes

$$\mathcal{I}_\omega = 4t^2 \times |\alpha|^2 \times \frac{(2 - g/\omega)^2}{4(1 - g/\omega)}, \quad (18)$$

which for $g \approx g_c$ can be approximated as

$$\mathcal{I}_\omega \approx 4t^2 \times |\alpha|^2 \times \frac{1}{4\left(1 - \frac{g}{g_c}\right)}. \quad (19)$$

The above quantum Fisher information can surpass the standard quantum limit, which can be achieved only with nonclassical resources [41,43]. The enhancement originates from the quantum effects of correlated virtual excitations which form the hybrid light-matter modes

$$\langle \hat{n}_c \rangle \equiv \langle \hat{e}^\dagger \hat{e} \rangle = \sinh^2 \xi_- = \frac{\left(1 - \sqrt{1 - \frac{g}{g_c}}\right)^2}{4\sqrt{1 - \frac{g}{g_c}}}. \quad (20)$$

Therefore, the quantum Fisher information can be expressed with the virtual excitations $\langle \hat{n}_c \rangle$ or, equivalently, the virtual squeezing parameter ξ_- as

$$\mathcal{I}_\omega \approx 16t^2 \times |\alpha|^2 \times \langle \hat{n}_c \rangle^2 = t^2 \times |\alpha|^2 \times \exp(-4\xi_-), \quad (21)$$

exhibiting Heisenberg scaling with respect to the number of virtual excitations [44]. Remarkably, suppressing noise using the same squeezing parameter ξ_- —which is equivalent to extracting the squeezed virtual excitations from the ground state and displacing it by α [see Fig. 1(a)]—yields

$$\mathcal{I}_\omega \approx 16t^2 \times |\alpha|^2 \times \langle \hat{n}_c \rangle = 4t^2 \times |\alpha|^2 \times \exp(-2\xi_-), \quad (22)$$

leading to quantum Fisher information that scales only linearly with the number of virtual excitations. In other words, for a fixed squeezing parameter ξ_- , modifying the oscillator frequency—thereby keeping the squeezing excitations virtual—enhances the quantum Fisher information more effectively than reducing quantum noise via real squeezing (see also Supplemental Material [29] for a more detailed comparison). This indicates that retaining squeezing in its virtual form can be more advantageous than extracting and directly converting the virtual excitations into real ones.

However, this enhancement requires operating very close to the critical point, in accordance with the principles of critical metrology [33]. Note that the above analysis assumes an isolated system, coherent evolution, and pre-loaded excitations and neglects the measurement process. A more realistic experimental approach is discussed in the next section.

Signal-to-noise ratio—The quantum Fisher information establishes the fundamental precision bound for parameter estimation, optimized over all measurements [45], even the most abstract ones. In essence, it is the maximum signal-to-noise ratio achievable for a given system [46]

$$\mathcal{I}_\omega = \max_{\hat{O}} S(\omega), \quad \text{where } S(\omega) = \frac{|\partial_\omega \langle \hat{O} \rangle|^2}{\Delta^2 \hat{O}}, \quad (23)$$

with $\hat{O}$ being a measured observable. Taking into account practical considerations, one cannot perform arbitrary measurements on the system. In particular, direct measurements on virtual excitations are not feasible. In the following case, however, we always deal with a Gaussian system and Gaussian states. Hence, all the information about the parameters of the system is contained in the quadrature operators [29]. Therefore, homodyne and heterodyne detection schemes constitute optimal measurements (saturating the quantum Cramér-Rao bound) in this scenario, for which the quantum Fisher information equals the signal-to-noise ratio [47].

Assuming that the two frequencies ω_- and ω_+ are well separated, we can consider only one mode of the system. We focus again on the lower normal mode described by $\hat{e}$.

Therefore, let us calculate the signal-to-noise ratio for the measurement of the $\hat{X} = (\hat{e} + \hat{e}^\dagger)/2$ quadrature for which we also assume that the system is open and loses excitations—a typical situation in cavity quantum electrodynamics experiments [48,49]. In the frame rotating with the frequency of the pumping laser (assumed to be in a coherent state), the Hamiltonian of the system is given by a driven harmonic oscillator:

$$\hat{H} = \delta \hat{e}^\dagger \hat{e} + \eta(\hat{e} + \hat{e}^\dagger), \quad (24)$$

where $\delta \equiv \omega_p - \omega_-$ is the pump-polariton detuning and η is the pump strength. Using the input-output relations—consistent with those considered in Refs. [7,50–52], ensuring the conservation of the total number of polaritons rather than photons—it is straightforward to calculate the quadrature properties of the transmission for the driven-dissipative system [53]

$$\langle \hat{X} \rangle = \sqrt{\kappa} |\alpha| \exp(-i\varphi) \quad \text{and} \quad \Delta^2 \hat{X} = \frac{1}{4t}, \quad (25)$$

where $|\alpha| = 2\sqrt{\eta^2/(\kappa^2 + 4\delta^2)}$ is the amplitude of the measurable photon field which depends on the unknown parameter ω through $\delta(\omega)$, κ is the cavity loss rate, $\varphi = \arctan[\kappa/2\delta]$ is the phase shift between the drive and the system response, and t is the time of the measurement. In the following, we focus on the amplitude measurement [29].

The signal-to-noise ratio of the quadrature amplitude after a measurement time t becomes [29]

$$S(\omega) = 4\kappa(\partial_\omega |\alpha|)^2 = \frac{256\kappa t \delta^2 \eta^2 (g/\omega - 2)^2}{(4\delta^2 + \kappa^2)^3 (1 - g/\omega)}, \quad (26)$$

which can be rewritten including the number of virtual excitations (20) close to the threshold point as

$$S(\omega) \approx \frac{64\kappa\delta^2}{(4\delta^2 + \kappa^2)^2} |\alpha|^2 \exp(4\xi_-) t \approx \frac{2^{10}\kappa\delta^2}{(4\delta^2 + \kappa^2)^2} |\alpha|^2 \langle \hat{n}_c \rangle^2 t. \quad (27)$$

The signal-to-noise ratio from the above equation is presented in Fig. 2. Although the largest amplitude requires the resonance condition $\delta = 0$, this is not the case for the signal, as at resonance the amplitude does not depend on detuning, $S(\delta = 0) = 0$. For small couplings the signal grows slowly [see Fig. 2(a)], and for large couplings the signal grows dramatically near to the threshold point [see Fig. 2(b)], reflecting macroscopic occupation of virtual excitations.

In contrast, reducing the noise using the same squeezing parameter ξ_- —which corresponds to extracting the squeezed virtual excitations from the ground state—leads to

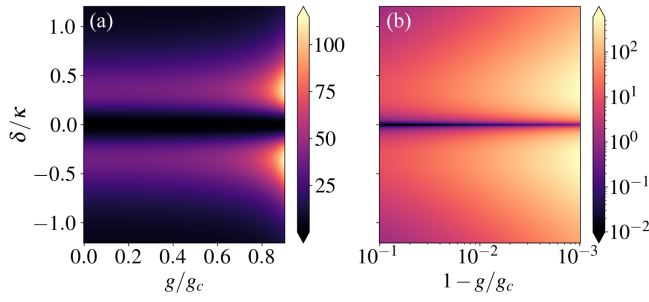


FIG. 2. The signal-to-noise ratio for amplitude measurements normalized to time t as a function of detuning δ/κ and the coupling strength g/g_c . At the resonance ($\delta/\kappa = 0$), the signal does not depend on the unknown parameters, so a slightly off-resonant drive is required for amplitude measurement. (a) At the optimal points, the enhancement is related to the derivative of the effective frequency with respect to the bare frequency. (b) A log-scale enlargement near the critical point, where the derivative is extremely large, illustrating the macroscopic occupation of virtual excitations. In the numerical simulation, we set $\kappa/\omega = 1$ and $\eta/\omega = 1$.

$$\Delta^2 \hat{X} = \frac{\exp(2\xi_-)}{4t} = \frac{\sqrt{1-g/g_c}}{4t} \approx \frac{1}{16\langle \hat{n}_c \rangle t} \quad (28)$$

and, thus, results in the signal-to-noise ratio that scales linearly with the number of virtual excitations, contrary to the quadratic scaling shown in Eq. (27).

Limitations—The calculations presented so far assume an infinite number of atoms, where the collective spin is approximated as a harmonic oscillator. This approximation, enabled by the Holstein-Primakoff transformation, is typically well justified in experiments with large atomic ensembles [54], but it has limitations. In the thermodynamic limit, the spectrum is fully harmonic. Reducing the number of atoms introduces anharmonicity, which can give rise to phenomena such as photon blockade [55], where increasing the driving power no longer leads to a proportional increase in output due to saturation and nonlinear effects. As a result, the magnitude of the coherent field at the system’s output is inherently limited. Furthermore, in finite systems, the energy gap cannot fully close, thereby constraining the maximum achievable derivative of the energy gap (frequency) and the virtual squeezing parameter (see Supplemental Material [29]).

Additionally, our calculations neglected the diamagnetic A^2 term in the Hamiltonian (2). Including this term results in the Hopfield model [56], where the Bogoliubov transformation shows that the cavity frequency increases as $\omega \rightarrow \tilde{\omega} > \omega$. This shift leads to an increase in the critical coupling strength, such that $g_c \rightarrow \tilde{g}_c = \sqrt{\tilde{\omega}\Omega} > \sqrt{\omega\Omega}$. However, as long as there is no upper limit on the coupling strength g , the diamagnetic term does not fundamentally prohibit the use of virtual excitations in quantum metrology. Note that the diamagnetic term is not inherent to light-matter

systems. For instance, it does not appear in optomechanical systems [14,17,57,58].

Another essential consideration is the precise knowledge of system parameters, apart from the unknown quantity being measured. Effective use of the signal-to-noise ratio (via the error propagation formula) requires exquisite control over one of the subsystems and the coupling strength [59]. This reveals a trade-off: Leveraging hybrid modes in precision metrology requires near-perfect control and knowledge of one mode to extract maximal information about the other.

We also emphasize that, while critical slowing down is a common feature near phase transitions, it does not universally limit the practical implementation of our scheme. In particular, for large coupling strengths $g \gg g_c$, the upper polariton mode becomes the relevant excitation, and the associated energy gap increases with g , avoiding divergent relaxation times. Moreover, in some strongly correlated platforms, such as polaritons in condensed matter [6], rapid thermalization can occur despite a vanishing gap due to many-body effects. These scenarios demonstrate that enhanced sensitivity near criticality can, in certain regimes, coexist with favorable dynamical properties, ensuring practical applicability of our approach.

Finally, it is important to emphasize that resonant conditions are not strictly required to exploit virtual excitations and similar results can be obtained in the dispersive regime. For instance, in the quantum Rabi model, large detunings ($\omega \ll \Omega$) are necessary to maximize virtual excitations in the normal phase [60]. In this regime, the mode with the lower frequency exhibits virtual squeezing [29]. Likewise, in the Dicke model, the ratio ω/Ω dictates which mode accumulates more virtual excitations [23,61]. In the extreme limit of $\omega/\Omega \rightarrow 0$, only the cavity mode becomes significantly populated with virtual excitations. This also implies that Ω can be treated as a control parameter when g is fixed and well known, in the case of ω being the to-be-measured parameter.

Conclusions—In this Letter, we demonstrate that virtual excitations—a hallmark of the ultrastrong coupling regime in light-matter systems—can be harnessed to enhance quantum precision measurements. Using the Dicke model as a test bed, we show that these virtual processes reshape the system’s eigenstructure and modify the curvature of the Hamiltonian’s parameter manifold, thereby boosting the quantum Fisher information [31,62] and increasing the information content of the quantum state.

In contrast to conventional squeezing, which requires precise external control, our approach leverages intrinsic squeezing that naturally arises in ultrastrongly coupled systems. The interaction between an unknown subsystem and a well-characterized one creates a hybrid mode whose renormalized frequency—shifted due to virtual excitations—acts as a robust and experimentally accessible signal carrier [63]. This frequency can then be measured using established quantum metrology techniques [64].

More broadly, this strategy provides a powerful sensing framework. By coupling an unknown quantum system to a controlled probe, the emergent hybrid mode encodes information in a way that is both sensitive and resilient. This is especially advantageous for macroscopic quantum systems, where traditional squeezing protocols face technical limitations. Our results suggest that ultrastrong coupling offers not just an alternative to engineered squeezing, but a complementary path to quantum-enhanced sensitivity. Importantly, our proposal is experimentally viable, with proof-of-concept implementations already within reach in platforms such as cavity QED [48,49,65,66], circuit QED [67,68], optomechanics [14,69], and trapped ions [70–73].

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Data availability—The data that support the findings of this Letter are not publicly available. The data are available from the authors upon reasonable request.

- [1] A. Frisk Kockum, A. Miranowicz, S. De Liberato, S. Savasta, and F. Nori, Ultrastrong coupling between light and matter, *Nat. Rev. Phys.* **1**, 19 (2019).
- [2] P. Forn-Díaz, L. Lamata, E. Rico, J. Kono, and E. Solano, Ultrastrong coupling regimes of light-matter interaction, *Rev. Mod. Phys.* **91**, 025005 (2019).
- [3] W. Qin, A. F. Kockum, C. S. Muñoz, A. Miranowicz, and F. Nori, Quantum amplification and simulation of strong and ultrastrong coupling of light and matter, *Phys. Rep.* **1078**, 1 (2024).
- [4] A. G. J. MacFarlane, J. P. Dowling, and G. J. Milburn, Quantum technology: The second quantum revolution, *Phil. Trans. R. Soc. A* **361**, 1655 (2003).
- [5] G. Kurizki, P. Bertet, Y. Kubo, K. Mølmer, D. Petrosyan, P. Rabl, and J. Schmiedmayer, Quantum technologies with hybrid systems, *Proc. Natl. Acad. Sci. U.S.A.* **112**, 3866 (2015).
- [6] C. Ciuti, G. Bastard, and I. Carusotto, Quantum vacuum properties of the intersubband cavity polariton field, *Phys. Rev. B* **72**, 115303 (2005).
- [7] C. Ciuti and I. Carusotto, Input-output theory of cavities in the ultrastrong coupling regime: The case of time-independent cavity parameters, *Phys. Rev. A* **74**, 033811 (2006).
- [8] A. A. Anappara, S. De Liberato, A. Tredicucci, C. Ciuti, G. Biasiol, L. Sorba, and F. Beltram, Signatures of the ultrastrong light-matter coupling regime, *Phys. Rev. B* **79**, 201303(R) (2009).
- [9] Y. Todorov, A. M. Andrews, R. Colombelli, S. De Liberato, C. Ciuti, P. Klang, G. Strasser, and C. Sirtori, Ultrastrong light-matter coupling regime with polariton dots, *Phys. Rev. Lett.* **105**, 196402 (2010).
- [10] S. Ashhab and F. Nori, Qubit-oscillator systems in the ultrastrong-coupling regime and their potential for preparing nonclassical states, *Phys. Rev. A* **81**, 042311 (2010).
- [11] F. Beaudoin, J. M. Gambetta, and A. Blais, Dissipation and ultrastrong coupling in circuit QED, *Phys. Rev. A* **84**, 043832 (2011).
- [12] A. Settineri, V. Macrì, A. Ridolfo, O. Di Stefano, A. F. Kockum, F. Nori, and S. Savasta, Dissipation and thermal noise in hybrid quantum systems in the ultrastrong-coupling regime, *Phys. Rev. A* **98**, 053834 (2018).
- [13] G. A. Peterson, S. Kotler, F. Lecocq, K. Cicak, X. Y. Jin, R. W. Simmonds, J. Aumentado, and J. D. Teufel, Ultrastrong parametric coupling between a superconducting cavity and a mechanical resonator, *Phys. Rev. Lett.* **123**, 247701 (2019).
- [14] F. Bibak, U. Delić, M. Aspelmeyer, and B. Dakić, Dissipative phase transitions in optomechanical systems, *Phys. Rev. A* **107**, 053505 (2023).
- [15] K. Gietka, C. Hotter, and H. Ritsch, Unique steady-state squeezing in a driven quantum Rabi model, *Phys. Rev. Lett.* **131**, 223604 (2023).
- [16] R. Stassi, M. Cirio, K. Funo, J. Puebla, N. Lambert, and F. Nori, Unveiling and veiling an entangled light-matter quantum state from the vacuum, *Phys. Rev. Res.* **5**, 043095 (2023).
- [17] K. Dare, J. J. Hansen, I. Coroli, A. Johnson, M. Aspelmeyer, and U. Delić, Ultrastrong linear optomechanical interaction, *Phys. Rev. Res.* **6**, L042025 (2024).
- [18] Y.-H. Chen, Z.-C. Shi, Y.-R. Zhang, F. Nori, and Y. Xia, Suppressed energy relaxation in the quantum Rabi model at the critical point, [arXiv:2411.03710](https://arxiv.org/abs/2411.03710).
- [19] X. Cao, J. Q. You, H. Zheng, and F. Nori, A qubit strongly coupled to a resonant cavity: Asymmetry of the spontaneous emission spectrum beyond the rotating wave approximation, *New J. Phys.* **13**, 073002 (2011).
- [20] K. Gietka, Vacuum Rabi splitting as a manifestation of virtual two-mode squeezing: Extracting the squeezing parameters from frequency shifts, *Phys. Rev. A* **110**, 063703 (2024).
- [21] B. M. Garraway, The Dicke model in quantum optics: Dicke model revisited, *Phil. Trans. R. Soc. A* **369**, 1137 (2011).
- [22] H. Ritsch, P. Domokos, F. Brennecke, and T. Esslinger, Cold atoms in cavity-generated dynamical optical potentials, *Rev. Mod. Phys.* **85**, 553 (2013).
- [23] C. Emary and T. Brandes, Chaos and the quantum phase transition in the Dicke model, *Phys. Rev. E* **67**, 066203 (2003).
- [24] J. Larson and E. K. Irish, Some remarks on ‘superradiant’ phase transitions in light-matter systems, *J. Phys. A* **50**, 174002 (2017).
- [25] J.-Y. Zhou, Y.-H. Zhou, X.-L. Yin, J.-F. Huang, and J.-Q. Liao, Quantum entanglement maintained by virtual excitations in an ultrastrongly-coupled-oscillator system, *Sci. Rep.* **10**, 12557 (2020).
- [26] D. N. Basov, A. Asenjo-Garcia, P. J. Schuck, X. Zhu, and A. Rubio, Polariton panorama, *Nanophotonics* **10**, 549 (2021).
- [27] C. M. Wilson, G. Johansson, A. Pourkabirian, M. Simoen, J. R. Johansson, T. Duty, F. Nori, and P. Delsing, Observation of the dynamical Casimir effect in a superconducting circuit, *Nature (London)* **479**, 376 (2011).

- [28] V. Dodonov, Fifty years of the dynamical Casimir effect, *Physics* **2**, 67 (2020).
- [29] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/cxvs-5pb1> for details on (i) the calculation of quantum Fisher information for a coherent state in the driven-dissipative ultrastrongly coupled Dicke model in the large N limit and (ii) a comprehensive analysis of quantum metrology based on the quantum Rabi model with virtual excitations.
- [30] S. De Liberato, Virtual photons in the ground state of a dissipative system, *Nat. Commun.* **8**, 1465 (2017).
- [31] L. Campos Venuti and P. Zanardi, Quantum critical scaling of the geometric tensors, *Phys. Rev. Lett.* **99**, 095701 (2007).
- [32] G. Salvatori, A. Mandarino, and M. G. A. Paris, Quantum metrology in Lipkin-Meshkov-Glick critical systems, *Phys. Rev. A* **90**, 022111 (2014).
- [33] M. M. Rams, P. Sierant, O. Dutta, P. Horodecki, and J. Zakrzewski, At the limits of criticality-based quantum metrology: Apparent super-Heisenberg scaling revisited, *Phys. Rev. X* **8**, 021022 (2018).
- [34] L. Garbe, M. Bina, A. Keller, M. G. A. Paris, and S. Felicetti, Critical quantum metrology with a finite-component quantum phase transition, *Phys. Rev. Lett.* **124**, 120504 (2020).
- [35] K. Gietka, F. Metz, T. Keller, and J. Li, Adiabatic critical quantum metrology cannot reach the Heisenberg limit even when shortcuts to adiabaticity are applied, *Quantum* **5**, 489 (2021).
- [36] K. Gietka and H. Ritsch, Squeezing and overcoming the Heisenberg scaling with spin-orbit coupled quantum gases, *Phys. Rev. Lett.* **130**, 090802 (2023).
- [37] R. Salvia, M. Mehboudi, and M. Perarnau-Llobet, Critical quantum metrology assisted by real-time feedback control, *Phys. Rev. Lett.* **130**, 240803 (2023).
- [38] L. Giannelli, E. Paladino, M. Grajcar, G. S. Paraoanu, and G. Falci, Detecting virtual photons in ultrastrongly coupled superconducting quantum circuits, *Phys. Rev. Res.* **6**, 013008 (2024).
- [39] V. Giovannetti, S. Lloyd, and L. Maccone, Quantum metrology, *Phys. Rev. Lett.* **96**, 010401 (2006).
- [40] V. Giovannetti, S. Lloyd, and L. Maccone, Advances in quantum metrology, *Nat. Photonics* **5**, 222 (2011).
- [41] L. Pezzè, A. Smerzi, M. K. Oberthaler, R. Schmied, and P. Treutlein, Quantum metrology with nonclassical states of atomic ensembles, *Rev. Mod. Phys.* **90**, 035005 (2018).
- [42] E. Polino, M. Valeri, N. Spagnolo, and F. Sciarrino, Photonic quantum metrology, *AVS Quantum Sci.* **2**, 024703 (2020).
- [43] L. Pezzè and A. Smerzi, Entanglement, nonlinear dynamics, and the Heisenberg limit, *Phys. Rev. Lett.* **102**, 100401 (2009).
- [44] R. Demkowicz-Dobrzański, J. Kołodyński, and M. Guţă, The elusive Heisenberg limit in quantum-enhanced metrology, *Nat. Commun.* **3**, 1063 (2012).
- [45] S. L. Braunstein and C. M. Caves, Statistical distance and the geometry of quantum states, *Phys. Rev. Lett.* **72**, 3439 (1994).
- [46] W. Górecki, F. Albarelli, S. Felicetti, R. D. Candia, and L. Maccone, Interplay between time and energy in bosonic noisy quantum metrology, *PRX Quantum* **6**, 020351 (2025).
- [47] C. Oh, C. Lee, C. Rockstuhl, H. Jeong, J. Kim, H. Nha, and S.-Y. Lee, Optimal gaussian measurements for phase estimation in single-mode gaussian metrology, *npj Quantum Inf.* **5**, 10 (2019).
- [48] K. Baumann, C. Guerlin, F. Brennecke, and T. Esslinger, Dicke quantum phase transition with a superfluid gas in an optical cavity, *Nature (London)* **464**, 1301 (2010).
- [49] R. Mottl, F. Brennecke, K. Baumann, R. Landig, T. Donner, and T. Esslinger, Roton-type mode softening in a quantum gas with cavity-mediated long-range interactions, *Science* **336**, 1570 (2012).
- [50] M. Bamba and T. Ogawa, Dissipation and detection of polaritons in the ultrastrong-coupling regime, *Phys. Rev. A* **86**, 063831 (2012).
- [51] M. Bamba and T. Ogawa, System-environment coupling derived by Maxwell's boundary conditions from the weak to the ultrastrong light-matter-coupling regime, *Phys. Rev. A* **88**, 013814 (2013).
- [52] S. De Liberato, Comment on System-environment coupling derived by Maxwell's boundary conditions from the weak to the ultrastrong light-matter-coupling regime, *Phys. Rev. A* **89**, 017801 (2014).
- [53] D. Barberena, R. J. Lewis-Swan, A. M. Rey, and J. K. Thompson, Ultra narrow linewidth frequency reference via measurement and feedback, *C.R. Phys.* **24**, 55 (2023).
- [54] F. Mivehvar, F. Piazza, T. Donner, and H. Ritsch, Cavity QED with quantum gases: New paradigms in many-body physics, *Adv. Phys.* **70**, 1 (2021).
- [55] K. M. Birnbaum, A. Boca, R. Miller, A. D. Boozer, T. E. Northup, and H. J. Kimble, Photon blockade in an optical cavity with one trapped atom, *Nature (London)* **436**, 87 (2005).
- [56] J. J. Hopfield, Theory of the contribution of excitons to the complex dielectric constant of crystals, *Phys. Rev.* **112**, 1555 (1958).
- [57] D. Marković, S. Jezouin, Q. Ficheux, S. Fedortchenko, S. Felicetti, T. Coudreau, P. Milman, Z. Leghtas, and B. Huard, Demonstration of an effective ultrastrong coupling between two oscillators, *Phys. Rev. Lett.* **121**, 040505 (2018).
- [58] K. Kustura, C. Gonzalez-Ballester, A. d. I. R. Sommer, N. Meyer, R. Quidant, and O. Romero-Isart, Mechanical squeezing via unstable dynamics in a microcavity, *Phys. Rev. Lett.* **128**, 143601 (2022).
- [59] G. Mihailescu, S. Campbell, and K. Gietka, Uncertain quantum critical metrology: From single to multi parameter sensing, *Phys. Rev. A* **111**, 052621 (2025).
- [60] M.-J. Hwang, R. Puebla, and M. B. Plenio, Quantum phase transition and universal dynamics in the Rabi model, *Phys. Rev. Lett.* **115**, 180404 (2015).
- [61] K. Gietka and T. Busch, Inverted harmonic oscillator dynamics of the nonequilibrium phase transition in the Dicke model, *Phys. Rev. E* **104**, 034132 (2021).
- [62] P. Zanardi, P. Giorda, and M. Cozzini, Information-theoretic differential geometry of quantum phase transitions, *Phys. Rev. Lett.* **99**, 100603 (2007).
- [63] T. W. Hänsch, Nobel lecture: Passion for precision, *Rev. Mod. Phys.* **78**, 1297 (2006).
- [64] C. Hotter, H. Ritsch, and K. Gietka, Combining critical and quantum metrology, *Phys. Rev. Lett.* **132**, 060801 (2024).

- [65] J. Klinder, H. Keßler, M. Wolke, L. Mathey, and A. Hemmerich, Dynamical phase transition in the open Dicke model, *Proc. Natl. Acad. Sci. U.S.A.* **112**, 3290 (2015).
- [66] A. J. Kollár, A. T. Papageorge, V. D. Vaidya, Y. Guo, J. Keeling, and B. L. Lev, Supermode-density-wave-polariton condensation with a Bose–Einstein condensate in a multi-mode cavity, *Nat. Commun.* **8**, 14386 (2017).
- [67] F. Yoshihara, T. Fuse, S. Ashhab, K. Kakuyanagi, S. Saito, and K. Semba, Superconducting qubit–oscillator circuit beyond the ultrastrong-coupling regime, *Nat. Phys.* **13**, 44 (2017).
- [68] X. Gu, A. F. Kockum, A. Miranowicz, Y. xi Liu, and F. Nori, Microwave photonics with superconducting quantum circuits, *Phys. Rep.* **718–719**, 1 (2017).
- [69] K. Dare, J. J. Hansen, I. Coroli, A. Johnson, M. Aspelmeyer, and U. Delić, Ultrastrong linear optomechanical interaction, *Phys. Rev. Res.* **6**, L042025 (2024).
- [70] J. Braumüller, M. Marthaler, A. Schneider, A. Stehli, H. Rotzinger, M. Weides, and A. V. Ustinov, Analog quantum simulation of the Rabi model in the ultra-strong coupling regime, *Nat. Commun.* **8**, 779 (2017).
- [71] A. Safavi-Naini, R. J. Lewis-Swan, J. G. Bohnet, M. Gärttner, K. A. Gilmore, J. E. Jordan, J. Cohn, J. K. Freericks, A. M. Rey, and J. J. Bollinger, Verification of a many-ion simulator of the Dicke model through slow quenches across a phase transition, *Phys. Rev. Lett.* **121**, 040503 (2018).
- [72] K. A. Gilmore, M. Affolter, R. J. Lewis-Swan, D. Barberena, E. Jordan, A. M. Rey, and J. J. Bollinger, Quantum-enhanced sensing of displacements and electric fields with two-dimensional trapped-ion crystals, *Science* **373**, 673 (2021).
- [73] M. L. Cai, Z. D. Liu, W. D. Zhao, Y. K. Wu, Q. X. Mei, Y. Jiang, L. He, X. Zhang, Z. C. Zhou, and L. M. Duan, Observation of a quantum phase transition in the quantum Rabi model with a single trapped ion, *Nat. Commun.* **12**, 1126 (2021).
- [74] S. Krämer, D. Plankensteiner, L. Ostermann, and H. Ritsch, Quantumoptics.jl: A JULIA framework for simulating open quantum systems, *Comput. Phys. Commun.* **227**, 109 (2018).