

Surface-Mediated Ultrastrong Cavity Coupling of Two-Dimensional Itinerant Electrons

Christian J. Eckhardt,^{1,2} Andrey Grankin,³ Dante M. Kennes^{1b},^{2,1} Michael Ruggenthaler,¹ Angel Rubio^{1b},^{1,4}
Michael A. Sentef^{1b},^{5,1} Mohammad Hafezi,³ and Marios H. Michael^{1b},^{*}

¹*Max Planck Institute for the Structure and Dynamics of Matter, Center for Free-Electron Laser Science (CFEL),
Luruper Chaussee 149, 22761 Hamburg, Germany*

²*Institut für Theorie der Statistischen Physik, RWTH Aachen University and
JARA-Fundamentals of Future Information Technology, 52056 Aachen, Germany*

³*Joint Quantum Institute, Department of Physics, University of Maryland, College Park, Maryland 20742, USA*

⁴*Initiative for Computational Catalysis (ICC), The Flatiron Institute, 162 Fifth Avenue, New York, New York 10010, USA*

⁵*Institute for Theoretical Physics and Bremen Center for Computational Materials Science,
University of Bremen, 28359 Bremen, Germany*



(Received 17 September 2024; accepted 10 September 2025; published 7 October 2025)

Engineering phases of matter in cavities requires effective light-matter coupling strengths that are on the same order of magnitude as the bare system energetics, coined the ultrastrong coupling regime. For models of itinerant electron systems, which do not have discrete energy levels, a clear definition of this regime has been lacking to date. Here, we argue that a change of the electronic mass exceeding 10% of its bare value may serve as such a definition. We propose a quantitative computational scheme for obtaining the electronic mass in relation to its bare vacuum value and show that coupling to surface polariton modes can induce such mass changes. Our results have important implications for cavity design principles that enable the engineering of electronic properties with quantum light.

DOI: [10.1103/2fw2-lbhy](https://doi.org/10.1103/2fw2-lbhy)

Introduction—Engineering the electromagnetic field inside optical resonators is emerging as a versatile tool to control materials properties [1–6]. In the corresponding field of cavity materials engineering, several intriguing effects have been demonstrated, such as alteration of the critical temperature of a charge density wave transition [7] and influence on the topologically protected conductance properties of quantum Hall systems [8]. The strong interaction between matter and photon fields in a cavity [2,9] further gives rise to the emerging field of polaritonic chemistry [3,10–19], promising to revolutionize the way we perceive materials science [1,2]. Theoretically, the change of the electromagnetic field inside cavities is often understood by considering one or few bosonic modes that have a modified light-matter coupling constant [1,2,20,21]. Such models have been successfully employed to describe situations where a material’s resonance is coupled to the resonance of a cavity [22–27]. If the effective coupling constant exceeds 0.1 times the energy of the mode’s

frequency, the system is said to be in the ultrastrong coupling regime [2].

For systems of itinerant electrons [26,28–36], like a gas of free electrons [37], where no intrinsic materials resonances are present that naturally *select* a particular cavity mode or energy scale, it is *a priori* unclear whether such few-mode models with an effectively enhanced light-matter coupling strength are applicable [38–41]. As a consequence, one may need to consider the full continuum of modes, whose coupling strength is generically given by the fine-structure constant α , hindering a straightforward extension of the concept of ultrastrong coupling to this situation.

Here, we make progress in this direction by proposing a computational scheme in which changes to the energetics of 2D electrons are computed, in particular their effective physical mass. When considering the effects of all electromagnetic modes, the change in the mass of the electron due to vacuum fluctuations diverges with the high energy cutoff, even in free space [37,42,43]. To address the issue of properly treating high- and low-frequency cutoffs, we adopt a perturbative renormalization approach, assuming the electron has the measurable physical mass m_e in free space and then computing the relative mass change between free space and different optical setups. We define a 10% change of the electronic mass relative to free space as the onset of ultrastrong coupling. We find that when the electrons are in the near field of an interface to an insulator hosting phonon polaritons, the ultrastrong coupling regime

^{*}Contact author: marios.michael@mpsd.mpg.de

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International](https://creativecommons.org/licenses/by/4.0/) license. Further distribution of this work must maintain attribution to the author(s) and the published article’s title, journal citation, and DOI. Open access publication funded by the Max Planck Society.

may indeed be reached when assuming system parameters from experiment. We show that this effect is due to coupling to the surface modes at the interface while the bulk modes are to a good approximation unaltered compared to the free-space vacuum. Furthermore we investigate the far-field to near-field crossover close to a Drude metal.

Our results introduce a computational tool to investigate the effects different cavity setups have on itinerant electron systems through the photonic density of states. Physically, changes in the electronic mass impact the density of states at the Fermi level, which directly controls the response of noninteracting electrons, such as heat capacity, conductivity, and magnetic susceptibility. For interacting systems, significant changes of the electronic mass underlie previous theoretical predictions in the field of cavity engineering of itinerant electrons [37,38,44–48]. Our findings reveal that cavities employing surface phonon polariton modes could indeed induce such substantial changes.

Electronic self-energy and partial local density of states— In this Letter we consider the full continuous electromagnetic field close to an interface with a material that we describe using a frequency dependent $\epsilon(\omega, \mathbf{r})$; see Fig. 1(a). We employ the Weyl gauge for which $\mathbf{E} = -\partial_t \mathbf{A}$ and expand the electromagnetic vector potential $\hat{\mathbf{A}}(\mathbf{r})$ in its eigenmodes,

$$\hat{\mathbf{A}}(\mathbf{r}) = \sum_q \sqrt{\frac{\hbar}{2\omega_q \epsilon_0}} \left(e^{iq_{\parallel} \cdot \mathbf{r}_{\parallel}} \mathbf{f}_q^*(z) \hat{a}_q^{\dagger} + e^{-iq_{\parallel} \cdot \mathbf{r}_{\parallel}} \mathbf{f}_q(z) \hat{a}_q \right). \quad (1)$$

Here, the multi-index q is given as $q = (\mathbf{q}_{\parallel}, q_z, \lambda)$ with $\mathbf{q}_{\parallel}$ the in-plane part of the wave vector, q_z its z component and λ denoting the polarization. $\hbar\omega_q$ is the eigenenergy of the modes where a photon is created by $\hat{a}_q^{\dagger}$ and annihilated by $\hat{a}_q$. The mode functions factorize into a plane-wave part that we have written explicitly and a perpendicular dependence $\mathbf{f}_{q_z, \lambda}(z)$ since we assume in-plane translational symmetry. We provide the explicit mode functions in Supplemental Material (SM), Sec. 2 [49]. With these one may compute the partial local density of states (PLDOS) of the electromagnetic field inside the cavity as [50]

$$\rho_n(\omega, \mathbf{r}) = \sum_q |\mathbf{f}_{q_z, \lambda}(\mathbf{r}) \cdot \mathbf{n}|^2 \delta(\omega - \omega_q), \quad (2)$$

where $\mathbf{n}$ is a vector of unit length. In free space the PLDOS is given as [50]

$$\rho_{n,0}(\omega, \mathbf{r}) = \frac{\omega^2}{3\pi^2 c^3}. \quad (3)$$

We consider a free electron in two dimensions with system size $L \times L$ that we couple to light in the minimal coupling scheme, assuming consistent in-plane periodic boundary conditions. The system size sets the longest length scale at

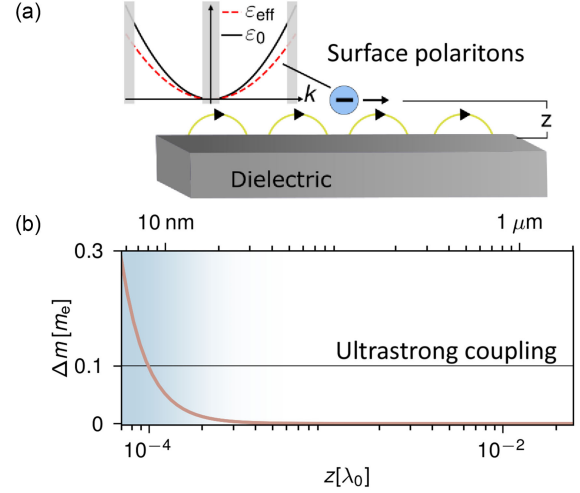


FIG. 1. Summary of results. (a) Illustration of a 2D electron (blue circle) with quadratic bare dispersion ϵ_0 (black line) traveling parallel to a surface hosting surface polaritons (yellow semicircles) at distance z . The coupling leads to a flattening of the electronic dispersion ϵ_{eff} (red dashed line) corresponding to an effectively increased electronic mass. The k^2 dispersion is not valid at infinitely long or short length scales, marked by the gray bars. (b) Change of the mass Δm (orange line) as computed via Eq. (8) in units of the physical electronic mass m_e in vacuum close to an interface with a material that hosts phonon polaritons as a function of distance z to that interface. The onset of strong coupling is marked by the thin black line and blue shading. The bottom x axis is in units of the length scale associated with the surface phonon polaritons $\lambda_0 = (2\pi c/\omega_\infty)$, where $\sqrt{2}\omega_\infty = \sqrt{\omega_{\text{LO}}^2 + \omega_{\text{TO}}^2}$. Parameters are chosen for the soft phonon of SrTiO₃—see main text.

which the k^2 dispersion is valid—also see Fig. 1(a). In order to also set a shortest length scale, we will later introduce a UV cutoff. We omit spin for brevity and neglect the diamagnetic term. Employing the dipole approximation, light and matter can be decoupled perturbatively using a Schrieffer-Wolf transformation [37,38], shown in SM Sec. 1 [49], which yields the effective Hamiltonian

$$\hat{H} = \sum_q \omega_q \hat{a}_q^{\dagger} \hat{a}_q + \left(\frac{\hbar^2 \mathbf{k}_{\parallel}^2}{2m_0} + \Sigma(\mathbf{k}_{\parallel}) \right) \hat{c}_{\mathbf{k}_{\parallel}}^{\dagger} \hat{c}_{\mathbf{k}_{\parallel}}, \quad (4)$$

where $\hat{c}_{\mathbf{k}_{\parallel}}^{\dagger}$ creates an electron with in-plane momentum $\hbar \mathbf{k}_{\parallel}$, while $\hat{c}_{\mathbf{k}_{\parallel}}$ annihilates such an electron and m_0 is the bare mass. We note that beyond the dipole approximation Eq. (4) would receive further corrections that may, however, be neglected by perturbative arguments. In Eq. (4), the effect of the photons on the electrons is incorporated via the self-energy Σ , which is defined as

$$\Sigma(\mathbf{k}_{\parallel}) = - \sum_q \frac{e^2 \hbar^2 |\mathbf{k}_{\parallel} \cdot \mathbf{f}_{q_z, \lambda}(z_0)|^2}{m_0^2 \hbar \omega_q} \frac{\hbar}{2\omega_q \epsilon_0}. \quad (5)$$

This self-energy explicitly depends on the continuous three-dimensional electromagnetic field through the mode functions and can be directly related to the PLDOS of the electromagnetic field using its definition in Eq. (2),

$$\Sigma(\mathbf{k}_{\parallel}, \Lambda_{\text{UV}}) = -\alpha \frac{4\pi\hbar^3 c |\mathbf{k}_{\parallel}|^2}{2m_0^2} \int_{\Lambda_{\text{IR}}}^{\Lambda_{\text{UV}}} d\omega \frac{\rho_{\mathbf{e}_{q_{\parallel}}}(\omega, z_0)}{\omega^2}, \quad (6)$$

where α is the fine-structure constant. We consider a high-energy cutoff that ensures the validity of the dipole approximation, $\Lambda_{\text{UV}} \ll c|k_{\parallel}|$. For typical momenta of interest in condensed matter, $k_{\parallel} \sim 1 \text{ \AA}$, this condition also guarantees that the cutoff we consider is below the Landau pole, Λ_{Landau} , for which perturbative nonrelativistic QED within the dipole approximation becomes ill-defined [43]. Additionally, we have included a low-energy cutoff Λ_{IR} that is not strictly necessary to regularize nonrelativistic QED [42,43] but is instead introduced for two reasons: (i) in order to be agnostic to specific boundary conditions that we set for the electromagnetic field discussed later; (ii) to account for the fact that the electronic k^2 dispersion is naturally not valid on infinite length scales [see Fig. 1(a)], which effectively necessitates a low-energy cutoff on the combined light-matter system. The latter point is further illustrated by the fact that in reality, light couples electrons of momentum k to $k + q$, and therefore the IR cutoff of the electronic system also creates a cutoff to the light modes that are able to couple to electrons. The range of validity of Eqs. (4) and (5) is discussed in SM Sec. 1.

Mass renormalization relative to free space—Incorporating the self-energy into an effective dispersion $\varepsilon_{\text{eff}}(\mathbf{k}_{\parallel}) = (\hbar^2 \mathbf{k}_{\parallel}^2 / 2m_0) + \Sigma(\mathbf{k}_{\parallel})$ naturally gives rise to a definition of the effective mass of an electron with an electromagnetic environment as $m_{\text{eff}} = \hbar^2 [\partial_{\mathbf{k}_{\parallel}}^2 \varepsilon_{\text{eff}}(\mathbf{k}_{\parallel})]^{-1}$. Note that due to the dipole approximation the self-energy is real and the mass is independent of the momentum of electrons. Corrections beyond the dipole approximation would introduce a momentum dependent mass. From Eq. (6) it becomes clear that the self-energy depends on the UV cutoff and diverges when removing the cutoff Λ_{UV} in free space, where the PLDOS is given by Eq. (3). This divergence in QED is usually mitigated perturbatively by using mass renormalization, where the coupling to the electromagnetic field leads to the replacement of a cutoff dependent bare mass m_0 by the physical mass m_e [57],

$$m_0(\Lambda_{\text{UV}}) = m_e + \frac{m_e^2}{\hbar^2} \partial_{\mathbf{k}_{\parallel}}^2 \Sigma_{\text{vac}}(\mathbf{k}_{\parallel}, \Lambda_{\text{UV}}) + \mathcal{O}(\alpha^2), \quad (7)$$

to leading order in the fine-structure constant. To fix the bare mass, we assume that the effective mass of the electron far away from the interface, where the density of states asymptotically approaches the free-space one, Eq. (3), is given by the physical mass m_e for any given cutoff Λ_{UV} .

This allows us to compute the change in the effective mass relative to this point of reference,

$$\Delta m := m_{\text{eff}} - m_e = -\frac{m_e^2}{\hbar^2} \partial_{\mathbf{k}_{\parallel}}^2 \Sigma_{\text{eff}}(\mathbf{k}_{\parallel}) + \mathcal{O}(\alpha^2), \quad (8)$$

which depends on an effective self-energy in which a reference PLDOS, $\rho_{n,r}$, is subtracted from the cavity one,

$$\Sigma_{\text{eff}}(\mathbf{k}_{\parallel}) = -\alpha \frac{4\pi\hbar^3 c |\mathbf{k}_{\parallel}|^2}{2m_e^2} \int_{\Lambda_{\text{IR}}}^{\Lambda_{\text{UV}}} d\omega \frac{\rho_{\mathbf{e}_{q_{\parallel}}}(\omega, z_0) - \rho_{n,r}(\omega)}{\omega^2}. \quad (9)$$

The above expression, Eq. (9), quantifies the effective enhancement or reduction of the electron-photon coupling inside a cavity through integrating the difference in the PLDOS compared to free space and constitutes the first major result of our Letter.

Redistribution of spectral weight in the PLDOS—We compute the change of mass close to a phonon-polariton system. We assume the lower half-space to be filled with the insulating material giving rise to a frequency dependent permittivity that we take as the one produced by a Lorentz-oscillator model for the phonons [51],

$$\varepsilon(\omega, \mathbf{r}) = \begin{cases} 1 & ; z > 0 \\ \varepsilon(\omega) = \frac{\omega_{\text{LO}}^2 - \omega^2}{\omega_{\text{TO}}^2 - \omega^2} & ; z < 0, \end{cases} \quad (10)$$

where ω_{LO} is the longitudinal phonon frequency while ω_{TO} the transversal one, $\omega_{\text{LO}} < \omega_{\text{TO}}$ [58]. In order to obtain concrete numbers we choose $\omega_{\text{LO}} = 2\pi \times 5.1 \text{ THz}$ and $\omega_{\text{TO}} = 2\pi \times 1.26 \text{ THz}$, which corresponds to the soft phonon of SrTiO_3 . We choose SrTiO_3 as a test case because it is a quantum paraelectric material characterized by a low-frequency, infrared-active phonon mode and is a suitable substrate for 2D electron gases hosted in atomically thin van der Waals materials [59]. As the point of reference we choose a distance to the interface z for which $z \gg (c/\omega_{\text{LO}})$, in particular $z = 0.25 \text{ m}$. In order to obtain discrete modes we assume perfect metal boundary conditions in z direction with a system extent $L_z \gg (c/\omega_{\text{LO}})$, in particular choosing $L_z = 1 \text{ m}$, and choose an IR cutoff of $\Lambda_{\text{IR}} = 50 \text{ GHz}$ considerably above the corresponding fundamental mode associated with the 1 m length scale. We have confirmed numerically that any cutoff Λ_{IR} far below the energy scales related to the surface phonon polariton, in particular ω_{TO} , does not change the results quantitatively. On the other hand we use a UV cutoff of $\hbar\Lambda_{\text{UV}} = 0.3 \text{ eV}$, such that $\Lambda_{\text{UV}} \gg \omega_{\text{LO}}$. Again we have checked that the results obtained via Eq. (8) do not change as long as the cutoff is far above the energies related to the surface phonon polariton, in particular ω_{LO} . We report the relative change of mass in Fig. 1(b), which shows that for realistic distances to the surface the change in mass may indeed

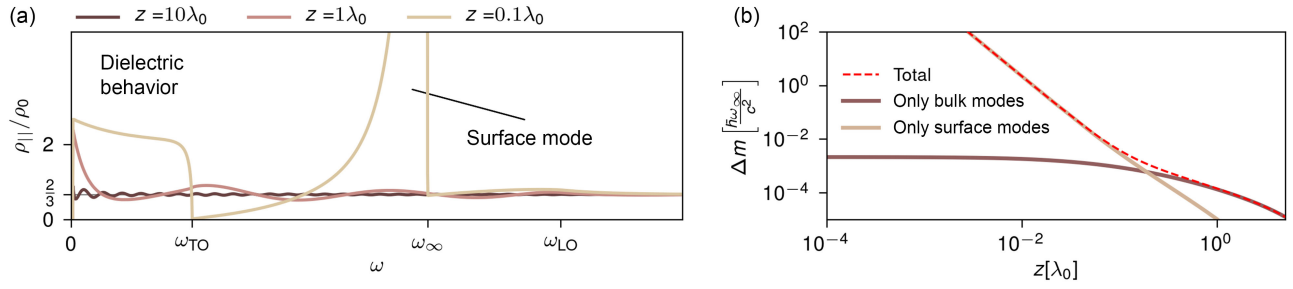


FIG. 2. Mass change due to surface modes. (a) In plane partial local density of states, $\rho_{||}$ in units of the free-space photonic density of states $\rho_0 = (\omega^2/\pi^2 c^3)$ as a function of frequency ω close to an interface hosting phonon-polaritons for different distances z to the interface (color lines) in units of the length scale associated with the surface phonon polaritons $\lambda_0 = (2\pi c/\omega_\infty)$, where $\sqrt{2}\omega_\infty = \sqrt{\omega_{LO}^2 + \omega_{TO}^2}$. Note that since we focus only on the in-plane component of the PLDOS, the answer away from the surface converges to $2/3$. (b) Mass change Δm of an electron relative to free space as a function of the distance to an interface z_0 with a material hosting phonon-polaritons only accounting for the surface modes [yellow line, Eq. (11)], only taking the bulk modes into account (brown line) and the full result (red dashed line). ω_{LO} and ω_{TO} are chosen as in Fig. 1(b) but the units are chosen such that a global rescaling of frequencies yields the same plot.

exceed 10% of the physical mass in free space. We coin such a drastic change in mass of a free electron as the ultrastrong coupling regime for this case [2]. Details of the calculation of the PLDOS are given in SM Sec. 4 [49]. We explain the results of the previous part in terms of a redistribution of spectral weight in the PLDOS, which we show in Fig. 2(a). At low frequencies, $\omega < \omega_{TO}$, spectral weight is added since the material acts as a dielectric in this regime. Within the restrahlenband and above, $\omega > \omega_{TO}$, spectral weight is depleted since $\epsilon(\omega)|_{\omega > \omega_{TO}} < 1$, with some oscillations at high frequencies. On top of this a strong peak around $\omega \lesssim \omega_\infty = \sqrt{(\omega_{LO}^2 + \omega_{TO}^2)/2}$ appears close to the interface, $z < (c/\omega_\infty)$, that can be attributed to the surface modes.

We compare the mass renormalization induced by surface and bulk modes separately. The result is shown in Fig. 2(b) using the same parameters as for Fig. 1(b) but choosing natural units such that a global rescaling of all frequencies does not alter the results. We find that the bulk mode contribution leads to a mass change only weakly depending on distance that makes up only a negligible portion of the overall mass change close to the interface. As a consequence, close to the interface, $z < (c/\omega_\infty)$, the bulk mode contribution nearly cancels the free-space contribution in Eq. (8), and the mass renormalization is dominated by the additional presence of the surface modes, i.e.,

$$\int_{\Lambda_{IR}}^{\Lambda_{UV}} d\omega \frac{\rho_{e_{q||}}(\omega, \mathbf{r}) - \rho_{n,r}(\omega)}{\omega^2} \approx \int_{\omega_{TO}}^{\omega_{LO}} \frac{\rho_{e_{q||}}^{SPP}(\omega, \mathbf{r})}{\omega^2}, \quad (11)$$

where $\rho_{e_{q||}}^{SPP}(\omega, \mathbf{r})$ denotes the density of states of the surface modes only. This allows us to derive an approximate analytic expression for the mass renormalization as the electron approaches the interface (derived in SM Sec. 5 [49]),

$$\Delta m_{SPP} = \alpha \frac{\hbar c (\omega_\infty^2 - \omega_{TO}^2)}{4\omega_\infty^4} \frac{1}{z^3}. \quad (12)$$

These findings justify a computational approach for quantifying effects from coupling to surface modes where bulk modes only enter implicitly in the calculation giving rise to the bare system parameters, while the surface modes are quantized and treated explicitly within the theory [60,61].

Drude metal—In this part we consider an electron moving close to a Drude metal. We model the Drude metal by a frequency dependent permittivity $\epsilon(\omega)$ that takes the same form as Eq. (10) but setting $\omega_P \equiv \omega_{LO}$ and $\omega_{TO} = 0$ and assume small losses such that we have for the scattering time $\tau^{-1} \ll \omega_P$. We show the in-plane PLDOS close to the interface in Fig. 3(a). A strong peak appears for small distances $z \ll (c/\omega_P)$ associated with the surface plasmon-polariton mode. At low frequency the density of states is reduced since the Drude metal acts as a perfect conductor in this regime.

We compute the effective change in mass for different distances to the interface using Eq. (8). Again we introduce cutoffs but note that the results do not change quantitatively as long as $\Lambda_{IR} \ll \omega_P \ll \Lambda_{UV}$ is fulfilled. The result is shown in Fig. 3(b). We find two distinct regimes: a far-field regime, $z \gg (2\pi c/\omega_P)$, where the effective mass of the electrons is decreased scaling as $\Delta m \sim -z^{-1}$, and a near-field regime, $z \ll (2\pi c/\omega_P)$, where the effective mass increases as $\Delta m \sim z^{-3}$. The increase close to the interface is due to coupling to surface modes, surface plasmon-polaritons in this case, similar to the case of the insulator presented in Fig. 2. In the far field, we attribute the decrease in mass to the perfect metal properties of the Drude metal at low frequencies signified by the depletion of spectral weight at low frequencies $\omega \ll \omega_P$.

The far-field effect of a Drude metal on the electronic mass has interesting implications for electrons in an idealized Fabry-Pérot cavity. One may model a

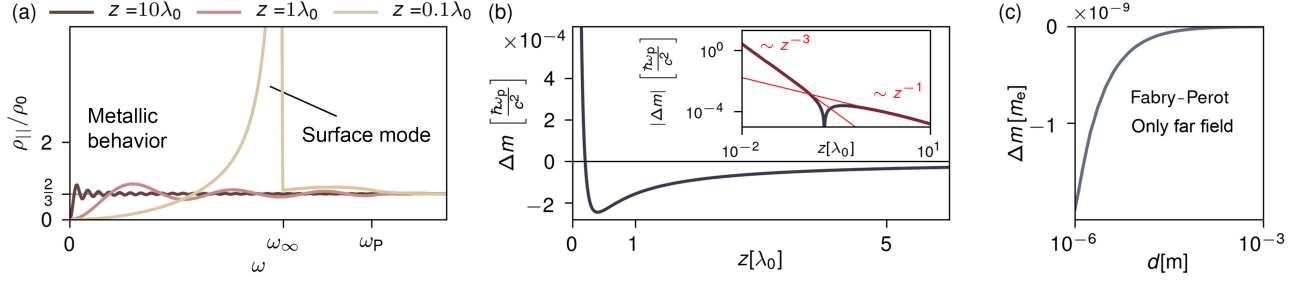


FIG. 3. Mass change close to a Drude metal. (a) In-plane partial local density of states $\rho_{||}$ (colored lines) in units of the free-space photonic density of states $\rho_0 = (\omega^2/\pi^2 c^3)$ as a function of frequency ω close to a Drude metal for different distances z to the interface (color lines) in units of the length scale associated with the surface plasmon-polaritons $\lambda_0 = (2\sqrt{2}\pi c/\omega_P)$. (b) Mass change Δm (blue line) as computed with Eq. (8) as a function of distance z to a Drude metal in units of λ_0 defined as in (a). The units of the plot are chosen such that a rescaling of the plasma frequency ω_P does not change the plot. In the inset the absolute value of the same quantity is plotted on a log-log scale with a scaling as z^{-3} in the near field and z^{-1} in the far field indicated (red lines). (c) Relative mass change of a free electron (blue line) in the middle of an idealized Fabry-Pérot cavity (perfect metal boundary conditions) assuming that the k^2 dispersion of the electron is valid at energy scales far below the fundamental frequency of the cavity.

Fabry-Pérot cavity by perfect mirror boundary conditions at $z = 0$ and $z = d$, giving rise to the fundamental frequency $\omega_0 = (\pi c/d)$. This reduces the effect of the mirrors to far-field effects of a metal and leads to the gapping out of all in-plane modes below the fundamental frequency ω_0 . We now assume a free electron at $z = (d/2)$ that has a k^2 dispersion that is valid far below the fundamental frequency of the cavity, as well as far above, $\Lambda_{IR} \ll \omega_0 \ll \Lambda_{UV}$. We find a decrease of the mass of the electron given by fundamental constants, $\delta m/m_e \sim -(\alpha\hbar/cdm_e)$, signifying a reduced effective light-matter coupling compared to free space—see Fig. 3(c). This change can be attributed to the gapping out of in-plane light modes below ω_0 to which the electron is sensitive when assuming $\Lambda_{IR} \ll \omega_0$. In the case of $\Lambda_{IR} \approx \omega_0$, which could be relevant for bound systems [42], a different result may be expected. More details on the calculation for the idealized Fabry-Pérot cavity can be found in SM [49].

Discussion—We have investigated the coupling of 2D free electrons to light, treating the full continuum of modes of the electromagnetic field instead of relying on single- or few-mode approximations. We extend the definition of ultrastrong coupling to this case as a 10% change in the effective electronic mass, which we compute using a method that we propose in this Letter. Near-field effects close to an insulator can lead to this ultrastrong coupling regime due to coupling to surface modes. Importantly we find that one may compute the mass change of the electron solely from the coupling to the surface modes without the need to renormalize with respect to the vacuum [60,61]. These results further emphasize the experimental prospects of engineering electronic properties via cavities employing near-field effects [8,52,62–72].

It is interesting to consider the applicability of sum rules to the change of mass. A sum rule derived in Ref. [73] directly applies to a perfect metal reproducing the $\Delta m \sim -z^{-1}$ scaling found in the present Letter in the far field

of a Drude metal, while close to interfaces sum rules derived for a frequency dependent permittivity $\epsilon(\omega)$, in Refs. [74,75], are consistent with the scaling $\Delta m \sim z^{-3}$ found here.

The mass of free electrons directly enters in their cyclotron frequency and is well documented in the literature for many materials. As a precise measurement of the predictions made in this Letter, we therefore propose to measure the cyclotron resonance in a dilute, two-dimensional gas of electrons close to an insulator that hosts suitable surface phonon polaritons, for example SrTiO₃. We expect our theory to be readily realizable in 2D van der Waals materials placed on top of appropriately chosen substrates hosting surface phonon polaritons [76], or at the interface of 3D materials with surface states [77].

We note that here we focus on the single-particle coupling. Strong coupling effects could be further amplified by collective strong coupling, where a macroscopic amount of particles couples to the photonic modes, such as in the case of polaritonic chemistry [3,78]. Future studies should investigate the complex behavior expected once the interparticle interactions of a macroscopic ensemble are included [14,79,80].

Acknowledgments—We acknowledge useful discussions with Junichiro Kono, Andrey Baydin, Eugene Demler, Jonathan Curtis, and Maximilian Daschner. M. H. M. is grateful for the support from the Alexander von Humboldt foundation. M. A. S. was funded by the European Union (ERC, CAVMAT, Project No. 101124492). D. M. K. acknowledge support by the DFG via Germany’s Excellence Strategy—Cluster of Excellence Matter and Light for Quantum Computing (ML4Q, Project No. EXC 2004/1, Grant No. 390534769), individual Grant No. 508440990. A. R. and M. R. acknowledge support from the Cluster of Excellence “CUI: Advanced Imaging of Matter”—EXC 2056—Project No. 390715994 and SFB-925 “Light induced dynamics and control of correlated

quantum systems”—Project No. 170620586 of the Deutsche Forschungsgemeinschaft (DFG), and Grupos Consolidados (IT1453-22), and from the Max Planck-New York City Center for Non-Equilibrium Quantum Phenomena at the Flatiron Institute. The Flatiron Institute is a division of the Simons Foundation. M. H. and A. G. acknowledge support by DARPA No. HR00112530313.

The views, opinions and/or findings expressed are those of the authors and should not be interpreted as representing the official views or policies of the Department of Defense or the U.S. Government.

- [1] F. Schlawin, D. M. Kennes, and M. A. Sentef, *Appl. Phys. Rev.* **9**, 011312 (2022).
- [2] A. Frisk Kockum, A. Miranowicz, S. De Liberato, S. Savasta, and F. Nori, *Nat. Rev. Phys.* **1**, 19 (2019).
- [3] F. J. Garcia-Vidal, C. Ciuti, and T. W. Ebbesen, *Science* **373**, eabd0336 (2021).
- [4] J. Bloch, A. Cavalleri, V. Galitski, M. Hafezi, and A. Rubio, *Nature (London)* **606**, 41 (2022).
- [5] M. Ruggenthaler, J. Flick, C. Pellegrini, H. Appel, I. V. Tokatly, and A. Rubio, *Phys. Rev. A* **90**, 012508 (2014).
- [6] H. Hübener, U. De Giovannini, C. Schäfer, J. Andberger, M. Ruggenthaler, J. Faist, and A. Rubio, *Nat. Mater.* **20**, 438 (2021).
- [7] G. Jarc, S. Y. Mathengattil, A. Montanaro, F. Giusti, E. M. Rigoni, R. Sergo, F. Fassioli, S. Winnerl, S. Dal Zilio, D. Mihailovic, P. Prelovšek, M. Eckstein, and D. Fausti, *Nature (London)* **622**, 487 (2023).
- [8] F. Appugliese, J. Enkner, G. L. Paravicini-Bagliani, M. Beck, C. Reichl, W. Wegscheider, G. Scalari, C. Ciuti, and J. Faist, *Science* **375**, 1030 (2022).
- [9] P. Forn-Díaz, L. Lamata, E. Rico, J. Kono, and E. Solano, *Rev. Mod. Phys.* **91**, 025005 (2019).
- [10] J. A. Hutchison, T. Schwartz, C. Genet, E. Devaux, and T. W. Ebbesen, *Angew. Chem., Int. Ed.* **51**, 1592 (2012).
- [11] T. W. Ebbesen, *Acc. Chem. Res.* **49**, 2403 (2016).
- [12] M. Ruggenthaler, N. Tancogne-Dejean, J. Flick, H. Appel, and A. Rubio, *Nat. Rev. Chem.* **2**, 1 (2018).
- [13] D. Sidler, M. Ruggenthaler, C. Schäfer, E. Ronca, and A. Rubio, *J. Chem. Phys.* **156**, 230901 (2022).
- [14] M. Ruggenthaler, D. Sidler, and A. Rubio, *Chem. Rev.* **123**, 11191 (2023).
- [15] K. Hirai, J. A. Hutchison, and H. Uji-i, *Chem. Rev.* **123**, 8099 (2023).
- [16] K. Nagarajan, A. Thomas, and T. W. Ebbesen, *J. Am. Chem. Soc.* **143**, 16877 (2021).
- [17] J. Flick, M. Ruggenthaler, H. Appel, and A. Rubio, *Proc. Natl. Acad. Sci. U.S.A.* **114**, 3026 (2017).
- [18] B. S. Simpkins, A. D. Dunkelberger, and J. C. Owrutsky, *J. Phys. Chem. C* **125**, 19081 (2021).
- [19] M. K. Svendsen, K. S. Thygesen, A. Rubio, and J. Flick, *arXiv:2305.02391*.
- [20] M. A. Sentef, M. Ruggenthaler, and A. Rubio, *Sci. Adv.* **4**, eaau6969 (2018).
- [21] M. K. Svendsen, M. Ruggenthaler, H. Hübener, C. Schäfer, M. Eckstein, A. Rubio, and S. Latini, *arXiv:2312.17374*.
- [22] C. Schäfer, J. Flick, E. Ronca, P. Narang, and A. Rubio, *Nat. Commun.* **13**, 7817 (2022).
- [23] D. N. Basov, R. D. Averitt, and D. Hsieh, *Nat. Mater.* **16**, 1077 (2017).
- [24] C. Ciuti, *Phys. Rev. B* **104**, 155307 (2021).
- [25] J. Fregoni, F. J. Garcia-Vidal, and J. Feist, *ACS Photonics* **9**, 1096 (2022).
- [26] V. Rokaj, J. Wang, J. Sous, M. Penz, M. Ruggenthaler, and A. Rubio, *Phys. Rev. Lett.* **131**, 196602 (2023).
- [27] G. Kipp *et al.*, *arXiv:2403.19745*.
- [28] V. Rokaj, M. Penz, M. A. Sentef, M. Ruggenthaler, and A. Rubio, *Phys. Rev. B* **105**, 205424 (2022).
- [29] A. Chakraborty and F. Piazza, *Phys. Rev. Lett.* **127**, 177002 (2021).
- [30] H. Gao, F. Schlawin, M. Buzzi, A. Cavalleri, and D. Jaksch, *Phys. Rev. Lett.* **125**, 053602 (2020).
- [31] F. Schlawin, A. Cavalleri, and D. Jaksch, *Phys. Rev. Lett.* **122**, 133602 (2019).
- [32] G. M. Andolina, A. De Pasquale, F. M. D. Pellegrino, I. Torre, F. H. L. Koppens, and M. Polini, *Phys. Rev. B* **109**, 104513 (2024).
- [33] E. Viñas Boström, A. Sriram, M. Claassen, and A. Rubio, *npj Comput. Mater.* **9**, 1 (2023).
- [34] A. Chiochetta, D. Kiese, C. P. Zelle, F. Piazza, and S. Diehl, *Nat. Commun.* **12**, 5901 (2021).
- [35] P. Fadler, K. P. Schmidt, J. Li, and M. Eckstein, *Phys. Rev. B* **109**, 085149 (2024).
- [36] K. Masuki and Y. Ashida, *Phys. Rev. B* **109**, 195173 (2024).
- [37] V. Rokaj, M. Ruggenthaler, F. G. Eich, and A. Rubio, *Phys. Rev. Res.* **4**, 013012 (2022).
- [38] C. J. Eckhardt, G. Passetti, M. Othman, C. Karrasch, F. Cavaliere, M. A. Sentef, and D. M. Kennes, *Commun. Phys.* **5**, 1 (2022).
- [39] G. Passetti, C. J. Eckhardt, M. A. Sentef, and D. M. Kennes, *Phys. Rev. Lett.* **131**, 023601 (2023).
- [40] O. Dmytruk and M. Schiró, *Phys. Rev. B* **103**, 075131 (2021).
- [41] J. Li, D. Golez, G. Mazza, A. J. Millis, A. Georges, and M. Eckstein, *Phys. Rev. B* **101**, 205140 (2020).
- [42] D. M. Welakuh, V. Rokaj, M. Ruggenthaler, and A. Rubio, *Phys. Rev. Res.* **7**, 013093 (2025).
- [43] F. Hiroshima and K. R. Ito, *Rev. Math. Phys.* **19**, 405 (2007).
- [44] O. Dmytruk and M. Schiró, *Commun. Phys.* **5**, 1 (2022).
- [45] J. B. Curtis, M. H. Michael, and E. Demler, *Phys. Rev. Res.* **5**, 043118 (2023).
- [46] S. Latini, D. Shin, S. A. Sato, C. Schäfer, U. De Giovannini, H. Hübener, and A. Rubio, *Proc. Natl. Acad. Sci. U.S.A.* **118**, e2105618118 (2021).
- [47] H. P. Ojeda Collado, M. H. Michael, J. Skulte, A. Rubio, and L. Mathey, *Phys. Rev. Lett.* **133**, 116901 (2024).
- [48] E. Vlasuk, V. K. Kozin, J. Klinovaja, D. Loss, I. V. Iorsh, and I. V. Tokatly, *Phys. Rev. B* **108**, 085410 (2023).
- [49] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/2fw2-lbhy> for (a) derivation of the Schrieffer-Wolff transformation that yields Eq. (4), (b) details on the calculation of effective mass in an idealized Fabry-Pérot cavity, (c) mode functions, (d) details for the computation of exact local density of states, and (e) derivation of the mass renormalization due to surface phonon polaritons, which includes Refs. [32,50–56].

- [50] W. L. Barnes, S. A. R. Horsley, and W. L. Vos, *J. Opt.* **22**, 073501 (2020).
- [51] C. R. Gubbin, S. A. Maier, and S. De Liberato, *Phys. Rev. B* **94**, 205301 (2016).
- [52] R. Riolo, A. Tomadin, G. Mazza, R. Asgari, A. H. MacDonald, and M. Polini, *Proc. Natl. Acad. Sci. U.S.A.* **122**, e2407995122 (2025).
- [53] C. Schäfer, M. Ruggenthaler, V. Rokaj, and A. Rubio, *ACS Photonics* **7**, 975 (2020).
- [54] R. J. Glauber and M. Lewenstein, *Phys. Rev. A* **43**, 467 (1991).
- [55] J. M. Pitarke, V. M. Silkin, E. V. Chulkov, and P. M. Echenique, *Rep. Prog. Phys.* **70**, 1 (2007).
- [56] H. Hörner, L. M. Rachbauer, S. Rotter, and U. Leonhardt, *Phys. Rev. B* **108**, 235430 (2023).
- [57] M. D. Schwartz, *Quantum Field Theory and the Standard Model* (Cambridge University Press, Cambridge, England, 2014).
- [58] G. D. Mahan, *Many-Particle Physics* (Springer US, Boston, MA, 2000).
- [59] J.-F. Ge, Z.-L. Liu, C. Liu, C.-L. Gao, D. Qian, Q.-K. Xue, Y. Liu, and J.-F. Jia, *Nat. Mater.* **14**, 285 (2015).
- [60] K. Lenk, J. Li, P. Werner, and M. Eckstein, *Phys. Rev. B* **106**, 245124 (2022).
- [61] C. J. Eckhardt, S. Chattopadhyay, D. M. Kennes, E. A. Demler, M. A. Sentef, and M. H. Michael, *Nat. Commun.* **15**, 2300 (2024).
- [62] G. Kipp *et al.*, [arXiv:2403.19745](https://arxiv.org/abs/2403.19745).
- [63] A. Thomas, E. Devaux, K. Nagarajan, T. Chervy, M. Seidel, D. Hagenmüller, S. Schütz, J. Schachenmayer, C. Genet, G. Pupillo, and T. W. Ebbesen, *J. Chem. Phys.* **162**, 134701 (2025).
- [64] A. Thomas, E. Devaux, K. Nagarajan, G. Rogez, M. Seidel, F. Richard, C. Genet, M. Drillon, and T. W. Ebbesen, *Nano Lett.* **21**, 4365 (2021).
- [65] G. Scalari, C. Maissen, D. Turčinková, D. Hagenmüller, S. De Liberato, C. Ciuti, C. Reichl, D. Schuh, W. Wegscheider, M. Beck, and J. Faist, *Science* **335**, 1323 (2012).
- [66] C. Maissen, G. Scalari, F. Valmorra, M. Beck, J. Faist, S. Cibella, R. Leoni, C. Reichl, C. Charpentier, and W. Wegscheider, *Phys. Rev. B* **90**, 205309 (2014).
- [67] A. Bayer, M. Pozimski, S. Schambeck, D. Schuh, R. Huber, D. Bougeard, and C. Lange, *Nano Lett.* **17**, 6340 (2017).
- [68] M. Halbhüser, J. Mornhinweg, V. Zeller, C. Ciuti, D. Bougeard, R. Huber, and C. Lange, *Nat. Photonics* **14**, 675 (2020).
- [69] H. Yu, Y. Peng, Y. Yang, and Z.-Y. Li, *npj Comput. Mater.* **5**, 1 (2019).
- [70] M. S. Tame, K. R. McEnery, S. K. Özdemir, J. Lee, S. A. Maier, and M. S. Kim, *Nat. Phys.* **9**, 329 (2013).
- [71] H. Herzig Sheinfux, L. Orsini, M. Jung, I. Torre, M. Ceccanti, S. Marconi, R. Maniyara, D. Barcons Ruiz, A. Hötger, R. Bertini, S. Castilla, N. C. H. Hesp, E. Janzen, A. Holleitner, V. Pruneri, J. H. Edgar, G. Shvets, and F. H. L. Koppens, *Nat. Mater.* **23**, 499 (2024).
- [72] F. Helmrich, I. Khanonkin, M. Kroner, G. Scalari, J. Faist, A. Imamoglu, and T. F. Nova, [arXiv:2408.00189](https://arxiv.org/abs/2408.00189).
- [73] S. M. Barnett and R. Loudon, *Phys. Rev. Lett.* **77**, 2444 (1996).
- [74] S. Scheel, *Phys. Rev. A* **78**, 013841 (2008).
- [75] S. Sanders and A. Manjavacas, *ACS Photonics* **5**, 2437 (2018).
- [76] A. K. Geim and I. V. Grigorieva, *Nature (London)* **499**, 419 (2013).
- [77] A. Goetzberger, E. Klausmann, and M. J. Schulz, *CRC Crit. Rev. Solid State Sci.* **6**, 1 (1976).
- [78] T. W. Ebbesen, A. Rubio, and G. D. Scholes, *Chem. Rev.* **123**, 12037 (2023).
- [79] D. Sidler, T. Schnappinger, A. Obzhairov, M. Ruggenthaler, M. Kowalewski, and A. Rubio, *J. Phys. Chem. Lett.* **15**, 5208 (2024).
- [80] D. Sidler, C. Schäfer, M. Ruggenthaler, and A. Rubio, *J. Phys. Chem. Lett.* **12**, 508 (2021).