Multielectron Ground State Electroluminescence

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The ground state of a cavity-electron system in the ultrastrong coupling regime is characterized by the presence of virtual photons. If an electric current flows through this system, the modulation of the light-matter coupling induced by this nonequilibrium effect can induce an extracavity photon emission signal, even when electrons entering the cavity do not have enough energy to populate the excited states. We show that this ground state electroluminescence, previously identified in a single-qubit system [Phys. Rev. Lett. **116**, 113601 (2016)] can arise in a many-electron system. The collective enhancement of the light-matter coupling makes this effect, described beyond the rotating wave approximation, robust in the thermodynamic limit, allowing its observation in a broad range of physical systems, from a semiconductor heterostructure with flatband dispersion to various implementations of the Dicke model.

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Introduction.—When the interaction between light and matter is stronger than the coupling to the environment, a variety of hybridization effects can be observed. In the context of cavity quantum electrodynamics, realizing this "strong-coupling" regime has been achieved in different ways, for example, by reducing the losses of the system [1], by enhancing the vacuum electromagnetic field in one-dimensional cavities [2], by increasing the dipole moment of the atom [3], or by taking advantage of collective properties [4]. Building upon these strategies, it has been possible to engineer light-matter couplings up to a significant fraction of the bare energies of the bare light and matter modes themselves [1,3,5–32].

This new cavity quantum electrodynamics (QED) "ultrastrong" regime has made possible the observation and study of a range of unique physical effects [33–46]. Among these phenomena are the ones originating from the hybridization of the ground state. This hybridization leads to a ground state photonic population that is sometimes called "virtual", as it is energetically forbidden from leaking into the environment. However, there are several proposals describing how these hybridized ground states can be observed, typically by modulating some system parameter [40–45,47], akin to the way the dynamical Casimir effect relies on amplifying vacuum fluctuations [48–53].

In particular, in Ref. [54] it was shown that the passage of an electronic current through a device where, within the device, electrons ultrastrongly couple to light in a cavity can result in extracavity emission, i.e., the conversion of virtual to real photons. In Ref. [54], such "ground state electroluminescence" was predicted for systems in which a *single* electron at a time interacts ultrastrongly with the cavity mode [55–63]. In this Letter, we analyze ground state electroluminescence in a much more general scenario in which *many* electrons at the same time are allowed to interact with the cavity mode [64–71]. This allows for

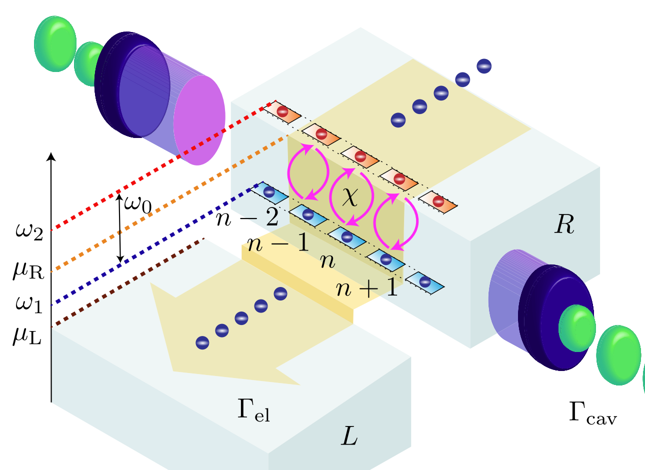

FIG. 1. A right lead (R) is connected to a left lead (L) via a middle region, the two elements kept at chemical potentials μ_R and μ_L, respectively, by applying an electrical bias, which induces an electron current quantified by a rate Γ_{el} for the free electrons (blue spheres) flowing out of the device. Sandwiched between the leads, a solid-state cavity (dark purple disks), enhances the electronic coupling to the photonic vacuum field (light purple disk cross section), at a strength quantified by χ for each electron. The bare frequency difference between the two electronic flatbands, $\omega_0 = \omega_2 - \omega_1$, separates the lower states (blue) and upper states (red). The presence of virtual photons inside the cavity induces an extracavity photon emission (green blobs) from the polaritonic ground state, at a rate Γ_{cav}.

stronger effective couplings through collective effects in a solid-state device [5,67,72–74], as sketched in Fig. 1.

As we will show, one could expect the electroluminescence effect to be washed out in a system containing many electrons because, while the coupling is enhanced by collective effects, the conversion of virtual to real photons relies on a process where an electron leaving the system effectively changes the light-matter coupling. In this many-electron system, such an effective modulation of the light-matter coupling is suppressed with the number of electrons, so one might expect that this negates the enhanced collective coupling.

However, surprisingly, we find here that the combination of collective coupling and the *multielectron* nature of the current combine to make the ground state electroluminescence macroscopically robust even in the thermodynamic limit. The transport-induced luminescent effect can be estimated by an intuitive bosonic theory that goes beyond the rotating wave approximation (RWA) by including counterrotating terms perturbatively. In the Supplementary Material (SM) [75], we test this model against a full bosonic model that includes non-RWA terms nonperturbatively, and a second-quantization fermionic theory, finding excellent agreement.

Light-matter system.—We consider a prototypical many-body fermionic system interacting with light in a solid-state quantum device. The model system can be generalized further due to the approximations that we will make, but, for definiteness, we begin by considering two electronic bands containing a maximum of $2N_T$ electrons which interact with a single electromagnetic mode confined in a cavity. We further neglect electron-electron interactions, band dispersion, and higher excitations. We thus consider a two flatband electronic model such that it can be described by the Hamiltonian ($\hbar = 1$ hereafter)

$$H = \omega_c a^\dagger a + \sum_n (\omega_1 c_{1,n}^\dagger c_{1,n} + \omega_2 c_{2,n}^\dagger c_{2,n}) + D(a + a^\dagger)^2 + \chi(a + a^\dagger)\sum_n (c_{2,n}^\dagger c_{1,n} + c_{1,n}^\dagger c_{2,n}), \tag{1}$$

where $c_{1,n}$ ($c_{2,n}$) represents the annihilation operator for the nth ($n = 1, \ldots, N_T$) fermion in the first (second) state with energy ω_1 (ω_2). Note that in Eq. (1) we are counting each fermion over the index n; in several solid-state systems, this can be shown to be equivalent to a model for flatbands, as in intersubband transitions with finite real in-plane momentum [7,10,13], in the limit of small photon momentum or strong magnetic confinement. In more general contexts, it may be required to include the photonic momentum, which can induce diagonal transitions [35,76–80]. The annihilation operator a is associated with a cavity mode of frequency ω_c. The light-matter interaction has strength χ, and the potential energy of the electromagnetic field is proportional to the frequency $D = N\chi^2/\omega_0$, relative to the diamagnetic term [25,26,46], where $\omega_0 = \omega_2 - \omega_1$.

To begin our analysis, we divide the Hilbert space in sectors closed under the Hamiltonian evolution. They are characterized by the set of sites occupied by a single electron $\{N\}$, the set of sites occupied by two electrons $\{N_2\}$, and the number of photons in the cavity.

Within each of these sectors, the coherent dynamics can be described by

$$H = \omega_c a^\dagger a + \omega_0 S^3 + \chi(a + a^\dagger)(S^- + S^+), \tag{2}$$

which takes the standard form of the Dicke Hamiltonian. The interaction of a cavity mode of frequency ω_c with a matter excitation of frequency ω_0, where $\omega_0 = \omega_2 - \omega_1$, is described beyond the RWA. Here we defined $\sigma_n^- = c_{1,n}^\dagger c_{2,n}$, $S_n^- = \sigma_n^-$, and $S_n^3 = \sigma_n^z/2$, where σ_n^α is the α-direction Pauli matrix operator. In Eq. (2), we performed a fermion-to-spin transformation that, with respect to Eq. (1), involves no approximations. The parameters have been renormalized following the Bogoliubov transformation needed to reabsorb the diamagnetic term proportional to D in Eq. (1) (see Sec. II of the SM [75]). For the sake of generality in Eq. (1), we neglect the Coulomb interaction, which would depend on microscopic details. In the case of parallel subbands, a theoretical description of electron-electron interactions in the bosonic approximation, directly applicable to our approach [81,82], can be completely captured by a renormalization of the system transition frequency, the so-called depolarization shift [83], and by a more complex functional dependency between the electronic operators c and the collective excitation operators S^α. While such more complex relations remain quadratic, and could thus be incorporated into our treatment, their deviations from Eq. (1) scale with the ratio between the plasma frequency, ω_p, and the bare excitation frequency, ω_0. Equation (1) thus remains quantitatively accurate while $\omega_0 \gg \omega_p$.

Environment.—We are interested in studying the effects of three environments on this model: a left (L) and right (R) electronic reservoir, which give rise to the electronic current, and the extracavity electromagnetic modes, into which the photons are emitted. The total environment-system interaction Hamiltonian is $H^I = H_{\rm el}^I + H_{\rm cav}^I = H_L^I + H_R^I + H_{\rm cav}^I$. Our aim is to compute the transition rates among eigenstates of the system induced by the interaction Hamiltonian, H^I, representing the physical interaction with the environmental degrees of freedom. We can model the interaction with the electronic reservoirs as $H_L^I = \lambda \sum_{n,\zeta}[(c_{1,n} + c_{2,n})c_{L;n,\zeta}^\dagger + \text{H.c.}]$, and identically for H_R^I (change $L \to R$), thus assuming that the energy scale, λ, is equal for the two fermionic reservoirs. The operators $c_{L(R);n,\zeta}$ label the annihilation operators for a fermion associated with a degree of freedom n and ζ in the left (right) reservoir.

Since we are interested in strong-coupling effects between light and matter only within the system, we treat all three environments perturbatively. They induce transitions only between eigenstates of the system, as given by Fermi's "golden rule." The total *electron-transport rates* can be calculated by summing over single electron scattering processes, $\Gamma_{\rm el}^{\alpha\to\beta}=\sum_n\Gamma_{{\rm el},n}^{\alpha\to\beta}$ (see Sec. I of the SM [75] for details), where

$$\Gamma_{{\rm el},n}^{\alpha\to\beta}\propto\Gamma_{\rm el}|M_{\alpha\beta}^n|^2=\Gamma_{\rm el}|\langle\beta|(c_{1,n}+c_{2,n})|\alpha\rangle|^2, \quad (3)$$

where α and β are the initial and final states for the system, $\Gamma_{\rm el}$ is the electron tunneling rate, and $M_{\alpha\beta}^n$ provides the *electron-current transition matrix element* for the nth electron site.

To calculate the ground state electroluminescence rate, we consider that, when N electrons are in the device and the device is in the hybridized light-matter ground state, $|\alpha\rangle=|G_N\rangle$, an electron within the device can leave, reducing the electron number to $(N-1)$. When an electron leaves, it can, due to the ground state light-matter hybridization, result in a transition to an excited state of the hybridized system with $(N-1)$ electrons, $|\beta\rangle=|E_{N-1}\rangle$, which contains a nonzero photonic population. We assume that the cavity loss rate $\Gamma_{\rm cav}$ is much faster than the electronic rates $\Gamma_{\rm el}$, such that this excited state immediately decays and emits an extracavity photon, decaying to the $(N-1)$ ground state, $|G_{N-1}\rangle$; this emission, arising only because the ground state itself contains photons, is the electroluminescence we want to produce. In addition, by imposing a chemical potential across the system which forbids electrons from entering directly into excited states of the coupled system, $\mu_L<\mu_R<\omega_2$, one can suppress "regular" electroluminescence and ensure that the observed photon emission arises only from the ground state.

Under the above assumptions ($\Gamma_{\rm cav}\gg\Gamma_{\rm el}$ and energetically forbidden regular electroluminescence), the overall rate of ground state–sourced photonic emission depends upon the electron-current transition matrix elements, $M_{\alpha\beta}^n$ of Eq. (3). This reduces to the problem of calculating the properties of the *ground state*, $|G_N\rangle$, and the various possible excited states, $|E_{N-1}\rangle$, that contribute to these transitions, and the overlap with the operators which destroy electrons. In the SM [75], we present a fully fermionic calculation of such rates, but it is much more instructive to first consider a simpler bosonic approximation which captures the essential physics.

Bosonic approximation.—To proceed further, we assume that thermalization effects are such that we can neglect double-occupied electron sites, $N_2\simeq 0$, and consider the following approximate Holstein-Primakoff transformation $S_+=\sqrt{N}b^\dagger+O(|b^\dagger b/\sqrt{N}|)$, and $S_z=b^\dagger b-j_N$ in terms of an effective bosonic mode b. In a dilute regime in which the number of electronic excitations is much smaller than the total number of electrons, we can neglect terms of order $|b^\dagger b/\sqrt{N}|$ and rewrite Eq. (2) as

$$\begin{aligned}H&\simeq\omega_c a^\dagger a+\omega_0 b^\dagger b+g_N(ab^\dagger+a^\dagger b)+g_N(ab+a^\dagger b^\dagger)\\&=H_{\rm JC}+g_N(ab+a^\dagger b^\dagger)=H_{\rm JC}+V,\end{aligned} \quad (4)$$

up to $\mathbb{C}$ numbers and terms of order $1/\sqrt{N}$, and where $g_N=\sqrt{N}\chi$ is the *bosonic* light-matter coupling. While the *full* bosonic Hamiltonian of Eq. (4) can be diagonalized analytically (see Secs. II and III of the SM [75]) to most clearly highlight the main idea behind the processes studied here, we will consider the counterrotating term V as a perturbation of the Jaynes-Cummings (JC) term, $H_{\rm JC}$, and rewrite Eq. (4) as $H\simeq\omega^- p_-^\dagger p_-+\omega^+ p_+^\dagger p_++V$, where $p_\pm^\dagger=\alpha_a^\pm a^\dagger+\alpha_b^\pm b^\dagger$ are the polaritonic excitations of the JC part of the original Hamiltonian, and where the explicit expression for the polariton energies $\omega^\pm$ and the dimensionless coefficients $\alpha_a^\pm$ and $\alpha_b^\pm$ are given in Sec. III of the SM [75]. First-order perturbation theory in V gives the following expression for the ground state and single-polariton states for the non-RWA system,

$$\begin{aligned}|G_N\rangle&=|G_N^{(0)}\rangle-\beta_{++}|++_N^{(0)}\rangle-\beta_{+-}|+-_N^{(0)}\rangle-\beta_{--}|--_N^{(0)}\rangle,\\|\pm_N\rangle&=|\pm_N^{(0)}\rangle+\cdots,\end{aligned} \quad (5)$$

where we introduced the perturbative coefficients $\beta_{\pm\pm}=-\sqrt{2}g_N(\alpha_a^\mp\alpha_b^\mp)/(2\omega^\pm)$ and $\beta_{+-}=g_N(\alpha_a^+\alpha_b^-+\alpha_a^-\alpha_b^+)/(\omega^++\omega^-)$, which are explicitly derived in Sec. III of the SM [75]. For the sake of clarity, we omitted higher-order terms in the expansion, which will not contribute to our results [indicated by the suspended dots in Eq. (5)]. Note that the ground state is a superposition of the Fermi sea for the system with N electrons in the first band and no cavity photons, $|G_N^{(0)}\rangle=\bigotimes_{n\in\mathbf{N}}c_{1,n}^\dagger|0_{\rm el}\rangle|0_{\rm ph}\rangle$, where $|0_{\rm el}\rangle$ and $|0_{\rm ph}\rangle$ represent the electronic and photonic vacuum states, respectively, and $\mathbf{N}$ is the set of occupied sites of cardinality N. The double-polariton states of the unperturbed basis, similar to all of the other excited states, can be defined by multiple applications of the JC polaritons, e.g., $|\pm_N^{(0)}\rangle=p_\pm^\dagger|G_N^{(0)}\rangle$ for the single-polariton states, and $|\pm\pm_N^{(0)}\rangle=(p_\pm^\dagger)^2|G_N^{(0)}\rangle/\sqrt{2}$ and $|+-_N^{(0)}\rangle=p_+^\dagger p_-^\dagger|G_N^{(0)}\rangle$ for the double-polariton states.

Ground state electroluminescence.—We assume that the system is initially in its ground state and, by emission of an electron, can transition to an excited state, which then decays by emitting photons, the process which constitutes ground state electroluminescence.

Setting $\mu_L<\mu_R<\omega_2$, we obtain that $\Gamma_{\rm el}^{G\to B}=\Gamma_{\rm el}\delta_{G,B}$, where G labels the ground state and B labels any state (see Sec. I of the SM [75] for details). This condition ensures that the regular direct electroluminescence is energetically forbidden and allows for the undiluted ground state process to occur. The ground state polariton creation leading to photon emission can be estimated as $\Gamma_{\rm GSE}=\sum_{E=\{\pm,\pm\pm,\pm\mp,\ldots\}}\Gamma_{\rm el}^{G\to E}$.

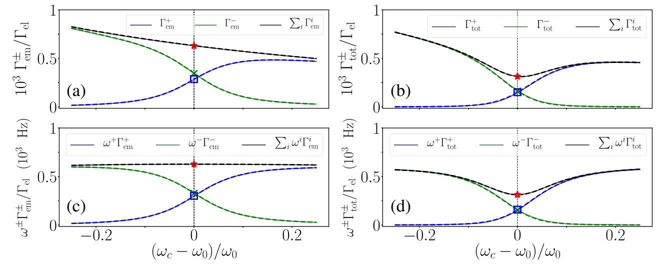

FIG. 2. (a),(c) Polariton emission rates $\Gamma_{\mathrm{em}}^{B'}$ and fluxes $\omega^{B'}\Gamma_{\mathrm{em}}^{B'}$, in units of the total electron-transport rate Γ_{el} for the upper polariton ($B' = +$, blue curves) and lower polariton ($B' = -$, green curves), and sum of the two signals (black curves) versus the normalized detuning for $g_N/\omega_0 = 0.05$. (b),(d) Total photon emission rates, from Eq. (6), and fluxes. Solid curves correspond to the bosonic RWA quantities, dashed curves to the full bosonic model developed in the SM [75].

We begin by calculating the transition from the ground state, $|G_N\rangle$, to the single-polariton states, $|\pm_{N-1}\rangle$. From Eq. (3), we have $M_{G\pm}^{n} = \langle \pm_{N-1}|(c_{1,n} + c_{2,n})|G_N\rangle$, where the state $|\pm_{N-1}\rangle$ is the state with $(N-1)$ electrons due to the tunneling of the nth electron. Here we use the perturbative expressions to expand the non-RWA contribution in these states, given in Eq. (5). To proceed further, one needs to calculate expectation values of fermionic operators onto light-matter many-body states intrinsically expressed in terms of polariton operators. This task can be crucially simplified by using Eqs. (1) and (2) and the Holstein-Primakoff mapping to rewrite $\sqrt{N}b = S_- = \sum_n c_{1,n}^{\dagger} c_{2,n}$, which, using the definition of the polariton modes, immediately gives $[p_{\pm}^{\dagger}, c_{1,n} + c_{2,n}] = \alpha_b^{\pm}[b^{\dagger}, c_{1,n} + c_{2,n}] = -(1/\sqrt{N})\alpha_b^{\pm} c_{1,n}$, which holds up to order $1/\sqrt{N}$ and holds linearly in any of the perturbative parameters (see Sec. III of the SM [75] for details). We then obtain an explicit expression for the matrix elements contained in Eq. (3), $M_{G\pm}^{n} = (\sqrt{2}\beta_{\pm\pm}\alpha_b^{\pm} + \beta_{+-}\alpha_b^{\mp})/\sqrt{N}$, which, together with the initial working condition $\Gamma_{\mathrm{cav}} \gg \Gamma_{\mathrm{el}}$, allows us to estimate the photon emission rate from the ground state to the single-polariton states, $\Gamma_{\mathrm{em}}^{\pm}$, as $\Gamma_{\mathrm{em}}^{\pm} \simeq \Gamma_{\mathrm{el}}^{G\to\pm} \propto \sum_n |M_{G\pm}^{n}|^2 = |\sqrt{2}\beta_{\pm\pm}\alpha_b^{\pm} + \beta_{+-}\alpha_b^{\mp}|^2$, that is, $\Gamma_{\mathrm{em}}^{\pm} = O(\eta^2) = O(N\chi^2/\omega_0^2)$, where $\eta = g_N/\omega_0$.

The contributions to Γ_{GSE} from the higher-excited states (which are double-polariton states, $|G_N\rangle \to |\pm\pm_{N-1}\rangle$ and $|G_N\rangle \to |+-_{N-1}\rangle$), are of $O(\eta^2/N)$, as detailed in Sec. III of the SM [75]. Thus the dominant contribution to the ground state electroluminescence (GSE) involves the single-polariton transitions, giving the total GSE rate $\Gamma_{\mathrm{GSE}} \simeq \Gamma_{\mathrm{em}}^{+} + \Gamma_{\mathrm{em}}^{-} = O(\eta^2)$.

Remarkably, this emission is of the same order of magnitude of the one predicted in systems containing a single electron [54] (but following the enhanced collective coupling rate, $\eta^2 = N\chi^2/\omega_0^2$). In the single electron case [54], the light-matter coupling was strongly modulated, as the single electron coupling was assumed to be ultrastrong, and the effective modulation of the coupling due to the emission of the electron was large. Here instead, the tunneling of a single electron (among N total) only minimally modulates the light-matter coupling, yet a *collective* enhancement occurs, to ensure the same η^2 scaling. This can be interpreted as a superradiant enhancement with respect to the single-particle light-matter coupling of the fermionic system, χ, and the overall large electron current.

In Fig. 2(a), we plot the GSE rates for the upper (blue curves) and lower (green curves) polariton channels versus the frequency detuning, as well as the total rate (black curves), calculating them using the JC polaritons (solid curves) and comparing them to the full bosonic model that retains the counterrotating terms (see the SM [75]). There is a clear inversion of the contribution to polariton creation versus detuning. In Fig. 2(c), we plot the energy flux of such an emission for $g_N = 0.05\omega_0$, which shows a peak at zero detuning and shows that for $\omega_c \ll \omega_0$ the extracted energy that is associated with the emission is limited, and that there is little dependency on the detuning, both results that are in accordance with recent predictions for dissipative systems [84]. The plots of Figs. 2(a) and 2(c) give indications for experiments based on electric current measurement.

In a photodetection spectroscopic experiment, the extracavity photonic emission is the product of two processes: First, there is the polariton scattering due to the extraction of an electron, calculated in Γ_{GSE} and which we have shown to be dependent on the $|G_N\rangle \to |\pm_{N-1}\rangle$ channel. Then there is a second relaxation process that involves the emission of a photon, $|\pm_{N-1}\rangle \to |G_{N-1}\rangle$, occurring with probability $|\alpha_{\mathrm{ph}}^{\pm}|^2$, proportional to the Hopfield coefficients associated with light (see the SM [75] for details). Thus the multi-electron GSE scattering first creates real polaritons in the cavity, which will then escape by emitting photons at their own eigenfrequencies. The spectrum of the system is thus made of two Lorentzian peaks at the polariton frequencies ω^- and ω^+. The effective total photon emission rate, $\Gamma_{\mathrm{tot}} = \Gamma_{\mathrm{tot}}^{+} + \Gamma_{\mathrm{tot}}^{-}$, needs to take into account also the efficiency of this conversion, which is determined by the cavity characteristic rate, Γ_{cav}, and by the rate of conversion of the bright polaritons into dark polaritons, $\Gamma_{\mathrm{dark}}^{\pm}$. For a leaky cavity, $\Gamma_{\mathrm{cav}} \gg \Gamma_{\mathrm{dark}}^{\pm}$, we have

$$\Gamma_{\mathrm{tot}}^{\pm} = |\alpha_{\mathrm{ph}}^{\pm}|^2 \frac{\Gamma_{\mathrm{em}}^{\pm}\Gamma_{\mathrm{cav}}}{\Gamma_{\mathrm{dark}}^{\pm} + \Gamma_{\mathrm{cav}}} \simeq |\alpha_{\mathrm{ph}}^{\pm}|^2 \Gamma_{\mathrm{em}}^{\pm}. \tag{6}$$

In Figs. 2(b) and 2(d), we plot such total photon emission rates and fluxes, which show that the decrease in photon collection is very modest, with respect to the electric signal measurement [Figs. 2(a) and 2(c)], by which they are normalized. In Fig. 3, we show the polariton emission rate and flux versus cavity-matter frequency detuning and of the

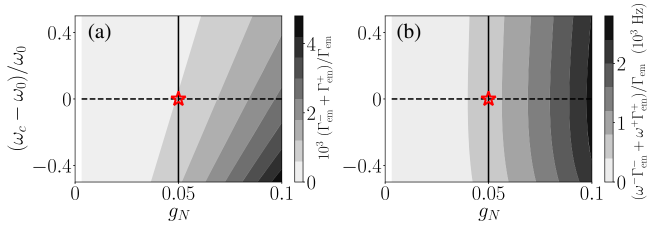

FIG. 3. (a) Polariton emission rate $\Gamma_{em} = \Gamma^-_{em} + \Gamma^+_{em}$ and (b) flux as a function of the frequency detuning and coupling strength, setting $\chi = 3\times10^{-3}\omega_0$ fixed and varying N, and thus $g_N = \sqrt{N}\chi$. The vertical solid black lines correspond to the cut shown in Fig. 2. The resonance condition is marked by dashed horizontal lines (and vertical lines in Fig. 2).

coupling, keeping fixed the single-particle coupling constant, χ, which characterizes solid-state structures with flatbands, so that we can span a wide range of effective light-matter coupling values up to $g_N = 0.1\omega_0$. Since $g_N = \sqrt{N}\chi$, moving rightward in the parameter space can be achieved simply by increasing the number of emitters in the system without requiring the light-matter coupling of the microscopic model to be ultrastrong. The contour plot relative to the emission rate ($\Gamma^{\pm}_{em}$) [Fig. 3(a)] shows an asymmetry in the detuning, favoring, at fixed g_N, the lower polariton emission. The flux ($\omega^{\pm}\Gamma^{\pm}_{em}$) [Fig. 3(b)] shows that the small emission frequency of the lower polariton curbs down the asymmetry, which is consistent with previous predictions [84].

Realizations.—The interplay of collective photonic effects in the presence of local dissipation and transport has recently been studied in several many-body fermionic systems [73,85–93]. Doped semiconductor quantum wells offer a many-body platform in which transport and ground state properties of cavity QED can be investigated [4,10,13,35,71,76,78,94–99].

Intersubband transitions in the conduction band of these systems (the first devices to reach the ultrastrong coupling regime [97]) allow us to dynamically control the electron density with external fields [4,94–96,98,100–104]. Thus multielectron GSE would be an effect relatively easy to explore in experiments compared to other features arising from vacuum fluctuations, such as the nonadiabatic modulation of the coupling strength.

Other candidate systems are superconducting circuits [105–114] and hybrid solid-state architectures [14, 55–57,115–126], especially quantum-dot based systems [55–58,60–63,127–136].

Conclusions.—In conclusion, we have described a novel mechanism for light emission controlled by an electric current, occurring from the ground state of a many-body dissipative light-matter system in the regime of ultrastrong light-matter coupling.

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