Ensemble-Induced Strong Light-Matter Coupling of a Single Quantum Emitter

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We discuss a technique to strongly couple a single target quantum emitter to a cavity mode, which is enabled by virtual excitations of a nearby mesoscopic ensemble of emitters. A collective coupling of the latter to both the cavity and the target emitter induces strong photon nonlinearities in addition to polariton formation, in contrast to common schemes for ensemble strong coupling. We demonstrate that strong coupling at the level of a single emitter can be engineered via coherent and dissipative dipolar interactions with the ensemble, and provide realistic parameters for a possible implementation with SiV^- defects in diamond. Our scheme can find applications, amongst others, in quantum information processing or in the field of cavity-assisted quantum chemistry.

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In the strong coupling regime of light-matter interactions, the coherent exchange of energy between matter and light occurs at a rate faster than all dissipative processes [1,2]. Strong coupling has been actively pursued in atomic [3–8] and molecular [9–17] physics, circuit quantum electrodynamics (QED) [18–22], and condensed-matter physics [23–26], with applications in, e.g., quantum information processing [27–33], transport [34–40], optomechanics [41–46], and quantum chemistry [47–53].

For a single quantum emitter, strong coupling is accompanied by nonlinear quantum effects leading to, e.g., the generation of nonclassical states of light used in photon quantum gates [54–56]. Reaching strong coupling, however, is often difficult as it relies on the integration of small volume resonators with large quality factors. For mesoscopic ensembles this requirement is relaxed as the coupling can be collectively enhanced by a factor $\sqrt{N}$ (for N emitters) [5,57,58]. This allows for the realization of collective polariton dynamics [59,60], however, it does not enhance nonlinear quantum effects for the cavity mode [61–64], limiting its use for, e.g., quantum optics and quantum information processing [65].

Here, we introduce an alternative approach where an ensemble of two-level quantum emitters with negligible populations modifies the cavity coupling of an individual beneficiary *target* two-level system. This allows for bringing the single target emitter from the weak to the strong coupling regime, even when the mediating ensemble is only weakly coupled to the cavity. The induced strong coupling is visible in collectively enhanced single emitter-cavity vacuum Rabi splittings. Remarkably, here the resulting photon nonlinearities are also enhanced by collective effects. We identify two distinct mechanisms for reaching strong coupling in the target-cavity system, based on an increase of coherent coupling strength or a modification of dissipation, respectively. In the first case, we provide a possible implementation of our scheme with realistic parameters considering silicon-vacancy (SiV^-) centers in diamond coupled to a microcavity.

Our results can provide a viable path towards the realization of quantum nonlinear optics experiments within cavity QED with moderate light-matter coupling strengths or cavity quality factors. In contrast to schemes using active dressing of, e.g., Rydberg states [66,67], our setup solely relies on virtual excitations of a nearby ensemble.

The basic scheme is shown in Fig. 1(a). The target emitter A and N ensemble emitters denoted as B are coupled to a cavity mode via a Jaynes-Cummings interaction [68] $H_{\rm JC} = a(g_A\sigma_A^+ + \sum_{\ell=1}^{N} g_\ell \sigma_\ell^+) + \text{H.c.}$ and to each other via coherent dipole-dipole interactions [69] $H_{\rm DD} = \sum_{i\neq j}\sqrt{\gamma_i\gamma_j}g(\vec{r}_{ij})\sigma_i^+\sigma_j^-$ (where $i,j = A,1,\ldots,N$). Here, $\sigma_i^\pm$ denote the raising and lowering operators for the emitters and $g_i = g_i^{(0)}\cos(ky_i)$ their position-dependent coupling strength to the cavity mode (wavelength λ, $k = 2\pi/\lambda$, annihilation operator a, and cavity mode profile constant along x and z). The coherent dipole-dipole coupling strengths depend on the relative positions of the emitters $\vec{r}_{ij} = \vec{r}_i - \vec{r}_j$ via the anisotropic function $g(\vec{r}_{ij})$ [70,79] and on their decay rate γ_A and $\gamma_\ell \equiv \gamma_B$ for all $\ell = 1,\ldots,N$. For significant collective enhancement of

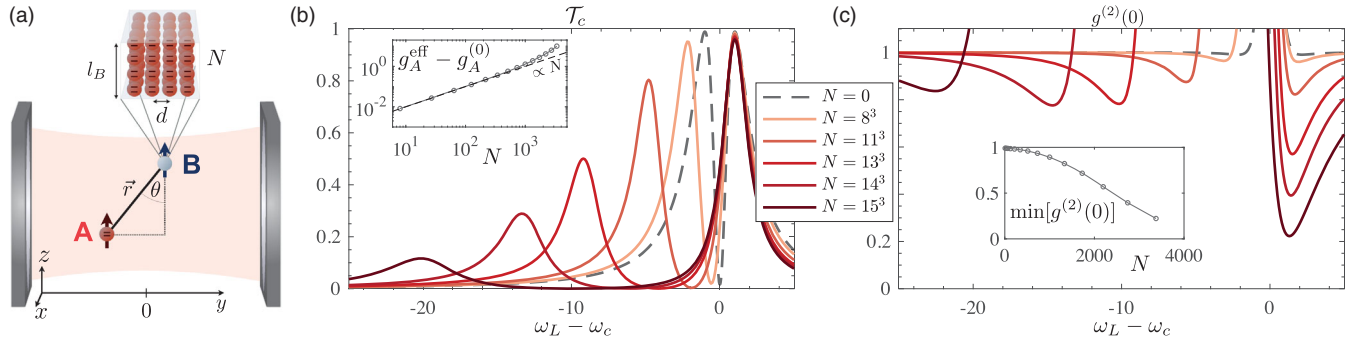

FIG. 1. (a) A single target emitter A is weakly coupled to a resonant ($\omega_A = \omega_c$) cavity with strength $g_A^{(0)}$, and close to a small ensemble B. As B we consider a cube of N emitters identical to A but with different transition frequencies $\omega_B > \omega_A$. When $l_B \ll r$, B can be approximated by a single point-dipole at distance $|\vec{r}| \equiv r$. (b) The normalized cavity transmission spectrum $\mathcal{T}_c$ probed by a weak laser field with frequency ω_L [70] shows a splitting increasing with N, which results from an enhanced effective cavity-coupling strength $g_A^{\rm eff} > g_A^{(0)}$ of the *single* target emitter A. We find $g_A^{\rm eff} - g_A^{(0)} \propto N$ (see inset and text). (c) Equal-time cavity photon-photon correlation function $g^{(2)}(0)$. The induced cavity-coupling strength occurs on a *single* emitter level, giving rise to photon blockade effects, $g^{(2)}(0) \ll 1$ and $\min[g^{(2)}(0)]$ decreasing with N (inset). [Parameters for (b) and (c): $g_A^{(0)} \equiv 1, \gamma_A = 0.01, \kappa = 2, \omega_B - \omega_A = 1000$, cavity-wavelength λ, $\vec{r}_A = (0,0,0)$, $\vec{r} = (0,0,\lambda/20)$, $d = 10^{-3}\lambda$ (lattice spacing)].

the A-cavity coupling, we require $|\vec{r}_{\ell A}| \lesssim \lambda$, so that the ensemble B is in the near field of A.

The dynamics of the system with density operator ρ is governed by a standard master equation

$$\partial_t \rho = -i[H_0 + H_{\rm JC} + H_{\rm DD}, \rho] + \mathcal{L}\rho, \tag{1}$$

where $\hbar \equiv 1$. In a frame rotating at the system A's frequency ω_A, $H_0 = \Delta_c a^\dagger a + \Delta_B \sum_\ell \sigma_\ell^+ \sigma_\ell^-$ represents the free Hamiltonian with detunings $\Delta_c = \omega_c - \omega_A$ and $\Delta_B = \omega_B - \omega_A$, where ω_c and ω_B are the bare cavity mode and ensemble emitter frequencies, respectively.

Dissipation is described by $\mathcal{L}\rho = -\kappa \mathcal{D}(a^\dagger, a)\rho - \sum_{i,j} \sqrt{\gamma_i \gamma_j} f(\vec{r}_{ij}) \mathcal{D}(\sigma_i^+, \sigma_j^-)\rho$, with superoperator $\mathcal{D}(X,Y)\rho = [X, Y\rho] + [\rho X, Y]$. The first term on the right-hand side governs photon losses at rate 2κ. Diagonal terms ($i = j$) in the second term correspond to standard spontaneous emission of the emitters. Off-diagonal ($i \neq j$) terms give rise to mutual decay into the environment depending on the separation of the emitters via the function $f(\vec{r}_{ij})$ [70,79], ranging from 0 at large separations to unity at zero separation. In this near-field limit, they are responsible for the emergence of super- and subradiant modes [80].

We are interested in obtaining an effective dynamics for the target A and the cavity, by adiabatically eliminating [70,79,81–83] the interacting ensemble of emitters B. When the latter is weakly populated, the equations of motion within the subsystem B become linear, $\partial_t \langle \sigma_\ell^- \rangle = -i \sum_m (\mathbf{M})_{\ell m} \langle \sigma_m^- \rangle$, with $(\mathbf{M})_{\ell m} = (\Delta_B - i\gamma_B)\delta_{\ell m} + (1 - \delta_{\ell m})\gamma_B [g(\vec{r}_{\ell m}) - i f(\vec{r}_{\ell m})]$ a $N \times N$ matrix [84,85]. Adiabatic elimination of B generally requires that the eigenvalues λ_η of $\mathbf{M}$ are the largest parameters (see Refs. [70,79] for details). This ensures that subsystem B evolves fast compared to the reduced system including A and the cavity. The coupling of the latter to B also has to be perturbatively small such that B remains weakly populated.

After elimination of B, the effective master equation for the reduced density operator $\rho^{\rm eff}$ for target A and cavity only reads

$$\partial_t \rho^{\rm eff} = -i[H_0^{\rm eff} + H_{\rm JC}^{\rm eff}, \rho^{\rm eff}] + \mathcal{L}^{\rm eff} \rho^{\rm eff}. \tag{2}$$

The effective free Hamiltonian $H_0^{\rm eff} = \Delta_A^{\rm eff} \sigma_A^+ \sigma_A^- + \Delta_c^{\rm eff} a^\dagger a$ contains renormalized detunings $\Delta_A^{\rm eff} = \omega_A^{\rm eff} - \omega_A = -{\rm Re}(\vec{v}^T \mathbf{M}^{-1} \vec{v})$ and $\Delta_c^{\rm eff} = \omega_c^{\rm eff} - \omega_A = \Delta_c - {\rm Re}(\vec{g}^T \mathbf{M}^{-1} \vec{g})$, which can be generally compensated. Here, $(\vec{g})_\ell = g_\ell$ and $(\vec{v})_\ell = \sqrt{\gamma_A \gamma_B}[g(\vec{r}_{\ell A}) - i f(\vec{r}_{\ell A})]$. The effective A-cavity interaction $H_{\rm JC}^{\rm eff} = g_A^{\rm eff}(a^\dagger \sigma_A^- + \sigma_A^+ a)$ has the modified coherent coupling

$$g_A^{\rm eff} = g_A - {\rm Re}(\vec{g}^T \mathbf{M}^{-1} \vec{v}). \tag{3}$$

The effective dissipator for the reduced system in Lindblad form reads $\mathcal{L}^{\rm eff} \rho^{\rm eff} = -\sum_{\nu=\pm} \gamma_\nu \mathcal{D}(L_\nu^\dagger, L_\nu)\rho^{\rm eff}$ with rates

$$\gamma_\pm = \frac{\kappa^{\rm eff} + \gamma_A^{\rm eff}}{2} \pm \sqrt{\frac{(\kappa^{\rm eff} - \gamma_A^{\rm eff})^2}{4} + \mu^2}. \tag{4}$$

The main results of this work follow from Eqs. (3) and (4). We define the system to be strongly coupled when $|g_A^{\rm eff}|$ exceeds the largest rate γ_+. The parameters $\kappa^{\rm eff} = \kappa + {\rm Im}[\vec{g}^T \mathbf{M}^{-1} \vec{g}]$ and $\gamma_A^{\rm eff} = \gamma_A + {\rm Im}[\vec{v}^T \mathbf{M}^{-1} \vec{v}]$ in Eq. (4) describe the modification of cavity decay and the target emitter linewidth, respectively, due to the interaction with the B ensemble. The rate $\mu = {\rm Im}[\vec{g}^T \mathbf{M}^{-1} \vec{v}]$ gives rise to mutual decay mechanisms, dissipatively coupling A and the cavity via Lindblad jump operators L_+ and L_- (see Refs. [70,79]).

Using the mode decomposition of $\mathbf{M}$ in terms of its eigenvectors $\vec{x}_\eta$, the effective coupling modification in Eq. (3) reads $\Delta g_A \equiv g_A^{\rm eff} - g_A = -\sum_\eta \mathrm{Re}(\vec{g}^T\vec{x}_\eta\vec{x}_\eta^T\vec{v}/\lambda_\eta)$, which depends both on the mode overlaps $\vec{g}^T\vec{x}_\eta$, $\vec{v}^T\vec{x}_\eta$, and the complex eigenvalues λ_η, and can increase with N. For example, when the sum $\sum_\eta$ is dominated by a nearly homogeneous ($l_B \ll |\vec{r}_{\ell A}| \ll \lambda$) coupling to a symmetric mode $\vec{x}_S = (1,1,\ldots,1)^T/\sqrt{N}$ of B with eigenvalue λ_S, the enhancement scales as $\sim N$ since both mode overlaps $\vec{g}^T\vec{x}_S$, $\vec{v}^T\vec{x}_S$ scale as $\sim\sqrt{N}$.

In Fig. 1(b) we show the normalized cavity transmission spectrum $\mathcal{T}_c$ for the crystalline cubic structure in Fig. 1(a), where the symmetric mode $\vec{x}_S$ is found to dominate [70]. Here, $\mathcal{T}_c(\omega_L)$ is computed numerically from Eq. (2) for a weak cavity drive (strength ϕ) at frequency ω_L, and is given by $\mathcal{T}_c(\omega_L) = \langle a^\dagger a\rangle(\kappa^2/\phi^2)$ [70]. We find a collective enhancement of the polariton splitting that scales approximately $\propto N$, with superlinear corrections for large N due to energy shifts in λ_S. We note that due to the small size of the cube ($l_B \ll |\vec{r}_{jA}| \ll \lambda$), the condition $\Delta_A^{\rm eff} \simeq \Delta_c^{\rm eff}$ is always satisfied. The polariton peaks display a marked asymmetry growing with N, which we attribute to the presence of the mutual decay rate μ [70,79].

Importantly, the collective enhancement leads to strong nonlinear effects at the single photon level. Figure 1(c) shows the equal-time photon-photon correlation function $g^{(2)}(\tau=0) = \langle a^\dagger a^\dagger aa\rangle/\langle a^\dagger a\rangle^2$, computed for the same parameters as in Fig. 1(b) and using the steady state of Eq. (2) with a weak laser drive [70]. The figure shows that $g^{(2)}(0)$ becomes $g^{(2)}(0) \lesssim 1$ at frequencies close to the polariton ones, and decreases with N. It can reach values $g^{(2)}(0) \ll 1$ for the upper polariton, reflecting the polariton asymmetry in $\mathcal{T}_c$ found above. This demonstrates that it is possible to achieve the "photon blockade" regime for a single target quantum emitter solely by increasing the number N of dipoles making up the mediating mesoscopic ensemble.

When $l_B \ll |\vec{r}_{\ell A}|$, the collective dipole of B can be approximated by a point dipole of strength $g_B^{(0)} = \sqrt{N}g_A^{(0)}$ and $\gamma_B = N\gamma_A$ positioned at $\vec{r}_B$ (each dipole moment in B is chosen identical to that of A). Analytical expressions for the parameters of Eq. (2) can then be obtained in the single-mode limit. In this case the effective coupling strength reads [70]

$$g_A^{\rm eff} = g_A^{(0)}\left\{\cos(ky_A) - \frac{\gamma_B\Delta_B g(\vec{r}) + {\gamma_B}^2 f(\vec{r})}{\Delta_B^2 + {\gamma_B}^2}\cos(ky_B)\right\}, \tag{5}$$

where $\vec{r} \equiv \vec{r}_B - \vec{r}_A$. We find an increased effective cavity decay, $\kappa^{\rm eff} = \kappa + \gamma_B g_B^2(\vec{r}_B)/(\Delta_B^2 + {\gamma_B}^2)$, as well as a modified linewidth

$$\gamma_A^{\rm eff} = \gamma_A\left\{1 - f(\vec{r})^2 + \frac{[g(\vec{r})\gamma_B - f(\vec{r})\Delta_B]^2}{\Delta_B^2 + {\gamma_B}^2}\right\}. \tag{6}$$

The mutual decay term entering Eq. (4) instead reads $\mu = g_B\sqrt{\gamma_A\gamma_B}(g(\vec{r})\gamma_B - f(\vec{r})\Delta_B)/(\Delta_B^2 + {\gamma_B}^2)$. The equations above show that the parameters in Eq. (2) are generally tuned by a combination of coherent (dispersive) [86] and incoherent (dissipative) terms [e.g., $\sim\Delta_B g(\vec{r})$ and $\sim\gamma_B f(\vec{r})$ in Eq. (5), respectively], and depend strongly on geometrical effects, due to the dipolar nature of the $A-B$ interaction. For a given geometry (i.e., positioning of A and B, and dipole orientations), $g_A^{\rm eff}$ can be significantly increased by modifying the detuning Δ_B, an effect that we term *coherent inheritance of coupling strength*. A competing effect that we dub *reduction of linewidth* in A can be similarly obtained for $\gamma_A^{\rm eff}$ in Eq. (6) (even in the absence of a cavity, e.g., by putting the third term on the rhs of the equation to zero, and for small separations where $f(\vec{r})^2 \sim 1$). Such linewidth narrowing effect can be related to the formation of a subradiant mode.

Figures 2(a) and 2(c) show contour plots of $|g_A^{\rm eff}/\gamma_+|$ as a function of the relative distance r between A and B, and the dipole-orientation angle θ [see Fig. 1(a)], for two choices of parameters where coherent inheritance of coupling strength [Fig. 2(a)] or reduction of linewidth [Fig. 2(c)] are the dominant effects. In the two panels, the bare parameters of Eq. (2) are chosen as $\gamma_A < g_A^{(0)} < \kappa$ and $\kappa < g_A^{(0)} < \gamma_A$, respectively. In Fig. 2(a) we further choose a large detuning $|\Delta_B| \gg \gamma_B$ (with $\gamma_B = 10^4\gamma_A$). This implies an effective coherent coupling $g_A^{\rm eff}/g_A^{(0)} \approx \cos(ky_A) - \{\gamma_B g(\vec{r})/\Delta_B\}\cos(ky_B)$, when the positions of A and B are such that $|g(\vec{r})| > |f(\vec{r})|$. Note that while larger effective couplings can be obtained for smaller $|\Delta_B|$, the latter has to be large enough to satisfy the conditions for adiabatic elimination. In Fig. 2(c), instead, A and B are resonant ($\Delta_B = 0$) and γ_B is large, such that the effective coherent coupling and the linewidth simplify to $g_A^{\rm eff}/g_A^{(0)} \approx \cos(ky_A) - f(\vec{r})\cos(ky_B)$ and $\gamma_A^{\rm eff} = \gamma_A[1 - f(\vec{r})^2 + g(\vec{r})^2]$, respectively. Furthermore, we position A such that $g_A^{\rm eff}/g_A^{(0)} \approx -f(\vec{r})\cos(ky_B)$ only stems from the dissipative dipole term $f(\vec{r})$. In both cases, the figure shows that a regime of strong coupling with $|g_A^{\rm eff}/\gamma_+| > 1$ between the target A and the cavity can be reached (area enclosed by the green line), e.g., in a "head-to-tail" configuration with $\theta \sim 0$ or close to the "magic angle" $\theta \sim 0.3\pi$ in Figs. 2(a) and 2(c), respectively. This is particularly surprising for Fig. 2(c), as coherent light-matter coupling is entirely due to dissipative effects.

Figures 2(b) and 2(d) confirm strong coupling with a simulation of the time evolution for the mean populations of A, B, and the cavity, starting with a single photon. The presence of B induces a clear coherent exchange of energy between the cavity and A (negligible population in B), in

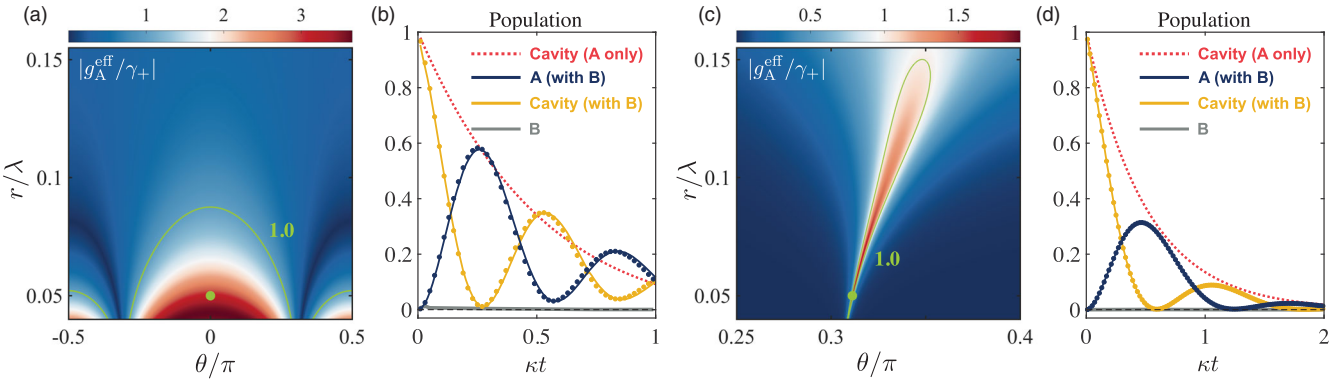

FIG. 2. (a),(c) Contour plots of $|g_A^{\mathrm{eff}}/\gamma_+|$ vs separation $\vec{r}(r,\theta)$ between A and B (point dipole, see Fig. 1). $|g_A^{\mathrm{eff}}/\gamma_+| > 1$ (inside the green contour line) indicates strong coupling. The parameters are chosen to illustrate either *coherent coupling inheritance* (a) or *dissipative linewidth narrowing* (c) for reaching strong coupling of A (see text). (b),(d) Time evolution of population in A, cavity (initial state with one photon) and B (barely visible). Oscillations exemplify strong coupling [parameters for green dots in (a) and (c), we adapt ω_c such that $\Delta_A^{\mathrm{eff}} \simeq \Delta_c^{\mathrm{eff}}$]. Lines: Full master equation simulation [Eq. (1) with two emitters], points: effective model [Eq. (2)], dashed line: cavity decay without B. [Parameters (a),(b): As in Fig. 1, but choosing $g_B^{(0)} = 100$ and thus $\gamma_B = \gamma_A(g_B^{(0)}/g_A^{(0)})^2 = 100$, A and B in the $x-z$ plane, $\vec{r} = r(\sin(\theta), 0, \cos(\theta))$. (c),(d): $\gamma_A = 2.5$, $\kappa = 0.1$, $\omega_B = \omega_A$, $g_B^{(0)} = 20$ and thus $\gamma_B = 1000$, $\vec{r}_A = (0, -1/4, 0)\lambda$, A and B in the $y-z$ plane, $\vec{r} = r(0, \sin(\theta), \cos(\theta))$].

contrast to the trivially damped behavior for the bare system. The validity of the effective master equation Eq. (2) is confirmed via a comparison to a full simulation of Eq. (1). We note that we have also demonstrated that our scheme is robust against inhomogeneous broadening and moderate positional disorder [70].

We now discuss a possible implementation of *coherent coupling inheritance* with realistic parameters considering SiV^- centers in diamond coupled to a microcavity of the kind described in Ref. [87] [Fig. 3(a)]. Two nanodiamonds separated by a distance $r = 60$ nm containing respectively a single (A) and an ensemble (B) of SiV^- defects are located above the cavity mirror. The nanodiamond with size l_B containing N dipoles B could be a fragment of a bulk diamond that was grown with a high percentage of the centers oriented along the y axis [88], which coincides with the $\langle 111 \rangle$ crystallographic direction [89].

The energy levels of SiV^- consist of electronic ground E_g and excited E_u state doublets with splittings of $\sim$50 GHz and $\sim$260 GHz, respectively [90]. At temperatures below 4 K and in the presence of a magnetic field H aligned along z, the states are Zeeman split and both A and B species can be prepared in the ground state $|1,\downarrow\rangle$ by spin-flipping optical pumping [91]. The splitting of the E_g doublet [92] then allows for microwave resolvable addressing to map the dipoles A to the state $|2,\downarrow\rangle$. A and B only couple to the optical transitions $|2\rangle \leftrightarrow |3\rangle$ and $|1\rangle \leftrightarrow |4\rangle$, respectively, [Fig. 3(b)] via photons polarized along z. A cavity mode with vacuum wavelength $\lambda = 737$ nm is resonant with the target A and detuned from B by $\Delta_B = 2\pi \times 310$ GHz.

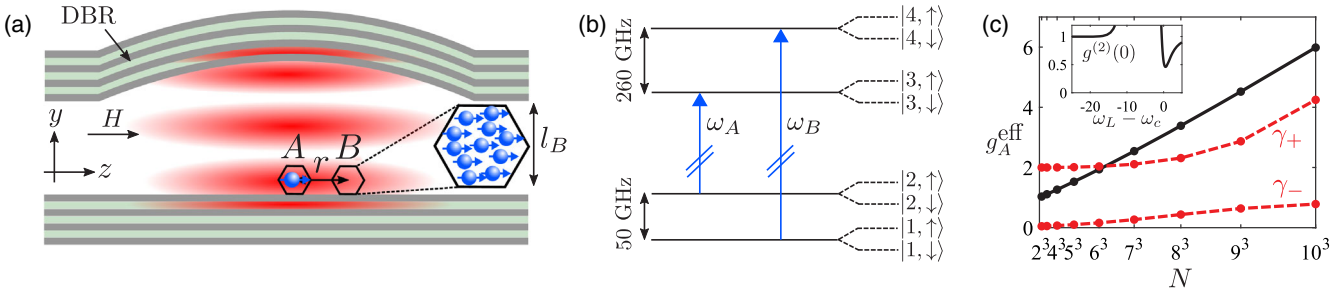

FIG. 3. Implementation with SiV^- centers coupled to a microcavity mode (red). (a) A single SiV^- center (A) is deposited above the bottom distributed Bragg reflector (DBR). A nanodiamond of size l_B (modeled as a cube with lattice spacing $d = 8$ nm $\simeq 0.01\lambda$) separated from A by a distance $\vec{r} = (0, 0, 60)$ nm contains an ensemble of N SiV^- centers (B). (b) Energy levels of a SiV^- center under a magnetic field H along z. The optical transitions $|2\rangle \to |3\rangle$ and $|1\rangle \to |4\rangle$ with frequencies ω_A (wavelength $\lambda = 737$ nm) and ω_B correspond to A and B, respectively. (c) Scaling of g_A^{eff} (black) showing that the system can be brought into strong coupling ($g_A^{\mathrm{eff}} > \gamma_+$) by increasing N. The $g^{(2)}(0)$ for $N = 1000$ is shown in the inset as a function of the detuning $\omega_L - \omega_c$. Other parameters are $g_A^{(0)} \equiv 1$, $\omega_A = \omega_c$, $\gamma_A = \gamma_B = 0.044$, $\kappa = 2$, and $\Delta_B = 265$, $\vec{r}_A = (0, 0, 0)$.

The bare couplings $g^{(0)}_{A/B}$ can be estimated by comparing our setup to that of Ref. [91], where pairs of SiV^- were coupled to a diamond photonic crystal cavity with strength $g_{\rm pcc} = 2\pi \times 7.3$ GHz and mode volume $V_{\rm pcc} = 0.5(\lambda/2.4)^3$ (2.4 is the refractive index of diamond). In our setup [Fig. 3(a)], we choose a microcavity with mode volume $V_{\rm mc} = 1.4\lambda^3$ [87], where the coupling strengths then read $g^{(0)}_{A/B} = g_{\rm pcc}\sqrt{V_{\rm pcc}/V_{\rm mc}} \approx 2\pi \times 1.17$ GHz. In Fig. 3 we model the nanodiamond B as a cube with lattice spacing $d = 8$ nm, which provides a density of emitters $\sim 2 \times 10^6 \mu{\rm m}^{-3}$ [93]. The decay rates are expected to be $\gamma_A = \gamma_B \sim 2\pi \times 51.5$ MHz [94]. We choose $\kappa = 2g^{(0)}_A$ corresponding to a moderate cavity quality factor $Q \approx 8.7 \times 10^4$ (see, e.g., Refs. [87,95,96]). This provides a finesse $F \approx 3 \times 10^4$ for a cavity of length $L = 3\lambda/2 \approx 1.1$ μm. Note that the nonradiative decay of the state $|2\rangle$ with rate $\sim 2\pi \times 13$ MHz [94] can be overcome using a repumping scheme.

The effective coupling strength $g^{\rm eff}_A$ is displayed as a function of N in Fig. 3(c), showing that the single SiV^- center A can be brought to strong coupling with the microcavity mode ($g^{\rm eff}_A > \gamma_+$) due to the presence of the nanodiamond containing the B emitters. Remarkably, this leads to strong nonlinear effects $g^{(2)}(0) \approx 0.46$ for $N = 1000$ and at a frequency close to that of the upper polariton.

In conclusion, we have investigated different mechanisms of how a *single* quantum emitter A, initially weakly coupled to a cavity mode, can be brought into strong coupling. The effect relies on both coherent and dissipative interactions with a nearby ensemble of identical emitters, and provides a viable path towards the realization of strong coupling at the level of a single SiV^- center, which has, to the best of our knowledge, never been achieved. In contrast to usual collective strong coupling, our scheme gives rise to collectively induced strong photon nonlinearities and can find important applications in quantum information technologies with SiV^- centers. Our work also emphasizes the important role played by optically active environments in cavity QED. In particular, it is an interesting prospect to explore whether our model could be extended to vibrational strong coupling in molecular ensembles [97], and to scalable microwave qubits magnetically coupled to a transmission line resonator on the verge of strong coupling [98].

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