Collective Effects in Casimir-Polder Forces

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We study cooperative phenomena in the fluctuation-induced forces between a surface and a system of neutral two-level quantum emitters prepared in a coherent collective state, showing that the total Casimir-Polder force on the emitters can be modified via their mutual correlations. Particularly, we find that a one-dimensional chain of emitters prepared in a super- or subradiant state experiences an enhanced or suppressed collective vacuum-induced force, respectively. The collective nature of dispersion forces can be understood as resulting from the interference between the different processes contributing to the surface-modified resonant dipole-dipole interaction. Such cooperative fluctuation forces depend singularly on the surface response at the resonance frequency of the emitters, thus being easily maneuverable. Our results demonstrate the potential of collective phenomena as a new tool to selectively tailor vacuum forces.

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Introduction.—Collections of atoms and solid-state quantum emitters coupled to waveguides and nanophotonic structures offer a promising platform for scalable quantum information processing [1–4]. The applications of such systems range from building long-ranged quantum networks [5,6], quantum memory devices [7–9], and metrology [10], to facilitating new experimental regimes with exotic light-matter interactions [11–13]. When interfacing small quantum systems and surfaces at nanoscales, fluctuation-induced phenomena such as vacuum forces [14], surface-modified dissipation [15,16], and decoherence [17] become important considerations. The need to achieve the control and coherence of photonic systems at that scale requires a detailed understanding of these phenomena, so as to determine the extent to which they can be tailored and controlled. In this work, we consider the possibility of using cooperative effects [18–20], as arising from having atoms or emitters in correlated quantum states, to modify fluctuation-induced forces, or Casimir-Polder (*CP*) forces [21,22].

The study of cooperative effects has a long history in the context of spontaneous emission from a collection of atoms in optical cavities and free space [23–31], and more recently near waveguides [7,32,33]. Considering that surface-modified spontaneous emission is the dissipative counterpart to the dispersive vacuum forces [34], one can expect to observe collective effects in dispersion forces as well. When considering vacuum forces, however, the role of quantum coherence within or between the interacting bodies is seldom discussed. While there have been some investigations into the effect of correlations on the van der Waals forces between two atoms [35,36], the effect of spatial wave function coherence on *CP* forces [37–39], the effect of surface modes on dipole-dipole correlations [40–44], and interference effects in vacuum forces in a three-level system [45], a general analysis of fluctuation-induced forces between an N-particle system prepared in a coherent collective state and a macroscopic body is yet to be explored in detail. The goal of this Letter is to analyze a proof of concept that illustrates cooperative effects in Casimir-Polder forces between a surface and a system of N two-level quantum emitters prepared in a Dicke state [23].

Model.—We consider a one-dimensional chain of N two-level quantum emitters with ground and excited levels $|g\rangle_n$ and $|e\rangle_n$ at a distance z_0 from the surface of a planar half-space medium, with the emitters separated by a distance x_0 from one another [see Fig. 1(a)]. We assume that the half-space $z<0$ is occupied by a medium of dielectric permittivity $\epsilon(\omega)$, while the upper half-space is vacuum. The levels $|g\rangle_n$ and $|e\rangle_n$ are coupled via an electric-dipole transition with resonance transition frequency ω_0 and spontaneous emission rate Γ_0, with $\hat{\sigma}_n^+ = (\hat{\sigma}_n^-)^\dagger = |e\rangle_n\langle g|_n$ being the ladder operators for the corresponding transition. Defining the collective spin operators $\hat{J}_k \equiv \sum_{n=1}^N \hat{\sigma}_n^k$ ($k \in \{x,y,z\}$), the Dicke states $|J,M\rangle$ correspond to [23]

$$\hat{\mathbf{J}}^2|J,M\rangle = J(J+1)|J,M\rangle \quad \text{and}$$
$$\hat{J}_z|J,M\rangle = M|J,M\rangle. \tag{1}$$

The total Hamiltonian for the system of emitters and the electromagnetic (EM) field is $\hat{H} = \hat{H}_S + \hat{H}_F + \hat{H}_{\text{int}}$, where

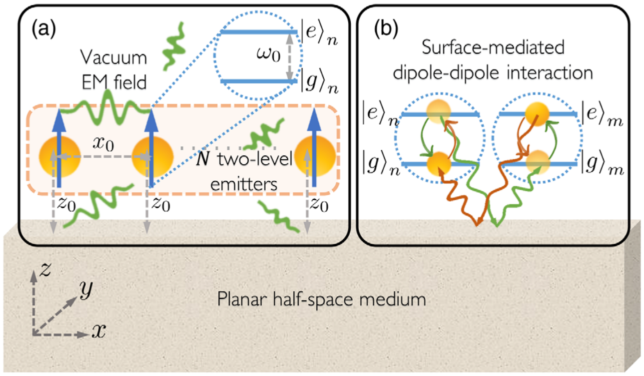

FIG. 1. (a) Schematic representation of N two-level quantum emitters prepared in a collective state, interacting with the vacuum EM field in the presence of a planar half-space medium. (b) Constructive (destructive) interference between the two processes shown in green and red leads to superradiance (subradiance) in the surface-mediated resonant dipole-dipole interactions.

$\hat{H}_S = \sum_{n=1}^{N} \hbar\omega_0 \hat{\sigma}_n^+ \hat{\sigma}_n^-$ is the Hamiltonian for the two-level emitters and $\hat{H}_F$ is the Hamiltonian for the medium-assisted EM field, which we assume to be in the vacuum state. The electric-dipole interaction Hamiltonian between the emitters and the EM field is $\hat{H}_{\rm int} = -\sum_{n=1}^{N} \hat{\mathbf{d}}_n \cdot \hat{\mathbf{E}}(\mathbf{r}_n)$, where $\hat{\mathbf{d}}_n = \mathbf{d}_n \hat{\sigma}_n^+ + \mathbf{d}_n^* \hat{\sigma}_n^-$ is the electric-dipole operator for the nth emitter and $\hat{\mathbf{E}}(\mathbf{r}_n)$ is the electric field at the position $\mathbf{r}_n$ of the nth emitter in the presence of the surface (see Ref. [46] for further details). We assume the dipole moments of all the emitters $\mathbf{d}_n \equiv d_0 \mathbf{e}_z$ to be equal in magnitude and aligned along the z direction [46].

The dynamics of the density matrix $\hat{\rho}_S$ of the emitters, after tracing out the EM field, is described by the Born-Markov master equation [48],

$$\frac{d\hat{\rho}_S}{dt} = -\frac{i}{\hbar}[\hat{H}'_S, \hat{\rho}_S] + \mathcal{L}'_S[\hat{\rho}_S], \tag{2}$$

where $\hat{H}'_S$ is the effective Hamiltonian for the emitters in the interaction picture:

$$\hat{H}'_S = \hbar\left(\sum_{n=1}^{N} \Omega_n^{(+)} \hat{\sigma}_n^+ \hat{\sigma}_n^- + \Omega_n^{(-)} \hat{\sigma}_n^- \hat{\sigma}_n^+ + \sum_{m>n} \Omega_{mn} \hat{\sigma}_m^- \hat{\sigma}_n^+\right). \tag{3}$$

Here, $\Omega_n^{(-)} = [\mu_0\omega_0/(\hbar\pi)] \int_0^\infty d\xi [\xi^2/(\xi^2+\omega_0^2)] \mathbf{d}_n^* \cdot \bar{\bar{G}}_{\rm sc}(\mathbf{r}_n, \mathbf{r}_n, i\xi) \cdot \mathbf{d}_n$ and $\Omega_n^{(+)} = -\Omega_n^{(-)} + \Omega_n^{(\rm res)}$ are the Casimir-Polder shifts for the ground and excited states of the nth emitter, respectively. These shifts correspond to processes wherein the nth dipole emits and reabsorbs a photon that is scattered off the surface, with the photon propagator given by the scattering Green's tensor $\bar{\bar{G}}_{\rm sc}(\mathbf{r}, \mathbf{r}', \omega)$ defined as the solution to the homogeneous Helmholtz equation [49,50]:

$$\nabla \times \nabla \times \bar{\bar{G}}_{\rm sc}(\mathbf{r}, \mathbf{r}', \omega) - \epsilon(\mathbf{r}, \omega)\frac{\omega^2}{c^2} \bar{\bar{G}}_{\rm sc}(\mathbf{r}, \mathbf{r}', \omega) = 0. \tag{4}$$

Here, $\epsilon(\mathbf{r}, \omega)$ is the space-dependent permittivity of the medium. Note that in addition to the broadband off-resonant contribution $\Omega_n^{(-)}$, the excited state has a resonant contribution [51],

$$\Omega_n^{(\rm res)} \equiv -\frac{\mu_0\omega_0^2}{\hbar} {\rm Re}[\mathbf{d}_n^* \cdot \bar{\bar{G}}_{\rm sc}(\mathbf{r}_n, \mathbf{r}_n, \omega_0) \cdot \mathbf{d}_n], \tag{5}$$

that depends on the response of the environment at the transition frequency ω_0 of the emitters.

The surface-modified resonant dipole-dipole interaction frequency Ω_{mn} between the emitters n and m can be expressed [52,53] as the sum of a contribution $\Omega_{mn}^{(\rm free)}$ from the resonant exchange of excitation between the two dipoles via a photon propagating in free space, and a contribution $\Omega_{mn}^{(\rm sc)}$ from a photon scattered off the surface, see Fig. 1(b), with

$$\Omega_{mn}^{(\rm sc,free)} = -\frac{\mu_0\omega_0^2}{\hbar} {\rm Re}[\mathbf{d}_m^* \cdot \bar{\bar{G}}_{\rm sc,free}(\mathbf{r}_m, \mathbf{r}_n, \omega_0) \cdot \mathbf{d}_n]. \tag{6}$$

Finally, the surface-modified Liouvillian is given by

$$\mathcal{L}'_S[\rho_S] = \sum_{m,n} \frac{\Gamma_{mn}}{2}(2\hat{\sigma}_m^- \rho_S \hat{\sigma}_n^+ - \hat{\sigma}_m^+ \hat{\sigma}_n^- \rho_S - \rho_S \hat{\sigma}_m^+ \hat{\sigma}_n^-), \tag{7}$$

where Γ_{nn} is the spontaneous emission rate for the excited state of the nth emitter, and $\Gamma_{mn} = \Gamma_{mn}^{(\rm free)} + \Gamma_{mn}^{(\rm sc)}$ is the dissipative coupling coefficient between emitters n and m, with

$$\Gamma_{mn}^{(\rm sc,free)} = \frac{2\mu_0\omega_0^2}{\hbar} {\rm Im}[\mathbf{d}_m^* \cdot \bar{\bar{G}}_{\rm sc,free}(\mathbf{r}_m, \mathbf{r}_n, \omega_0) \cdot \mathbf{d}_n]. \tag{8}$$

From Eqs. (5) and (8) we see that the dissipative coefficients $\Gamma_{nn}^{(\rm sc)}$ and $\Gamma_{mn}^{(\rm sc,free)}$ are related to the resonant dispersive shift $\Omega_n^{(\rm res)}$ and the dipole-dipole interactions $\Omega_{mn}^{(\rm sc,free)}$, respectively, being the real and imaginary parts of the same response function [58,59]. As we show below, this implies that a collective enhancement or suppression of resonant van der Waals forces is concomitant with the cooperative behavior of spontaneous emission.

Results.—We define the total *CP* force for the system of emitters in a state $\hat{\rho}_S$ as $F_{\rm CP}[\hat{\rho}_S] = -(\partial/\partial z){\rm Tr}[\hat{H}'_S \hat{\rho}_S]$, so that

$$F_{\rm CP}[\hat{\rho}_S] = -\hbar \sum_{n=1}^{N}\left[\frac{\partial}{\partial z}\Omega_n^{(+)} \langle \hat{\sigma}_n^+ \hat{\sigma}_n^- \rangle + \frac{\partial}{\partial z}\Omega_n^{(-)} \langle \hat{\sigma}_n^- \hat{\sigma}_n^+ \rangle\right] - \hbar \sum_{m>n} \frac{\partial}{\partial z}\Omega_{mn}^{(\rm sc)} \langle \hat{\sigma}_m^- \hat{\sigma}_n^+ + \hat{\sigma}_n^- \hat{\sigma}_m^+ \rangle, \tag{9}$$

where all the averages are taken over the density operator $\hat{\rho}_S$. The first term corresponds to the *CP* forces on the individual emitters and the second term to the contribution from surface-modified dipole-dipole interactions.

Focusing on that term we observe that while the operator average $(\langle\hat{\sigma}_m^-\hat{\sigma}_n^+ + \hat{\sigma}_n^-\hat{\sigma}_m^+\rangle)$ depends on the correlations between the dipoles in the state $\hat{\rho}_S$, the surface-modified dipole-dipole frequency $\Omega_{mn}^{(\mathrm{sc})}$ depends on the distance of the emitters from the surface. Hence, by preparing the emitters in a suitable collective state $\hat{\rho}_S$, the *CP* force on an ensemble can be modified. Since this modification depends only on the resonant frequency response of the surface, as evident from Eq. (6), it can thus be tailored easily by engineering surface resonances around the resonance frequency of the emitters. This is one of the main results of this Letter.

As a first illustration consider two emitters prepared near a metal surface in one of the four internal states $|\Psi_g\rangle \equiv |gg\rangle$, $|\Psi_e\rangle \equiv |ee\rangle$, $|\Psi_{\mathrm{sup}}\rangle \equiv (|eg\rangle + |ge\rangle)/\sqrt{2}$, or $|\Psi_{\mathrm{sub}}\rangle \equiv (|eg\rangle - |ge\rangle)/\sqrt{2}$. We assume the surface to be described by the Drude model with permittivity $\epsilon(\omega) = 1 - \omega_p^2/(\omega^2 + i\omega\gamma)$, where ω_p and γ are the plasma frequency and loss parameter for the metal, respectively. From Eq. (9) it follows that the force $F_{g(e)}$ for the state $|\Psi_{g(e)}\rangle$ is the sum of the forces on the individual emitters in the ground (excited) state:

$$F_{eg} = -\hbar\frac{\partial}{\partial z}[\Omega_1^{(\pm)} + \Omega_2^{(\pm)}] \approx -\frac{9\omega_p\hbar\Gamma_0 k_0}{32(\omega_p \mp \sqrt{2}\omega_0)\tilde{z}_0^4}. \quad (10)$$

Here the approximate second expression corresponds to the nonretarded, or near-field, limit of the *CP* force valid in the emitters-surface distance regime $\tilde{z}_0 \equiv k_0 z_0 \ll 1$, with $k_0 \equiv \omega_0/c$ [22,46,49].

In contrast, the force on the super- and subradiant states,

$$F_{\mathrm{sup\,sub}} = -\frac{\hbar}{2}\frac{\partial}{\partial z}\left[\Omega_1^{(\mathrm{res})} + \Omega_2^{(\mathrm{res})} \pm 2\Omega_{12}^{(\mathrm{sc})}\right], \quad (11)$$

includes a contribution that depends on the surface-mediated dipole-dipole interaction in addition to the resonant *CP* shifts of the individual emitters. In the nonretarded limit, it can be written as

$$F_{\substack{\mathrm{sup}\\ \mathrm{sub}}} \approx F_\infty[1 \pm f(\tilde{x}_0, \tilde{z}_0)], \quad (12)$$

where we have introduced the asymptotic force for infinitely separated emitters,

$$F_\infty \equiv -\frac{9\omega_p^2\hbar\Gamma_0 k_0}{16(\omega_p^2 - 2\omega_0^2)\tilde{z}_0^4}, \quad (13)$$

and

$$f(\tilde{x}_0, \tilde{z}_0) \equiv \frac{8\tilde{z}_0^4}{3}\int_0^\infty d\kappa\kappa e^{-2\kappa\tilde{z}_0}(\kappa^2+1)J_0\left(\tilde{x}_0\sqrt{\kappa^2+1}\right) \quad (14)$$

quantifies the cooperativity due to geometric configuration of the dipoles, with $\tilde{x}_0 \equiv k_0 x_0$. For coincident dipoles and to lowest order in $\tilde{z}_0$, $\lim_{x_0\to 0} f(\tilde{x}_0, \tilde{z}_0) \approx 1$.

As illustrated in Fig. 2(a), at small emitter separations $(x_0 \lesssim z_0)$ the cooperative contribution leads to an enhanced and suppressed *CP* force for the super- and subradiant state, respectively. For larger separations, $\lim_{x_0\to\infty} f(\tilde{x}_0, \tilde{z}_0) \approx 0$ and the interference effect in the resonant dipole-dipole interaction is attenuated, such that the super- and subradiant states experience an incoherent average of the ground and excited state forces, i.e., $F_{\mathrm{sup,sub}} \approx (F_g + F_e)/2 = F_\infty$. This is generally true for a state $|\Psi_{\theta,\phi}\rangle \equiv \cos\theta|eg\rangle + e^{i\phi}\sin\theta|ge\rangle$ with a single shared excitation between the emitters. We note that the total force

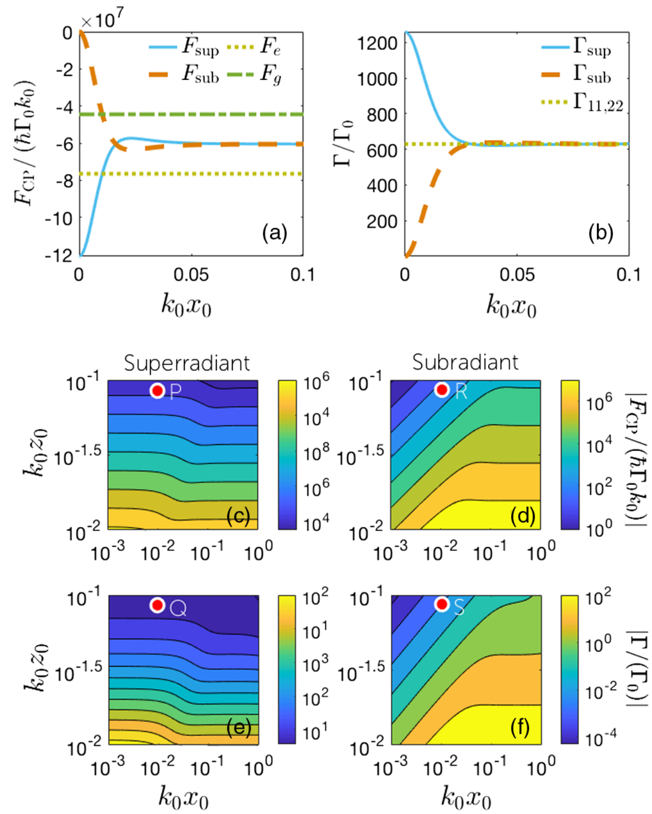

FIG. 2. (a) Collective Casimir-Polder force (in units of $\hbar\Gamma_0 k_0$) and (b) spontaneous emission (in units of Γ_0), on a system of two emitters near a gold surface, as a function of the separation between the emitters. Here the distance of the emitters from the surface is assumed to be $k_0 z_0 = 0.01$. (c) [(d)] Collective Casimir-Polder force and (e) [(f)] spontaneous emission on two emitters as a function of their distance from the surface and their mutual separation, for the dipoles prepared in the superradiant [subradiant] state $|\Psi_{\mathrm{sup}}\rangle$ [$|\Psi_{\mathrm{sub}}\rangle$]. For reference, the ground state Casimir-Polder force on the two emitters at point P is roughly $|F_g| \sim 10^4\hbar\Gamma_0 k_0$. The surface is described by the Drude model with a plasma frequency $\omega_p \approx 1.37 \times 10^{16}$ Hz (9 eV) and loss parameter $\gamma \approx 5.31 \times 10^{13}$ Hz (35 meV) for gold [60].

on the state $|\Psi_{\theta,\phi}\rangle$ is given by $F_{\theta,\phi} = -\hbar(\partial/\partial z)[\Omega_{1,2}^{(\mathrm{res})} + \Omega_{12}^{(\mathrm{sc})}\sin(2\theta)\cos\phi]$, which can vary between the super- and subradiant values in Eq. (12), depending on the relative amplitudes ($\tan\theta$) and phase ($\cos\phi$) between the states $|eg\rangle$ and $|ge\rangle$. The collective spontaneous emission for the superradiant (subradiant) state, given by $\Gamma_{\mathrm{sup}} = 1/2[\Gamma_{11} + \Gamma_{22} + 2\Gamma_{12}]$ ($\Gamma_{\mathrm{sub}} = 1/2[\Gamma_{11} + \Gamma_{22} - 2\Gamma_{12}]$) is depicted in Fig. 2(b) [46].

Figures 2(c)–2(f) give a more comprehensive picture of the collective *CP* forces and spontaneous emission as a function of the geometrical configuration of the dipoles. Assuming the emitter resonant wavelength to be $\lambda_0 \equiv 2\pi c/\omega_0 \sim 700$ nm, we see from the points R and S in Figs. 2(d) and 2(f), respectively, that a subradiant state of two emitters separated by $x_0 \sim 1$ nm, and at a distance $z_0 \sim 10$ nm from a gold surface, experiences a total force that is suppressed by a factor of $F_{\mathrm{sub}}/F_g \sim 10^{-2}$ relative to the ground state van der Waals force, with a spontaneous emission $\Gamma_{\mathrm{sub}}/\Gamma_0 \sim 10^{-2}$. Hence, subradiant *CP* forces provide a potential way to avoid both dissipation and undesirable *CP* attraction.

For a system of N dipoles the *CP* force on the Dicke superradiant state $|J = N/2, M = 0\rangle$ can be written as

$$F_{\mathrm{sup}} = -\frac{\hbar}{2}\sum_{n=1}^{N}\frac{\partial\Omega_n^{(\mathrm{res})}}{\partial z} - 2\hbar\frac{\binom{N-2}{-1+N/2}}{\binom{N}{N/2}}\sum_{m>n}\frac{\partial\Omega_{mn}^{(\mathrm{sc})}}{\partial z}, \quad (15)$$

where $\binom{N}{k}$ is a binomial coefficient. In the limit of superposed dipoles, $x_0 \to 0$, it reduces to

$$\lim_{x_0\to 0} F_{\mathrm{sup}} = -\frac{9\omega_p^2\hbar\Gamma_0 k_0}{32(\omega_p^2 - 2\omega_0^2)\tilde{z}_0^4}\left(N + \frac{N^2}{2}\right), \quad (16)$$

which demonstrates the characteristic N^2 scaling of the collective *CP* force on the superradiant state, depicted in the inset of Fig. 3, similar to free-space superradiant spontaneous emission at small emitter separations ($\tilde{x}_0 \ll 1$) [24]. We also remark that, for $N > 2$, multiple states in the degenerate subspace of subradiant Dicke states with $|J = 0, M = 0\rangle$ exhibit a suppressed *CP* force; see Ref. [46].

Discussion.—We have identified collective effects in vacuum-induced dispersion forces that result from the interference between the different channels contributing to the surface-modified resonant dipole-dipole interaction, as sketched in Fig. 1(b). Such cooperative enhancement or suppression of fluctuation forces occurs for the resonant contribution to the total *CP* force, and can be physically understood as the dispersive counterpart to super- or subradiance in spontaneous emission [see Eq. (12)]. In addition to the quantum correlations [63] in the state of the emitters this contribution to the total *CP* force depends only on the surface response at the resonance frequency of the emitters; see Eq. (6). It can thus be controlled by suitably tailoring the response of the surface around the resonant frequency of the emitters.

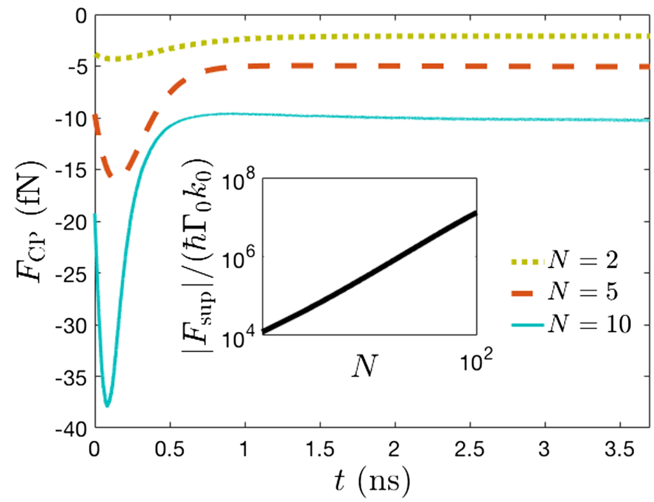

FIG. 3. Superradiant boost of the time-dependent total attractive *CP* force on a linear chain of N SiV emitters initially prepared in the excited level of the 737 nm transition with a lifetime of 1.7 ns [61], placed $z_0 \approx 10$ nm from a gold surface [47,62]. The inset shows the absolute value of the maximum boost as a function of the number of emitters, illustrating the N^2 scaling of the force for the superradiant state $|J = N/2, M = 0\rangle$.

Given that cooperative effects in optical dipole forces on solid-state emitters in nanodiamonds have been discussed both theoretically and experimentally [67–69], we suggest that it should be possible to observe a boost in the cooperative vacuum-induced forces by placing a similar nanodiamond doped with emitters near a surface. To estimate the feasibility of observing the collectively enhanced *CP* force, we consider a system of N silicon-vacancy (SiV) centers embedded in a cantilever near a metal surface [70,71]. We assume that they are initially prepared in the excited state, and solve the superradiance master equation Eq. (2) numerically [72]. As the system decays in a collective manner, it occupies the superradiant manifold transiently and experiences an enhanced *CP* force, as shown in Fig. 3. For a system of $N = 10$ SiV centers at a distance of $z_0 \approx 10$ nm from a gold surface, we find a superradiant boost in the collective *CP* force of $\Delta F_{\mathrm{CP}} \approx 20$ fN over a timescale of $\Delta\tau \approx 0.5$ ns [73,75]. The numerical results for the $N = 10$ case are from a trajectory simulation averaged over 1000 trajectories, whereas for the smaller $N \leq 6$, a direct simulation of the master equation Eq. (2) was performed. While the magnitude of the enhanced force is large enough to be observable with current technologies [76], the time resolution required to sense the enhancement would appear to pose an experimental challenge.

While we have concentrated here on the specific case of a fully inverted ensemble, a thorough analysis of different initial states will be considered in the future (see Ref. [68] for a pulsed excitation scheme used to observe superradiant emission). We also anticipate that the effects described here will be relevant and possibly observable in a variety of other platforms such as van der Waals shifts on atoms placed near optical fibers [32], quantum dot systems [77,78], and superconducting qubits [31].

In terms of potential applications one can speculate that superradiant states could be used to boost hard to observe weak fluctuation forces such as the ones discussed in Ref. [79], as well as a probe of surface properties [80]. More interestingly perhaps, given that subradiant states suppress undesirable Casimir-Polder attraction and exhibit long lifetimes with respect to the single-atom excited states [Fig. 2(f)], they can be a useful resource for trapping particles near surfaces, with potential applications in terms of quantum computation given their robustness to environmental decoherence [81,82].

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Note added.—The related works in Refs. [83,84] were brought to our attention recently.

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[53] We have also ignored here the long-range van der Waals tail contributions as discussed in Refs. [54–56]. Noting that in the distance regime of interest $a_0 \ll k_0 x_0 \ll \hbar c/\mathcal{L}$, where a_0 is the Bohr radius and $\mathcal{L}$ is the Lamb energy (~1 GHz [57]), either the van der Waals potential or the Wick-rotated contribution [55] dominates. We see that both these contributions scale roughly as $\sim [E_h (a_0 k_0)^6/(k_0 x_0)^6]$, where $E_h = \hbar c\alpha/a_0$ refers to the Hartree energy (with α as the fine-structure constant), such that they can be ignored in comparison to the resonant second-order contribution in the pertinent distance regime $k_0 x_0 \gtrsim \{[c\alpha(a_0 k_0)^6]/a_0\Gamma_0\}^{1/3}$.

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