Nonadditivity of Optical and Casimir-Polder Potentials

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(Received 28 November 2017; revised manuscript received 30 May 2018; published 22 August 2018)

An atom irradiated by an off-resonant laser field near a surface is expected to experience the sum of two fundamental potentials, the optical potential of the laser field and the Casimir-Polder potential of the surface. Here, we report a new nonadditive potential, namely, the laser-induced Casimir-Polder potential, which arises from a correlated coupling of the atom with both the laser and the quantum vacuum. We apply this result to an experimentally realizable scenario of an atomic mirror with an evanescent laser beam leaking out of a surface. We show that the nonadditive term is significant for realistic experimental parameters, transforming potential barriers into potential wells, which can be used to trap atoms near surfaces.

DOI: 10.1103/PhysRevLett.121.083603

In experiments involving an applied electromagnetic field and particles trapped in or near material objects, two forces of very different origin arise. One is the Casimir-Polder (CP) force [1] arising from the interaction between an atom and the electromagnetic vacuum field, which is restricted and modified by the presence of nearby surfaces [2]. The other is the optical force stemming from the direct interaction between the applied electromagnetic field and the atom. The latter was first applied in experiments where micron-size dielectric spheres were trapped by two laser beams [3], eventually leading to the first observation of optical trapping of atoms by a single strongly focused Gaussian laser beam [4]. Atoms can also be reflected by an evanescent laser field at a surface, as shown, for example, in Ref. [5] where state-selective reflection of Na atoms was demonstrated using an evanescent field.

The CP force has been studied extensively. It is well described by many established theoretical approaches, ranging from a quantum-mechanical linear-response formalism [6,7] to macroscopic extensions of quantum electrodynamics that incorporate material properties, see, e.g., Ref. [8]. There are also several experimental techniques to investigate this force. Almost fifty years after its theoretical prediction, the CP force was first measured by observing the deflection of atoms passing through a V-shaped cavity [9]. Later, it was demonstrated that the scattering of slowly moving atoms from a surface can also be used to deduce the CP potential [10]. Recently, approaches involving a single Rb atom optically trapped close to a surface have been employed [11]. There, the sum of the trapping potential and the CP potential determines the equilibrium position of the atom, so that for a known trapping potential the CP potential can be determined.

This technique has recently been used to make a first direct measurement of CP forces between solid surfaces and atomic gases in the transition regime between short distances (nonretarded) and long distances (retarded) [11,12]. In this experiment, ultracold ground-state Rb atoms are reflected from an evanescent wave barrier at a glass prism.

The question arises whether the two fundamentally different ingredients, the CP potential (U_{CP}) and the light-induced optical potential (U_{L}), can be simply added to obtain the total potential. Being inherently related to universal scaling laws of dispersion potentials [13], this question is of fundamental importance. In this Letter we report a new nonadditive laser-induced CP potential. We show that it plays an important role under specific experimental conditions.

The CP potential can be viewed as the modification of a fluctuating dipole moment due to the electric field those fluctuations induce. It is clear that this is a recursive process; the dipole induces a field, which in turn interacts with the dipole, changing the field it produces, which then interacts again with the dipole and so on. Normally, only the first step in this process is considered, where the fluctuations of the dipole moment are taken to come from the field that would be present if the dipole were not there. This electric field could come from the vacuum field or a laser field depending on the specific context. This leads to the usual approach taken when considering atoms subject to vacuum and laser fields; the two potentials are simply added (cf. Ref. [14]). Here we consider and take into account the effect of the laser field on the CP potential.

Quantum vacuum forces between an atom in its ground state and a surface are attractive and in general not suitable to create a stable position for the atom [15]. Nevertheless, Ref. [16] describes the possibility of trapping ground-state atoms by dressing them with an excited state whose potential is repulsive in a laser field. This method is similar in spirit to our work, especially because the laser-induced

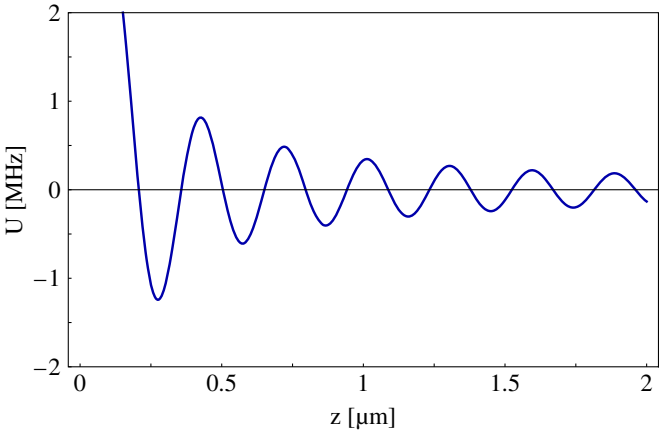

FIG. 1. Difference $\Delta U = U_{\mathrm{CP}} - U_{\mathrm{LCP}}$ for a perfectly conducting mirror with reflective coefficients $r_s = -1$ and $r_p = 1$. The laser intensity is $I = 5\ \mathrm{W/cm^2}$ used for Rb atoms with a transition frequency of $\tilde{\omega}_{10} = 2.37 \times 10^{15}$ rad/s. The detuning is $\Delta = 2\pi \times 10^8$ rad/s. The angle between the z axis and the orientation of the field $\mathbf{E}(\mathbf{r}_A)$ is $\theta = \pi/2$.

CP potential for the atomic ground state resembles the excited-state CP potential (cf. Fig. 1), but has a major difference in its assumption of a position-independent Rabi frequency.

We describe the system using a Hamiltonian that governs the coupling of the atom to the electric field and consequently consists of a field part, an atomic part, and a multipolar dipole-field coupling. The field Hamiltonian $\hat{H}_F$ can be expressed as

$$\hat{H}_F = \sum_{\lambda=\mathrm{e,m}} \int d^3r \int_0^\infty d\omega \hbar\omega \hat{\mathbf{f}}_\lambda^\dagger(\mathbf{r},\omega) \cdot \hat{\mathbf{f}}_\lambda(\mathbf{r},\omega), \quad (1)$$

where $\hat{\mathbf{f}}_\lambda$ and $\hat{\mathbf{f}}_\lambda^\dagger$ are creation and annihilation operators for composite field-matter excitations. Here we model a driving laser as being a result of a source occupying a volume V_S in space. This leads us to represent field states as a product

$$|\psi\rangle_F = |\{\mathbf{f}_\lambda(\mathbf{r},\omega)\}\rangle_{\mathbf{r}\in V_S} \otimes |\{0\}\rangle_{\mathbf{r}\notin V_S} \quad (2)$$

of coherent excitations $|\{\mathbf{f}_\lambda(\mathbf{r},\omega)\}\rangle$ in the source region and the vacuum state $|\{0\}\rangle$ for all other regions. If the annihilation operator $\hat{\mathbf{f}}_\lambda(\mathbf{r},\omega)$ acts on the state (2), there are, consequently, two contributions

$$\hat{\mathbf{f}}_\lambda(\mathbf{r},\omega)|\psi\rangle_F = \begin{cases} \mathbf{f}_\lambda(\mathbf{r},\omega)|\psi\rangle_F & \text{if } \mathbf{r} \in V_S \\ 0 & \text{if } \mathbf{r} \notin V_S \end{cases}. \quad (3)$$

The atom-field Hamiltonian $\hat{H}_{\mathrm{AF}} = -\hat{\mathbf{d}} \cdot \hat{\mathbf{E}}(\mathbf{r}_A)$ in the multipolar coupling scheme is determined by the electric field at the atom's position $\mathbf{r}_A$ and the dipole operator $\hat{\mathbf{d}}$. The electric field is given by the respective classical Green's tensor $\mathbf{G}(\mathbf{r},\mathbf{r}_A,\omega)$ and the field operator $\hat{\mathbf{f}}_\lambda(\mathbf{r},\omega)$. Solving the Heisenberg equation of motion for the field operator using the field Hamiltonian $\hat{H}_F$ (1) and the coupling Hamiltonian $\hat{H}_{\mathrm{AF}}$ and inserting this back into the electric field, the final expression for the time-dependent electric field operator yields

$$\begin{aligned} \hat{\mathbf{E}}(\mathbf{r},\omega,t) &= \hat{\mathbf{E}}_{\mathrm{free}}(\mathbf{r},\omega,t) + \hat{\mathbf{E}}_{\mathrm{ind}}(\mathbf{r},\omega) \\ &= \hat{\mathbf{E}}(\mathbf{r},\omega)e^{-i\omega(t-t_0)} \\ &\quad + \frac{i\mu_0}{\pi}\omega^2 \int_{t_0}^{t} dt' e^{-i\omega(t-t')} \mathrm{Im}\mathbf{G}(\mathbf{r},\mathbf{r}_A,\omega) \cdot \hat{\mathbf{d}}(t'), \end{aligned} \quad (4)$$

where μ_0 is the permeability of free space. The induced contribution represents the inhomogeneous part of the solution and couples the Green's tensor to the atomic dipole moment as shown in Ref. [17]. The state (3) can be inserted into Eq. (4) where the free component is modeled as a classical laser driving field of frequency ω_L at the atom's position, $\mathbf{E}(\mathbf{r}_A,t) = \mathbf{E}(\mathbf{r}_A)\cos(\omega_L t)$. In a similar way, one can compute the Heisenberg equation of motion for the atomic flip operator $\hat{A}_{mn}(t)$ [18] defined in such a way that the atomic part of the Hamiltonian is $\hat{H}_{\mathrm{A}} = \sum_n E_n \hat{A}_{nn}$.

We first assume that the atom-field coupling is weak and that the field is far-detuned from the atomic resonance. This assumption is later relaxed to obtain a more general result, as described in Ref. [19]. The electric field (4) is evaluated using the Markov approximation for weak atom-field coupling and we discard slow nonoscillatory dynamics of the flip operator by setting $\hat{A}_{mn}(t') \simeq e^{i\tilde{\omega}_{mn}(t'-t)}\hat{A}_{mn}(t)$ for the time interval $t_0 \le t' \le t$. To apply the Markov approximation we have assumed that the atomic transition frequency $\tilde{\omega}_{10}$ is not close to any narrowband resonance mode of the medium. If there were such a mode, the atom would mostly interact with it, similar to a cavity. In this case the mode could be modeled by a Lorentzian profile [18,20,21].

The parameters entering the dynamics are the shifted frequency $\tilde{\omega}_{mn} = \omega_{mn} + \delta\omega_{mn}$, where ω_{mn} is the atom's pure eigenfrequency and the CP frequency shift $\delta\omega_{mn}$ due to the presence of the surface, and the rate of spontaneous emission Γ_{mn}. The fast-oscillating nondiagonal parts $\hat{A}_{mn}(t)$ can be decoupled from the slowly oscillating diagonal operator terms $\hat{A}_{mm}(t)$ by assuming that the atom does not have quasidegenerate transitions. Moreover the atom is unpolarized in each of its energy eigenstates so that $\mathbf{d}_{mm} = \mathbf{0}$, which is guaranteed by atomic selection rules [18]. Finally, we assume the atom stays in its initial state with $\langle\hat{A}_{kl}(t')\rangle \approx \langle\hat{A}_{kl}(t)\rangle \approx \delta_{kn}\delta_{ln}$. Consequently, to compute the dipole moment we only need the nondiagonal elements of the atomic flip operator

$$\langle \dot{\hat{A}}_{mn}(t)\rangle = \mathrm{i}\tilde{\omega}_{mn}\langle \hat{A}_{mn}(t)\rangle - \frac{1}{2}[\Gamma_n + \Gamma_m]\langle \hat{A}_{mn}(t)\rangle + \frac{\mathrm{i}}{\hbar}\sum_k [\langle \hat{A}_{mk}(t)\rangle \mathbf{d}_{nk} - \langle \hat{A}_{kn}(t)\rangle \mathbf{d}_{\mathrm{km}}]\cdot \mathbf{E}(\mathbf{r}_\mathrm{A}, t). \tag{5}$$

Similar to the electric field (4), the equation for the time evolution of the dipole operator $\hat{\mathbf{d}}(t)$ can be split into free and induced (electric field dependent) parts $\hat{\mathbf{d}}_{\mathrm{free}}(t)$ and $\hat{\mathbf{d}}_{\mathrm{ind}}(t)$, given, respectively, by the first and second lines of Eq. (5).

In the Markov approximation, we use the expression of the complex atomic polarizability for an atom in a spherically symmetric state with negligible damping

$$\alpha_n(\omega) = \frac{2}{3\hbar}\sum_k \frac{\tilde{\omega}_{kn}|\mathbf{d}_{nk}|^2}{\tilde{\omega}_{kn}^2 - \omega^2}\mathbf{1}, \tag{6}$$

where $\mathbf{1}$ is the unit matrix. The dipole moment in the time domain reads

$$\langle \hat{\mathbf{d}}_{\mathrm{ind}}(t)\rangle_n = \frac{1}{2}[\alpha_n(\omega_L)\mathbf{E}(\mathbf{r}_A)e^{-i\omega_L t} + \mathrm{H.c.}], \tag{7}$$

which oscillates with the laser frequency ω_L. An equivalent expression for the induced electric field $\hat{\mathbf{E}}_{\mathrm{ind}}(\mathbf{r}, t)$ is obtained in a similar way.

The quantity we are interested in is the total potential $U = -\frac{1}{2}\langle \hat{\mathbf{d}}(t)\cdot\hat{\mathbf{E}}(\mathbf{r}_A, t)\rangle$, which consists of various contributions as shown in Table I. The order of the terms is determined by the number of dipole moments. At leading order, the induced part of the field (dipole moment) depends on the free part of the dipole moment (field), given by the first iteration of Eq. (4). This leads to two well-known potentials, namely, the laser-light potential and the standard, undriven CP potential. Going one iteration further yields an additional contribution containing four dipole moments in total, the CP potential under the influence of the driving laser field. This contribution is nonadditive; i.e., the total potential experienced by the atom can no longer be obtained simply by summing the Casimir-Polder potential and the laser-induced potential. We will show that this nonadditive contribution can be significant under certain circumstances.

TABLE I. Summary of contributions to the total potential U. The contributions to lowest order (yellow) form the ordinary CP potential U_{CP} and the laser-light potential U_{L}. Higher-order terms (green) build up the nonadditive laser-induced CP potential U_{LCP}, with certain terms vanishing as explained in the main text.

	$\hat{\mathbf{d}}^{(1)}_{\mathrm{free}}$	$\hat{\mathbf{d}}^{(2)}_{\mathrm{ind}}$	$\hat{\mathbf{d}}^{(3)}_{\mathrm{ind}}$	$\hat{\mathbf{d}}^{(4)}_{\mathrm{ind}}$	...
$\hat{\mathbf{E}}^{(0)}_{\mathrm{free}}$	0	U_{L}	0	$\frac{1}{2}U_{\mathrm{LCP}}$	
$\hat{\mathbf{E}}^{(1)}_{\mathrm{ind}}$	U_{CP}	0	$U^{(2)}_{\mathrm{CP}}$		
$\hat{\mathbf{E}}^{(2)}_{\mathrm{ind}}$	0	$\frac{1}{2}U_{\mathrm{LCP}}$			
$\hat{\mathbf{E}}^{(3)}_{\mathrm{ind}}$	$U^{(2)}_{\mathrm{CP}}$				
⋮					

The term of lowest order $\langle \hat{\mathbf{d}}_{\mathrm{free}}(t)\cdot\hat{\mathbf{E}}_{\mathrm{free}}(\mathbf{r}_\mathrm{A}, t)\rangle$ contains the free dipole moment and the free electric field. For $\mathbf{r}' \in V_\mathrm{S}$, this expression leads to the vanishing expectation value of the free dipole moment $\langle \hat{\mathbf{d}}_{\mathrm{free}}(t)\rangle = 0$. In the case of $\mathbf{r}' \notin V_\mathrm{S}$, this term vanishes as well according to Eqs. (2) and (3). The occurrence of a free dipole moment is the reason for the terms of odd order to vanish. The standard, undriven CP potential is obtained from the term $\langle \hat{\mathbf{d}}_{\mathrm{free}}(t)\cdot\hat{\mathbf{E}}_{\mathrm{ind}}(\mathbf{r}_\mathrm{A}, t)\rangle$ [22]

$$U_{\mathrm{CP}}(\mathbf{r}_A) = \frac{\hbar\mu_0}{2\pi}\int_0^\infty d\xi\xi^2\alpha_n(i\xi)\mathrm{Tr}\mathbf{G}^{(S)}(\mathbf{r}_A, \mathbf{r}_A, \mathrm{i}\xi) - \frac{\mu_0}{3}\sum_{k<n}\tilde{\omega}_{nk}^2|\mathbf{d}_{nk}|^2\mathrm{Tr}[\mathrm{Re}\mathbf{G}^{(\mathrm{S})}(\mathbf{r}_A, \mathbf{r}_A, \tilde{\omega}_{nk})]. \tag{8}$$

$\mathbf{G}^{(S)}(\mathbf{r}_\mathrm{A}, \mathbf{r}_A, \tilde{\omega}_{nk})$ is the scattering part of the classical Green's tensor $\mathbf{G}(\mathbf{r}_A, \mathbf{r}_A, \tilde{\omega}_{nk})$. The CP potential can be split into resonant and off-resonant contributions, where the ground state only shows the latter. The third term of the total potential U results in the laser-light potential $\langle \hat{\mathbf{d}}_{\mathrm{ind}}(t)\cdot\hat{\mathbf{E}}_{\mathrm{free}}(\mathbf{r}_A, t)\rangle = \langle \hat{\mathbf{d}}_{\mathrm{ind}}(t)\rangle\cdot\mathbf{E}(\mathbf{r}_A, t)$ from the coherent time-averaged electric field ($\mathbf{r}' \in V_S$),

$$U_{\mathrm{L}}(\mathbf{r}_A) = -\frac{1}{4}\alpha_n(\omega_L)\mathbf{E}^2(\mathbf{r}_A), \tag{9}$$

and the fourth term of the total potential $\langle \hat{\mathbf{d}}_{\mathrm{ind}}(t)\cdot\hat{\mathbf{E}}_{\mathrm{ind}}(\mathbf{r}_A, t)\rangle$ vanishes again both for $\mathbf{r}' \in V_S$ and $\mathbf{r}' \notin V_S$.

We are interested in a higher-order iteration, where the induced part of the dipole moment itself depends on the induced part of the electric field, while the induced electric field itself contains the induced dipole moment. Combining the two high-order perturbative expressions for $\hat{\mathbf{d}}$ and $\hat{\mathbf{E}}$ leads to the final result for the driven CP potential

$$U_{\mathrm{LCP}}(\mathbf{r}_A) = -\frac{1}{2}\langle \hat{\mathbf{d}}_{\mathrm{ind}}(t)\cdot\hat{\mathbf{E}}^{(2)}_{\mathrm{ind}}(\mathbf{r}_A, t)\rangle - \frac{1}{2}\langle \hat{\mathbf{d}}^{(2)}_{\mathrm{ind}}(t)\cdot\hat{\mathbf{E}}_{\mathrm{free}}(\mathbf{r}_A, t)\rangle = -\frac{\mu_0\omega_L^2}{2}\alpha_n^2(\omega_L)\mathbf{E}(\mathbf{r}_A)\cdot\mathrm{Re}\mathbf{G}^{(\mathrm{S})}(\mathbf{r}_A, \mathbf{r}_A, \omega_L)\cdot\mathbf{E}(\mathbf{r}_A). \tag{10}$$

This expression contains induced dipole moments and induced electric fields, both of second order. The CP force corresponding to the potential (10) is computed by taking

the gradient of the potential $\mathbf{F}_{\mathrm{LCP}}(\mathbf{r}_A) = -\nabla_A U_{\mathrm{LCP}}(\mathbf{r}_A)$ and can be expressed using the two contributions $\nabla\langle\hat{\mathbf{d}}^{(2)}_{\mathrm{ind}} \cdot \hat{\mathbf{E}}_{\mathrm{free}}(\mathbf{r})\rangle_{\mathbf{r}=\mathbf{r}_A} + \nabla\langle\hat{\mathbf{d}}_{\mathrm{ind}} \cdot \hat{\mathbf{E}}^{(2)}_{\mathrm{ind}}(\mathbf{r})\rangle_{\mathbf{r}=\mathbf{r}_A}$, where one can use the relation $\nabla\mathbf{E}(\mathbf{r}_A) \cdot \mathbf{E}(\mathbf{r})|_{\mathbf{r}=\mathbf{r}_A} = \frac{1}{2}\nabla_A \mathbf{E}^2(\mathbf{r}_A)$ and the symmetry of the Green's tensor $\nabla\mathbf{G}^{(S)}(\mathbf{r}, \mathbf{r}_A)|_{\mathbf{r}=\mathbf{r}_A} = \frac{1}{2}\nabla_A \mathbf{G}^{(S)}(\mathbf{r}_A, \mathbf{r}_A)$.

Figure 1 shows the difference between the laser-induced CP potential and the standard CP potential U_{CP} for a perfectly conducting mirror, whose Green's tensor is well known. The laser-induced potential for a constant laser potential is similar to the CP potential of an excited atom. The strongest addition effect of U_{LCP} is seen in the nonretarded regime, which we investigate further for a more realistic evanescent laser field.

For simplicity we first consider a two-level, isotropically polarized atom a distance z away from a dielectric half-space described by s (transverse electric: TE) and p

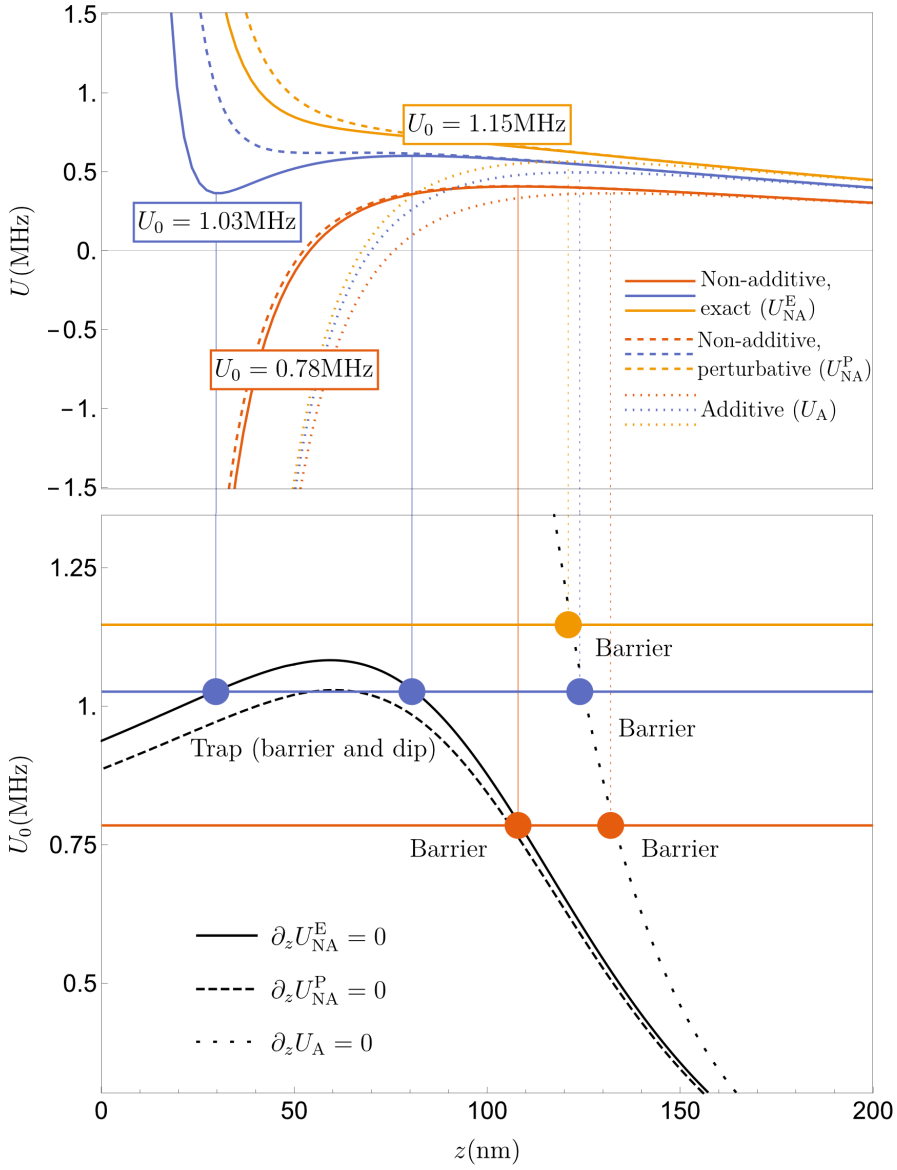

FIG. 2. Upper figure: Total potential with (solid lines) and without (dotted lines) the nonadditive term calculated here, for three different values of the potential scale U_0. The dashed lines are the total potential using the perturbative result, as opposed to the nonperturbative one. The parameters taken from Refs. [11,12] are as follows: decay length $z_0 = 430$ nm, detuning $\Delta = 2\pi \times 10^8$ Hz, dipole moment $d = 2.53 \times 10^{-29}$ Cm, transition frequency $\tilde{\omega}_{10} = 2.37 \times 10^{15}$ rad/s and quality factor $Q = 37.5 = -\mathrm{Re}(r_p)$. As reported in Ref. [11] the potential scale is given by $U_0 \sim P(\mathrm{W}) \times 10^{-23}$ J/W, meaning the range of values chosen for this quantity are commensurate with milliwatt laser power. The lower figure shows the positions of the minima and maxima of the additive potential U_A (dotted), the perturbative potential U^P_{NA} (dashed) and exact nonadditive potential U^E_{NA} (solid) as a function of the potential scale U_0 and the distance. This demonstrates that the nonadditive potential undergoes a transition from a traplike to barrierlike as a function of the atom-surface distance, while the additive potential remains barrierlike for all parameters shown here.

(transverse magnetic: TM) polarized reflection coefficients r_s and r_p. We display our formulas in the nonretarded (small z) limit, but all plots are produced including the effects of retardation. We find,

$$U_{\mathrm{LCP}} = -\frac{|\mathbf{d}_{nk}|^2}{192\pi\epsilon_0^2 z^3}\frac{\Omega^2}{\Delta^2}\mathrm{Re}r_p = U_{\mathrm{CP}}U_{\mathrm{L}}\frac{\mathrm{Re}r_p}{\hbar\Delta}, \qquad (11)$$

where we have defined $\Omega \equiv \mathbf{E}(\mathbf{r}_{\mathrm{A}})\cdot\mathbf{d}_{nk}/\hbar$ and $\Delta = \omega_L - \tilde{\omega}_{10}$, and on the right-hand side we have written the result in a suggestive form. This tells us that the laser-assisted Casimir-Polder potential is given by the product of the laser potential and the Casimir-Polder potential, modulated by the detuning Δ and the reflection coefficient r_p.

So far we have used a perturbative approach that is valid for a large detuning Δ. As we shall see later, this is not reliably satisfied for the particular application to the evanescent laser field, since the detuning and Rabi frequency Ω turn out to be of the same order of magnitude. Repeating the calculation by solving the optical Bloch equations, we find a more general nonperturbative result (see Ref. [19] for details);

$$U_{\mathrm{LCP}}^{\mathrm{NP}} = -\frac{|\mathbf{d}_{nk}|^2}{192\pi\epsilon_0^2 z^3}\frac{\Omega^2}{\Delta^2+\Omega^2}\mathrm{Re}r_p, \qquad (12)$$

which is seen to reduce to the perturbative result (11) in the large-detuning limit. The potential arising from an evanescent laser field can be written as $U_{\mathrm{L}} = U_0 e^{-2z/z_0}$, where U_0 is a constant that sets the scale of the interaction, and z_0 is its range. The electric field that is associated with this potential satisfies $E_L^2 = 6\hbar\Delta U_{\mathrm{L}}/d^2$, which gives the Rabi frequency as $\Omega_L^2 = 6\Delta U_{\mathrm{L}}/\hbar$.

The final ingredient we require is the electromagnetic response of the surface, described by a permittivity ε, which enters into Eq. (12) through the nonretarded limit of the p-polarized reflection coefficient: $r_{\mathrm{p}} = (\varepsilon-1)/(\varepsilon+1)$. By taking the Drude-Lorentz model as a basis, the reflection coefficient becomes $r_p \approx \omega_S/2\gamma$ with the plasmon resonance frequency $\omega_S = \sqrt{\omega_0^2 + \omega_P^2/2}$, which in turn depends on the absorption line frequency ω_0, the plasma frequency ω_P and the damping constant γ, cf. Ref. [16]. It is important to note that the real part of the reflection coefficient has both positive and negative wings around its resonance frequency, so that when viewed as a quality factor $\mathrm{Re}(r_p) = \pm Q$. As shown in Refs. [23–26], surface plasmon resonances, connected in the context of CP potentials, can be used to produce a maximum Purcell enhancement factor (or quality factor Q) of up to 60.

In order to produce concrete predictions, we consider the setup of Refs. [11,12], where an evanescent wave is created close to a surface creating a repulsive dipole potential into which a rubidium atom is placed. In order to assess the contribution of our nonadditive term U_{LCP} we compare the "additive" potential $U_A = U_{\mathrm{CP}} + U_{\mathrm{L}}$ to that with the exact nonadditive term included: $U_{\mathrm{NA}}^E = U_{\mathrm{CP}} + U_{\mathrm{L}} + U_{\mathrm{LCP}}^{\mathrm{NP}}$. Our results are shown in Fig. 2, and demonstrate that the nonadditive term calculated here has a drastic effect on even the qualitative character of the potential. The perturbative version $U_{\mathrm{NA}}^P = U_{\mathrm{CP}} + U_{\mathrm{L}} + U_{\mathrm{LCP}}$ of the laser-assisted Casimir-Polder potential is shown for reference, it is seen that the qualitative character of the effect is captured in the perturbative regime, but precise predictions require the nonperturbative result.

In this Letter we have derived and theoretically evaluated a nonadditive laser-induced CP potential. The electric field and the dipole moment were each split into a free contribution and an induced contribution, each of which depends on the other. In this way the laser light potential and the standard CP potential are reproduced as lowest-order terms. The higher-order correction term leads to the nonadditive potential which we have derived in both a perturbative and nonperturbative approach. We have shown that this term makes a significant contribution under certain experimental conditions. If the laser power is high enough and in combination with an additional enhancement by a surface plasmon resonance, the occurrence and position of barriers and minima in the total potential can significantly change, leading to local minima, which can be used to trap atoms near surfaces.

We acknowledge helpful discussions with Diego Dalvit, Francesco Intravaia, and Ian Walmsley. This work was supported by the German Research Foundation (DFG, Grants No. BU 1803/3-1 and No. GRK 2079/1). S. Y. B is grateful for support by the Freiburg Institute of Advanced Studies.

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