

Modifying atom-surface interactions with optical fields

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The ability to control matter on the nanometer scale is greatly influenced by the van der Waals (vdW) interaction. Therefore, understanding and manipulating the vdW interaction is of interest to the fields of nanotechnology and atom optics. We show that near-resonant light can significantly modify atom-surface vdW interactions in the nonretarded regime. A theory based on quantized electromagnetic fields is used to calculate (1) the ordinary vdW interaction, (2) corrections to the ordinary vdW interaction due to thermal radiation, and (3) modifications to the ordinary vdW interaction that result from monochromatic (laser) radiation. Near-resonant laser light with an intensity of 5 W/cm^2 is predicted to double the vdW interaction strength for sodium atoms, and possible experiments to detect this effect are discussed.

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I. INTRODUCTION

The strength of the van der Waals (vdW) interaction between matter is usually determined by the atomic or bulk material properties of a system. Therefore, the vdW force can typically only be altered by changing the atomic constituents of the matter involved in the interaction, limiting the ability to control nanometer scale forces. This circumstance resonates with the point of view in which the vdW interaction can be thought of as being caused by the fluctuating dipole moment of a quantum-mechanical atom [1]. These fluctuating dipoles can then interact with each other, leading to the vdW force between the atoms. However, there exists an alternative perspective in which the origin of the vdW interaction is due to electromagnetic-field fluctuations in the vacuum, which in turn induce a fluctuating dipole moment in the atoms [2]. This suggests that the vdW interaction can be affected by the radiation environment that the matter resides in, opening the exciting possibility of using light to control or even inhibit the vdW interaction.

One way to change the radiation environment of matter is through temperature, which exposes the atoms to a thermal electromagnetic field. In principle, this should modify the vdW or Casimir-Polder interaction. Just such an effect has recently been observed for an atom and a nearby surface in the retarded regime [3], where the atom-surface distance z is much larger than the principle transition wavelength λ of the atom. However, thermal modifications are predicted to be negligible in the nonretarded regime [4,5], where distances are such that $z \ll \lambda$. This is unfortunate since the ability to change the force between atoms and a surface would have an impact on atom chips, atom optics, and quantum reflection experiments.

Milonni and Smith [6] have predicted that the vdW interaction between two atoms can be significantly changed with only a few W/cm^2 of near-resonant laser radiation. A similar field-modified particle-particle interaction called optical binding has been observed for micron scale dielectric spheres [7] and has even been used to organize the spheres into crystalline lattices [8]. However, to our knowledge a field modi-

fied atom-surface interaction in the nonretarded regime has not been observed or proposed. In this paper, a theory for a field-modified atom-surface vdW interaction in the nonretarded regime will be described, along with a survey of possible experiments which must satisfy several critical requirements to observe such an effect.

This paper begins by introducing some familiar calculations for the strength of vdW interactions in the nonretarded regime, along with an estimate for when field-modified vdW interactions will become important. A model in which the atom and field are treated classically will then be described, and used to calculate the optical field intensity required to double the vdW interaction. Next, the electromagnetic field will be treated quantum mechanically, leading to an expression which can be used to compare the vacuum, thermal, and field-modified contributions to the atom-surface vdW interaction. Finally, a discussion of the theoretical results and candidate experiments will be presented.

II. BACKGROUND

As background knowledge, we first summarize that the Lifshitz formula [9,10] gives the nonretarded potential for atoms near a surface at zero temperature as

$$U^{(\text{vdW},0)} = -\frac{C_3^{(0)}}{z^3}, \quad z \ll \lambda, \quad (1)$$

with

$$C_3^{(0)} = \frac{\hbar}{4\pi} \int_0^\infty \frac{\alpha(i\omega) \epsilon(i\omega) - 1}{4\pi\epsilon_0 \epsilon(i\omega) + 1} d\omega, \quad (2)$$

where $\alpha(i\omega)$ is the polarizability of the atom and $\epsilon(i\omega)$ is the permittivity of the surface [11,12]. For a perfect conductor ($\epsilon = \infty$) and for an undamped atom modeled as a Lorentz oscillator [$\alpha(i\omega) = \alpha_{\text{DC}} / [1 + (\omega/\omega_0)^2]$], the Lifshitz formula gives

$$C_3^{(0)} = \frac{\alpha_{\text{DC}} \hbar \omega_0}{4\pi\epsilon_0 8}, \quad (3)$$

where $\alpha_{\text{DC}} = e^2/m\omega_0^2$ is the static atomic polarizability, $\alpha_{\text{DC}}(4\pi\epsilon_0)^{-1}$ is in SI units of volume, and ω_0 is the principal transition frequency of the atom [13]. We will refer to this ordinary result as $C_3^{(0)}$ since it is derived using quantum mechanics in the absence of any applied fields.

Comparing the different ways to derive $C_3^{(0)}$ leads to some important insight into the physical mechanisms behind the vdW interaction. To this end one can show that Eq. (3) is equivalent to the Lennard-Jones result, which is often written in terms of the atomic dipole moment p as $C_3 = \langle p^2 \rangle / 12(4\pi\epsilon_0)$, by utilizing the fact that $\alpha_{\text{DC}} = 2\langle p^2 \rangle / (3\hbar\omega_0)$ for a two-level atom [13–15]. Since the vdW coefficient in this expression depends on the fluctuating dipole moment $\langle p^2 \rangle$, this result is reminiscent of a source field picture [1]. Later on in this paper we will show that the vdW coefficient can also be expressed in terms of $C_3 \sim \sum_k \hbar \omega_k / 2$, where $\hbar \omega_k / 2$ is associated with the zero-point energy of a mode k of the electromagnetic vacuum, which is essentially a vacuum field picture [2]. It is quite surprising, and perhaps most important, that both pictures yield the same physical consequence: the vdW interaction. Both of these viewpoints are valid from the perspective of quantum electrodynamics and the choice of which one to subscribe to is largely a matter of taste or convenience. In fact, it has been shown that by changing the ordering of atomic and field operators between symmetric or normal, one can emphasize the role of the vacuum field or source field [16,17].

It is also instructive to recall that $p = \alpha(\omega)E$, which allows us to estimate what intensity of monochromatic radiation would be required to have a comparable influence on an atom as the vdW interaction. By comparing the Lennard-Jones result to Eq. (3) we can calculate the effective size of the fluctuating dipole $\langle p^2 \rangle = 3\alpha_{\text{DC}}\hbar\omega_0/2 = \alpha(\omega)^2 \langle E^2 \rangle$ induced by the vacuum. After substituting $I \equiv \epsilon_0 c E^2 / 2$ for E^2 (as defined in Ref. [18]), we find that an optical field far below resonance would need to have an intensity of $I = 3\epsilon_0 c \hbar \omega_0 / 2 \alpha_{\text{DC}} \sim 10^{14}$ W/cm² in order to induce a dipole moment comparable to the vdW interaction. Clearly, this indicates that a field-modified vdW interaction far detuned from resonance would be negligible for experimentally attainable CW intensities of ~ 1 W/cm². However, near resonance it is well known that the polarizability can be enhanced by a factor of $\omega_0/\gamma \sim 10^7$, where γ is the width of an atomic resonance at ω_0 . This will reduce the intensity requirement by a factor of $(\gamma/\omega_0)^2$. Thus, an optical field near resonance will have an influence comparable to the vdW interaction when $I \sim 1$ W/cm². A more quantitative and rigorous calculation of the required field will be put forth in the subsequent sections.

III. CLASSICAL MODEL

To investigate how the atom-surface interaction potential changes when the atom is in the presence of monochromatic light, we first explore a simple model with classical electric fields and polarizable atoms. In the Lorentz oscillator model,

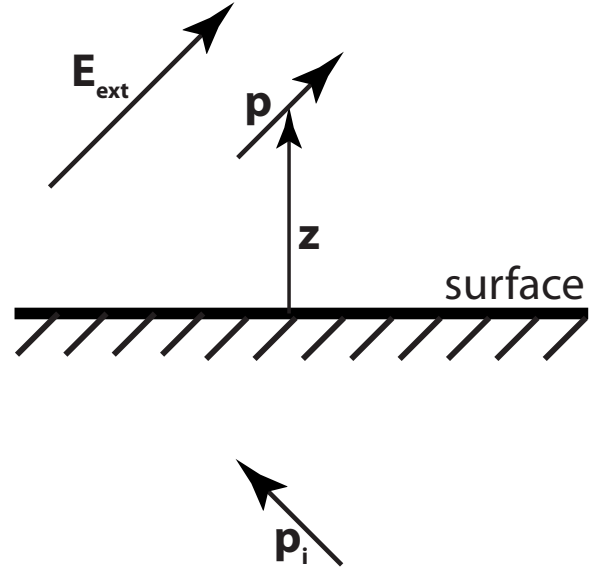


FIG. 1. Depiction of the field-modified atom-surface interaction. An external field \mathbf{E}_{ext} induces a dipole moment \mathbf{p} in an atom some distance z from a surface. The atom also experiences a field from its image dipole \mathbf{p}_i within the surface, leading to a force between the atom and surface.

atoms have a polarizability defined by $\mathbf{p} = \alpha\mathbf{E}$, where \mathbf{E} is the electric field that induces an electric dipole moment \mathbf{p} in the atom. The magnitude of polarizability depends on frequency as [19]

$$|\alpha(\omega)| = \frac{e^2}{m\sqrt{(\omega_0^2 - \omega^2)^2 + \gamma^2\omega^2}}, \quad (4)$$

where ω_0 is the resonance frequency and γ describes the damping for the oscillator. The phase of \mathbf{p} lags behind the phase of \mathbf{E} by an amount $\phi = \tan^{-1}[\gamma\omega/(\omega_0^2 - \omega^2)]$, so in general the polarizability can be written as a complex quantity [14]

$$\alpha(\omega) = \frac{e^2}{m(\omega_0^2 - \omega^2 + i\gamma\omega)}. \quad (5)$$

Away from resonance where $|\omega - \omega_0| \gg \gamma$,

$$\alpha(\omega) \approx \frac{e^2}{m(\omega_0^2 - \omega^2)} \approx -\alpha_{\text{DC}} \frac{\omega_0}{2\Delta}, \quad (6)$$

where $\alpha_{\text{DC}} \equiv e^2/m\omega_0^2$. The very last expression is valid in the limit where the detuning $\Delta \equiv (\omega - \omega_0)$ satisfies both $\omega_0 \gg |\Delta| \gg \gamma$ so the approximation $\omega_0 + \omega \approx 2\omega_0$ can be used.

If we describe the electric field also with classical physics (i.e., without including any fluctuations), then in the absence of an externally applied electric field this Lorentz oscillator has no reason to be polarized. Even its mean square dipole moment $\langle p^2 \rangle$ is zero. Therefore, the oscillator will not interact with a surface, a point discussed by Spruch in Ref. [20].

Let us consider the situation in Fig. 1, where an atom with some dipole moment p induced by an external field \mathbf{E}_{ext} is located a distance z from a perfectly conducting surface. Using the method of images [18], we can replace the perfectly

conducting surface with an image dipole p_i . Therefore the potential experienced by the atom is given by

$$\begin{aligned} U &= -\frac{1}{2}\alpha(\omega)|\mathbf{E}_{\text{ext}} + \mathbf{E}_{\text{image}}|^2 \\ &= -\frac{1}{2}\alpha(\omega)[|\mathbf{E}_{\text{ext}}|^2 + |\mathbf{E}_{\text{image}}|^2 + 2\mathbf{E}_{\text{ext}} \cdot \mathbf{E}_{\text{image}}], \end{aligned} \quad (7)$$

where the atom interacts with the superposition of the external field \mathbf{E}_{ext} and the field from its image dipole $\mathbf{E}_{\text{image}}$.

The first term in Eq. (7) is just the ac Stark shift which does not depend explicitly on the atom-surface distance, and plays no role in modifying the vdW interaction. However, the importance of distinguishing this term from a field-modified vdW interaction will be mentioned later. The second term in the brackets can be neglected since it is of order $[\alpha(\omega)z^{-3}]^2$ and will be small compared to the cross-term which is of order $\alpha(\omega)z^{-3}$, so long as $\alpha(\omega)z^{-3}/4\pi\epsilon_0 \ll 1$. Thus, we identify the third term as the explicit and leading-order atom-surface interaction potential given by

$$U_{\text{AS}} = -\alpha(\omega)[\mathbf{E}_{\text{ext}} \cdot \mathbf{E}_{\text{image}}]. \quad (8)$$

The external field at the location of the atom can be defined as

$$\mathbf{E}_{\text{ext}} = \frac{1}{2}\mathbf{E}_0 e^{-i\omega t} + \text{c.c.}, \quad (9)$$

where the factor of 1/2 is required to make the electric field have the form $\mathbf{E}_{\text{ext}} = \mathbf{E}_0 \cos(\omega t)$. Therefore, the field at the location of the atom due to the image dipole can be written as

$$\mathbf{E}_{\text{image}} = \frac{1}{2}[3(\hat{\mathbf{p}}_i \cdot \hat{\mathbf{z}})\hat{\mathbf{z}} - \hat{\mathbf{p}}_i] \frac{\alpha(\omega)E_0}{4\pi\epsilon_0(2z)^3} e^{-i[\omega t - (2z\omega/c)]} + \text{c.c.}, \quad (10)$$

where \vec{z} is the vector from the closest point on the surface to the atom and \mathbf{p}_i is the image dipole. Dipole radiation terms which vary as z^{-2} and z^{-1} are neglected since the atom-surface distances considered here are in the near-field limit [21]. From the method of images with a perfectly conducting surface, $\mathbf{p}_{i,z} = \mathbf{p}_z$ and $\mathbf{p}_{i,\perp} = -\mathbf{p}_\perp$ as indicated in Fig. 1. The phase factor $2z\omega/c$ comes from retardation.

Hence the atom-surface interaction potential is

$$U_{\text{AS}} = -\frac{\alpha(\omega)^2 |E_0|^2}{64\pi\epsilon_0 z^3} [(\hat{\mathbf{E}}_0 \cdot \hat{\mathbf{z}})^2 + 1] \cos(2z\omega/c), \quad (11)$$

where we have averaged over rapidly oscillating terms such as $e^{\pm i2\omega t}$. The expression in brackets ranges from 1 to 2 depending on the field polarization with respect to $\hat{\mathbf{z}}$.

We now consider the nonretarded limit so the cosine factor in Eq. (11) equals 1. The field-induced atom-surface interaction potential (11) has the form of $U_{\text{AS}} = -C_3/z^3$, where the coefficient C_3 for this all-classical model is

$$C_3^{(\text{class})} = \left(\frac{\alpha_{\text{DC}}}{4\pi\epsilon_0}\right)^2 \left(\frac{\omega_0}{\Delta}\right)^2 [(\hat{\mathbf{E}}_0 \cdot \hat{\mathbf{z}})^2 + 1] \frac{\pi I}{8c}, \quad (12)$$

where $I = \epsilon_0 c |E_0|^2 / 2$ is the optical intensity of a plane wave and we have used Eq. (6) for the polarizability near resonance.

Let us compare the strength of this potential to the benchmark obtained with the Lifshitz formula. The two are equal for the case of polarization perpendicular to $\hat{\mathbf{z}}$ when

$$I = \frac{\hbar c \Delta^2}{\left(\frac{\alpha_{\text{DC}}}{4\pi\epsilon_0}\right) \omega_0 \pi}. \quad (13)$$

For sodium atoms $\frac{\alpha_{\text{DC}}}{4\pi\epsilon_0} = 24 \times 10^{-30} \text{ m}^3$ [22], and if we consider $\Delta = 10\gamma = 2\pi \times 100 \text{ MHz}$, we find that an optical intensity of $I = 5 \text{ W/cm}^2$ is needed to make $C_3^{(\text{class})} = C_3^{(0)}$. In other words, at this intensity the field-modified atom-surface interaction will be equal to the vdW interaction of an atom in vacuum according to Eq. (3). In the following section we consider the same problem but with a quantized electric field, which will enable a comparison of vacuum, thermal, and monochromatic field contributions to C_3 .

IV. QUANTUM MODEL

Milonni and Smith used a field operator method to show how externally applied electromagnetic fields affect atom-atom interactions [6]. We will apply this approach to the atom-surface problem and arrive at some of the same conclusions: (1) Modifications to the potential are enhanced when the frequency of the applied field is near an atomic resonance. (2) Such modifications could be observed with optical intensities easily obtainable in laboratories. (3) The modifications also depend on the polarization of the applied fields. However, unlike the atom-atom system the image dipole depends on the atomic polarization in a way that makes the nonretarded atom-surface interactions always attractive, and can only be increased by illuminating the system with laser light. Although, this situation changes if the atomic and image dipoles are allowed to be out of phase due to retardation effects or when the surface response is such that $\epsilon(\omega) < 0$ [23,24].

We now allow the externally applied electric field at the location of the atom to have many frequency components

$$\mathbf{E}_{\text{ext}} = \frac{1}{2} \sum_k [\mathbf{E}_{k,0} e^{-i\omega_k t} + \text{c.c.}], \quad (14)$$

and thus rewrite the atom-surface interaction (11) as

$$U_{\text{AS}} = -\text{Re} \sum_k \frac{\alpha(\omega_k)^2}{64\pi\epsilon_0 z^3} \mathbf{E}_{k,0}^* \cdot [\hat{\mathbf{z}}\hat{\mathbf{z}} + 1] \cdot \mathbf{E}_{k,0} e^{i2kz} \quad (15)$$

with $k = \omega/c$. In writing Eq. (15) we have assumed that different frequency components are uncorrelated, and we have averaged over rapidly oscillating terms. We can use a plane-wave mode decomposition for the electric field where each mode index k corresponds to a wave vector \mathbf{k} and a polarization index $\lambda (= 1, 2)$. This can be expressed as

$$\mathbf{E}_{k,0} = \mathbf{e}_{k\lambda} A_{k\lambda,0} e^{i\mathbf{k}\cdot\mathbf{x}}, \quad (16)$$

where $\mathbf{e}_{k\lambda}$ is a unit polarization vector ($\mathbf{k}\cdot\mathbf{e}_{k\lambda}=0$, and $\mathbf{e}_{k\lambda}\cdot\mathbf{e}_{k\lambda'}=\delta_{\lambda\lambda'}$). Then the atom-surface interaction (15) becomes

$$U_{AS} = -\text{Re} \sum_{k\lambda} \frac{\alpha(\omega_k)^2}{64\pi\epsilon_0 z^3} [(\mathbf{e}_{k\lambda}\cdot\hat{\mathbf{z}})^2 + 1] A_{k\lambda,0}^* A_{k\lambda,0} e^{i2kz}. \quad (17)$$

Equation (17) is formulated in classical electromagnetic theory. As described in Ref. [2], when the electromagnetic field is quantized the amplitude takes the form $A_{k\lambda,0} = (\hbar\omega_k/2\epsilon_0 V)^{1/2}(a_{k\lambda} - a_{k\lambda}^\dagger)$, where $[a_{k\lambda}, a_{k\lambda}^\dagger] = 1$. Using a similar prescription as Ref. [6], we can then make the following replacement for the mode amplitudes in Eq. (17)

$$A_{k\lambda,0}^* A_{k\lambda,0} \rightarrow \left(\frac{\hbar\omega_k}{2\epsilon_0 V} \right) [a_{k\lambda} a_{k\lambda}^\dagger + a_{k\lambda}^\dagger a_{k\lambda} - a_{k\lambda}^2 - a_{k\lambda}^{\dagger 2}], \quad (18)$$

where $a_{k\lambda}^\dagger$ and $a_{k\lambda}$ are the field mode creation and annihilation operators for the field mode ($\mathbf{k}\lambda$) and V is the quantization volume. After taking the expectation value of Eq. (17) over the field state, the last two terms of Eq. (18) vanish. Following these operations, the potential becomes

$$U_{AS} \rightarrow - \sum_{k\lambda} \left(\frac{\hbar\omega_k}{\epsilon_0 V} \right) \frac{\alpha^2(\omega_k) \cos(2kz)}{64\pi\epsilon_0 z^3} \times \left(N_{k\lambda} + \frac{1}{2} \right) [(\mathbf{e}_{k\lambda}\cdot\hat{\mathbf{z}})^2 + 1], \quad (19)$$

where $N_{k\lambda} = \langle a_{k\lambda}^\dagger a_{k\lambda} \rangle$ is the average number of photons in mode $\mathbf{k}\lambda$. It is important to note that in the steps leading up to Eq. (19) we have again replaced the surface with an image dipole, which incorporates the boundary conditions on the modes of the electric field according to Maxwell's equations. This image dipole method can lead to more intuition about the physics of the problem, and it avoids the mathematically complicated task of finding an expression for the allowed modes of the electric field for an atom near a surface.

A. Ordinary C_3

In order to calculate the atom-surface potential in the quantum vacuum $U_{AS}^{(\text{vac})}$, we consider Eq. (19) in the absence of any external radiation, which implies that $N_{k\lambda}=0$. Indeed, we find that the remaining interaction is attributed to the zero-point electromagnetic field energy of $\frac{1}{2}\hbar\omega_k$ per mode. An expression for $U_{AS}^{(\text{vac})}$ can be derived by going to the mode continuum limit in which $\sum_{k\lambda} \rightarrow (\frac{V}{8\pi^3}) \int d^3k \sum_\lambda$, leading to

$$U_{AS}^{(\text{vac})} = - \frac{\hbar}{(4\pi\epsilon_0)^2} \frac{1}{64z^3\pi^2} \int d^3k \omega_k \alpha^2(\omega_k) \cos(2kz) \times \sum_\lambda [(\hat{\mathbf{e}}_{k\lambda}\cdot\hat{\mathbf{z}})^2 + 1]. \quad (20)$$

With the identity $\sum_{\lambda=1}^2 (\hat{\mathbf{e}}_{k\lambda}\cdot\hat{\mathbf{z}})^2 = 1 - (\hat{\mathbf{k}}\cdot\hat{\mathbf{z}})^2$ this can be written as

$$U_{AS}^{(\text{vac})} = - \frac{\hbar}{(4\pi\epsilon_0)^2} \frac{1}{64z^3\pi^2 c^3} \int d\omega \omega^3 \alpha^2(\omega) \cos(2z\omega/c) \times \int d\Omega_k [3 - (\hat{\mathbf{k}}\cdot\hat{\mathbf{z}})^2], \quad (21)$$

where $d\Omega_k$ is an element of solid angle about $\mathbf{k}=k\hat{\mathbf{k}}=\omega\hat{\mathbf{k}}/c$. Using $\int d\Omega_k [3 - (\hat{\mathbf{k}}\cdot\hat{\mathbf{z}})^2] = 32\pi/3$ and the full expression (4) for the atomic polarizability yields the following expression:

$$U_{AS}^{(\text{vac})} = - \left(\frac{\alpha_{\text{DC}}}{4\pi\epsilon_0} \right)^2 \frac{\hbar\omega_0^4}{6z^3\pi c^3} \int_0^\infty d\omega \frac{\omega^3 \cos(2z\omega/c)}{(\omega_0^2 - \omega^2)^2 + \gamma^2 \omega^2}. \quad (22)$$

The integral in Eq. (22) diverges logarithmically at the upper limit of integration in the absence of the cosine factor that came from retardation. A careful use of cutoffs [20] leads to the value $\pi\omega_0/2\gamma$ for the integral. Recalling $\gamma = \alpha_{\text{DC}}\omega_0^4/(6\pi\epsilon_0 c^3)$ [14,15], we find that

$$U_{AS}^{(\text{vac})} = - \frac{\alpha_{\text{DC}}}{4\pi\epsilon_0} \frac{\hbar\omega_0}{8} \frac{1}{z^3} \equiv - \frac{C_3^{(0)}}{z^3}. \quad (23)$$

This confirms that the calculation method presented here, in which the electromagnetic field is quantized while the atom is regarded as a classical polarizable body, can correctly describe the ordinary vdW interaction.

One may interpret this result as a statement that vacuum fluctuations in the electromagnetic field cause the ordinary vdW interaction. Although other points of view (i.e., quantum mechanical atoms and classical fields) also work to describe the ordinary vdW interaction. One virtue of the theory described here is that it can readily be adapted to describe modifications to $U_{AS}^{(\text{vac})}$ due to thermal radiation, or other externally applied electromagnetic fields, by adding the appropriate additional terms for $N_{k\lambda}$ in Eq. (19).

B. Thermal fields

According to Ref. [6], in the case of a thermal field at temperature T the photon number distribution is given by $N(\omega) = (e^{\hbar\omega/k_B T} - 1)^{-1}$, where k_B is Boltzmann's constant, implying that $N(\omega) + \frac{1}{2} = \frac{1}{2} \coth(\hbar\omega/2k_B T)$. Neglecting retardation we find that the van der Waals energy at temperature T is

$$U_{AS}^{(\text{therm})} = - \left(\frac{\alpha_{\text{DC}}}{4\pi\epsilon_0} \right)^2 \frac{\hbar\omega_0^4}{6z^3\pi c^3} \int_0^\infty d\omega \frac{\omega^3 \coth(\hbar\omega/2k_B T)}{(\omega_0^2 - \omega^2)^2 + \gamma^2 \omega^2}. \quad (24)$$

The integral in Eq. (24) is dominated by values near $\omega \sim \omega_0$, and we can classify the temperature regimes with respect to $\hbar\omega_0$. Thus, in the low-temperature limit $\hbar\omega_0/k_B T \gg 1$, and $\coth(\hbar\omega/2k_B T) \sim 1$ implying $U_{AS}^{(\text{therm})} = -C_3^{(0)}/z^3$. In the high-temperature limit, $\hbar\omega_0/k_B T \ll 1$, $\coth(\hbar\omega/2k_B T) \sim 2k_B T/\hbar\omega$, and

$$U_{AS}^{(\text{therm})} = - \left(\frac{\alpha_{\text{DC}}}{4\pi\epsilon_0} \right)^2 \frac{\hbar\omega_0^4}{6z^3\pi c^3} \frac{2k_B T}{\hbar} \times \int_0^\infty d\omega \frac{\omega^2}{(\omega_0^2 - \omega^2)^2 + \gamma^2\omega^2}. \quad (25)$$

The integral evaluates to $\pi/2\gamma$, implying

$$U_{AS}^{(\text{therm})} = \frac{2k_B T}{\hbar\omega_0} U_{AS}^{(\text{vac})}, \quad (26)$$

which is also a well-known result, for example, written as Eq. (20) in Ref. [5]. For sodium atoms in a room-temperature environment $2k_B T/\hbar\omega_0 \approx 0.02$, which unfortunately indicates that thermal contributions to atom-surface interactions in the nonretarded regime are usually negligible.

C. Light-enhanced C_3

Now we return to look at terms in Eq. (19) that arise when monochromatic light is applied so that $N_{\mathbf{k}\lambda} \neq 0$ for one particular mode. If one particular mode ($\mathbf{k}\lambda$) is occupied by many photons (i.e., $N_{\mathbf{k}\lambda} \gg 1$), it follows from Eq. (18) that $N_{\mathbf{k}\lambda}/V = \epsilon_0 |E_0|^2 / \hbar\omega_k$. Equation (19) then becomes

$$U_{AS}^{(\text{CP,field})} = U_{AS}^{(\text{CP},0)} - \frac{|E_0|^2 \alpha^2(\omega_k)}{64\pi\epsilon_0 z^3} [(\mathbf{e}_{\mathbf{k}\lambda} \cdot \hat{\mathbf{z}})^2 + 1] \cos(2kz). \quad (27)$$

As we foresaw with the classical model, if the populated photon mode has polarization perpendicular to $\hat{\mathbf{z}}$ then the term in brackets is unity. If we consider the nonretarded limit then the cosine term is also unity, and we are left with

$$U_{AS}^{(\text{vdW,field})} = U_{AS}^{(\text{vdW},0)} - \left(\frac{\alpha(\omega_k)}{4\pi\epsilon_0} \right)^2 \frac{\pi I}{2cz^3} = U_{AS}^{(\text{vdW},0)} - \left(\frac{\alpha_{\text{DC}}}{4\pi\epsilon_0} \right)^2 \left(\frac{\omega_0}{\Delta} \right)^2 \frac{\pi I}{8cz^3}. \quad (28)$$

Hence we find that 5 W/cm² will be required to double the ordinary $C_3^{(0)}$ for sodium atoms by applying light with $\Delta = 2\pi \times 100$ MHz and polarization perpendicular to $\hat{\mathbf{z}}$, as was found from Eq. (13).

V. DISCUSSION

Equation (27) offers several testable predictions. For example, experiments could verify how the laser intensity, frequency, and polarization affect $U_{AS}^{(\text{vdW,field})}$. The qualitative dependence on the laser parameters is similar to the atom-atom system considered in Ref. [6]. However, unlike the atom-atom case, the field-modified atom-surface interaction considered here is always attractive. This is because the image dipole is entirely determined by the atomic dipole in the nonretarded regime. Although, the image dipole could be out of phase with the atomic dipole if one allows for the dynamic response of a realistic surface, described by $\epsilon(\omega)$. If $\epsilon(\omega_k) < 0$ then repulsive interactions can also be caused in the nonretarded limit [23,24]. In addition, the atomic and image

dipole could be out of phase in the retarded regime, possibly leading to repulsive interactions.

Scattering may always be a problem for experiments because

$$\frac{U_{AS}^{(\text{vdW,field})}}{U_{AS}^{(\text{vdW},0)}} = 1 + \frac{\alpha_{\text{DC}}}{4\pi\epsilon_0} \frac{\pi\omega_0}{\hbar c} \frac{I}{\Delta^2}. \quad (29)$$

Equation (29) emphasizes how the relative enhancement factor depends on I/Δ^2 . We note that the spontaneous emission rate depends on the same factor. Hence the enhancement of C_3 due to optical fields cannot be arbitrarily large without also causing spontaneous emission. As another point of view, one may use this fact to see that the enhancement we are discussing requires a significant atomic population in the excited state. This superposition of ground and excited states is necessary to induce an electric dipole moment in the atom.

Candidate experiments

Many experiments in the last two decades have measured the strength of atom-surface interactions, so it is worth investigating why this effect has not yet been observed. We will catalog several experiments here according to two criteria. First, are they sensitive only in the nonretarded regime? Second, can the intensity of near-resonant light be varied from 0 to 10 W/cm² without jeopardizing the experimental sensitivity to C_3 ?

The first question is important because the enhancement we are discussing will be very different in the retarded regime. In particular, the $\cos(2kz)$ term in Eq. (19) and the long-range terms which were neglected in Eq. (10) can cause the interaction to change sign as a function of atom-surface distance on the scale of an optical wavelength. If an experiment is sensitive in the retarded regime, it may average over a range of distances such that the modification to the atom-surface interactions due to laser light will be obscured. For example, during quantum reflection experiments [25–27], Talbot Lau interferometer experiments [28], atom beam transmission through cavities [29], and BEC oscillation near an atom-chip surface [30] the atom-surface interactions are probed at distances larger than an optical wavelength. The second criterion is important because the signal used to measure C_3 may require laser light. Then the intensity of the light itself may alter the signal even if the vdW interaction is not directly affected by the light. For example, experiments with atoms reflected from evanescent wave potentials [31,32], spectroscopic shifts of atomic resonances with atomic beams [33], and atomic vapors in nanocells [34] all need light in order to measure C_3 . Hence, an attempt to change C_3 by modifying the intensity of the light may cause changes in the experimental observations simply due to ac Stark shifts or power broadening.

One candidate experiment that we suggest for measuring the difference between $U_{AS}^{(\text{vdW,field})}$ and $U_{AS}^{(\text{vdW},0)}$ uses atom diffraction from nanostructure transmission gratings (nanogratings). The atoms are transmitted through a 100 nm period grating with $\sim 50\%$ open fraction so that each atom passes within 25 nm of a grating bar surface. In these experiments the atom-surface interaction is measured by observing

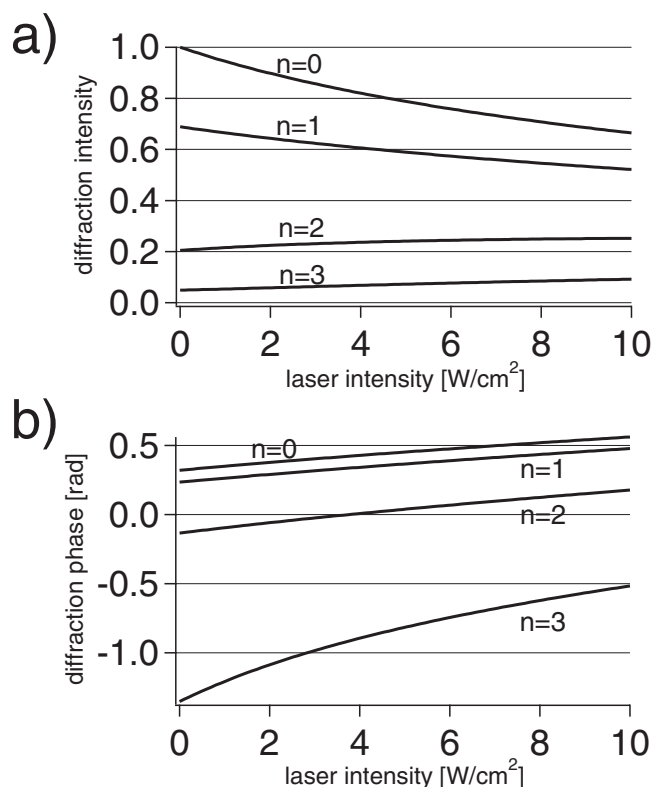


FIG. 2. Searching for a field-modified vdW interaction with atomic diffraction from a material grating. The intensity (a) and phase (b) of atom-wave diffraction orders n are plotted as a function of the intensity of an applied laser field, using Eq. (28) and the diffraction theory from Ref. [35] for typical experimental parameters [36]. As the atoms pass by the grating walls their de Broglie wave phase is influenced by the vdW interaction with the grating surface, altering the far-field atom-wave diffraction pattern [35–37]. The grating surface is assumed to be perfectly conducting so that when the laser field is off (zero intensity) Na atoms will experience a vdW coefficient of $C_3=6.3 \text{ meV nm}^3$, which is doubled when the laser intensity is 5 W/cm^2 .

changes in the diffraction order intensity [36,38] or phase [35,37]. The experiments probe the atom-surface interaction in the range of 9 to 20 nm and are sensitive without laser light, avoiding the effects of spatial averaging and ambiguity of light-induced changes to the experimental signal. Therefore, the atomic diffraction experiments satisfy both of the criteria mentioned at the beginning of this section. Figure 2 shows how the intensity and phase of atom-wave diffraction patterns can be significantly modified by external laser fields. For example, if we consider Na atoms and a perfectly conducting grating the ratio of second to zeroth order diffraction intensity increases by 50% when a 5 W/cm^2 laser field illuminates the atoms. In addition, a 5 W/cm^2 laser field will cause a ~ 0.2 radian phase shift to the zeroth diffraction order and a ~ 0.5 rad shift to the third order. These changes are well within the detection sensitivity of the diffraction experiments [35–38].

Since the nanogratings are only 120 nm thick, and made of dielectric silicon nitride material, the optical intensity inside each grating slot will not have nodes as in a macroscopic waveguide. Thus, although more calculations are

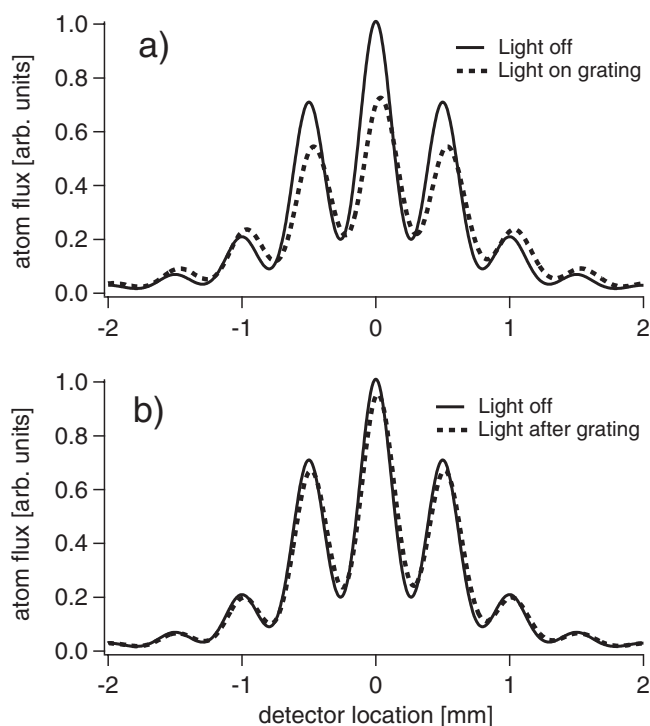


FIG. 3. Simulated laser-modified atom diffraction patterns. Light is predicted to change the diffraction patterns because of modified atom-surface interactions and also photon scattering. Plot (a) shows the effect of a laser beam, with intensity $I=5 \text{ W/cm}^2$ and detuning $\Delta=100 \text{ MHz}$, illuminating the atoms inside a nanograting. Plot (b) shows the effect of the same laser illuminating atoms after passing through a nanograting. In each simulation atoms exposed to the laser have an equal probability of spontaneously scattering 1 or 0 photons.

needed to verify this claim, the intensity may be homogeneous enough throughout each grating slot so that the ac Stark shift has a negligible gradient compared to the atom-surface interaction potential. Furthermore, the atoms in a beam only interact with the grating bar surfaces for a time on the order of 10^{-10} s —much shorter than a spontaneous decay lifetime. Hence, a light beam may be designed to illuminate the atoms for such a short time that only one or zero photons will be scattered by each atom.

The signature of a change in C_3 in these experiments [36,38] is a change in the relative intensities of coherently diffracted orders. Scattering can be distinguished from this effect because the momentum gained by scattering a photon (in the range 0 to $2 \hbar k_{\text{ph}}$) is smaller than the momentum difference $\Delta p = \hbar/d \approx 6 \hbar k_{\text{ph}}$ due to first order diffraction from a $d=100 \text{ nm}$ period nanograting, where the number 6 is valid for sodium atoms with $\lambda_0=590 \text{ nm}$. The effect of scattering could also be isolated by comparing diffraction patterns obtained when the laser is illuminating the atoms before, during, and after their passage through the grating. To further illustrate these concepts, Fig. 3 shows how laser light will modify atom diffraction patterns. Spontaneous emission will broaden and shift each diffraction peak. On the other hand, if the laser affects how atoms interact with the nanograting bars, then modified atom-surface interactions will change the total atom flux in each order [given by the area of

each diffraction peak and plotted in Fig. 2(a)]. Figure 3(a) shows a diffraction pattern with both effects. Note how the zeroth order is reduced, but the second order is slightly increased as a result of the laser light despite the scattering effect. Figure 3(b) shows atom diffraction modified by scattering alone. In this case the area of each diffraction order does not change. An additional control experiment could be operated by comparing diffraction of different atomic species, one for which the light is near resonance and one that is not resonant with the light. Also we note that the ac Stark shift can be reversed with the sign of Δ , whereas the enhancement factor in Eq. (29) is independent of the sign of Δ .

In addition to diffraction of ground state atoms from material gratings there are other candidate experiments that satisfy the selection criteria previously mentioned. Equation (29) indicates that the relative magnitude of the field-modified vdW interaction depends linearly on the static polarizability α_{DC} of the atom. Atoms prepared in an excited state generally have a much larger polarizability than those in the ground state, which could enhance the size of this effect and reduce the intensity requirement for the external laser field. In fact, using excited state atoms to achieve a larger vacuum vdW coefficient has already been used in atomic deflection [39] and spectroscopy experiments [33]. Atom scattering from magnetized surfaces may also provide a means to search for a field modified vdW interaction [40]. Reflection-mode diffraction of atomic beams from crystal surfaces also satisfies the two criteria for a candidate experiment [41]. Bound-state resonance features in these experiments can reveal information about the atom-surface potential [42] and, in principle, could be used to search for the

field-modified vdW interaction. However, a remaining practical challenge for using reflection-mode diffraction is that these experiments primarily work with ground-state noble-gas atoms. Experiments which can utilize alkali atoms have the advantage of larger polarizabilities (e.g., $\alpha_{Na}/\alpha_{He} \sim 125$) and near-resonant light which can be generated conveniently from diode lasers.

VI. CONCLUSION

We used a quantized electric field approach to show how the atom-surface interaction potential can be modified by applying near-resonant electromagnetic (optical) fields. In addition, we have demonstrated that the quantum electrodynamics approach used here also serves to derive two well-known results (the low and high temperature vdW interaction potential), and this lends credence to the method of calculation. The derivations presented here also show how electromagnetic vacuum fluctuations and fluctuations of the atomic dipole moment provide complementary descriptions of the atom-surface van der Waals interaction. Finally, we discussed candidate experiments that may serve to test the prediction that near-resonant laser light can modify the interaction between atoms and surfaces.

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