Lateral Casimir-Polder forces by breaking time-reversal symmetry

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We examine the lateral Casimir-Polder force acting on a circular rotating emitter near a dielectric plane surface. As the circular motion breaks time-reversal symmetry, the spontaneous emission in a direction parallel to the surface is in general anisotropic. We show that a lateral force arises which can be interpreted as a recoil force because of this asymmetric emission. The force is an oscillating function of the distance between the emitter and the surface, and the lossy character of the dielectric strongly influences the results in the near-field regime. The force exhibits also a population-induced dynamics, decaying exponentially with respect to time on timescales of the inverse of the spontaneous decay rate. We propose that this effect could be detected measuring the velocity acquired by the emitter, following different cycles of excitation and thermalisation. Our results are expressed in terms of the Green’s tensor and can therefore easily be applied to more complex geometries.

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I. INTRODUCTION

Casimir-Polder (CP) forces are forces between atoms and magneto-dielectric bodies originating from the quantum fluctuations of the electromagnetic field and the atomic charges [1, 2]. Despite their generally low magnitude they have to be taken into account when working on nanoscales which today is common in experiments as well as applications [3].

Lateral Casimir forces are a relatively young topic, having first been measured in 2002 [4]. They are characterised by their direction which is parallel to the surface instead of the usual normal direction, and have been suggested to facilitate contactless force transmissions [5]. Lateral forces are typically achieved by breaking the translational symmetry of the surface, using for example periodically structured surfaces [4, 6–9] or corrugated surfaces [10–13]. A lateral force has also been realised by breaking the mirror symmetry, using chiral particles near a surface [14, 15]. It is discriminatory since it pushes chiral particles with opposite handedness in opposite directions.

In this article we show that a lateral force can arise by breaking time-reversal symmetry via a rotating dipole which emits asymmetrically. A dipole moment is created with left-handed or right-handed circularly polarised light with spin parallel to the surface. More specifically we excite a Cesium atom from the hyperfine ground-state \( |6^2S_{1/2}, F = 4, M_F = 4 \rangle \) to the excited state \( |6^2P_{3/2}, F' = 5, M'_{F} = 5 \rangle \) using a resonant right-handed circularly polarised laser beam that propagates along the y-direction, creating a dipole moment rotating in the \( x - z \) plane, where \( \hat{z} \) is the direction normal to the surface. The excitation of guided and radiation modes that propagate in the \( +x \) and \( -x \) directions may be expected to be asymmetric in this case [16, 17]. The asymmetric emission of guided modes has been extensively investigated in the literature and relies on spin-orbit coupling of light mediated by a particle near a surface [18–24]. The conservation of total momentum in the system in conjunction with the asymmetric emission suggests the existence of a lateral force opposite to the direction of stronger emission. This lateral force could be measured observing the asymmetric emission distribution. A similar effect has been studied for the same atomic system close to an optical nanofiber [16]. It is the circular polarisation of the illuminating light that creates a rotating dipole-moment which breaks time-reversal symmetry. Using linearly polarised light, the time-reversal symmetry is conserved and the force is purely normal to the surface.

Note that this system has been analysed previously

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by classical methods, studying the interaction of the rotating dipole with the reflected field \[25\]. However the excited atom will unavoidably decay to the ground state, for which the lateral force is forbidden by energy conservation. Hence the lateral force has a population-induced dynamics which can be captured only by quantum approaches to the atom-field coupling. It is the aim of this article to investigate this effect.

The problem can be solved by expanding the electric field into guided and radiation modes. However here we use a different approach which relies on the Green’s tensor and permits to study geometries different from the planar configuration as well the impact of dissipation \[26, 27\].

The article is organised as follows. In section II we study the asymmetric emission of a multilevel circular emitter near a surface and investigate how the lateral force arises from this asymmetric emission. We also develop a dynamical approach to the atom–field coupling which gives the same expression for the lateral force. In section III we consider a two-level Cesium atom near a surface and investigate how the lateral force has a population-induced excitation. Hence the lateral force is forbidden by energy conservation: for which the lateral force is forbidden by energy conservation: the excited atom will unavoidably decay to the ground state, leaving the atom and the field in the stationary state, if it existed, one could accelerate the atom along the surface, leaving the atom and the field in the stationary ground-states. A lateral force can arise only from the asymmetric emission of the photon as we because of the asymmetric emission of the photon as we assume the atomic eigenstate \(|n\rangle\) to exhibit a well-defined \(y\)-component of angular momentum. It is then a circular emitter whose transitions matrix elements are not real, in general, but obey \(d_{kn} = d_{nk}^* \[26\].

**II. LATERAL FORCE FROM THE ASYMMETRIC EMISSION**

We consider a multi-level atom, with atomic frequencies \(\omega_{nk} = (E_n - E_k)/\hbar\), and dipole moments \(d_{nk}\), near a homogeneous and isotropic dielectric plate (see Fig. 1). The atom is prepared in an incoherent superposition of energy eigenstates \(|n\rangle\) with occupation \(p_n(t)\) and is placed at position \(r_A = z_A \hat{z}\) from the dielectric plate. We assume the atomic eigenstate \(|n\rangle\) to exhibit a well-defined \(y\)-component of angular momentum. It is then a circular emitter whose transitions matrix elements are not real, in general, but obey \(d_{kn} = d_{nk}^* \[26\].

![Vacuum](image)

**FIG. 1.** Sketch of the system, expected lateral force and correlated asymmetric decay due to rotational dielectric moment.

We first derive the lateral force using a dynamical approach to the atom–field coupling. According to the dynamical approach the Casimir-Polder force between the atom and the surface reads \[27 29\]:

\[
F(r_A, t) = \frac{\mu_0}{\pi} \lim_{\varepsilon \to 0^+} \left( \sum_{n} p_n(t) \sum_{k} \int_{0}^{\infty} d\omega \omega^2 \frac{1}{\omega - \omega_{nk} - i\varepsilon} \times \nabla d_{nk} \cdot \text{Im} G(r, r_A, \omega) \cdot d_{kn} \biggr|_{r=r_A + c.c.} \right). 
\]

This integral can be rotated to the imaginary axis leading to the following non-resonant and resonant contributions:

\[
F^\text{nr}(r_A, t) = \frac{\mu_0}{\pi} \sum_{n} p_n(t) \sum_{k} \int_{0}^{\infty} d\xi \xi^2 \left( \frac{1}{\omega_{nk} - i\xi} 
+ \frac{1}{\omega_{nk} + i\xi} \right) \text{Re}\left\{ \nabla d_{nk} \cdot G(r, r_A, i\xi) \cdot d_{kn} \biggr|_{r=r_A} \right\},
\]

\[
F^r(r_A, t) = 2\mu_0 \sum_{n} p_n(t) \sum_{k<n} \omega_{nk}^2 
\times \text{Re}\left\{ \nabla d_{nk} \cdot G(r, r_A, \omega_{nk}) \cdot d_{kn} \biggr|_{r=r_A} \right\}. 
\]

We first note that the lateral component of the non-resonant contribution vanishes, since \(\partial_x G^{(1)}(r, r_A, i\xi)\biggr|_{r=r_A}\) is an antisymmetric tensor and \(G^{(1)}(r, r_A, i\xi)\) is real. In fact the non-resonant term survives also for ground-state atoms, where the existence of a lateral force is forbidden by energy conservation: if it existed, one could accelerate the atom along the surface, leaving the atom and the field in the stationary ground-states. A lateral force can arise only from the resonant contribution, which describes the atom recoil because of the asymmetric emission of the photon as we shall show later.

Most importantly the force shows a population induced dynamics. For example if the atom is a two-level system initially excited in the state \(|1\rangle\) the populations of the ground-state and excited-state satisfy the rate equations:

\[
p_0 = \Gamma p_1, \\
p_1 = -\Gamma p_1, 
\]

where \(\Gamma\) is the spontaneous decay rate for the transition \(1 \to 0\) (0 is the ground-state). The solutions, subject to the initial conditions \(p_0(0) = 0\) and \(p_1(0) = 1\), are:

\[
p_0(t) = 1 - e^{-\Gamma t}, \\
p_1(t) = e^{-\Gamma t}. 
\]

The lateral force for this two-level system hence reads:

\[
F_x(r_A, t) = 2\mu_0 e^{-\Gamma t} \omega_{10}^2 
\times \text{Re}\left\{ d_{10} \cdot \partial_x G^{(1)}(r, r_A, \omega_{10}) \cdot d_{01} \biggr|_{r=r_A} \right\}, 
\]

and decays exponentially with respect to time. Because of this force the particle will acquire some velocity in the
lateral direction. Supposing that the initial velocity is zero, the mean-velocity acquired reads:

\[
v = \int_{0}^{\infty} \frac{d}{dt} F_{x}(r_{A}, t) \frac{m}{m} = \frac{F_{x}(r_{A}, 0)}{m}.
\]  

(7)

We next investigate the asymmetric emission of the circular emitter and how the lateral force arises from this asymmetry. The total emission rate \( \Gamma \) is the sum of the free-space emission rate \( \Gamma^{(0)} \) and the surface-assisted emission rate \( \Gamma^{(1)} \) where the photon is reflected by the dielectric surface. If the atom is prepared in the eigenstate \( |n\rangle \) it will decay to lower lying energy levels, and the decay rate can be expressed in terms of the scattering Green’s tensor \( G^{(1)} \) \cite{27,29}:

\[
\Gamma^{(1)}_{n}(z_{A}) = 2\mu_{0} \frac{1}{\hbar} \sum_{k<n} \omega_{nk}^{2}
\]

\[
\times \text{Im}\left\{ \mathbf{d}_{kn} \cdot G^{(1)} \left( r_{A}, r_{A}, \omega_{nk} \right) \cdot \mathbf{d}_{dkn} \right\},
\]

(8)

where we have used the property \( \mathbf{d}_{nk} \cdot \text{Im} G \left( r, r, \omega \right) \cdot \mathbf{d}_{kn} = \text{Im} \left\{ \mathbf{d}_{nk} \cdot G \left( r, r, \omega \right) \cdot \mathbf{d}_{kn} \right\} \).

The scattering Green’s tensor for a dielectric reads \((z > 0, z' > 0)\) \cite{26,27}:

\[
G^{(1)} \left( r, r', \omega \right) = \int_{0}^{2\pi} \frac{d\varphi}{\pi} \int_{0}^{\infty} \frac{dk^{\parallel}}{k^{\parallel}} G^{(1)} \left( r, r', \omega, k^{\parallel} \right),
\]

\[
G^{(1)} \left( r, r', \omega, k^{\parallel} \right) = \frac{i}{8\pi^{2}} \frac{1}{k^{\perp}} e^{i\mathbf{k}^{\parallel} \cdot \left( \mathbf{r} - \mathbf{r}' \right)} e^{+i\left( z + z' \right)}
\]

\[
\times \sum_{\sigma = s, p} r_{\sigma} e_{\sigma+} e_{\sigma-},
\]

(9)

where \( k^{\parallel} = k^{\parallel} \left( \cos \phi, \sin \phi, 0 \right) \) is wave vector parallel to the surface and \( k^{\perp} = \sqrt{\omega^{2} / c^{2} - k^{\parallel}^{2}} \), \( k^{\perp} = \sqrt{\epsilon(\omega) \omega^{2} / c^{2} - k^{\parallel}^{2}} \) are the perpendicular components of the wave-vector in vacuum and in the dielectric plate. The polarisation vectors and Fresnel reflection coefficients for s-polarised and p-polarised waves read:

\[
\mathbf{e}_{s \pm} = (\sin \varphi, -\cos \varphi, 0),
\]

\[
\mathbf{e}_{p \pm} = \frac{\epsilon}{\epsilon} \left( \mp k^{\perp} \cos \varphi, \mp k^{\perp} \sin \varphi, k^{\parallel} \right),
\]

\[
r_{s} = \frac{k^{\perp} - k_{m}^{\perp}}{k^{\perp} + k_{m}^{\perp}},
\]

\[
r_{p} = \frac{\epsilon(\omega) k^{\perp} - k_{m}^{\perp}}{\epsilon(\omega) k^{\perp} + k_{m}^{\perp}}.
\]

Using Eq. (9) the assisted rate reads:

\[
\Gamma^{(1)}_{n}(z_{A}) = \int_{0}^{2\pi} \frac{d\varphi}{\pi} \int_{0}^{\infty} dk^{\parallel} \omega_{nk}^{2} \gamma_{n} \left( z_{A}, k^{\parallel} \right),
\]

(10)

where \( \gamma_{n} \left( z_{A}, k^{\parallel} \right) \) is the emission rate density:

\[
\gamma_{n} \left( z_{A}, k^{\parallel} \right) = \frac{2\mu_{0}}{\hbar} \sum_{k<n} \omega_{nk}^{2}
\]

\[
\times \text{Im}\left\{ \mathbf{d}_{kn} \cdot G^{(1)} \left( r_{A}, r_{A}, \omega_{nk} \right) \cdot \mathbf{d}_{kn} \right\}.
\]

(11)

A lateral force may result from an unbalanced spontaneous emission into the +x and −x directions. According to the conservation of the total momentum the force is opposite to the momentum of the emitted photon. If the atom is prepared in an incoherent superposition of energy eigenstates \( |n\rangle \) with population \( p_{n} \) the force reads:

\[
F_{x}(z_{A}, t) = -\int_{0}^{2\pi} d\varphi \int_{0}^{\infty} dk^{\parallel} \frac{1}{\hbar} \sum_{n} p_{n} \gamma_{n} \left( z_{A}, k^{\parallel} \right)
\]

\[
\times \text{Re}\left\{ \mathbf{d}_{kn} \cdot \partial_{\mathbf{k}^{\parallel}} G^{(1)} \left( r, r_{A}, \omega_{nk} \right) \cdot \mathbf{d}_{kn} \right\}_{r=r_{A}},
\]

(12)

where we have used the relation \( \text{Im} \left( ix \right) = -\text{Re} x \). The lateral force is associated with the recoil of the atom because of asymmetric emission. It vanishes if the atom is in the ground-state or if the atomic dipole moments are real since \( \partial_{\mathbf{k}^{\parallel}} G^{(1)} \left( r, r_{A}, \omega_{nk} \right) \bigg|_{r=r_{A}} \) is an antisymmetric tensor. Note that the bulk part of the Green’s tensor gives no contribution to the lateral force since in the absence of the dielectric plate the emission is obviously symmetric.

### III. APPLICATION: CESIUM ATOM

The aim of this section is to investigate the lateral force for a Cesium atom near a dielectric plate and compare the orders of magnitude of the force to previous works in literature using Cesium atoms near a optical nanofiber \cite{16}.

The Cesium atom in its ground-state can be excited to the excited state \( |6^{2}P_{3/2}, F' = 5, M_{F} = 5\rangle \) by using a right-handed circularly polarised laser beam. Since the beam is propagating along the y direction the resulting electric dipole moment is rotating in the \( x - z \) plane:

\[
\mathbf{d}_{10} = d \left( i, 0, 1 \right),
\]

(13)

where \( i \) denotes the excited state and 0 the ground-state. The magnitude of the dipole moment is \( d = 1.9 \times 10^{-29} \text{Cm} \) and the wavelength of the emitted photon is \( \lambda_{10} = 852 \text{nm} \) \cite{19}. Furthermore \( \mathbf{d}_{10} = \mathbf{d}_{10}^{\ast} \). The emitted photon carries the momentum and is responsible for the lateral force. There is only one decay channel to the ground state with the emission of a \( \sigma^{+} \) photon, the atom hence can be treated as a effective two-level system.
Substituting the Green’s tensor [9] into (6) and performing the trivial angular integrals we find:

\[ F_x(r_A, t) = -\frac{e^{-i\Omega t}d^2}{2\pi\varepsilon_0} \text{Im} \left\{ \int_0^\infty dk \|k\|^3 e^{2ik\cdot z_A} F_p \right\}. \quad (14) \]

In the limit of short times the force does not vanish in the classical limit \( h \to 0 \), showing that it could be described by classical methods [25]. Only the \( p \)-polarised components of the field affect the force, since the \( s \)-polarised components lead to an emission equal in the \( +x \) and \( -x \) half spaces.

The first interesting special case is a perfect conducting body, described by a unity reflection coefficient of \( p \)-polarised waves. Remarkably we note that the force can be performed analytically for any given distance:

\[ F^\text{PC}_x(r_A, t) = e^{-i\Omega t} \left\{ \frac{3d^2}{4\pi\varepsilon_0\lambda_{10}z_A} \cos \left( \frac{4\pi z_A}{\lambda_{10}} \right) + \frac{d^2}{\varepsilon_0} \left( \frac{1}{\lambda_{10}^2 z_A} - \frac{3}{16\pi z_A^2} \right) \sin \left( \frac{4\pi z_A}{\lambda_{10}} \right) \right\}. \quad (15) \]

where \( k = \omega_{10}/c \) is the transition wave-vector. This is depicted in Fig. 2 for short times.

\[ \text{FIG. 2. LCP force for a perfectly conducting half space.} \]

We see that the force shows characteristic Drexhage-type oscillations. Remarkably we note that the force can change sign depending on the distance between the emitter and the surface. This shows the importance of radiation modes to describe this effect. In fact, guided modes are excited strongly only in one direction giving a lateral force of constant sign.

Another interesting feature of this force is that it is not conservative. In fact if we take the curl of the total force we obtain a non-vanishing result:

\[ \nabla \times F = \frac{\partial F_x}{\partial z_A} \hat{y} = \frac{d^2}{\varepsilon_0} \left\{ \left( \frac{4\pi^2}{\lambda_{10}^2 z_A} - \frac{3}{\lambda_{10} z_A^2} \right) \cos \left( \frac{4\pi z_A}{\lambda_{10}} \right) + \left( \frac{3}{4\pi z_A^2} - \frac{5\pi}{\lambda_{10}^2 z_A} \right) \sin \left( \frac{4\pi z_A}{\lambda_{10}} \right) \right\} \hat{y}. \quad (16) \]

\[ \text{Hence the Casimir-Polder force is not conservative and can not be generally derived from a potential, by just taking its gradient. Such class of forces, called curl forces, have been explored in the past [20].} \]

We next consider the case the case of a dissipative dielectric medium described by a complex relative permittivity. An analytical expression can be derived in the non-retarded regimes when the distance between the emitter and the surface is much smaller than the atomic wavelength \( z_A \ll \lambda_{10} \):

\[ F_x^{\text{ret}}(r_A, t) = -e^{-i\Omega t} \frac{3d^2}{8\pi\varepsilon_0 z_A^2} \text{Im} \frac{\varepsilon}{[\varepsilon + 1]^2}, \quad (17) \]

where \( \varepsilon = \varepsilon(\omega_{10}) \) and we have used the following Fresnel coefficient for \( p \)-polarised waves \( r_p = (\varepsilon - 1)/(\varepsilon + 1) \). It shows a divergence when the emitter approaches the surface because of the lossy nature of the medium. Note that an analogous divergence is observed in the spontaneous emission of a real dipole moment near a lossy surface, where the emitter is excited with a linear-polarised resonant laser beam [31, 32].

In the retarded regime, namely when the distance between the emitter and the surface is much greater than the atomic wavelength \( z_A \gg \lambda_{10} \) we have:

\[ F_x^{\text{ret}}(r_A, t) = e^{-i\Omega t} \frac{d^2\pi}{\varepsilon_0\lambda_{10} z_A} \text{Re} \left\{ \sqrt{\varepsilon - 1} \right\} \times \text{Im} \left\{ \sqrt{\varepsilon - 1} \right\} \cos \left( \frac{4\pi z_A}{\lambda_{10}} \right), \quad (18) \]

where the following Fresnel coefficient for \( p \)-polarised wave has been used: \( r_p = (\sqrt{\varepsilon - 1})/(\sqrt{\varepsilon + 1}) \).

The distance between two roots of the retarded case is a quarter of the emission wavelength. This agrees with intuition. The atom emits at some time an electromagnetic field while having a specific dielectric moment orientation. The field is partially reflected by the surface and interacts with the dipole after some delay, when the dipole moment has a different orientation. If the distance between the emitter and the surface is incremented by \( \lambda_{10}/4 \) the delay is incremented by \( \Delta t = \lambda_{10}/2c \), leading to an opposite orientation of the dipole \( d \rightarrow d e^{i\omega_{10}\Delta t} = -d \). This translates to a change of sign of the force.

In Fig. 3 we display the lateral force at the initial time when the emitter is near a gold dielectric (\( \varepsilon_{Au} \approx 1.40 + 1.35i \)) or silica (\( \varepsilon_{Si} \approx 1.45 + 2.05 \times 10^{-7}i \)). The force shows a strong increment in the near-field regime, having a magnitude of around \( 10^{-21} \) Newton. The same orders of magnitude have been obtained for a circular emitter near a nanofiber [16]. To judge the strength of the effect we consider the lateral velocity acquired by the atom due to the recoil, which is in the nm/s range.

The atom in the ground-state can be excited also to the hyperfine level \( |6^2P_{1/2} \rangle, F = 5, M'_F = -5 \rangle \) by using a left-handed circularly polarised laser beam. The atom
we consider the emission spectrum:

\[ \bar{\Gamma}(z_A, \varphi) = A(z_A) + B(z_A) \cos \varphi + C(z_A) \cos^2 \varphi + D(z_A) \sin^2 \varphi, \]  

(20)

where:

\[ A(z_A) = -\frac{d^2}{4\pi^2 \varepsilon_0} \text{Re} \left[ \int_0^\infty dk^0 k^4 e^{2ik^0 z_A} \right], \]

\[ B(z_A) = \frac{d^2}{2\pi^2 \varepsilon_0} \text{Im} \left[ \int_0^\infty dk^0 k^3 e^{2ik^0 z_A} \right], \]

\[ C(z_A) = \frac{d^2}{4\pi^2 \varepsilon_0} \text{Re} \left[ \int_0^\infty dk^0 k^2 e^{2ik^0 z_A} \right], \]

\[ D(z_A) = -\frac{d^2}{4\pi^2 \varepsilon_0} \text{Re} \left[ \int_0^\infty dk^0 k^2 \frac{\sigma^2}{c^2} e^{2ik^0 z_A} \right]. \]

(21)

The emission is in general asymmetric along the x direction because of the factor B(z_A):

\[ \int_{-\pi/2}^{3\pi/2} d\varphi \bar{\Gamma}(z_A, \varphi) - \int_{-\pi/2}^{\pi/2} d\varphi \bar{\Gamma}(z_A, \varphi) = 4B(z_A) \]

\[ = \frac{2d^2}{\pi^2 \varepsilon_0} \text{Im} \left[ \int_0^\infty dk^0 k^2 \frac{\sigma^2}{c^2} e^{2ik^0 z_A} \right]. \]

(22)

The comparison with Eq. (14) shows that if the emission is asymmetric along the x-direction then the lateral force is finite.

In Fig. 4 and Fig. 5 we show the emission spectrum for two atom-plate distances: z_A = 264 nm, 302 nm. Figure 4 shows an asymmetric emission with stronger emission in the positive x direction; this suggests a negative lateral force as shown in Fig. 3. For z_A = 302 nm the emission spectrum is symmetric along the x-direction suggesting that the lateral force is zero, as shown in Fig. 3.

V. CONCLUSION

We have predicted a lateral Casimir-Polder force for a circular excited emitter near a planar surface. The underlying reason for a non-vanishing force is the breaking of time-reversal symmetry. The sign of the force depends on the polarization of the emitted light and can be controlled by changing the quantum excited state of the emitter. We have also shown that the lateral force is an atom recoil force stemming from the asymmetric emission of the photon.

Our dynamical approach of the field-atom coupling shows that the lateral force has a population-induced dynamics, decaying exponentially with time, on time scales
FIG. 4. Asymmetric emission spectrum for $z_A = 264$ nm along the $x$ direction.

FIG. 5. Symmetric emission spectrum for $z_A = 302$ nm.

of the inverse of the spontaneous decay rate. Moreover it exhibits characteristic Drexhage-type oscillations when changing the distance between the emitter and the surface. The near field regime is strongly influenced by the dissipative character of the medium: for short distances it converges to zero for a lossless medium, while it diverges if the medium has a complex dielectric permittivity. This effect could be detecting measuring the velocity acquired by the atom after the recoil.

Our formalism which uses the scattering Green’s tensor can be adopted for more complex geometries, like spheres, cylinders or resonating planar cavities which can enhance the spatial oscillations of the force.

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