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Coupled surface polaritons and the Casimir force

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The Casimir force between metallic plates made of realistic materials is evaluated for distances in the nanometer range. A spectrum over real frequencies is introduced and shows narrow peaks due to surface resonances (plasmon polaritons or phonon polaritons) that are coupled across the vacuum gap. We demonstrate that the Casimir force originates from the attraction (repulsion) due to the corresponding symmetric (antisymmetric) eigenmodes, respectively. This picture is used to derive a simple analytical estimate of the Casimir force at short distances. We recover the result known for Drude metals without absorption and compute the correction for weakly absorbing materials.

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I. INTRODUCTION

Van der Waals and Casimir forces are among the few macroscopic manifestations of vacuum fluctuations. Since the seminal paper by Casimir [1] showing the existence of an attraction between two perfect conductors separated by a vacuum gap, an abundant literature has been devoted to this effect. In particular, the relevance of retardation, finite conductivity, and finite temperature have been studied (see, e.g., [2]). Exhaustive lists of references can be found in several review papers such as [3, 4, 5].

In the last five years, the interest in Casimir forces has increased due to the existence of new measurements with improved accuracy [6, 7]. This has challenged theoreticians to quantify the corrections to the ideal case (zero temperature, perfect conductors, flat interfaces) that must be taken into account for an accurate comparison with experiments [8, 9, 10, 11, 12, 13].

Furthermore, the developments of micro-electromechanical systems (MEMS), for example, have shown that the Casimir effect is becoming an issue in nano-engineering [14, 15]. Indeed, these short-range forces could seriously disturb the performances of MEMS [16].

From a theoretical point of view, different methods exist to calculate Casimir forces. Casimir himself [1] determined the electromagnetic eigenfrequencies of the system and summed them in order to obtain the system's zero-point energy. The force is found by differentiation of this energy with respect to the geometrical distance separating the bodies [1, 17]. Ingenious subtraction procedures

are often required to obtain a finite value for the Casimir energy, and realistic dispersive or absorbing materials can be dealt with using contour integrals over complex frequencies [18]. Another method, used by Lifshitz [19], considers fluctuating currents driven by thermal or vacuum fluctuations in the whole space. These currents, whose spatial correlations are known through the fluctuation dissipation theorem, radiate electromagnetic fields. The force is obtained by calculating the flux of the Maxwell stress tensor across a surface separating the bodies. One thus gets an integral over all possible plane wave contributions. For two bodies separated by a vacuum gap, for example, the plane waves can be labelled by their frequency, wave vector parallel to the interface, and polarization. By using clever deformation contour methods, Lifshitz greatly simplified the calculation of the Casimir force integral. The principal drawback of this approach is that the integrand can no longer be interpreted as a force spectrum.

In this paper, we use an alternative approach and study the force integral over real frequencies and wave vectors. We show for generic materials (semiconductors and real metals) that in the near-field regime (separation distance small compared to the wavelengths considered), the frequency spectrum of the force exhibits peaks located close to surface-polariton frequencies. These peaks give the essential contribution to the Casimir force in this regime. We identify two types of resonant surface modes, binding and antibinding, that contribute respectively with an attractive and a repulsive term to the force. This substantiates early suggestions [20, 21] that the Casimir force is due to surface modes.

We finally focus on materials whose dielectric constant is modeled by a Drude formula, including a nonzero absorption. We are able to use the qualitative suggestions mentioned above and propose a quantitative estimation of the Casimir force in terms of coupled surface resonances. The dominant contribution of these resonances at nanometer distances allows to integrate exactly over the mode frequencies, whereas the integral over the

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wave vector is performed to first order in the absorption. We show that the respective contributions of the binding/antibinding modes give a simple and accurate analytical estimate for the short-distance Casimir force that coincides with previous results for nonabsorbing Drude materials [9]. For the first time, corrections due to material losses are incorporated. The paper concludes with a discussion of possibilities to “tune” the Casimir force that are suggested by our approach.

II. SURFACE RESONANCES IN THE FREQUENCY SPECTRUM

The starting point for our calculation of the Casimir force is Rytov’s theory of fluctuating electrodynamics in absorbing media [22] that has first been used by Lifshitz in his seminal paper [19]. This scheme applies to dispersive or absorbing materials, as long as their dielectric response is linear. It has also been shown to provide a suitable framework for a consistent quantization procedure of the macroscopic Maxwell equations (see [23] and references therein).

In the following, we focus on the standard geometry of two planar half-spaces made from identical material (of local dielectric constant $\varepsilon(\omega)$) and separated by a vacuum gap of width d . In the Rytov-Lifshitz method, the Casimir force is computed from the expectation value of the Maxwell stress tensor at an arbitrary position in the gap. At zero temperature and after subtraction of divergent contributions, Lifshitz gets a force per unit area given by [19]

$$F = \int_0^\infty \frac{d\omega}{2\pi} \int_0^\infty \frac{du}{2\pi} F(u, \omega) \quad (1)$$

$$F(u, \omega) = -\frac{2\hbar\omega^3}{c^3} \text{Im} uv \sum_{\mu=s,p} \frac{r_\mu^2(u, \omega) e^{-2\frac{\omega}{c}vd}}{1 - r_\mu^2(u, \omega) e^{-2\frac{\omega}{c}vd}} \quad (2)$$

where $v = (u^2 - 1)^{1/2}$ ($\text{Im} v \leq 0$), and r_μ is the Fresnel reflection coefficient for a plane wave with polarization μ and wavevector $K = \frac{\omega}{c}u$ parallel to the vacuum-medium interface. We use the convention that an attractive force corresponds to $F < 0$. We note that Rytov’s approach allows for an easy generalization to different media held at different nonzero temperatures. The radiation force on a small polarizable sphere above a heated surface has been discussed previously in [24]. Results for the non-equilibrium Casimir force will be reported elsewhere.

Lifshitz evaluated the integrals (1) by deforming frequency ω and wavevector u into the complex plane, to arrive at an integral over imaginary frequencies $\omega = i\xi$. The integration then requires the continuation of the dielectric function from real-frequency data to $\varepsilon(i\xi)$, using analyticity properties, see e.g. [9]. We follow here a different route and continue to work with real ω and u , taking advantage of the fact that Lifshitz’ results provides us with an expression for the frequency spectrum $F(\omega) = \int 4\pi^2 F(u, \omega) du$ of the Casimir force (as defined

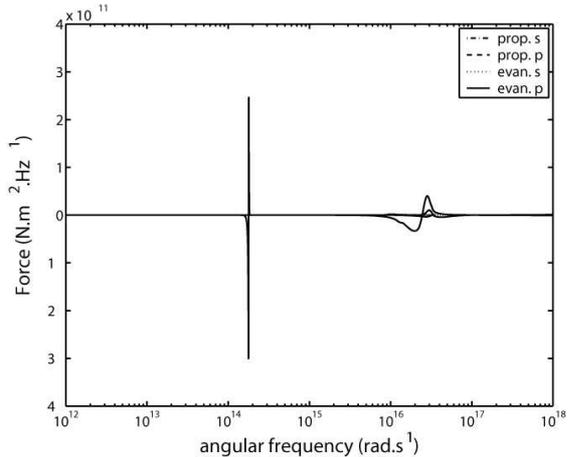


FIG. 1: Contributions of s and p polarized, propagating and evanescent modes to the force spectrum (Eq. (2)) integrated on the wavevector u . Distance $d = 10$ nm. Material: SiC, dielectric function with two resonances. The angular frequencies of the corresponding surface resonances are $1.78 \times 10^{14} \text{ s}^{-1}$ in the IR and $2.45 \times 10^{16} \text{ s}^{-1}$ in the UV [25]

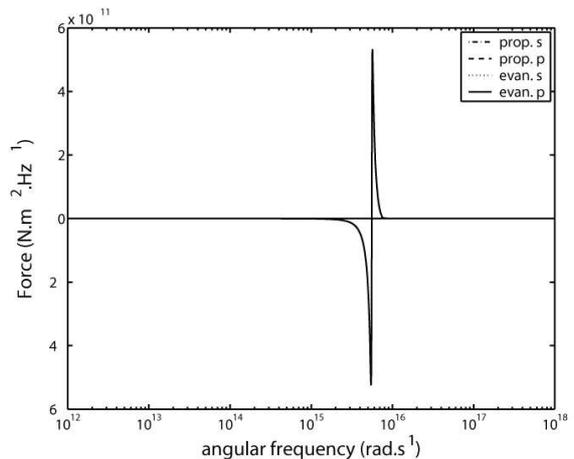


FIG. 2: Contributions of s and p polarized, propagating and evanescent modes to the force spectrum (Eq. (2)) integrated on the wavevector u . Distance $d = 10$ nm. Material: Chromium. Parameters of a Drude model (Eq. (7)): $\Omega = 5.57 \times 10^{15} \text{ s}^{-1}$ and $\gamma = 1.41 \times 10^{14} \text{ s}^{-1}$ [26].

in Eq. (2)). Note that the force spectrum is more difficult to define in a calculation based on the summation over modes (see, e.g., [27, 28]).

The spectrum of the force shows positive and negative peaks for very well defined frequencies (Fig. 1) when the distance d is reduced to the nanometer range. For a

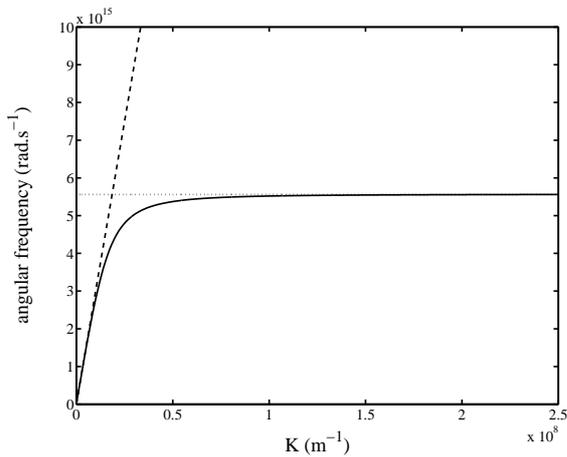


FIG. 3: Dispersion relation of the SPP on a flat interface vacuum/chromium. The real part of ω is represented versus the real part of the parallel wavevector $K = u\omega/c$.

polar material like SiC, the force is dominated by narrowpeaks in the UV and in the IR. These peaks can be ascribed to the surface phonon-polaritons in the IR and to surface plasmon polaritons in the UV. The largest contribution comes from the UV surface plasmon polariton even though it is broader due to larger losses than for the IR polariton. The large difference between the UV and the IR contributions in Fig. 1 is due to the factor ω^3 in Eq. (2). In Fig. 2, we plot the spectrum of the force between two chromium half-spaces. The dominant contribution to the force is clearly due to the surface plasmon polaritons. Indeed, the frequency of the peaks corresponds to the frequency Ω of the asymptote of the surface plasmon polaritons dispersion relation [29]:

$$u_{\text{SPP}} = \sqrt{\frac{\epsilon(\omega)}{\epsilon(\omega) + 1}}, \quad (3)$$

where the sign of the square root is chosen such that $\text{Re } u_{\text{SPP}} > 1$. It is seen in Eq. (3) that the frequency Ω is given by the condition $\text{Re } \epsilon(\Omega) = -1$. This corresponds to a large increase of the density of states and therefore of a peak of density of energy [30]. The dispersion relation suggests that the main contribution to the force comes from the surface waves with large wave vector u . We have separated the contribution according to the modes polarization (s or p). The modes in the cavity can also be separated into surface mode ($u > 1$) and propagating (guided) modes ($0 \leq u \leq 1$). The normal component of the wavevector is then either imaginary or real.

Among the four contributions (s or p, guided waves or surface waves) it is seen that the leading contributions comes from the p-polarized surface modes. This is a further indication that the leading contribution to the force is due to the SPP. It is worthwhile pointing out that if we were using a perfectly conducting metal, the spectrum of the force would be completely different because of the lack of SPP. Note also that the usual cal-

ulation yields the total contribution of all the modes called surface modes. Our detailed analysis shows that among all the modes, only the p-polarized surface modes contribute. These surface modes are ignored by the usual simple picture of the Casimir effect based on a modified density of states due to a quantification of the propagating states between two perfectly conducting planes. This simple picture includes only what we have called guided modes.

While the total force is always negative, we have noted in Fig. 1,2 that the contribution of the force is either positive or negative depending on the frequency. We shall analyse this behaviour in the next section.

III. BINDING AND ANTIBINDING RESONANCES

In order to further analyse the role of SPP in the Casimir force, we plot in Fig. 4a the integrand $F(u, \omega)$ as given by Eq. (2) for two chromium half-spaces separated by a distance of 10 nm. It is seen that the upper branch yields a positive contribution whereas the lower branch yields a negative contribution. These two branches are reminiscent of the dispersion relation of a SPP on a two interfaces system. It is given by the complex poles of the reflection factor of the two interfaces system in the (u, ω) plane.

$$1 - r_p^2 e^{-2\frac{u}{c}vd} = 0 \quad (4)$$

In order to illustrate the influence of the complex pole (i.e. of the SPP) on the force, we plot the quantity $1/|1 - r_p^2 e^{-2\frac{u}{c}vd}|^2$ in Fig. 4b in the real (u, ω) plane. Upon comparison between Fig. 4b and Fig. 4a, it is clearly seen that the main contribution to the force can be attributed to the SPP. In addition, we observe on Fig. 4b a dark line which corresponds to minima of $1/|1 - r_p^2 e^{-2\frac{u}{c}vd}|^2$. The minima can be attributed to very large values of the reflection factor of a plane interface r_p . Thus, the locus of the minimum is the dispersion relation of the SPP on a single flat interface.

In Fig. 5, we plot the force for a spacing $d = 100$ nm. It is seen that the two branches tend to merge with the flat interface dispersion relation. It is thus clear that the overlapping of the two SPP leads to a splitting of the polariton frequencies [31, 32]. The frequency splitting can be found by looking at the poles of the reflection coefficient for the cavity formed by the two interfaces [33]. This coefficient is precisely the integrand of Eq. (2) so that we get two (complex) poles given by

$$r_p(u, \omega) = \pm e^{\frac{u}{c}vd}. \quad (5)$$

The signs correspond to either symmetric or antisymmetric mode functions (for the magnetic field), as shown in Appendix A and sketched in Fig. 4b. The symmetric (antisymmetric) branch corresponds to a lower (higher) resonance frequency, respectively, similar to molecular orbitals and tunneling doublets [34].

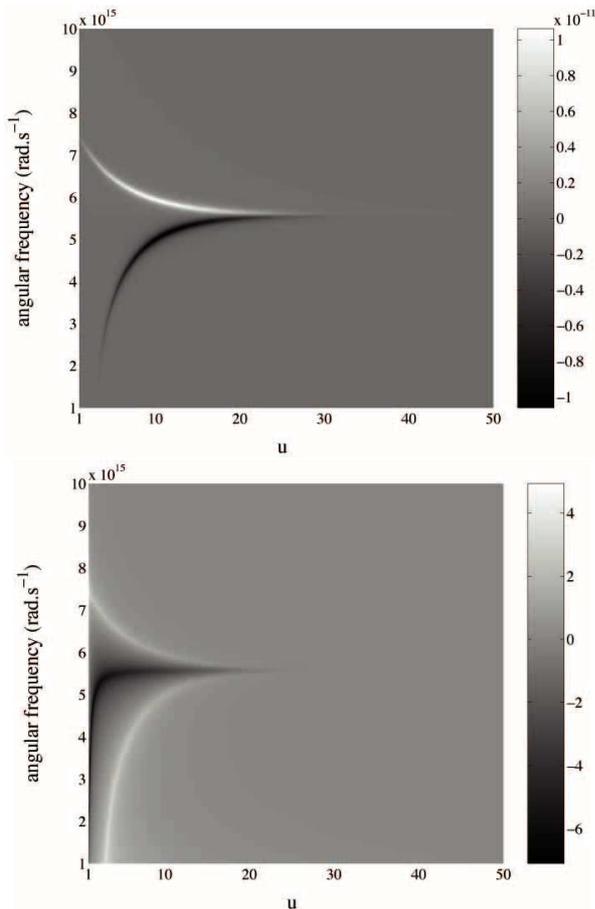


FIG. 4: (a) Wavevector resolved spectrum of the Casimir force (Eq. (2)) in the (u, ω) plane between two chromium half spaces separated by a distance of 10 nm. The frequency of the flat asymptote corresponds to the peaks of the force spectrum Fig. 2. (b) The mode functions corresponding to the polariton branches are sketched.

These branches contribute with opposite signs to the Casimir force, due to the following identity

$$\frac{2r_p^2(\omega, u)e^{-2\omega vd}}{1 - r_p^2(\omega, u)e^{-2\omega vd}} = \frac{r_p(\omega, u)e^{-\omega vd}}{1 - r_p(\omega, u)e^{-\omega vd}} - \frac{r_p(\omega, u)e^{-\omega vd}}{1 + r_p(\omega, u)e^{-\omega vd}}, \quad (6)$$

where the first (second) term is peaked at the symmetric (antisymmetric) cavity mode. The symmetry of the resonance mode function hence determines the attractive or repulsive character of its contribution to the Casimir force. We show in Appendix A by evaluating explicitly the Maxwell stress tensor, that *symmetric modes are binding* as in molecular physics. This behavior is already apparent in Fig. 4: the two surface-plasmon polariton branches in the u, ω plane contribute with opposite signs to the Casimir force.

We note that the splitting in Eq. (6) of the force spectrum gives meaningful results also after integration be-

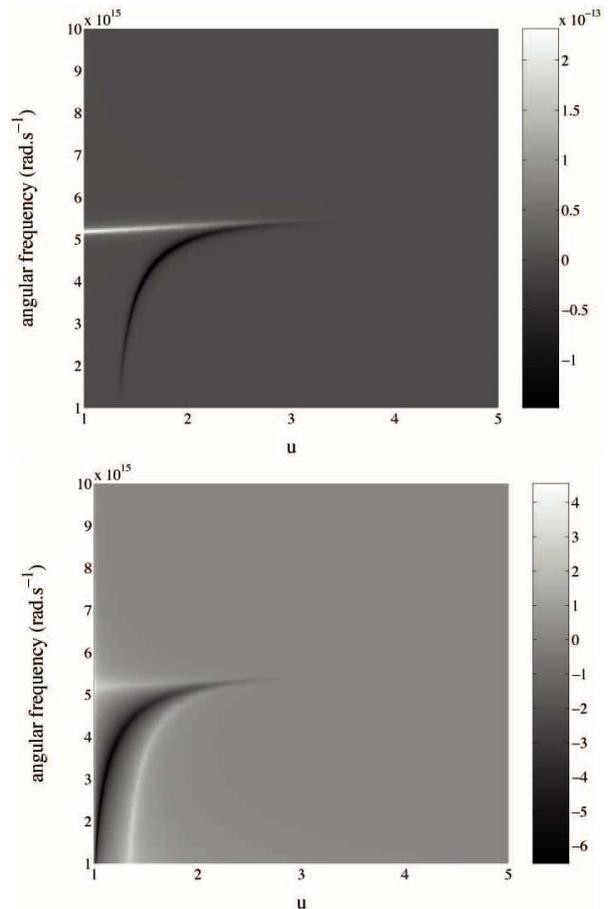


FIG. 5: (a) Wavevector resolved spectrum of the Casimir force (Eq. (2)) in the (u, ω) plane between two chromium half spaces separated by a distance of 100 nm. (b) The mode functions corresponding to the polariton branches are sketched.

cause for evanescent waves, both terms converge separately. We also point out that for a complex permittivity $\varepsilon(\omega)$ (as required by the Kramers-Kronig relations for a dispersive material), the polariton dispersion relation necessarily moves into the complex plane and never occurs on the real u - or ω -axis in the integral (1), thus excluding any singularities of the integrand.

IV. SHORT-DISTANCE LIMIT

The short-distance behaviour of the Casimir force between non-perfect metals has been computed in [9] using an integration over imaginary frequencies. We show here that these results can also be recovered with a real frequency calculation. In fact, we prove that the interaction between surface polariton resonances across the vacuum gap quantitatively accounts for the short-distance Casimir force derived in [9], thus completing qualitative discussions put forward by Gerlach [21] and the Lambrecht group [33].

For definiteness and for the ease of comparison with the

literature, let us focus on a dielectric function of Drude type

$$\varepsilon(\omega) = 1 - \frac{2\Omega^2}{\omega(\omega + i\gamma)} \quad (7)$$

where $\sqrt{2}\Omega$ is the plasma frequency and γ the damping coefficient. With this convention, the large u asymptote of the surface plasmon resonance (Eq. (3)) occurs at $\omega \approx \Omega$ ($\varepsilon(\Omega) \approx -1$). In the region of large wavevectors, the p-polarized reflection coefficient indeed has a pole there,

$$u \gg 1: \quad r_p(\omega, u) \approx \frac{\Omega^2}{\Omega^2 - i\gamma\omega - \omega^2}. \quad (8)$$

From Eq. (5), we thus get the following dispersion relation for the (anti)symmetric surface plasmon resonances, neglecting for the moment the damping coefficient γ :

$$\omega_{\pm}^2 \approx \Omega^2 (1 \mp e^{-\omega_{\pm} u d}), \quad (9)$$

where we have used $v \approx u$ for $u \gg 1$. For large u , we solve by iteration and find that $\omega_{\pm} \lesssim \Omega$. As announced above, the symmetric mode thus occurs at a lower resonance frequency.

To derive an analytical estimate for the Casimir force, we retain in Eq. (2) only the contribution of p-polarized, evanescent waves. Introducing the new variables ω and $x = \frac{\omega}{c}vd$, we get

$$F = -\frac{\hbar}{4\pi^2 d^3} \text{Im} \int_0^{\infty} d\omega \int_0^{\infty} x^2 dx e^{-x} \times \sum_{\lambda=\pm 1} \frac{\lambda}{r_p^{-1}(\omega, x/(\omega/c)d) - \lambda e^{-x}}, \quad (10)$$

where $\lambda = \pm 1$ corresponds to symmetric (antisymmetric) modes, respectively. The integral is dominated by the range $x \sim 1$ and $\omega \sim \Omega$. To leading order in $\Omega d \rightarrow 0$, we can thus use the asymptotic form of r_p valid for large u given by Eq. (8). Performing the integral over ω analytically and including damping to first order in γ/Ω yields

$$F = -\frac{\hbar\Omega}{4\pi d^3} \int_0^{\infty} dx x^2 e^{-x} \times \sum_{\lambda=\pm 1} \left(\frac{\lambda}{2\sqrt{1 - \lambda e^{-x}}} - \frac{\lambda\gamma}{2\pi\Omega(1 - \lambda e^{-x})} \right). \quad (11)$$

This result shows clearly that anti/symmetric modes give Casimir forces of opposite sign. The first term in the parenthesis can be computed by expanding the square root in a power series in λe^{-x} , leading to an infinite series given in [9, 33]. The second term, the correction due to damping, can be explicitly integrated in terms of the Riemann Zeta function ζ , so that we finally have

$$F = -\frac{\hbar\Omega}{4\pi d^3} \left(\alpha - \frac{\gamma\zeta(3)}{4\pi\Omega} \right), \quad (12)$$

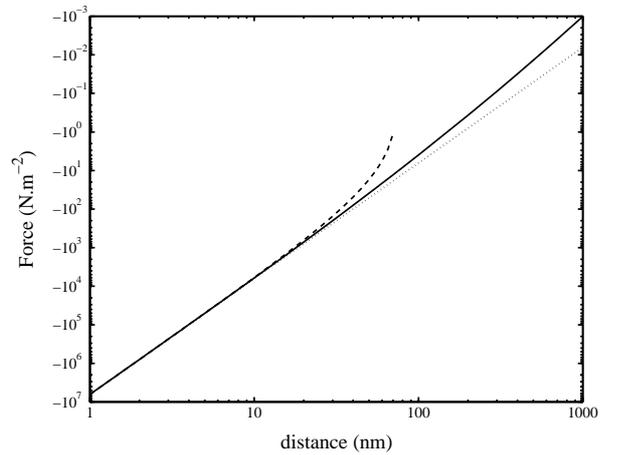


FIG. 6: Casimir force (2) (plain), p-polarization evanescent contribution (dashed) and analytical short-distance limit (12) for Cr (dotted).

where

$$\alpha = \frac{1}{4} \sum_{n=1}^{\infty} \frac{(4n-3)!!}{n^3(4n-2)!!} \approx 0.1388, \quad (13)$$

and

$$\zeta(3) = \sum_{n=1}^{\infty} \frac{1}{n^3} \approx 1.202. \quad (14)$$

Our result Eq. (12) for the short-distance Casimir force is compared in Fig. 6 to the full integral Eq. (2) and to the evanescent p-polarization contribution: it turns out to be quite accurate for distances $d \leq 0.03 \lambda_{\text{SPP}}$ where $\lambda_{\text{SPP}} = 338$ nm is the wavelength of the surface plasmon polariton with the largest frequency [35]. In the case of chromium, the first order correction in γ/Ω is 1.75% the zeroth order value of the force.

We finally note that the correction of order γ/Ω derived here must not be confused with the one for the finite conductivity of the media. Finite conductivity (i.e., a non-perfect mirror) is already taken into account by the finite value of the plasma frequency Ω and is responsible for the appearance of the short-distance regime where the Casimir force $\sim 1/d^3$ [19]. At large distances, a finite Ω leads to a small correction to the well-known Casimir force $\sim 1/d^4$ between perfect conductors [2, 9].

V. CONCLUSION

We have pointed out that the Casimir attraction between realistic materials can be quantitatively understood, at short distances, in terms of the interaction between electromagnetic surface plasmon (or phonon) polaritons. The modes overlap the vacuum gap and split into symmetric and antisymmetric combinations which

contribute with different signs to the Maxwell stress tensor and hence to the Casimir force. We derived an analytic formula that recovers previous results for nonabsorbing materials and evaluated a correction due to absorption. Our results have been validated by comparing to a numerical calculation based on Lifshitz theory.

The approach presented here has the advantage of linking in a transparent way the Casimir force to the actual physical properties of the material surface. This suggests the possibility of engineering the surface plasmon polariton dispersion relation to modify the Casimir force. Indeed, as it has been shown, the Casimir force at short distances is entirely due to the interaction between surface plasmon polaritons. Magnetic materials which exhibit Casimir repulsion and support s-polarized surface waves when $\text{Re } \mu < -1$ [37] are good candidates. The folding of the dispersion relation in the k-space by a grating, known to change the surface wave behaviour [38] could also lead to a substantial modification of the Casimir force.

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APPENDIX A: ANGULAR SPECTRUM ANALYSIS

In this appendix, we compute the Casimir force in terms of an angular spectrum representation of the electromagnetic fields that is adapted to the planar geometry at hand.

Letting the vacuum gap occupy the region $-d < z < 0$, we can expand the electric field in the gap as

$$\mathbf{E}(\mathbf{x}, \omega) = \sum_{\mu=s, p} \int d^2 K (E_{-}^{\mu}(\mathbf{K}) e^{-ik_z z} \mathbf{e}_{\mu}^{-} +$$

$$E_{+}^{\mu}(\mathbf{K}) e^{ik_z(z+d)} \mathbf{e}_{\mu}^{+}) e^{i\mathbf{K}\cdot\mathbf{x}} \quad (\text{A1})$$

where $\mathbf{K} = (k_x, k_y)$ is the component of the wavevector parallel to the interfaces and $k_z = \sqrt{(\omega/c)^2 - K^2}$ its perpendicular component. \mathbf{e}_{μ}^{\pm} ($\mu = s, p$) are the unit polarization vectors, and $E_{\pm}^{\mu}(\mathbf{K})$ are the amplitudes of down- and upwards propagating plane waves. A similar expansion holds for the magnetic field $\mathbf{H}(\mathbf{x}, \omega)$ with amplitudes $H_{\pm}^{\mu}(\mathbf{K})$. We get the averaged Maxwell stress tensor by integrating incoherently over the contributions $T_{zz}^{\mu}(\mathbf{K})$ of individual modes. For the particular case of a p-polarized evanescent mode ($K > \omega$), we get by straightforward algebra

$$T_{zz}^p(\mathbf{K}) = -2\mu_0 v^2 \text{Re} [H_{+}^{p*}(\mathbf{K}) H_{-}^p(\mathbf{K})], \quad (\text{A2})$$

The up- and downward propagating amplitudes are of course related via the reflection coefficient from the upper interface. Taking the phase references in Eq. (A1) into account, we have

$$H_{-}^p = r_p H_{+}^p e^{ik_z d} = r_p e^{-\frac{\omega}{c} v d} H_{+}^p \approx \pm H_{+}^p, \quad (\text{A3})$$

where the last equality applies in the vicinity of the coupled surface resonances defined by Eq. (5). The condition $r_p e^{-\frac{\omega}{c} v d} = +1$ thus corresponds to a symmetric magnetic field distribution on both interfaces, because $H_{+}^p = H_{-}^p$. In addition, with our sign convention, this mode gives an attractive contribution proportional to $-|H_{-}^p|^2$ to the stress tensor (A2). The opposite is true for antisymmetric modes.

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