Exact surface impedance formulation of the Casimir force: application to spatially dispersive metals

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We obtain exact expressions for the Casimir forces between arbitrary materials using the concept of surface impedance. We verify their consistency with the well known expressions for perfect conductors and with Lifshitz formula for semi-infinite local homogeneous media. As an application we present a full and rigorous calculation of the Casimir force between two metallic half-spaces described by a hydrodynamic non-local dielectric response.

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I. INTRODUCTION

In recent years experimental studies of Casimir vacuum forces between metallic surfaces have reached the necessary accuracy to test in detail the theoretical predictions [1–6] posed originally by Casimir in 1948 for perfectly-conducting parallel plates [7]. Indeed, detailed properties of the materials such as absorptivity, rugosity, or finite temperature effects [8–10] have become very relevant due to the accuracy reached in the experiments. The study of vacuum forces between realistic materials was pioneered by Lifshitz in 1956, who proposed a macroscopic theory for semi-infinite slabs described by a frequency ω dependent dielectric function $\varepsilon(\omega)$ [11]. Lifshitz formula, which reduces to the Casimir result in the limiting case of perfect conductors, has been successfully employed in a number of experimental situations. Different authors have elaborated alternative derivations of Lifshitz formula that permit a simpler, more transparent approach to vacuum forces in realistic materials [12]. Among several proposals to calculate the Casimir force, the impedance approach was employed for the first time by Mostepanenko and Trunov [13] to derive the Lifshitz formula in an approximate fashion. In a series of investigations of vacuum forces in realistic materials, both at zero and finite temperature, Klimchitskaya [14] et al. have found small discrepancies in the spatial behaviour of the forces when calculated according to the impedance or the Lifshitz approach. In those papers it is argued that the discrepancies arise from the approximate nature of the impedance concept. Indeed, within the particular context of the optics of metals, the surface impedance is usually derived by taking approximations valid only for good conductors below their plasma frequency, and it is interpreted in terms of induced surface currents [15]. However, such approximations are unnecessary and the concept of surface impedance can be straightforwardly applied to arbitrary materials [16–18]. In an equivalent framework, Kats [22] introduced the reflection amplitude coefficients r_a^α ($\alpha = s, p$) for metallic media to derive also an approximate version of Lifshitz formula to study the influence of non-locality on van der Waals interactions in a semi-quantitative way. Noticeably, he stated incorrectly that the reflection coefficient cannot be expressed merely in terms of the surface impedance for dielectrics. Recently, a more rigorous derivation of the Casimir force in terms of the reflection amplitudes was discussed by Reynaud and collaborators [23] using a *S*-matrix formalism.

The surface impedance Z of any planar surface may be *defined* as the ratio of the complex electric and magnetic tangential field components at the surface [16, 17]. For an *s*-polarized wave incident on a planar surface $z = z_0$ moving at an angle θ towards the positive z direction, the impedance boundary condition is

$$\mathbf{E}(z_0^-) = Z^s(\mathbf{H}(z_0^-) \times \hat{\mathbf{z}}), \tag{1}$$

while for a *p*-polarized wave, the corresponding definition is

$$Z^p\mathbf{H}(z_0^-) = \hat{\mathbf{z}} \times \mathbf{E}(z_0^-), \tag{2}$$

where $\hat{\mathbf{z}}$ is a normal vector pointing inside the surface, and z_0^- denotes a position immediately before the $z = z_0$ interface. A main advantage of this concept is that it relates only the tangential fields outside of the material, without the need of involving the internal degrees of freedom of the material, which are taken into account through the value of Z^s and Z^p. Equations (1) and (2) are functional definitions of the surface impedance, so that they are *exact* and valid, not only for perfect conductors, but also for real metals and insulators. Indeed, from these definitions the exact reflection coefficients for light impinging on a system (a) from vacuum (0) are

$$r_a^s = \frac{Z_a^s - Z_0^s}{Z_a^s + Z_0^s}, \quad r_a^p = \frac{Z_0^p - Z_a^p}{Z_0^p + Z_a^p}, \tag{3}$$

with Z_a^α and Z_0^α the corresponding surface impedances for $\alpha =$ s or p polarization. For a local homogeneous semi-infinite medium with dielectric response ε_a, $Z_0^s = q/k$, $Z_0^p = k/q$, $Z_a^s = q/k_a$ and $Z_a^p = k_a/\varepsilon_a q$ exactly, yielding the classical Fresnel coefficients [16].

$$r_a^s = \frac{k - k_a}{k + k_a}, \quad r_a^p = \frac{\varepsilon_a k - k_a}{\varepsilon_a k + k_a}, \tag{4}$$

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where $q=\omega/c$, ω is the frequency, c the speed of light, and $k_a=\sqrt{\varepsilon_a q^2-Q^2}$ and $k=k_0=\sqrt{q^2-Q^2}$ are the wavevector projections normal to the surface. However, the surface impedances are more general and are useful even for spatially dispersive systems, for which the Fresnel relations are not applicable [17, 18]. Spatial dispersion (or non-locality) is an important part of a full optical characterization of a system needed for a precise comparison between theory and experimental data [6]. The non-locality of the dielectric response means that the polarization induced at a given position $\mathbf{r}$ depends not only on the electric field applied at that point $\mathbf{r}$, but also at nearby points $\mathbf{r}'$. Thus, the non-local dielectric response of a material is characterized by a function of both space and time coordinates, $\varepsilon(\mathbf{r},\mathbf{r}',t-t')$, where $t-t'$ is the delay between the time t' at which an excitation is applied and the time t at which the response is observed. Manifestations of non-local behavior are the anomalous skin effects [19], additional bulk waves in semiconductors [20], excitonic effects [21], and plasma oscillations in conductors. Spatial dispersion effects add further energy-transport mechanisms that have an effect on the optical properties of materials.

In section II of this paper we obtain exact resuts for the Casimir forces valid for a wide class of materials using Eqs.(1) and (2). As a non-trivial application of the surface impedance approach, in section III we calculate the Casimir force between two semi-infinite slabs with a non-local dielectric response. Finally, we present our conclusions in section IV.

II. THEORY

Consider a system S consisting of two slabs $a=1,2$ parallel to the $x-y$ plane within free space and separated by a vacuum cavity V of length L along the z-direction, with inner boundaries at $z_1=0$ and $z_2=L$ as shown in Fig. 1. We assume that the slabs are non-chiral, translational invariant and isotropic within the $x-y$ plane, but otherwise they may be arbitrary. A given photon within V impinging upon a slab a may be reflected with a probability amplitude r_a^α which depends on its polarization α, acquiring a phase kL as it moves on to the other slab. Otherwise, it may be transmitted into the material with a probability $T_a^\alpha=1-|r_a^\alpha|^2$ where it can be absorbed, exciting electronic or vibrational degrees of freedom, or it can be transmitted into the vacuum beneath the slab, in any case, becoming lost forever (multiple reflections within the slab are implictly accounted for by the reflection amplitudes r_a^α). In thermodynamic equilibrium there would be photons coming from the outer vacuum and photons radiated by the materials themselves that would compensate exactly for the photons from V lost through absorption and transmission, appearing with a probability T_a^α with no definite phase relation to the lost photons. Thus, the equilibrium radiation within the cavity V depends exclusively on its geometry, characterized by L, and on its reflection amplitudes r_a^α. Therefore, if we construct an auxiliary fictitious system S' made of two infinitesimal sheets at $z=z_a$, and we postulate that their reflection amplitudes are given exactly by the same amplitudes as those of the original slabs, we assure that the radiation field within the fictitious cavity V' corresponds to the real one. We further assume that in S' a photon may be transmitted from V' into the vacuum outside with an amplitude t_a^α. By choosing $|r_a^\alpha|^2+|t_a^\alpha|^2=1$ we make certain that energy is conserved without having to account for any internal degrees of freedom of the fictitious sheets, even when the real system is dissipative.

We study first the case of s-polarized waves. The translational invariance implies that the wavevector projection $\mathbf{Q}$ onto the $x-y$ plane is a conserved quantity. Thus, the incident electric field of a normal mode can be written as $\mathbf{E}(\mathbf{r},t)=\mathbf{E_0}e^{i(Qx-\omega t)}\phi(z)$, with amplitude $\mathbf{E_0}=(0,E_y,0)$, where we have chosen, temporarily and without loss of generality, $x-z$ as the plane of incidence The magnetic field $\mathbf{B}=(B_x,0,B_z)$, is determined by Maxwell equations: $-iqB_x=\partial_z E_y$, and $qB_z=QE_y$. The field component E_y satisfies a one-dimensional wave equation within V'. The solution satisfying the boundary condition (1) at the interface of slab 2 ($z=L^-$) is

$$E_y^>(z)=e^{ik(z-L)}+r_2^s e^{-ik(z-L)}. \tag{5}$$

The corresponding solution satisfying the boundary condition at slab 1 ($z=0^+$) is

$$E_y^<(z)=e^{-ikz}+r_1^s e^{ikz}, \tag{6}$$

where the reflection amplitudes r_a^α are given by Eqs.(3) and k^2 plays the role of eigennumber.

The *electric* Green's function can now be calculated as

$$G_{k^2}^E(z,z')=\frac{E_y^<(z_<)E_y^>(z_>)}{W}, \tag{7}$$

where $z_<$ and $z_>$ are the smaller and larger of z and z', respectively, and W is their Wronskian. The magnetic Green's function is obtained by replacing $E_y\to B_x$ and $r_a^s\to -r_a^s$ in Eqs. (5)-(7). We do not consider B_z separately, as it is simply proportional to E_y. Therefore, for each $\mathbf{Q}$, the local density of states per unit k^2 is given by [24]

$$\rho_{k^2}^s(z)=-\frac{1}{2\pi}\mathrm{Im}\left(G_{\tilde{k}^2}^E(z,z)+G_{\tilde{k}^2}^B(z,z)\right), \tag{8}$$

with ($\tilde{k}\equiv k+i0^+$). By substituting Eqs. (5),(7) and its magnetic analogues we obtain

$$\rho_{k^2}^s=\frac{1}{2\pi\tilde{k}}\mathrm{Re}\left[\frac{1+r_1^s r_2^s e^{2i\tilde{k}L}}{1-r_1^s r_2^s e^{2i\tilde{k}L}}\right], \tag{9}$$

which is independent of z. Given the symmetries of the problem, the p-polarization density of states $\rho_{k^2}^p$ may be derived similarly, by replacing $B_x\to -E_x$, and $E_y\to B_y$. The final expression is simply given by Eq.(9) after replacing all the superscripts $s\to p$. Finally, the total density of states is $\rho_{k^2}=\rho_{k^2}^s+\rho_{k^2}^p$.

Take now given values of α and $\mathbf{Q}$ and consider a photon in a state characterized by k^2. Its momentum and velocity along the z direction are $\pm\hbar k$ and $\pm ck/q$, where the upper sign applies if it moves towards slab 2 and the lower if it moves towards slab 1. In either case, its contribution to the

momentum flux towards the slab 2 is $\hbar c k^2/q$. Denoting by $f(k) = N_(k) + 1/2$ the photon occupation number of state k, the total momentum flux is obtained by summing over k^2, $\mathbf{Q}$ and α, yielding

$$\sum_{\alpha} \int Q dQ dk^2 \hbar \frac{k^2}{q} f \rho^{\alpha}_{k^2}. \tag{10}$$

There is a similar contribution coming from the semi-infinite vacuum on the other side of the slab, obtained by substituting $r_2^{\alpha} \to 0$ above and reversing the flux direction $z \to -z$. The total force per unit area is obtained by adding the contributions from both sides,

$$\frac{F(L)}{A} = \frac{\hbar c}{2\pi^2} \int_0^{\infty} dQ Q \int_{q \geq 0} dk \frac{k^3}{q} f \mathrm{Re} \frac{1}{\tilde{k}} \left[\frac{1}{\xi^s - 1} + \frac{1}{\xi^p - 1} \right], \tag{11}$$

with $\xi^{\alpha} = (r_1^{\alpha} r_2^{\alpha} \exp(2i\tilde{k}L))^{-1}$. The integral over k runs from iQ to 0 and then to ∞, so that q remains real and positive. This expression depends only on the reflection coefficients or, equivalently, on the surface impedances of the system and the slab separation. Lifshitz formula is recovered upon substitution of the *local* Fresnel amplitudes Eqs.(4), whereas for perfect mirrors ($r_a^{\alpha} = \pm 1$) Eq.(11) yields the expected Casimir expression.

III. PLASMA OSCILLATIONS IN CONDUCTORS

A non-trivial application of the impedance approach is the calculation of the Casimir force between two spatially dispersive conductors. Kats [22] studied non-local effects in an approximate way, and as he stated, it is necessary to specify correctly the dependence of the dielectric function on the wave vector and frequency. Furthermore, he studied the spatial dispersion due to the anomalous skin effect. Here, we are interested in the non-local screening effects due to excitation of plasma oscillations in metallic systems. In a conductor, the normal component of an incident *p*-polarized wave pushes the conduction charge towards or away from the surface creating an excess charge. When the frequency of the electromagnetic wave is above the plasma frequency ω_p, this charge propagates as a longitudinal wave (plasmon). Thus, at a surface, *p*-polarized waves couple to bulk plasmons which carry away energy from the inicident wave, reducing the reflection amplitudes. Thus, conductors support both longitudinal and transverse oscillations.

Translation invariance within the bulk of homogeneous systems implies that its non-local dielectric function ε depends on the separation $\mathbf{r} - \mathbf{r}'$ between the excitation and observation positions $\mathbf{r}$ and $\mathbf{r}'$ only, and not on $\mathbf{r}$ and $\mathbf{r}'$ separately. Thus, the Fourier transformed response $\varepsilon(\mathbf{k}_a, \omega)$ depends on the wavevector $\mathbf{k}_a$ and frequency ω of the field within the medium. Within an isotropic system there is no prefered direction beyond that of $\mathbf{k}_a$, so that the non-local response may be characterized by two scalar functions, the longitudinal $\varepsilon_l(\mathbf{k}_a, \omega)$ and transverse $\varepsilon_t(\mathbf{k}_a, \omega)$ dielectric functions, describing the response to longitudinal $\mathbf{E}_l$ and transverse $\mathbf{E}_t$ fields, parallel and perpendicular to $\mathbf{k}_a$ respectively. The total field is the sum of the longitudinal and transverse parts $\mathbf{E} = \mathbf{E}_l + \mathbf{E}_t$. Taking $x - z$ as the plane of propagation, within a metal we can have transverse waves with wave vectors $\mathbf{k}_t = (Q, 0, k_t)$ satisfying the dispersion relation $k_t^2 = \varepsilon_t(\mathbf{k}_t, \omega) q^2 - Q^2$ and longitudinal waves with wavevectors $\mathbf{k}_l = (Q, 0, k_l)$ that obey the dispersion relation $\varepsilon_l(\mathbf{k}_l, \omega) = 0$. The former relation is consequence of the electromagnetic wave equation while the latter may be obtained from Gauss law in the absence of external charges. From our discussion above, we expect *p*-polarized waves to couple to plasmons at surfaces, where translational invariance is lost. Notice that in order for plamons to propagate, i.e., for their frequency to depend on their wavevector, the longitudinal dielectric function has to be spatially dispersive. On the other hand, transverse waves do not require spatial dispersion to propagate. Thus, to illustrate the effects of plasmon excitation at surfaces it is enough to choose a non-local $\varepsilon_l(\mathbf{k}_l, \omega)$ and a local transverse dielectric function $\varepsilon_t(\omega)$.

The most simple model of the response of a metal exhibiting spatial dispersion and coupling to plasmons at surfaces is the hydrodinamic model[25]. In it, the transverse dielectric function is chosen as a local Drude function $\varepsilon_t = 1 - \omega_p^2/(\omega(\omega + i\gamma))$, where ω_p is the plasma frequency and γ the damping. For the longitudinal part we consider a hydrodynamic dielectric function [25] given by

$$\varepsilon^l(\mathbf{k}_a, \omega) = 1 - \frac{\omega_p^2}{\omega^2 + i\omega\gamma - \beta^2(Q^2 + k_a^2)}, \tag{12}$$

where $\beta^2 = 3v_f^2/5$ with v_f the Fermi velocity of the metal [25]. In this model, the compressibility of the electron gas β is responsible for the spatial dispersion [26].

To calculate the surface impedance for *p*-polarized waves we have to account for the non-homogeneity of the surface. Since we now have an additional longitudinal field, this can be done most simply by postulating an additional boundary condition (ABC) [27], not derivable from Maxwell equations. Charge conservation and the physical impossibility of infinite charge densities suggest the ABC $j_{zl} + j_{zt} = 0$, i.e., the normal component of the total current should vanish at the surface. Although much more realistic models of a metallic surface which don't require ABC's are available [28], our choice is adequate to illustrate the effects of spatial dispersion. Within our model, the surface impedance is simply

$$\begin{aligned} Z_p &= \frac{E_{x\perp} + E_{x||}}{B_y} \\ &= \frac{k_t c}{\varepsilon_t \omega} + \frac{(\varepsilon_t - 1) Q^2 c}{\varepsilon_t k_l \omega}, \end{aligned} \tag{13}$$

and substitution into (3) yields

$$r_p = \frac{\varepsilon_t k - k_t - (\varepsilon_t - 1) Q^2/k_l}{\varepsilon_t k + k_t + (\varepsilon_t - 1) Q^2/k_l}, \tag{14}$$

while r_s agrees with its local counterpart [Eq.(4)].

The Casimir force between two semi-infinite conductors described by the HD model can now be calculated again using Eq. (11). In Fig. 2 we present the normalized difference

between the Casimir force for the local case(F_L) with that obtained with non-local effects (F_{NL}), i.e. $\Delta = (F_L - F_{NL})/F_L$ calculated at zero temperature ($f = 1/2$) for three different metals (K, Au, Al) as a function of the separation between their surfaces in units of the plasma wavelength of the metal $L/\lambda_p = \omega_p L/2\pi c$. The parameters for each metal (ω_p, γ, v_f) are taken from the literature [29, 30]. As expected the non-local effects become important for separations of the order of the plasma wavelength λ_p, and the difference between them can be significant for separations less than λ_p. For these separations the vacuum modes that contribute mainly to the Casimir force have a frequency larger than ω_p thus exciting propagating modes withinin the metals. Furthermore the better the conductor the higher the difference between the local and non-local cases.

The disagreement reported between the Casimir forces associated to the Lifshitz and the impedance approach Refs.[14] can be traced back to the introduction of an approximate expression for the impedance. In the original work of Mostepanenko and Trunov [13], they employed an expression for the impedance $Z(\omega) \approx 1/\sqrt{\varepsilon(\omega)}$, valid only for small values of the electromagnetic wavevector $\mathbf{Q}$, parallel to the surface of incidence. However, by introducing the definitions (1) and (2) in the same formalism, then the *exact* expression for the Casimir force (Eq.(11)) is obtained by just replacing $Z \to Z^p$ in Eq.(41) of Ref.([14]), and $Z \to Z^s$ in Eq.(42) of the same reference.

IV. CONCLUSIONS

We have shown that with the correct definitions for the surface impedances, an exact expression for the Casimir force is obtained. This result is very general as no assumption about the nature of the system was made, beyond that of translational invariance and isotropy along the surfaces. Thus, it may be applied to homogeneous semi-infinite systems, finite layered heterostructures, metals, insultators, etc. As a case study, we calculated the surface impedance and Casimir force between homogeneous semiinfinite conductors with a non-local optical response. We employed an approximate hydrodinamic model for the dynamical response of the metals, as it is the simplest one that describes the coupling at surfaces of *p*-polarized light to bulk plasmons. Within this model, non-local effects arise from Pauli's principle, which yields a finite compressibility of the conduction electron gas due to the finite velocity of its electrons even at zero temperature, of the order of the Fermi velocity $v_f \approx 10^{-2}c$. Correspondingly, we found that non-local effects become relevant for separations of the order of the plasma wavelength of the slab, for which they produce corrections of the order of 1%, and they grow substantially as the separation is further diminished. The non-local effects allow plasmons to take away energy from the incident electromagnetic wave, reducing the values of the reflectance and of the Casimir force. Correspondingly, a local calculation predicts a higher value of the force. As seen in recent experiments, it is at close separations that the difference between theory and experiment become larger. The hydrodinamic model employed here ignores many aspects of the dynamics of electrons close to metal surfaces such as the presence of Landau damping through the decay of plasmons into electron-hole pairs, the fact that the self-consistent electronic profile is smooth, and that the dynamic electronic density peaks outside instead of within the nominal metal surface. However, these additional effects may be accounted for within our formalism by replacing within our formulae (Eqs. (3) and (11) the approximate hydrodinamic expressions for the surface impedances (13) by more accurate expressions [28]. This constitutes work in progress.

In summary, we have derived expressions for the Casimir force of a very large classs of systems employing an exact surface impedance formalism. Our expressions allowed us to make a calculation of nonlocal effects in the Casimir force between metals. We used a simple hydrodinamic model whichs accounts qualitatively for the excitation of plasmons by electromagnetic fields at surfaces. We found that non-local corrections to the Casimir force are of the order of 1% at separations $L \sim \lambda_p$ and increase as L decreases, so they should not be ignored at small distances.

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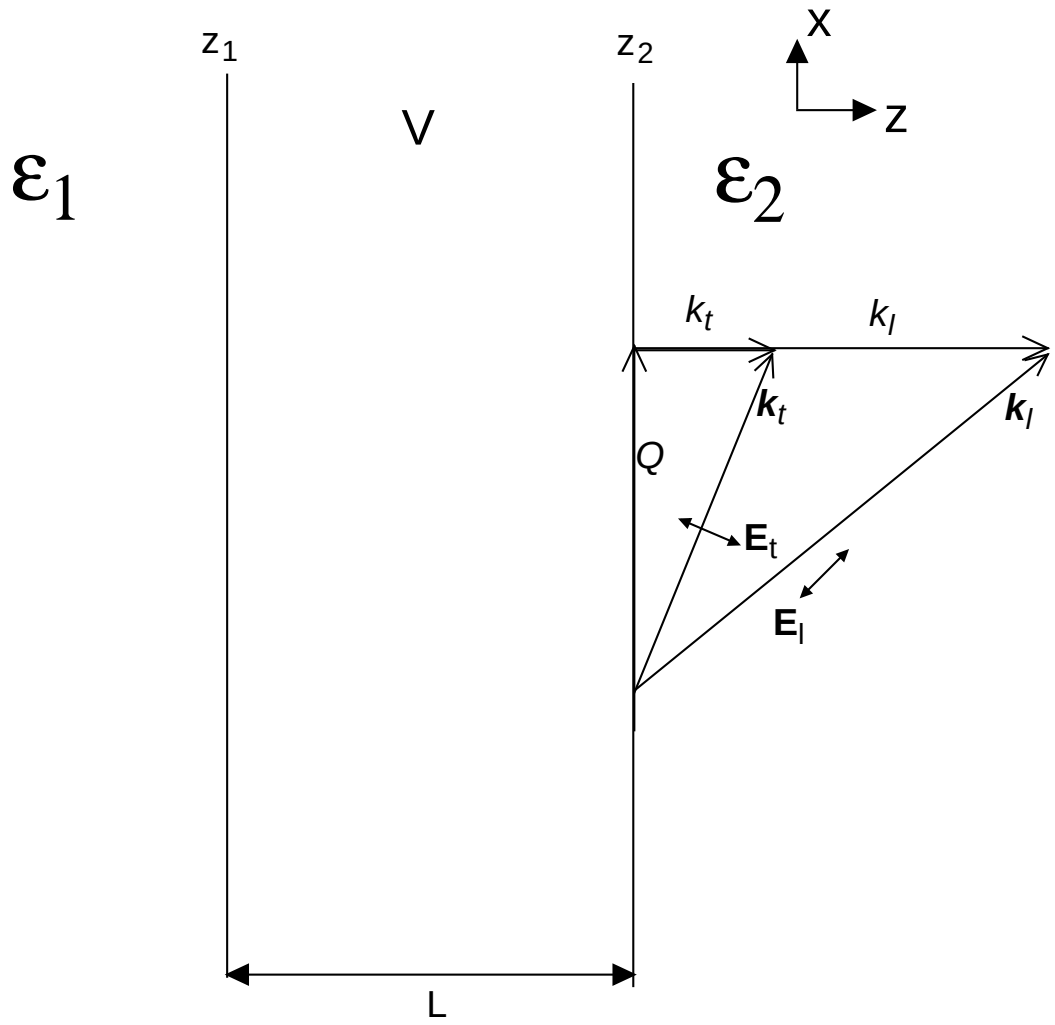

FIG. 1: Cavity V of width L between two metals with dielectric response functions ε_1 and ε_2 with surfaces at z_1 and z_2. Transverse modes with wavevector $\mathbf{k}_t = (Q, 0, k_t)$ and p-polarized electric field $\mathbf{E}_t$ are coupled at the surface of non-local conductors to longitudinal modes with field $\mathbf{E}_l$ and wavevector $\mathbf{k}_l = (Q, 0, k_l$.

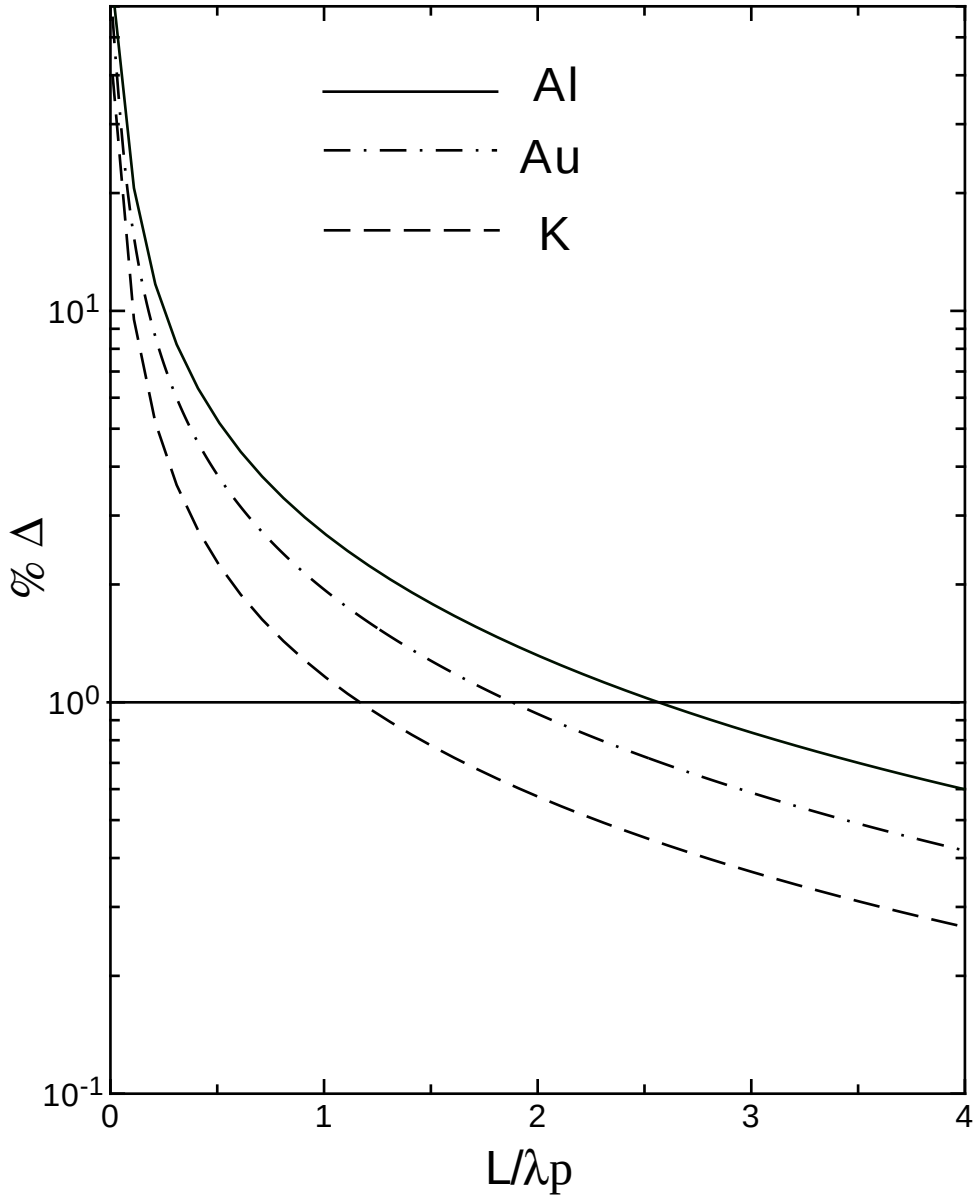

FIG. 2: Normalized difference $\Delta = (F_L - F_{NL})/F_L$ between the Casimir force in the local and the non-local cases. The separation between the slabs is normalized to the plasma wavelength λ_p of the metals. The curves correspond to Au, Al and K. The horizontal line is a visual aid to indicate where both calculations differ by 1%.