Modify Plasmon Modes in a Cavity to Create Repulsive Casimir /vdw Forces

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Abstract

The Current paper describes the influence of surface plasmons on the Casimir effect between two uncharged, conducting parallel plates at arbitrary separating distances. This quantum electrodynamics (QED) phenomena has renewed the interest for application to nanomechanical devices for the last decade. In order to write Casimir energy, two sorts of modes can be identified: evanescent surface modes and propagative modes. Using the plasma model to describe the optical response of the metal, we can express the Casimir energy as the sum over these two modes. In contrast to ordinary propagative cavity modes (photonic modes) which have a real longitudinal wavevector, the evanescent modes (plasmonic modes) have an imaginary longitudinal wavevector.

The contribution of the plasmonic modes to the Casimir energy will be analytically evaluated. This contribution is essential not only at short but also at large distances. One of the two plasmonic modes gives rise to a repulsive contribution. The Casimir interaction changes from attractive to repulsive with distance. Hence, it becomes repulsive for intermediate and large distances. Recent experiments have been able to detect and accurately measure this effect, However, for the first time, a repulsive force was seen.

Contents

I. INTRODUCTION

"I mentioned my results to Niels Bohr, during a walk. That is nice, he said,

that is something new. I told him that I was puzzled by the extremely simple

form of the expressions for the interaction at very large distances and he mumbled

something about zero-point energy. That was all, but it put me on a new track."

-H. B. G. Casimir (Milonni 1994)

Dutch physicist Hendrik B. G. Casimir (figure 1) proposed the existence of a force between neutral mirrors in vacuum in 1948 ([1], [2] and [3]). While Casimir was employed at Philips Research Laboratories in Eindhoven, the Netherlands, he showed that two perfectly conducting parallel plates attract each other. The interaction between two electrically neutral parallel plates placed a few micrometers apart without an external electromagnetic field, is dominated by the force arising from quantum fluctuations of the electromagnetic field, known as the Casimir force. Since electromagnetic fluctuations depend on the dielectric function of the surfaces, different materials might be used to reveal new aspects of the Casimir force and suggest novel solutions for the design of micro and nanofabricated devices [4,5]. In quantum field theory the Casimir effect is a physical force arising from a quantized field. From classical point of view, the lack of an external filed means that there is no field between plates and no force would be measured between them. So, one needs to consider quantum electrodynamics (QED) for this remarkable prediction. In the quantum electrodynamics description it is seen that the plates do affect the virtual photons which constitute the field and generate a net force. This force could be either attractive or repulsive, depending on the specific arrangement of the two plates. According to current descriptions of a quantum vacuum, all fields in a vacuum are never completely eliminated, especially the electromagnetic fields. This allows the system to have a certain amount of energy that is always present, commonly referred to as the zero-point or vacuum energy. Casimir effect has been witnessing a renewed interest because of its application in micro- and nanotechnology. At the moment, many theorists and experimentalists are still intrigued by the Casimir effect and recent advances in both precision measurements of the Casimir force and the downscaling of devices, has resulted in many publications and a growing interest in this topic. Also, more people are working on possible applications of the Casimir force.

In the past few years, the interest in surface plasmons has increased due to the connection with a broad range of topics such as near-field spectroscopy, sub wavelength resolution [6,7] and extraordinary optical transmission through sub wavelength metallic hole arrays [8,9,10]. The contribution of surface plasmons to the Casimir energy will be evaluated in section 4. This contribution is essential not only at short distances but also at large distances. It should be mentioned that surface plasmons give a large contribution at large distances, and even a repulsive contribution [11-13]. We investigate the role of surface plasmons on the Casimir effect. The role of surface plasmons is important in many fields of physics, such as plasmon-associated light transmission through metallic structures [14,15,16] or dispersion forces between electronic Wigner crystals that are relevant for biomolecular physics [17].

The Casimir force could be explained by the radiation pressure due to the quantum fluctuations of the quantized electromagnetic field when in the presence of boundaries. This is illustrated in figure 2.

Because only the electromagnetic modes that have nodes on both walls can exist within the cavity, the zero-point energy depends on the separation between the plates, giving rise to an attractive force [18,19]. Calculating the Casimir force for perfect conductors is simple, However for real materials it is more complicated. In 1956, the Russian physicist Evgeny Lifshitz extended Casimir's formulation to a more general treatment of the problem for real materials described by their dielectric properties outside a vacuum [20]. This led to what we know now as the Casimir-Lifshitz forces. Later on, the Casimir effect was confirmed experimentally. In 1978, Blokland and Overbeek presented strong experimental evidence for the Casimir effect. Later on, K. Lamoureaux reported the first high precision measurements of the Casimir force in 1996 [21]. He presented an experiment analogous to Casimir's, where he used a spherical conductor and a flat plate conductor, thus altering the force equation. This adjustment allows modifying the dependence on the plate's surface area to be eliminated. It was followed by several other experimental studies, which have produced evidence further supporting the Casimir effect. A similar experiment was conducted by Jeremy N. Munday and colleagues in 2009, However, the aim of their study was to verify that the Casimir force not only exhibits an attractive but also a repulsive force [22].

Fig 1. Hendrik "Henk" Brugt Gerhard Casimir (1909-2000)

II. Casimir Force

In fact Casimir's original goal was to compute the van der Waals force between polarizable molecules of the metallic plates. He [3] noticed that the energy should be a finite expression. So, he summed the electromagnetic eigenfrequencies of the system in order to obtain the system's zero-point energy. The force is calculated by a differentiation of this energy with respect to the geometrical distance separating the bodies. He considered an ideal setting with perfectly reflecting mirrors in vacuum.

The corresponding interaction energy E depends only on geometrical quantities, distance between the mirrors L and their surfaces $A \rangle\rangle L^2$,and two fundamental constants, the speed of light c and reduced Planck constant ħ.

$$E = E_{cas} = -\frac{\hbar c \pi^2 A}{720 L^3} \qquad (1)$$

According to the thermodynamics, a negative energy corresponds to a binding energy. This remarkable universal feature corresponds to the fact that the optical response of perfect mirrors is saturated : mirrors cannot reflect more than 100 % of the incoming light. However, the experiments are performed with real reflectors, typically metallic mirrors, as metals are good reflectors at frequencies much smaller than the plasma frequency ω_p or at wavelengths much larger than $\lambda_p = 2\pi c / \omega_p$.

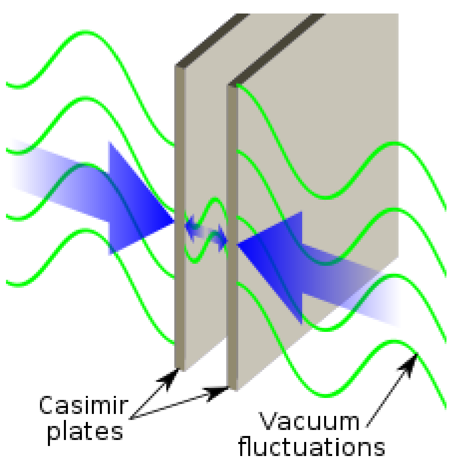

Fig. 2. Inside the cavity less electromagnetic modes can exist than in the vacuum.

III. Surface Plasmon modes

We investigate the influence of surface Plasmon modes on the Casimir effect. We study this role between two plane parallel metallic mirrors at arbitrary distances and write the Casimir energy as a sum of contributions associated with evanescent surface Plasmon modes and propagative cavity modes. By choosing some parameters, one of the surface Plasmon modes becomes propagative sector and it is counted with the surface Plasmon contribution. This contribution is called plasmonic. To achieve a correct result for the Casimir energy, the contribution of plasmonic modes is essential at all distances. In contrast to propagative cavity modes which give rise to an attractive contribution, one of the two plasmonic modes gives rise to a repulsive contribution. This contribution balances out the attractive contributions. As the mirror separation is increased, the combined plasmonic contribution to the Casimir energy changes sign [23,24].

Lifshitz recovered expression (1) for metallic mirrors with separations much larger than the plasma wavelength λ_p associated with the metals. For short separations $L\langle\langle\lambda_p$, the Casimir energy can be expressed in terms of the coulomb interaction between surface plasmons. In 1968, Van Kampen and co-workers [25] showed that in the limit of small separations $L\langle\langle\lambda_p$, indeed the Casimir effect is dominated by the coulomb interaction between the surface plasmons. This includes all of the electron excitations propagating on two metallic mirrors. The corresponding electromagnetic field modes are evanescent inside the cavity. In contrast to ordinary propagating cavity modes, which have a real longitudinal wavevector, plasmonic modes have an imaginary longitudinal wavevector [25,27,28].

Van Kampen and co-workers computed the Casimir energy in the case of small separations, $L\langle\langle\lambda_p$. In this limit the Casimir energy becomes[2]

$$E \approx \alpha \frac{L}{\lambda_p} E_{cas} \quad \text{with } \alpha = 1.790 \tag{2}$$

Which is reduced with respect to Eq. (1) [29,30]. The formula shows that the energy is dependent on the material parameter, λ_p .

At large separations, retardation has to be taken into account. In the limit of short distances, plasmonic modes dominate the interaction. But, they do not vanish for large separations. They give rise to a contribution which has a negative sign simultaneously in separations larger than $\lambda_p/4\pi$. The magnitude of this contribution is too large with respect to expression (1).

For simplicity in quantum field theory of the Casmir effect, ordinary propagating modes are called photonic modes and they determine the Casimir effect at large distances.

IV. The Role of surface plasmon modes in the Casimir effect

The Casimir free energy of two metallic plates is obtained by summing the zero point energies over the electromagnetic modes vibrating inside the cavity.

Here, the restriction of two infinitely large plane mirrors at zero temperature is considered and the Casimir formula (1) is modified by considering the metals finite conductivity. This modification is done by evaluating the radiation pressure of a vacuum field upon the two mirrors [13].

$$E = -\sum_{\mu}\sum_{k}\sum_{\omega}\frac{i\hbar}{2}\ln\left(1 - r_k^{\mu}[\omega]^2 e^{2ik_zL}\right) + c.c \qquad (3)$$

The energy E is calculated by summing over the polarization $\varepsilon = (TE, TM)$ and the wavevector $k \equiv (k_x, k_y)$ parallel to the mirrors and with frequency ω; k_z is the perpendicular wavevector associated with the mode. The components of the type r_k^{μ} ,corresponds to the reflection amplitudes, that which are assumed to be the same for both mirrors. They are casual retarded functions obeying high-frequency transparency.

Now, Casimir energy can be calculated as a sum over the cavity modes by using the plasma model for the mirrors dielectric function

$$\varepsilon[\omega] = 1 - \frac{\omega_p^2}{\omega^2} \qquad (4)$$

where ω_p signifies the plasma frequency and $\lambda_p = 2\pi c / \omega_p$ the plasma wavelength for metals used in modern experiments. λ_p lies in the sub-micron range (107 nm for Al and 137 nm for Cu and Au) [13].

In this case the zeros of the argument of the integrand in (3) lie on the real axis. Now, (3) can be rewritten as the sum over the solutions $[\omega_k^{\mu}]_n$ of the equation labeled by an integer index n

$$r_k^{\mu}[\omega]^2 e^{2ik_zL} = 1 \qquad (5)$$

Simple algebraic manipulations exploiting residues theorem, combined with complex integration techniques then lead to the Casimir energy expressed as sums over these modes (6).

$$E=\sum_{\mu,k}\left[\sum_{n}{}^{'}\frac{\hbar}{2}\omega_n^{\mu}(k)\right]_{L\to\infty}^{L}=\mathrm{Im}\sum_{\mu,k}\int_0^{\infty}\frac{d\omega}{2\pi}\ln\left(1-r_k^{\mu}[\omega]^2e^{2ik_zL}\right)\tag{6}$$

The prime in the sum over n signifies that the term m=0 has to be multiplied by $1/2$.

The upper mathematical identity was proven by Schram in 1973. The left-hand side corresponds to Casimir's sum over the zero point energies, but in this case the relevant modes are those of the real cavity. The notation $[...]_{L\to\infty}^{L}$ expresses the difference for finite and infinite mirror distance L. The right hand side corresponds to the Lifshitz formula for the Casimir energy. In 1955, Lifshitz used a different method and obtained the force by computing the expectation value of the Maxwell stress tensor inside the cavity [4]. He considered the electromagnetic fields as being radiated by fluctuating sources in the medium composing the mirrors, similar to the Van der Waals force between atoms and molecules which was derived by London.

Lifshitz simplified the calculation of the Casimir force integral by considering that mirrors are non-dissipative. He showed that this approach gives the same result as the Casimir sum over the zero-point energies.

Expression (6) has its limitation though. For large distances it requires perfect mirrors, whereas for short distances surface plasmon resonances. For arbitrary distances, not only the plasmonic modes, but also the photonic modes are important. The mode frequencies $\omega_n^{\mu}(k)$ are related to the zeros and branch cuts of

$$D_{\mu}[\omega;k]=1-r_k^{\mu}[\omega]^2e^{2ik_zL}\tag{7}$$

The Fresnel formulas are used for the reflection amplitudes. For thick mirrors we have

$$r^{TE}=\frac{k-k_m}{k+k_m},\qquad r^{TM}=\frac{k_m-\varepsilon[\omega]k}{k_m+\varepsilon[\omega]k}\tag{8}$$

Where

$$k_z=ik=i\sqrt{|k|^2-\omega^2/c^2}\tag{9a}$$

$$k_m=\sqrt{|k|^2-\varepsilon[\omega]\omega^2/c^2}=\sqrt{k^2+\omega_p^2/c^2}\tag{9b}$$

Signs for the square roots are chosen such that $\mathrm{Re}[k_i]\rangle 0$ and $\mathrm{Im}[k_i]\langle 0$. This shows that Eq (7) has no solution in the upper half plane.

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Eq (7) shows that the dielectric function for a metal, $\varepsilon[\omega]$, can be described by the plasma model, with ω_p being the plasma frequency. This constant is related to the specific physical properties of the metal. Up to $\omega \sim \omega_p$ the dielectric constant is different from Unity, such that the behavior of the metals is different from the surrounding vacuum. In the case of $\omega \rangle\rangle \omega_p$ the dielectric constant approaches Unity and the metal is transparent. In this way the plasma model implements the high-frequency cut-off for the mirror reflectivity. When we introduce the dielectric properties of the mirrors, it gives rise to a number of important modifications of the field modes. First of all, we cannot write dispersion relations $\omega_n^\mu(k)$ in terms of elementary functions. Figs. 3 and 4 show the results of numerical calculations from these figures. It is clear that imperfect reflection modifies the dispersion relation (solid lines) as, compared to a perfect reflector (dashed lines).

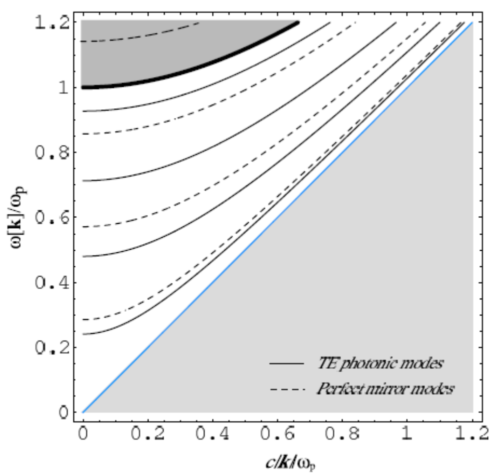

Fig. 3: Dispersion relations for TE-polarized modes between two metallic mirrors described by the plasma model (solid line), compared perfect conductors (dashed line). Mode frequency $\omega\ (k)$ and wavevector in the mirror plane, $|k|$, are normalized to the plasma frequency ω_p . Mirror distance $L = 1.75\lambda_p$. The (blue) diagonal line is the light cone below which the field is evanescent in the cavity (evanescent modes). Above the thick solid line, the field propagates through in the mirror material (bulk modes) [13].

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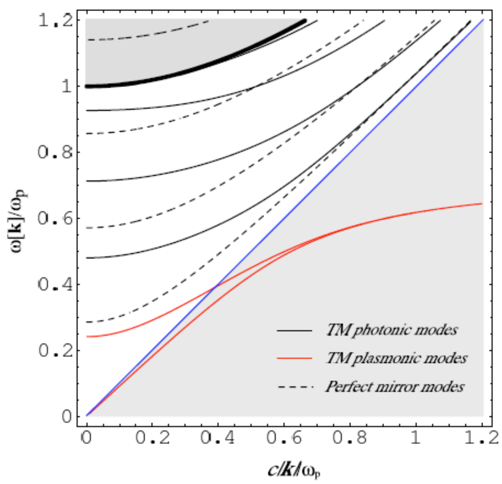

Fig. 4: Dispersion relations for TM-polarized modes. The solid red curves represent the plasmonic modes and the black curves the photonic modes. Note that one of the plasmonic modes crosses the light cone [13].

There are tree observable regions, starting from the top :

1- Bulk modes occur for $\omega \rangle \omega_B(k) = \left(\omega_p^2 + c^2|k|^2\right)^{1/2}$ (shaded above the thick line); their propagation will be in the cavity and inside the mirrors. These modes form a continuum that is mathematically represented by a branch cut of Eq (6) in the complex ω-plane.

   To solve these problems, Schram worked instead with a mirror of finite thickness, d. Therefore, the problem is restricted to thick mirrors.

2- Propagating (ordinary) cavity modes: these modes lie in the region above the light cone and below the bulk continuum. $c|k|\langle \omega \langle \omega_B(k)$. These modes propagate inside the cavity between the mirrors and they behave like a medium optically thinner than vacuum, $0\langle \varepsilon[\omega]\langle 1$. For a given k they lead to a discrete set of mode frequencies. In this region, the reflection coefficients (8) have unit modulus and a frequency-dependent phase.

   Figs. 3, 4 show that the cavity modes are displaced compared to the perfect cavity modes as a direct consequence of the phase shift δ acquired by vacuum field upon reflection.

3- Evanescent modes: they occur in the region below the light cone, $\omega\langle c|k|$ (shaded below the diagonal). Their electromagnetic field decreases exponentially when going away from the vacuum-mirror interface.

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An analytical continuation procedure can be used to obtain optical properties of evanescent modes (reflection and transmission) from ordinary modes. Dispersion relation (7) can be solved for the evanescent region and obtain two non-degenerate mode frequencies in only one polarization, considering a non-magnetic media. These modes are called surface plasmons. The field amplitude related to surface plasmons decreases exponentially away from the interface.

By considering an isolated interface, surface plasmons occur when $\varepsilon[\omega] < -1$. For two interfaces, there are two surface plasmons which are coupled through their corresponding evanescent field tails within the cavity. The zeros of Eq (7) for real k and k_m determines these obtained modes.

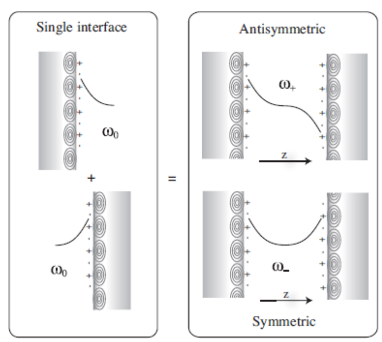

Fig. 5: A surface plasmon mode is associated with oscillating surface charge and surface current densities. The electric field related to a single body is evanescent and the excitation of the plasmon for a plane interface can be done by approaching from the vacuum side a medium with higher index refraction and illuminating the latter in total reflection [13].

The plasmonic modes are created due to frequency splitting. The antisymmetric (ω^+) and symmetric mode (ω^-) have respectively higher and lower energy than the isolated mode. In contrast to ω^- modes which contribute an attractive (binding) force, ω^+ modes contribute a repulsive (antibinding) Casimir force. (Figure 5)

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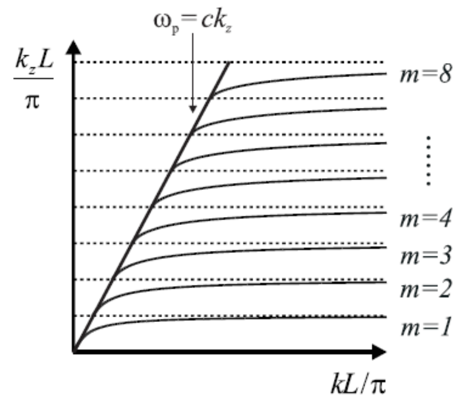

Fig. 6: Mode plot of the first photonic TE modes (m =1, 2, ...8) with the plasma model for ck = 0.5 ω_p. Modes are presented through their longitudinal wavevector as a function of kL/π. The dotted lines correspond to the cavity modes with perfect mirrors [12].

We can say that TE and TM modes inside the cavity are formed by the two mirrors. By solving Eq (5) and using the standard expression for the reflection coefficients, the different modes can be obtained.

From figure 6 we can see the phase shift caused by the TE modes through the influence of imperfect reflection. They can be represented through their perpendicular wavevector as a function of kL. The TE polarization admits only photonic modes which can be written under the standard form $k_z L = m\pi - \delta$, where the integer $m = 1,2...\infty$ is the order of cavity modes and δ the phase shift of the mode on a mirror. Dotted lines caused by perfect mirrors correspond to $\delta^{TE} = 0$. By using the plasma model the photonic modes will have shifted relative to perfectly reflecting mirrors. The limit of perfect reflection corresponds to the large distances limit. Metallic mirrors are transparent at high frequencies, which imposes an upper bound to their perpendicular wavevector $ck_z \langle \omega_p$,where all photonic modes coincide. We can obtain similar photonic modes for TM polarization which are labeled by a positive integer n.

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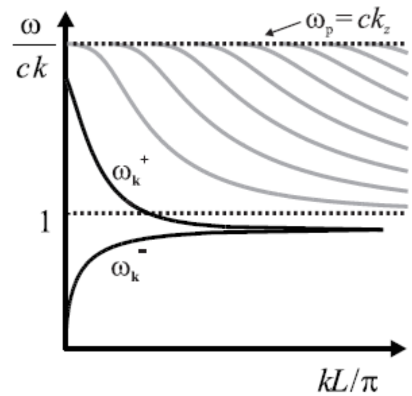

Fig. 7: Mode plot of the two plasmonic modes ω_k^- and ω_k^+ (black) in the sector $\omega\langle ck$ and of photonic modes (gray) in the sector $\omega\rangle ck$ for ck = 0.5 ω_p. Modes are presented through their frequency as a function of kL/π [12].

It is worthwhile pointing out that for the TE-polarization, all modes lie above the light cone, while for TM-polarization two modes which are called "plasmonic" enter the evanescent region in at least part of wavevectors' range. These two additional modes $\left[\omega_k^{pl}\right]_\pm$ are labelled in the limit of small distances. Figure 7 shows the plasmonic modes as solid black line, while gray lines are related to photonic modes. The modes are represented through their frequency as a function of kL. So, plasmonic modes with their imaginary wavevector will be visible. In contrast to photonic modes which lie in the sector $\omega\rangle ck$, plasmonic modes lie in the sector $\omega\langle ck$. When the separation between mirrors is infinite, plasmonic modes are given by the usual dispersion relation for the surface plasmons in a metallic bulk

$$\left[\omega_k^{pl}\right]_\pm \xrightarrow{L\to\infty} \frac{\omega_p^2 + 2|k|^2 - \sqrt{\omega_p^4 + 4|k|^4}}{2} \tag{10}$$

Finally, the total Casimir energy is rewritten by a sum of these terms

$$E = \underbrace{\sum_k \left[\frac{\hbar\omega_+}{2} + \frac{\hbar\omega_-}{2}\right]_{L\to\infty}^{L}}_{plasmonic\,modes(E_{pl})} + \underbrace{\overbrace{\sum_{\mu,k}\left[\sum_{\omega\langle\omega_B} \frac{\hbar\omega_n^\mu}{2}\right]_{L\to\infty}^{L}}^{cavity\,modes} + \overbrace{\lim_{d\to\infty}\sum_{\mu,k}\left[\sum_{\omega\geq\omega_B}\frac{\hbar\omega_n^\mu}{2}\right]_{L\to\infty,d}^{L}}^{bulk\,modes}}_{photonic\,modes(E_{ph})} \tag{11}$$

We cannot measure these terms separately because they don't have physical meaning on their own.

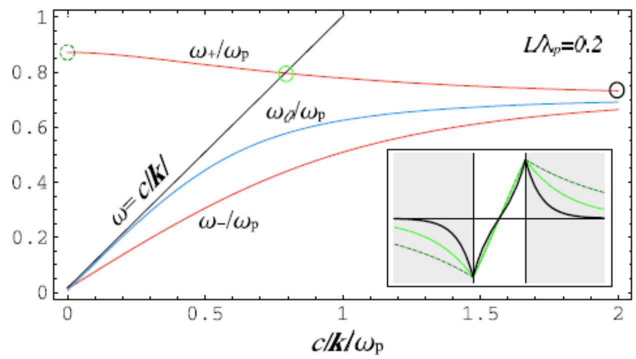

Fig. 8: A plot of the plasmonic dispersion relations $\omega_+(k)$, $\omega_0(k)$, $\omega_-(k)$, as function of $|k|$ for $L = 0.2\lambda_p$ ($\lambda_p = {}^{2\pi c}/_{\omega_p}$). Frequencies and wavevectors are scaled to the plasma frequency ω_p and ${}^{\omega_p}/_{c}$, respectively. Inset: magnetic field amplitude for chosen points along the branch $\omega_+(k)$, as labelled by the circles [13].

The dispersion relation is shown in Fig. 8 for two modes in the first sum of Eq.(11). The field related to these modes is evanescent both in vacuum and in the mirrors. $\omega_-(k)$ is restricted to the plasmonic mode sector (below the light cone), while $\omega_+(k)$ lies in the plasmonic mode sector for large distances, but it crosses the barrier $\omega\langle ck$ and dies out in the photonic mode sector for ${}^{kL}/_{\pi} \to 0$.

First sum of the Eq.(11) shows the plasmonic mode contribution

$$E_{pl} = \underbrace{\sum_k \left[\frac{\hbar\omega_+}{2} + \frac{\hbar\omega_-}{2}\right]_{L\to\infty}^{L}} \tag{12}$$

Both modes approach $\omega_0(k)$ for $L \to \infty$ such that zero-point energy for two isolated surface plasmons is subtracted. The phase shift δ for photonic modes at infinite distances approaches zero and their dispersion relation for perfect mirrors is given by

$$\left[\omega_k^{\varepsilon}\right]_n = \sqrt{|k|^2 + k_z^2} \tag{13}$$

with wavevectors $k_z = {}^{n\pi}/_{L}$.

In the limit of small distance L, the plasmonic mode ω_k^+ acquires a phase shift with the same sign as the TM photonic modes below the plasma frequency. The corresponding frequency at short distance is increased compared to long distances, while the frequency of ω_k^- for short distances is smaller than the one in the large distances.

V. Conclusion

In this paper the contribution of surface plasmon modes in the Casimir effect was investigated. Based on what we have seen in this article we know that Plasmonic modes are much more important not only in the short distances limit, but also at large distances. In the short distance limit, Casimir energy is very well approximated as the coulomb interaction energy between the two surface plasmons. As the mirror separation is increased, the photonic mode contribution increases while plasmonic modes give a larger repulsive contribution. This behavior attributed to the behavior of ω_k^+, which gives a repulsive contribution at all distances. This is balanced in the total Casimir energy by the contributions of photonic modes (cavity and bulk modes). So, the whole Casimir energy is the result of balance between the large attractive photonic contribution and the large repulsive plasmonic contribution. The outcome of this balance keeps the sign of a binding energy. However, this result relies heavily on the symmetry of the Casimir geometry with two plane mirrors. Based on the works that are done in this field until now, it would be interesting to investigate changing the balance (the contribution of plasmonic modes) and therefore the value or even the sign of the Casimir force by changing the field mirror coupling (geometry), using a non-planar mirror, metallic surfaces with nanostructures graved into it or even using the hole arrays used to enhance the transmission of light through metallic structures [15] or nano structured metallic surfaces.

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REFERENCES

References

[1] H. B. G. Casimir, D. Polder, Phys. Rev. **73**, 360 (1948)

[2] H. B. G. Casimir, J. Chim, Phys **46**, 407 (1949)

[3] H. B. G. Casimir, Proc. K. Ned. Akad. Wet 51, 793 (1948)

[4] E.M. Lifshitz, Sov. Phys. JETP 2 (1956) 73.

[5] I.E. Dzyaloshinskii, E.M. Lifshitz, L.P. Pitaevskii, Adv. Phys. **10** (1961) 165.

[6] M. L. Brongersma and P. G. Kik, eds., *Surface Plasmon Nanophotonics*, Springer Series in Optical Sciences (Springer, Berlin/Heidelberg/New York, 2007).

[7] S. A. Maier, *Plasmonics: Fundamentals and Applications* (Springer, Berlin/Heidelberg/New York, 2007).

[8] F. J. Garcia-Vidal, L. Martin-Moreno, T. W. Ebbesen, and L. Kuipers, Rev. Mod. Phys. **82**, 729 (2010).

[9] E. Altewischer, M. van Exter, and J. Woerdman, Nature (London) **418**, 304 (2002).

[10] S. Fasel, F. Robin, E. Moreno, D. Erni, N. Gisin, and H. Zbinden, Phys. Rev. Lett. **94**, 110501 (2005).

[11] F. Intravaia, Ph.D. thesis, Universit´e Paris VI, 2005.

[12] F. Intravaia and A. Lambrecht, Phys. Rev. Lett. **94**, 110404 (2005).

[13] F. Intravaia, C. Henkel, and A. Lambrecht, Phys. Rev. A **76**, 033820 (2007).

[14] T. Ebbesen, H. Lezec, H. Ghaemi, T. Thio, and P.Wolff, Nature 391, 667 (1998).

[15] L. Martin-Moreno, F. J. Garcia-Vidal, H. J. Lezec, K. M. Pellerin, T. Thio, J. B. Pendry, and T. W. Ebbesen, Phys. Rev. Lett. **86**, 1114 (2001).

[16] E. Altewischer, M. van Exter, and J. Woerdman, Nature 418, 304 (2002).

[17] A. W. C. Lau, D. Levine, and P. Pincus, Phys. Rev. Lett. 84, 4116 (2000).

[18] H. B. G. Casimir, Proc. Koninkl. Ned. Akad. Wetenschap. **51**, 793 ~1948!.

[19] P. W. Milloni, *The Quantum Vacuum: An Introduction to Quantum Electrodynamics* ~Academic, London, 1994!.

[20] E. M. Lifshitz, J. Exp. Theor. Phys. 2, 73 (1956)

[21] Lamoreaux, S. K. 1996, Phys: Rev: Lett:, 78, 1

[22] Munday, J. N., Capasso, F., & Parsegian, V. A., 2009, Nature, 457, 170-173

[23] C. Genet, F. Intravaia, A. Lambrecht and S. Reynaud, Ann. Fond. L. de Broglie 29 311 (2004); arXiv:quant-ph/0302072

[24] C. Henkel, K. Loulain, J-Ph. Mulet, and J.-J. Greffet, Phys. Rev. A69 023808 (2004).

[25] N. V. Kampen, B. Nijboer, and K. Schram, Phys. Lett. A 26, 307 (1968).

(1975); F. Forstmann and H. Stenschke, ibid. B17 1489 (1978).

[27] G. Barton, Rep. Prog. Phys. 42 65 (1979).

[28] K. Schram, Phys. Lett. A 43, 282 (1973).

[29] M.T. Jaekel and S. Reynaud, J. Physique I-1 1395 (1991).

[30] C. Genet, A. Lambrecht and S. Reynaud, Phys. Rev. A67 043811 (2003).