Diamagnetic effect on the Casimir force

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The Casimir force between a diamagnetic plate and a magnetodielectric plate at finite temperature is considered. Under the condition that the permittivity of the magnetodielectric plate is sufficiently small, we show that the diamagnetic property dominantly determines the asymptotic behavior of the repulsive Casimir force for large separations. On the basis of this simple property, we present numerical results showing that if an effective permeability of a superconductor is much less than one, its diamagnetic response can be indirectly detected by measuring the Casimir force.

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I. INTRODUCTION

The recent interest in the repulsive Casimir force [1,2] is inspired by the development of new magnetic materials. Although a variety of methods [3–10] to generate the repulsive Casimir force have been proposed, neither the repulsive Casimir force nor the magnetic effect on the Casimir force have been observed in vacuum, yet. Recently, Geyer *et al.* reported a very important understanding of the Casimir force between magnetodielectric plates [11,12]: at finite temperature, if the permeability along the imaginary axis at the nonzero Matsubara frequencies is very small, then the static permittivity and the static permeability mainly determine the dependence of the magnetic property on the Casimir force. Thus, to observe the magnetic effect on the Casimir force, we must increase either the permeability at the nonzero Matsubara frequencies or the contribution to the repulsive Casimir force at zero frequency. The former approach of increasing the permeability at the nonzero Matsubara frequencies has been already studied extensively [4,13]. Since the first nonzero Matsubara frequency at room temperature is of the order of 10^{14} rad/s, a novel optical material, namely a metamaterial, must be synthesized. In this paper, we focus on the latter approach of increasing the repulsive contribution of the Casimir force at zero frequency due to the diamagnetic effect. Although the permeability of existing diamagnetic materials is slightly less than one, Wood and Pendry designed a metamaterial having arbitrary effective permeability between 0 and 1 at zero frequency by using highly conductive materials [14]. The diamagnetic metamaterials have potential to be useful materials for investigating the Casimir effect.

Superconductors are also useful materials for examining the Casimir effect [15–18]. However, we do not have sufficient knowledge about the interaction of vacuum fluctuations with a superconductor, in particular magnetic interactions. Although it is often remarked that a superconductor is a perfect diamagnetic material with zero permeability, we generally regard superconductors as conductive materials having an infinite static conductivity. The most significant difference between a perfectly conductive material and a superconductor is that only the superconductor can exclude any magnetic field from it, i.e., the Meissner effect [19]. If the magnetic property of the superconductors affects the Casimir force, we may have to consider the Casimir and the Meissner effects simultaneously.

Since the Casimir force depends on electrical and magnetic properties, we expect that the material properties can be derived from the force-distance curve. Unfortunately, this is not possible for permittivity [20]. However, the dc magnetic permeability, which is a very important quantity to consider for the Meissner effect, may be indirectly measured by measuring the Casimir force. To examine this possibility, we calculate the Casimir force between a diamagnetic plate and a magnetodielectric plate. Our numerical results show that the effective dc permeability of a superconductor for the Casimir effect can be measured by current experiments, if the permeability is much less than one. Thus, the question of whether a superconductor should be considered as a diamagnetic or a nonmagnetic material can be answered by the experiments.

The structure of this paper is as follows. In Sec. II we briefly introduce the Lifshitz formula at finite temperature and consider the dependence of the direction of the Casimir force on the permeability of the diamagnetic metamaterial designed by Wood and Pendry. In Sec. III the repulsive Casimir force between a high-T_c superconductive plate and a porous magnetic nanocomposite is considered. We chose $Bi_2Sr_2CaCu_2O_{8+\delta}$ as a material of a high-T_c superconductor and show the contributions to the Casimir force at zero frequency, which determine the magnitude of the repulsive Casimir force for large separations. In Sec. IV we present our conclusions and discuss the possibility of the indirect measurement of the effective permeability of the superconductor.

II. CASIMIR FORCE ACTING ON DIAMAGNETIC METAMATERIAL

The Casimir force between infinite parallel plates at temperature T is determined by the separation length between the plates a, the complex permittivity $\epsilon(\omega)$, and the complex permeability $\mu(\omega)$ of the plates. According to the Lifshitz theory [21], the Casimir force per unit area of the plate can be expressed by adding the following four components:

$$P(a,T) = P_0^{\mathrm{TM}}(a,T) + P_0^{\mathrm{TE}}(a,T) + P_p^{\mathrm{TM}}(a,T) + P_p^{\mathrm{TE}}(a,T). \quad (1)$$

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Here, $P_0^{\rm TM}$ and $P_0^{\rm TE}$ are contributions at zero frequency to the Casimir force on the transverse magnetic (TM) and transverse electric (TE) modes, respectively and $P_p^{\rm TM}$ and $P_p^{\rm TE}$ are contributions at positive frequencies to the Casimir force of the TM and TE modes, respectively. In this study, the total contributions at zero frequency play an important role in determining the sign of the Casimir force for large separations, and they are given by

$$P_0^{(\rm mode)}(a,T) = -\frac{k_B T}{2\pi}\left(\frac{\alpha}{c}\right)^3 \int_0^\infty \kappa^2 d\kappa\, G_0^{(\rm mode)}(\kappa), \quad (2)$$

where k_B is the Boltzmann constant, $\alpha = 2\pi k_B T/\hbar$, and κ is the modulus of the wave-vector projection on the plate normalized by α/c. The coefficient of the integral $\alpha^3 k_B T/2\pi c^3$ is evaluated as 0.74 $(T/300\text{ K})$ (m Pa). The function $G_0^{(\rm mode)}(\kappa)$ is defined by

$$G_0^{(\rm mode)}(\kappa) = \left[\frac{e^{\beta a\kappa}}{r_{\rm mode}^{(1)}(\kappa) r_{\rm mode}^{(2)}(\kappa)} - 1\right]^{-1}, \quad (3)$$

where $\beta \equiv 2\alpha/c = 1.65(T/300\text{ K})\ \mu\text{m}^{-1}$. The electric and magnetic dependences of the Casimir force at zero frequency are determined by the reflection coefficients

$$r_{\rm TM}^{(n)}(\kappa) = \frac{\epsilon^{(n)}(0)\kappa - \sqrt{\kappa^2 + \left(\kappa_0^{(n)}\right)^2}}{\epsilon^{(n)}(0)\kappa + \sqrt{\kappa^2 + \left(\kappa_0^{(n)}\right)^2}}, \quad (4)$$

$$r_{\rm TE}^{(n)}(\kappa) = \frac{\mu^{(n)}(0)\kappa - \sqrt{\kappa^2 + \left(\kappa_0^{(n)}\right)^2}}{\mu^{(n)}(0)\kappa + \sqrt{\kappa^2 + \left(\kappa_0^{(n)}\right)^2}}, \quad (5)$$

where n is the index of the plate, and $\epsilon^{(n)}(0) \equiv \lim_{\xi\to 0}\epsilon(i\xi)$ and $\mu^{(n)}(0) \equiv \lim_{\xi\to 0}\mu(i\xi)$ denote the permittivity and permeability of the plate, respectively, labeled by n at zero frequency. The constant $\kappa_0^{(n)}$ is defined by

$$\kappa_0^{(n)} = \lim_{\xi\to 0}\sqrt{\epsilon^{(n)}(i\xi)\mu^{(n)}(i\xi)}\,\xi. \quad (6)$$

We consider a metamaterial designed by Wood and Pendry as an example of a diamagnetic plate, in the form of a cubic array of a highly conductive material (see the inset in Fig. 1). The effective permittivity $\epsilon_{\rm eff}$ and permeability $\mu_{\rm eff}$ at zero frequency are approximately expressed using the ratio of the cubic size to the lattice constant $\gamma(<1)$ as $\epsilon_{\rm eff} = \gamma^2/(1-\gamma)$ and $\mu_{\rm eff} = 1-\gamma^2$ near $\gamma = 1$. As the ratio approaches 1, the permittivity diverges and the permeability converges to zero. However, the product of the permittivity and the permeability does not exceed 2 for any γ. If the product of the permittivity and the permeability is finite, the contributions to the Casimir force at zero frequency can be obtained using the following analytical formulas:

$$P_0^{\rm TM}(a,T) = -\frac{k_B T}{8\pi a^3}{\rm Li}_3\left(\frac{\epsilon^{(1)}(0)-1}{\epsilon^{(1)}(0)+1}\frac{\epsilon^{(2)}(0)-1}{\epsilon^{(2)}(0)+1}\right), \quad (7)$$

$$P_0^{\rm TE}(a,T) = -\frac{k_B T}{8\pi a^3}{\rm Li}_3\left(\frac{\mu^{(1)}(0)-1}{\mu^{(1)}(0)+1}\frac{\mu^{(2)}(0)-1}{\mu^{(2)}(0)+1}\right), \quad (8)$$

where ${\rm Li}_3(z)$ is the polylogarithm function, which can be expressed by $z + z^2/8 + O(z^3)$ near zero and changes the sign at zero. Since the permittivities $\epsilon^{(1)}(i\xi)$ and $\epsilon^{(2)}(i\xi)$ are always greater than one, the contribution of the TM mode at zero frequency to the Casimir force is always attractive. However, $\mu^{(1)}(i\xi)$ and $\mu^{(2)}(i\xi)$ can be less than one without violating the Kramers-Kronig relation [22]. Thus, if $\mu^{(1)}(0) < 1$ and $\mu^{(2)}(0) > 1$, then the contribution of the TE mode at zero frequency to the Casimir force is always repulsive. If we assume that the plate labeled by $n = 1$ is a metamaterial having $\gamma = 1$, then the contribution at zero frequency to the Casimir force is expressed by

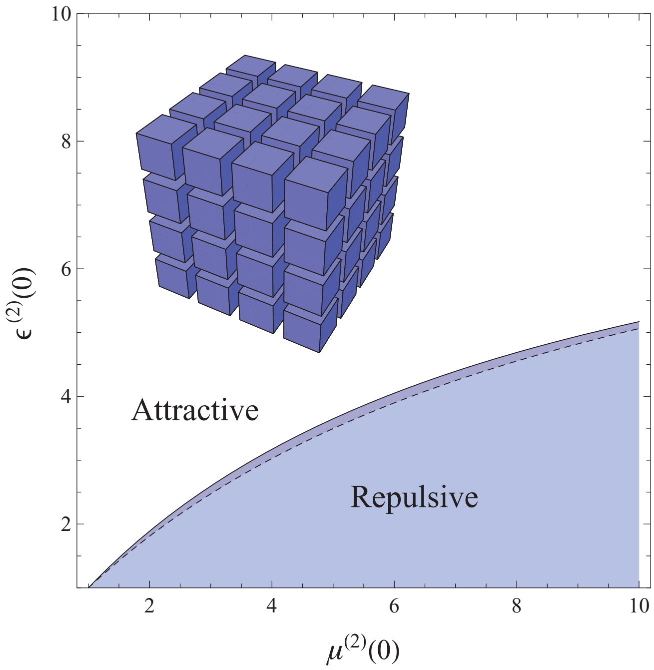

FIG. 1. (Color online) Direction of the Casimir force between a diamagnetic metamaterial and a magnetodielectric plate with $\mu^{(2)}(0)$ and $\epsilon^{(2)}(0)$ for large plate separations. The solid and dashed curves show boundaries corresponding to $\gamma = 1$ and $\gamma = 0.8$, respectively. The inset illustrates the structure of a cubic array of a metamaterial designed by Wood and Pendry.

$$P_0(a,T) = -C_0\frac{k_B T}{8\pi a^3}, \quad (9)$$

where $C_0 = {\rm Li}_3\{[\epsilon^{(2)}(0)-1]/[\epsilon^{(2)}(0)+1]\} + {\rm Li}_3\{-[\mu^{(2)}(0)-1]/[\mu^{(2)}(0)+1]\}$. Figure 1 shows the sign of the coefficient C_0 for $\epsilon^{(2)}(0)$ and $\mu^{(2)}(0)$. If the static permittivity of the magnetodielectric plate is much less than its static permeability, then the contribution at zero frequency to the Casimir force is always repulsive, independent of the separation length. On the other hand, if $\epsilon^{(2)}(0) > 8.67$, the contribution at zero frequency is attractive for any $\mu^{(2)}(0)$.

The contributions at positive frequencies are given by

$$P_p^{(\rm mode)}(a,T) = -\frac{k_B T}{2\pi}\left(\frac{\alpha}{c}\right)^3 \sum_{l=1}^{\infty}\int_0^\infty \sqrt{\kappa^2 + l^2}\,d\kappa \times \left[\frac{e^{\beta a q_l}}{r_{\rm mode}^{(1)}(i\xi_l,\kappa) r_{\rm mode}^{(2)}(i\xi_l,\kappa)} - 1\right]^{-1}, \quad (10)$$

where $\xi_l = 2\pi k_B T l/\hbar$ with positive integers l are the Matsubara frequencies and

$$q_l^2 \equiv q_l^2(l,\kappa) = \kappa^2 + l^2. \tag{11}$$

The reflection coefficients for positive frequencies are given by

$$r_{\rm TM}^{(n)}(i\xi_l,\kappa) = \frac{\epsilon_l^{(n)}\sqrt{\kappa^2+l^2} - \sqrt{\kappa^2+\epsilon_l^{(n)}\mu_l^{(n)}l^2}}{\epsilon_l^{(n)}\sqrt{\kappa^2+l^2} + \sqrt{\kappa^2+\epsilon_l^{(n)}\mu_l^{(n)}l^2}}, \tag{12}$$

$$r_{\rm TE}^{(n)}(i\xi_l,\kappa) = \frac{\mu_l^{(n)}\sqrt{\kappa^2+l^2} - \sqrt{\kappa^2+\epsilon_l^{(n)}\mu_l^{(n)}l^2}}{\mu_l^{(n)}\sqrt{\kappa^2+l^2} + \sqrt{\kappa^2+\epsilon_l^{(n)}\mu_l^{(n)}l^2}}, \tag{13}$$

where $\epsilon_l^{(n)} = \epsilon^{(n)}(i\xi_l)$ and $\mu_l^{(n)} = \mu^{(n)}(i\xi_l)$. For large separations, the exponential function in Eq. (8) rapidly decreases as l increases; thus, the sign of the Casimir force for large separations is determined by the contribution at zero frequency.

III. CASIMIR FORCE BETWEEN SUPERCONDUCTOR AND NANOCOMPOSITE

Let us consider the Casimir force between a high-T_c superconductor $Bi_2Sr_2CaCu_2O_{8+\delta}$ (BSCCO) and a porous magnetic nanocomposite plate. We can now evaluate the Casimir force if explicit formulas of the permittivity and permeability are given. Romanowsky and Capasso studied the Casimir force acting on BSCCO [18]. They showed that the anisotropy of the permittivity of BSCCO gives rise to a difference in the Casimir force. Accordingly, the permittivity of the superconductor must be expressed by the component parallel to the optical axis $\epsilon_\parallel^{(1)}$ and the component perpendicular to the optical axis $\epsilon_\perp^{(1)}$. In this paper, we define the optical axis as perpendicular to the conducting copper-oxide planes. Both of these components were described in the framework of the oscillator model,

$$\epsilon(i\xi) = 1 + \sum_{j=1}^{K} \frac{g_j}{\omega_j^2 + \xi^2 + \gamma_j \xi}. \tag{14}$$

We took the above parameters (oscillator frequencies ω_j, oscillator strengths g_j, and relaxation parameters γ_j) from Refs. [18,23]. These parameters are based on the experimental measurement at room temperature. The optical data of BSCCO below T_c is also measured in Ref. [23], and the spectrum does not show drastic change. We note that the temperature dependence of the permittivity in the framework of the Drude model is theoretically studied using a nested Fermi-liquid analysis [24]. Although the optical properties of the superconductor depend also on other properties such as the amount of dope and the properties of the surfaces [25], the most important property of BSCCO in this study is that the product $\sqrt{\epsilon(i\xi)}\xi$ vanishes in the limit $\xi \to 0$; the parameter $\kappa_0^{(1)}$ defined in Eq. (6) is zero in both the oscillator model and the Drude model.

Figure 2(a) shows $\epsilon_\parallel^{(1)}$ and $\epsilon_\perp^{(1)}$ along the imaginary frequency axis. Note that $\epsilon_\parallel^{(1)}$ is much less than $\epsilon_\perp^{(1)}$, and this implies suppression of the contribution of the TM mode to the attractive Casimir force for short separations. For the ac permeability of a high-T_c superconductor, we do not have sufficient experimental data [26], and we set $\mu_l = 1$ for $l \geqslant 1$. As shown below, if $\mu_l < 1$ for $l \geqslant 1$, its modification makes the repulsive force increase, but the final result is not sensitive to ac permeability.

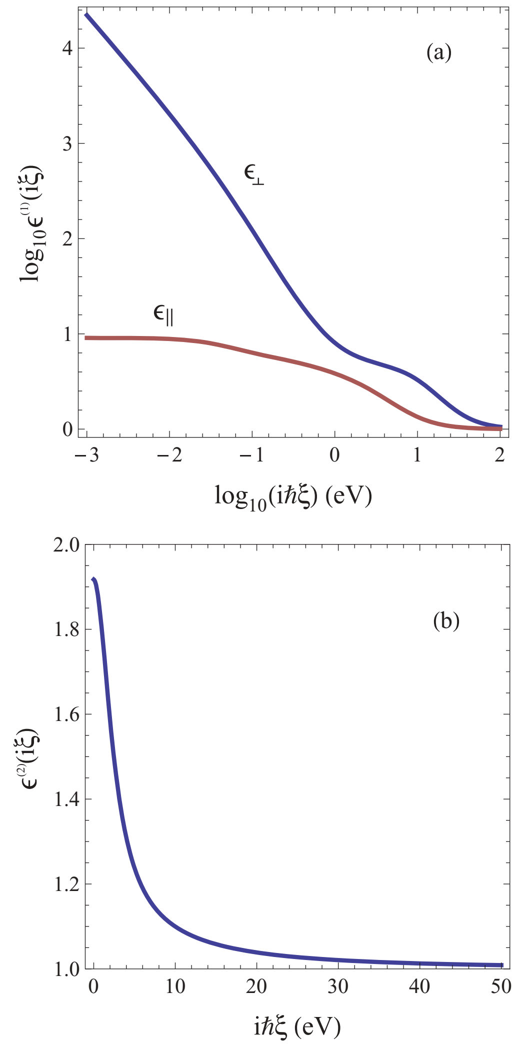

FIG. 2. (Color online) Permittivity of (a) the superconductor and (b) the nanocomposite along the imaginary frequency axis.

To contrast the contribution to the Casimir force at zero frequency with that at positive frequencies, we have to reduce the permittivity of the magnetodielectric plate without destroying its magnetism. A porous nanocomposite that consists of nickel nanoparticles and polystyrene is a well-known magnetodielectric material with a small permittivity [27]. Assuming that the permittivity of a porous nanocomposite $\epsilon_d(i\xi)$ can be described by the Maxwell-Garnett theory, the permittivity of the nanocomposite is a solution of the equation

$$\frac{\epsilon_d(i\xi)-1}{\epsilon_d(i\xi)-2} = f_1\frac{\epsilon_{\rm Ni}(i\xi)-1}{\epsilon_{\rm Ni}(i\xi)-2} + f_2\frac{\epsilon_{\rm poly}(i\xi)-1}{\epsilon_{\rm poly}(i\xi)-2}, \tag{15}$$

where $\epsilon_{\rm Ni}(i\xi)$ and $\epsilon_{\rm poly}(i\xi)$ are the dielectric functions along the imaginary frequency axis for nickel and polystyrene, respectively, and f_1 and f_2 are the volume fractions of nickel and polystyrene, respectively. For nickel, we used the plasma model with plasma frequency 3.94 eV [28], and $\epsilon_{\rm poly}(i\xi)$ is

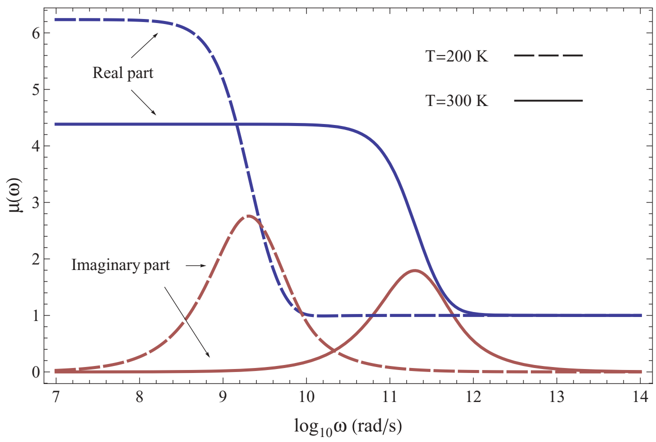

FIG. 3. (Color online) Real and imaginary parts of the permeability of a nanocomposite at $T = 300$ K (solid line) and $T = 200$ K (broken line).

expressed by the oscillator model with $K = 4$. Figure 2(b) shows the permittivity of a porous nanocomposite at $f_1 = 0.2$ and $f_1 = 0.1$ [29–31]. We find that the permittivity of the nanocomposite at zero frequency, $\epsilon^{(2)}(0) = 1.9$, is less than the above-mentioned threshold 8.67.

The permeability of a nanocomposite strongly affects the contribution to the Casimir force at zero frequency. According to Onsager's theory [32,33], if the size of a nanoparticle a_{Ni} is much smaller than the skin depth of nickel, then the complex permeability of the nanocomposite $\mu_d(\omega)$ is given by

$$\mu_d(\omega) = \frac{1}{4}\left[1 + \frac{4\pi Rf}{1 - i\omega\tau} + \sqrt{8 + \left(1 + \frac{4\pi Rf}{1 - i\omega\tau}\right)^2}\right]. \quad (16)$$

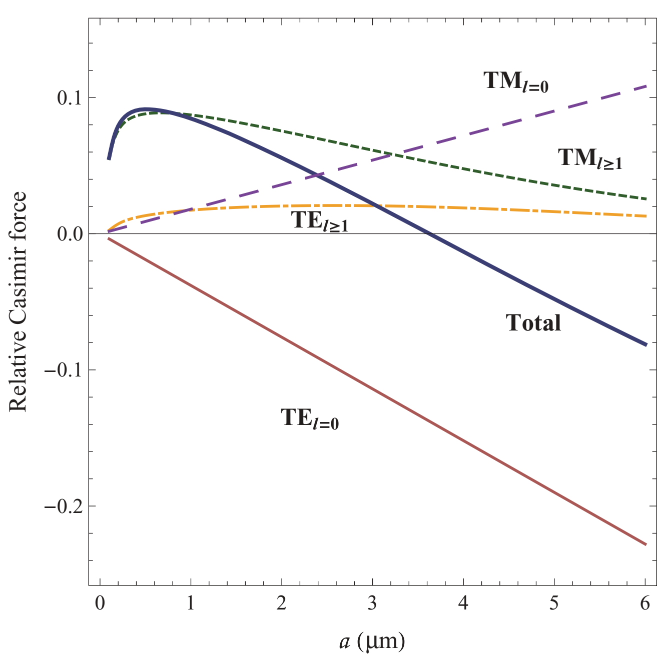

FIG. 4. (Color online) Plot of four contributions of the Casimir force between a superconductor and a nanocomposite at $T = 130$ K to the Casimir force between perfectly conductive plates: P_0^{TM} (long dashed line), P_0^{TE} (short dashed line), P_p^{TM} (dotted line), and P_p^{TE} (dotted dashed line), as well as the total Casimir force (solid line).

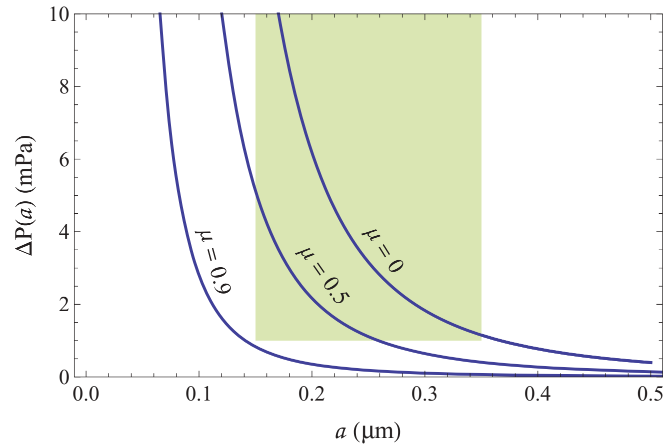

FIG. 5. (Color online) Absolute difference $\Delta P(a)$ between the Casimir force acting on a superconductor with the static permeability $\mu(0)$ and on a nonmagnetic superconductor. The rectangular shaded area indicates the region in which the Casimir force could be measured by current experiments.

Here τ is the Debye relaxation time and R is defined by $4\pi a_{\mathrm{Ni}}^3 \bar{m}_0^2 / k_B T$, where $\bar{m}_0$ is the saturation magnetization. The solid lines in Fig. 3 show the permeability of nanocomposite for $a_{\mathrm{Ni}} = 5$ nm, $\tau = 4.8$ ps, and $\bar{m}_0 = 480$ G at $T = 300$ K. Note that the static permeability depends on the temperature, and changes from 4.2 to 6.1 where the temperature decreases from 300 to 200 K. Also, the frequency where the imaginary part of the permeability has a maximum value is much smaller than the Matsubara frequency with $l = 1$ ($\approx 10^{14}$ rad/s). Thus the magnetic effect on the contribution to the Casimir force at positive frequencies is very small.

Figure 4 shows the contributions of the relative Casimir force between a superconductor having $\mu^{(1)}(0) = 0$ and a nanocomposite having $\mu^{(2)}(0) = 7$ to the Casimir force between perfectly conductive plates given by $-\pi^2\hbar c/240a^4$ at $T = 130$ K, which is a typical transition temperature of high-T_c superconductors. To take the anisotropic property of the superconductor into account, we use the generalized Lifshitz formula for uniaxial crystals [12]. The total Casimir force changes from attractive to repulsive near $a = 3.6$ μm, and the contribution of the TE mode at zero frequency dominantly determines the asymptotic behavior for large plate separations. The magnitude of the repulsive Casimir force is very small and difficult to measure with current experiments. However, the absolute value of the contribution at zero frequency rapidly increases as the plate separation decreases, as shown in Eq. (9). Figure 5 shows the difference between the Casimir force acting on magnetic superconductors with three different permeabilities and on a nonmagnetic superconductor. According to the measurement criterion for the Casimir force given in Ref. [13], the Casimir force can be measured in the rectangular shadow region in Fig. 5.

IV. CONCLUDING REMARKS

One of the necessary conditions when the contribution to the Casimir force between a magnetic plate and a nonmagnetic plate at zero frequency is repulsive is that one of the reflection coefficients for the TE mode $r_{\mathrm{TE}}^{(n)}(\kappa)$, defined in Eq. (5), is negative. Since the denominator of $r_{\mathrm{TE}}^{(n)}(\kappa)$ is always positive, the numerator must be negative if $r_{\mathrm{TE}}^{(n)}(\kappa)$ is negative. This means that $\kappa^2 + [\kappa_0^{(n)}]^2$ is greater than $[\mu^{(n)}(0)\kappa]^2$. If the magnetic plate is not diamagnetic, the permeability is greater than 1. Thus the value of $[\kappa_0^{(n)}]^2$ must be positive if the reflection coefficient is negative. Since the parameter $\kappa^{(n)}$ is defined by $\sqrt{\epsilon(i\xi)}\xi$, the permittivity along the imaginary frequency must diverge at the zero frequency. The metallic plate in which the permittivity is described by the plasma model satisfies this condition; this possibility was studied by Geyer *et al.* [11].

Another possibility of the negative reflection coefficient is considered in this paper. Even if the permittivity is finite at zero frequency, i.e., $\kappa_0^{(n)} = 0$, the reflection coefficient becomes negative if $\mu^{(n)}(0) < 1$. In contrast to the former case, the sign of the reflection coefficient is independent of the wave number κ and is always negative.

We calculated contributions to the Casimir force between a diamagnetic plate and a magnetodielectric plate at finite temperature and showed that if the permeability of the diamagnetic material at zero frequency $\mu(0)$ satisfies $\mu(0) \ll 1$, then the diamagnetic effect determines the asymptotic behavior of the repulsive Casimir force for large plate separations. Based on this result, we showed that for the combination of a superconductive plate and a porous magnetic nanocomposite, if the effective permeability $\mu(0) < 0.9$, current experiments could detect the diamagnetic effect of the superconductor on the Casimir force, suggesting that a superconductor should be considered as a diamagnetic material.

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