Van der Waals Interactions in a Magneto-Dielectric Medium

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(Dated: April 1, 2022)

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The continuing interest in Casimir effects and the related effects of vacuum field fluctuations in the case of dielectric media has stimulated some interest in the van der Waals interactions of atoms embedded in magneto-dielectric media. Unlike previous calculations which infer the atom-atom interaction from the dilute-medium limit of the macroscopic, many-body van der Waals interaction, the interaction is calculated directly for the system of two atoms in a magneto-dielectric medium. Two approaches are presented, the first based on Green functions that allow for absorption. We show that the correct van der Waals interactions are obtained regardless of whether absorption in the host medium is taken into account.

PACS numbers: 12.20.-m, 42.50.Nn, 78.67.-n

I. INTRODUCTION

The van der Waals interaction between ground-state atoms in vacuum is often regarded as a consequence of the vacuum fluctuations of the electromagnetic field. It is well known that the interaction between two macroscopic, uncharged dielectric bodies cannot be obtained by pairwise addition of this interatomic van der Waals interaction except in the dilute-medium limit. In this limit Lifshitz, for instance, obtained the retarded and unretarded pairwise van der Waals interaction between electrically polarizable atoms from a more general expression for the interaction energy of two dielectric half-spaces\textsuperscript{1}. The van der Waals interaction obtained in this way is identical to that obtained more directly by Casimir and Polder\textsuperscript{2} for the system of two atoms in vacuum.

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II. VAN DER WAALS INTERACTIONS FROM QUANTIZED FIELD WITHOUT ABSORPTION

We consider two identical atoms in a homogeneous and isotropic magneto-dielectric medium which we regard as a continuum with real electric permittivity $\epsilon(\omega)$ and magnetic permeability $\mu(\omega)$, and therefore real refractive index $n(\omega)$. The calculation of the van der Waals interactions in this section will be based on the following expressions for the electric and magnetic fields in the non-absorbing magneto-dielectric medium:\textsuperscript{11}

\begin{equation}
E(r, t) = \sum_{k\lambda} \left( \frac{2\pi \hbar \omega k}{n_k \gamma_k \epsilon V} \right)^{1/2} \left[ a_{k\lambda}(t)e^{ikr} - a_{k\lambda}^\dagger(t)e^{-ikr} \right] \epsilon_{k\lambda},
\end{equation}

\begin{equation}
H(r, t) = i \sum_{k\lambda} \left( \frac{2\pi \hbar c^2}{\omega_k n_k \gamma_k \mu V} \right)^{1/2} \left[ a_{k\lambda}(t)e^{ikr} - a_{k\lambda}^\dagger(t)e^{-ikr} \right] k \times \epsilon_{k\lambda}.
\end{equation}
We employ a standard notation in which \( a_{k\lambda}(t), a_{k\lambda}^\dagger(t) \) are Heisenberg-picture photon annihilation and creation operators for the plane-wave mode with wave vector \( k \) \(|k| = k = n_k \omega_k/c, n_k = n(\omega_k)\) and (linear) polarization unit vector \( e_{k\lambda}(\mathbf{r} \cdot e_{k\lambda} = 0, \lambda = 1,2) \). The refractive index is \( n_k = (\epsilon_k \mu_k)^{1/2} \), where \( \epsilon_k \) and \( \mu_k \) are the electric permittivity and magnetic permeability, respectively, at frequency \( \omega_k \). The group index is \( \gamma_k = n_k + \omega_k d\mu_k/d\omega_k \), and \( V \) is the quantization volume for the box normalization of the plane-wave modes. The Hamiltonian is

\[
H = H_A + H_B + H_F - d_A(t) \cdot \mathbf{E}(\mathbf{r}_A,t) - d_B(t) \cdot \mathbf{E}(\mathbf{r}_B,t) - m_A(t) \cdot \mathbf{H}(\mathbf{r}_A,t) - m_B(t) \cdot \mathbf{H}(\mathbf{r}_B,t),
\]

where \( H_A, H_B, \) and \( H_F \) are the Hamiltonian operators for atom A, atom B, and the electromagnetic field, respectively. The electric dipole moment operators are denoted by \( \mathbf{d}(t) \) and the magnetic dipole moment operators by \( \mathbf{m}(t) \).

### A. Van der Waals interaction between electrically polarizable particles

We first consider the van der Waals interaction between two electrically polarizable atoms separated by a distance \( R \) in the magneto-dielectric medium. The approach we will take follows closely that used by various authors for the derivation of the van der Waals interaction in the case in which the particles are in vacuum \( (n = 1) \). The basic idea of the method is that quantum vacuum fluctuations of the electromagnetic field in the medium induce in the atoms fluctuating electric and magnetic dipoles that interact with each other. The electric dipole moment induced in an atom at \( \mathbf{r} \) by an electric field is \( \mathbf{d}(t) = \alpha_e(t) \mathbf{E}(\mathbf{r},t) \), where \( \alpha_e \) is the (real) electric polarizability of the atom. (Since it is only the real part of the polarizability that determines shifts in energy levels, we can assume without loss of generality throughout this paper that the polarizabilities of the guest atoms are real.) The dipole interaction energy between the two induced, fluctuating electric dipoles is then

\[
W_{ee}(R) = \sum_{k\lambda} \alpha_e^A(\omega)\alpha_e^B(\omega) \langle \mathbf{E}_i(\mathbf{r}_A,k\lambda)\mathbf{E}_j(\mathbf{r}_B,k\lambda) \rangle V_{ij}^{ee}(\omega,R),
\]

where \( \mathbf{R} = \mathbf{r}_B - \mathbf{r}_A \), \( R = |\mathbf{R}| \). The two-point vacuum electric-field correlation function, summed over polarization states, follows easily from (6):

\[
\sum_{\lambda} \langle \mathbf{E}_i(\mathbf{r}_A,k\lambda)\mathbf{E}_j(\mathbf{r}_B,k\lambda) \rangle = \frac{2\pi\hbar\omega_k\mu_k}{n_k k} V(\delta_{ij} - \hat{k}_i\hat{k}_j) e^{-i\mathbf{k} \cdot \mathbf{R}},
\]

with \( \hat{k}_i = k_i/k \). The interaction potential \( V_{ij}^{ee}(\omega,R) \) between two oscillating electric dipoles embedded in the magneto-dielectric is calculated as follows. From the Hamiltonian (3) and the Heisenberg equations of motion for the annihilation and creation operators, one obtains a formal expression for the electric field generated by the electric dipole B at the position of the electric dipole A:

\[
\mathbf{E}_e(\mathbf{r}_A,t) = \frac{i}{\pi c^3} \int_0^\infty d\omega n(\omega)\mu(\omega)\omega^3 \left[ a \frac{\sin kR}{kR} - b \left( \frac{\sin kR}{k^3R^3} - \frac{\cos kR}{k^2R^2} \right) \right] dt' p_B(t') e^{i\omega(t'-t)} + h.c.,
\]

where \( a = d_B - (d_B \cdot \mathbf{R}) \mathbf{R}, b = d_B - 3(d_B \cdot \mathbf{R}) \mathbf{R}, \) and \( \mathbf{R} = \mathbf{R}/R \). We have used the notation \( \mathbf{d}(t) = \mathbf{d}(t) \), with \( \mathbf{d} \) a unit vector specifying the direction of the electric dipole moment. We are interested in the case of two electric dipoles, both oscillating at frequency \( \omega' \): \( p_B(t) = C_B e^{-i\omega't} + C_B' e^{i\omega't} \) and \( p_A(t) = C_A e^{-i\omega't} + C_A' e^{i\omega't} \), where \( C_A, C_B \) are arbitrary constants. We define the interaction \( V_{ee}(\mathbf{R}) = -d_A p_A(t) \cdot \mathbf{E}_e(\mathbf{r}_A,t) \). After performing the time integration for times \( t \gg 1/\omega' \), and taking the time average of the resulting expression, we get

\[
V_{ee}(\mathbf{R}) = -\frac{2\pi c^3}{\pi c^3} \text{Re} \int_0^\infty d\omega n(\omega)\mu(\omega)\omega^3 \left[ a \frac{\sin kR}{kR} - b \left( \frac{\sin kR}{k^3R^3} - \frac{\cos kR}{k^2R^2} \right) \right] \left[ \frac{C_A' C_B}{\omega - \omega' - i\eta} + \frac{C_A C_B'}{\omega + \omega' - i\eta} \right],
\]

where as usual \( \eta \to 0^+ \). Here \( a = d_A \cdot d_B - (d_A \cdot \mathbf{R})(d_B \cdot \mathbf{R}) \) and \( b = d_A \cdot d_B - 3(d_A \cdot \mathbf{R})(d_B \cdot \mathbf{R}) \). We can also define the interaction as \( V_{ee}(\mathbf{R}) = -d_A p_B(t) \cdot \mathbf{E}_e(\mathbf{r}_A,t) \), which amounts to interchanging A and B above. This implies that we can take \( C_A' C_B = C_A C_B' \). Note also that the time average \( p_A(t)p_B(t) = C_A C_B' + C_A' C_B \), so that we can write

\[
V_{ee}(\mathbf{R}) = -\frac{1}{\pi c^3} p_A p_B \text{Re} \int_0^\infty d\omega n(\omega)\mu(\omega)\omega^3 \left[ a \frac{\sin kR}{kR} - b \left( \frac{\sin kR}{k^3R^3} - \frac{\cos kR}{k^2R^2} \right) \right] \left( \frac{1}{\omega - \omega' - i\eta} + \frac{1}{\omega + \omega' - i\eta} \right).
\]

Next we use the fact that \( n \) and \( \mu \) (or actually their real parts which are implicit here) are even functions of \( \omega \) to rewrite this as

\[
V_{ee}(\mathbf{R}) = -\frac{1}{\pi c^3} p_A p_B \int_{-\infty}^\infty d\omega n(\omega)\mu(\omega)\omega^3 \left[ a \frac{\sin kR}{kR} - b \left( \frac{\sin kR}{k^3R^3} - \frac{\cos kR}{k^2R^2} \right) \right] \left( \frac{1}{\omega - \omega' - i\eta} + \frac{1}{\omega - \omega' + i\eta} \right).
\]
Performing the trivial contour integrations, writing \( \omega \) instead of \( \omega' \) for the dipole frequencies, and using again \( k = n(\omega) \omega / c \), we obtain the electric dipole-dipole interaction tensor needed in Eq. (1):

\[
V_{ij}^{ee}(\omega, \mathbf{R}) = \frac{1}{\epsilon_k} \frac{1}{R^3} \left[ (\delta_{ij} - \mathbf{3} \hat{\mathbf{R}}_i \hat{\mathbf{R}}_j)(\cos kR + kR \sin kR) - (\delta_{ij} - \hat{\mathbf{R}}_i \hat{\mathbf{R}}_j) k^2 R^2 \cos kR \right].
\] (9)

Using Eqs. (4) and (7), and passing to the continuum limit \( \sum_k \rightarrow (V/8\pi^3) \int_0^\infty dk k^2 \int d\Omega_k \), we obtain the van der Waals interaction energy between electrically polarizable particles:

\[
W_{ee}(R) = -\frac{\hbar}{16\pi R^6} \int_0^\infty du \alpha_e^A(iu) \alpha_e^B(iu) \frac{1}{\epsilon^2(iu)} F \left[ \frac{2n(iu)uR}{c} \right] e^{-2n(iu)uR/c},
\] (11)

where \( F(x) = x^4 + 4x^3 + 20x^2 + 48x + 48 \). Recall that along the imaginary frequency axis the electric permittivity and refractive index are real and positive. The electric-electric van der Waals force that results from Eq. (11) is therefore always attractive, regardless of the frequency dependence of \( \epsilon(\omega) \) and \( n(\omega) \).

Our calculation based on the quantized fields (11) and (2) gives a van der Waals interaction (11) in full agreement with that obtained recently by Tomaš \[4\], for instance. The same is true for the other van der Waals interactions we calculate in this paper. The main point of this section is to show that correct results for van der Waals interactions involving ground-state atoms in dispersive media can be obtained straightforwardly, without having to go to a dilute-medium limit of an interaction between macroscopic bodies, and without having to introduce complexities arising from absorption. We discuss this further in Section IV.

### B. Van der Waals interaction between magnetically polarizable particles

We next use the same approach to calculate the van der Waals interaction between magnetic dipoles induced in the atoms by fluctuations of the zero-point magnetic field. For this the relation between an induced magnetic dipole moment at position \( \mathbf{r} \) and the magnetic field is \( \mathbf{m}(t) = \alpha_m(t) \mathbf{H}(\mathbf{r}, t) \), where \( \alpha_m \) is the (real) magnetic polarizability of the atom. The dipole interaction between the two induced, fluctuating magnetic dipoles is

\[
W_{mm}(R) = \sum_{k\lambda} \alpha^A_m(\omega) \alpha^B_m(\omega) \langle \mathbf{H}_i(r_A, k\lambda) \mathbf{H}_j(r_B, k\lambda) \rangle V_{ij}^{mm}(\omega, \mathbf{R}),
\] (12)

where the two-point vacuum magnetic field correlation function summed over polarizations is found from (2) to be

\[
\sum_{\lambda} \langle \mathbf{H}_i(r_A, k\lambda) \mathbf{H}_j(r_B, k\lambda) \rangle = \frac{2\pi \hbar n_k \omega_k}{\mu_k \gamma_k V} (\delta_{ij} - \hat{\mathbf{k}}_i \hat{\mathbf{k}}_j) e^{-ik \cdot R}.
\] (13)

One can derive the magnetic dipole-dipole interaction tensor following steps similar to those above for the electric dipole-dipole interaction:

\[
V_{ij}^{mm}(\omega, \mathbf{R}) = \frac{1}{\mu_k} \frac{1}{R^3} \left[ (\delta_{ij} - \mathbf{3} \hat{\mathbf{R}}_i \hat{\mathbf{R}}_j)(\cos kR + kR \sin kR) - (\delta_{ij} - \hat{\mathbf{R}}_i \hat{\mathbf{R}}_j) k^2 R^2 \cos kR \right],
\] (14)

which differs from (11) simply by the replacement of \( \epsilon_k \) by \( \mu_k \). The details of the evaluation of (12) are essentially the same as for the electric van der Waals interaction, and lead straightforwardly to the expression

\[
W_{mm}(R) = -\frac{\hbar}{16\pi R^6} \int_0^\infty du \alpha^A_m(iu) \alpha^B_m(iu) \frac{1}{\mu^2(iu)} F \left[ \frac{2n(iu)uR}{c} \right] e^{-2n(iu)uR/c}.
\] (15)

Recall that along the imaginary frequency axis the magnetic permeability is real and positive. The magnetic-magnetic van der Waals force that results from Eq. (15) is always attractive, regardless of the frequency dependence of \( \mu(\omega) \) and \( n(\omega) \).
C. Van der Waals interaction between an electrically polarizable particle and a magnetically polarizable particle

In calculating $W_{ee}(R)$ and $W_{mm}(R)$ it has not been necessary to account for the fact that the field operators in (14) and (12) do not commute. Because of this noncommutativity, it is more appropriate to write $W_{ee}(R)$, for instance, in the symmetrized form

$$W_{ee}(R) = \frac{1}{2} \sum_{k\lambda} \alpha^A_A(\omega)\alpha^B_B(\omega) \left| \langle E_i(r_A, k\lambda)|E_j(r_B, k\lambda) \rangle + \langle E_j(r_B, k\lambda)|E_i(r_A, k\lambda) \rangle \right| V_{ij}^{ee}(\omega, R)$$

$$= \text{Re} \sum_{k\lambda} \alpha^A_A(\omega)\alpha^B_B(\omega) \langle E_i(r_A, k\lambda)|E_j(r_B, k\lambda) \rangle V_{ij}^{ee}(\omega, R).$$

(16)

The forms (5) and (9), however, show that symmetrization is actually not required because the summation over $k$ does not require us to distinguish between $\langle E_i(r_A, k\lambda)|E_j(r_B, k\lambda) \rangle$ and $\langle E_j(r_B, k\lambda)|E_i(r_A, k\lambda) \rangle$.

The situation in the case of the van der Waals interaction between an electrically polarizable particle and a magnetically polarizable particle, however, is different because the electric-magnetic correlation function summed over polarization states,

$$\sum_{\lambda} \langle E_i(r_A, k\lambda)|H_j(r_B, k\lambda) \rangle = \frac{2\pi\hbar \omega_k}{\gamma_k} \epsilon_{ijl} \int e^{-ikR}$$

(17)

is not purely real when summed over $k$. ($\epsilon_{ijl}$ is the Levi-Civita tensor.) Moreover the interaction tensor in this case, which we calculate to be

$$V_{ij}^{em}(\omega, R) = \frac{\omega^3}{c^3} n^2(\omega) \epsilon_{ijl} \int \left[ \frac{\sin kR}{k^2 R^2} - \frac{\cos kR}{kR} \right]$$

in a manner directly analogous to the electric-electric and magnetic-magnetic tensors, is antisymmetric.

Let $E_m(r_A, t)$ be the electric field (operator) at $r_A$ due to a magnetic dipole at $r_B$. We write the interaction between the fluctuating electric and magnetic dipole moments in the symmetrized form

$$W_{em}(R) = -\text{Re} \sum_{k\lambda} \alpha^A_A(\omega) \langle E_i^{(+)}(r_A, k\lambda)|E_m^{(-)}(r_A, k\lambda) \rangle,$$

(19)

where $E^{(+)}(r_A, t)$ is the positive-frequency (photon annihilation) part of the source-free (“vacuum”) electric field operator at $r_A$, and $E_i^{(-)}(r_A, t)$ is the negative-frequency (photon creation) part of the electric field produced by the magnetic dipole moment at $r_B$. This electric field is induced by the source-free magnetic field $H(r_B, t)$, so that the evaluation of (19) involves the electric-magnetic correlation function (17). The calculation is essentially just the same as that presented by Farina et al. [13] for the case where the two particles are in free space, except of course that in our case the refractive index $n(\omega)$ appears:

$$W_{em}(R) = \frac{\hbar}{4\pi c^2 R^4} \int_0^\infty du \omega^2 \alpha^A_A(i\omega)\alpha^B_B(i\omega) G \left[ \frac{2n(i\omega)uR}{c} \right] e^{-2n(i\omega)uR/c},$$

(20)

where $G(x) = (x + 2)^2$. $n(i\omega)$ is real and positive, so that $W_{em}(R)$ is always repulsive, regardless of the frequency dependence of the refractive index.

III. VAN DER WAALS INTERACTIONS FROM QUANTIZED FIELD WITH ABSORPTION

In this section we will calculate the van der Waals interactions considered in the previous section for two atoms embedded in a magneto-dielectric, but now taking absorption in the host medium into account. We use the quantization procedure for the EM field in a dispersive and absorbing medium based on the Green-function formulation [10]. The dyadic Green function $G(r, r', \omega)$ satisfies [10]

$$\left[ \nabla \times \kappa(r, \omega) \nabla \times -\frac{\omega^2}{c^2} \epsilon(r, \omega) \right] G(r, r', \omega) = \delta(r, r', \omega),$$

(21)
as well as the appropriate boundary conditions. Here \( \kappa(r, \omega) = \mu^{-1}(r, \omega) \). In an infinite, homogeneous material,

\[
G_{ij}(r, r', \omega) = \frac{\mu(\omega)}{4\pi k^2(\omega)} \left[ k^2(\omega)(\delta_{ij} - \hat{R}_i \hat{R}_j) - (\delta_{ij} - 3\hat{R}_i \hat{R}_j) \left( \frac{1}{R} - \frac{ik(\omega)}{R} \right) \right] \frac{e^{ik(\omega)R}}{R},
\]

(22)

where \( \mathbf{R} = r - r' \), \( R = |\mathbf{R}| \), \( \hat{\mathbf{R}} = \mathbf{R}/R \), and \( k(\omega) = n(\omega)c/\omega \). The refractive index of the medium is given by \( n^2(\omega) = \epsilon(\omega)\mu(\omega) \), with \( \epsilon(\omega) = \epsilon' + i\epsilon''(\omega) \) the complex electric permittivity and \( \mu(\omega) = \mu' + i\mu''(\omega) \) the complex magnetic permeability.

The quantized electric field in dispersive, absorbing media may be written in the form

\[
\mathbf{E}(r, \omega) = \sum_{\lambda=e,m} \int d^3r' \mathbf{G}_\lambda(r, r', \omega) \cdot \mathbf{f}_\lambda(r', \omega) + h.c.,
\]

(23)

where the operators \( \mathbf{f}_\lambda \) are bosonic operators satisfying the usual commutation relations:

\[
[f_{\lambda,i}(r, \omega), f_{\lambda',j}^\dagger(r', \omega')] = \delta_{\lambda,\lambda'}\delta_{ij}\delta(\omega - \omega') \quad ; \quad [f_{\lambda,i}(r, \omega), f_{\lambda',j}(r', \omega')] = 0.
\]

(24)

These operators may be regarded as being variables of the system composed of the EM field and the medium including the dissipative system. The electric and magnetic dyadic Green functions are defined in terms of the full Green function as

\[
\mathbf{G}_e(r, r', \omega) = \frac{i\omega^2}{c^2} \sqrt{\frac{\hbar}{\pi}} \Im \epsilon(\omega) \mathbf{G}(r, r', \omega),
\]

\[
\mathbf{G}_m(r, r', \omega) = -\frac{i\omega^2}{c^2} \sqrt{\frac{\hbar}{\pi}} \Im \kappa(\omega) [\mathbf{G}(r, r', \omega) \times \nabla r'].
\]

(25)

(26)

Note that for an absorbing medium \( \Im \epsilon(\omega) > 0 \), \( \Im \mu(\omega) > 0 \), and \( \Im \kappa(\omega) < 0 \). The quantized magnetic field, similarly, may be written as

\[
\mathbf{H}(r, \omega) = \sum_{\lambda=e,m} \frac{c}{i\omega\mu(\omega)} \int d^3r' \mathbf{r} \times \mathbf{G}_\lambda(r, r', \omega) \cdot \mathbf{f}_\lambda(r', \omega) + h.c.,
\]

(27)

and the total Hamiltonian for the free field is

\[
\hat{H} = \sum_{\lambda=e,m} \int d^3r \int_0^\infty d\omega \hbar \omega \mathbf{f}_\lambda^\dagger(r, \omega) \cdot \mathbf{f}_\lambda(r, \omega).
\]

(28)

We will require the following two-point vacuum field correlation functions obtained from these expressions:

\[
\langle \mathbf{E}_i(r, \omega) \mathbf{E}_j^\dagger(r', \omega') \rangle = \frac{\hbar \omega^2}{\pi c^2} \delta(\omega - \omega') \Im [\mathbf{G}(r, r', \omega)]_{ij},
\]

\[
\langle \mathbf{H}_i(r, \omega) \mathbf{H}_j^\dagger(r', \omega') \rangle = \frac{\hbar}{\pi} \frac{1}{|\mu(\omega)|^2} \delta(\omega - \omega') \Im [\nabla_r \times \nabla r' \times \mathbf{G}(r, r', \omega)]_{ij},
\]

\[
\langle \mathbf{E}_i(r, \omega) \mathbf{H}_j^\dagger(r', \omega') \rangle = -\frac{1}{i\omega\mu(\omega)} \frac{\hbar \omega^2}{\pi} \delta(\omega - \omega') \Im [\nabla_r \times \mathbf{G}(r, r', \omega)]_{ij}.
\]

(29)

(30)

(31)

The electric and magnetic fields at the position of the atom A are given by the sum of the “vacuum” contributions, \( \mathbf{E}_0(r_A, \omega) \) and \( \mathbf{H}_0(r_A, \omega) \), plus the fields generated by the atom B which contain both electric dipole and magnetic dipole components. An electric dipole located at position \( r_B \) generates fields at position \( r_A \) given by

\[
\mathbf{E}_e(r_A, \omega) = \omega^2 \alpha_e^B(\omega) \mathbf{G}(r_A, r_B, \omega) \cdot \mathbf{E}_0(r_B, \omega) + h.c.,
\]

\[
\mathbf{H}_e(r_A, \omega) = -i\omega \kappa(\omega) \alpha_e^B(\omega) [\nabla_{r_A} \times \mathbf{G}(r_A, r_B, \omega)] \cdot \mathbf{E}_0(r_B, \omega) + h.c.,
\]

(32)

(33)

while the fields generated by a magnetic dipole at position \( r_B \) are

\[
\mathbf{E}_m(r_A, \omega) = -i\omega \kappa(\omega) \alpha_m^B(\omega) [\nabla_{r_A} \times \mathbf{G}(r_A, r_B, \omega)] \cdot \mathbf{H}_0(r_B, \omega) + h.c.,
\]

\[
\mathbf{H}_m(r_A, \omega) = c\kappa(\omega) k^2(\omega) \alpha_m^B(\omega) \mathbf{G}(r_A, r_B, \omega) \cdot \mathbf{H}_0(r_B, \omega) + h.c.
\]

(34)

(35)
The vacuum expectation value of the van der Waals energy may be written as the sum of three contributions, one purely electric, one purely magnetic, and one mixed. Using the above expressions for the two-point correlation functions of the EM field, one can easily find each of these terms. The purely electric part stems from the $p \cdot E$ interaction, and is found to be

$$W_{ee}(R) = \frac{1}{2} \int_0^\infty d\omega d\omega' \alpha_e^A(\omega)e^{i(\omega-\omega')t}\langle \hat{E}_0(r_A, \omega) \cdot \hat{E}_1^\dagger(r_A, \omega') \rangle + h.c.$$  

$$= -\frac{\hbar}{2\pi} \int_0^\infty d\omega \alpha_e^A(\omega)\alpha_e^B(\omega)\omega^4 Re [G(r_A, r_B, \omega)]_{ij} \text{Im} [G(r_A, r_B, \omega)]_{ij}.$$  

Comparing this expression with Eq. (4) we see that $\omega^2 \text{Im} [G(r_A, r_B, \omega)]_{ij}$ is related to the electric dipole-dipole interaction tensor $V_{ij}(\omega, R)$, and that $\omega^2 \text{Re} [G(r_A, r_B, \omega)]_{ij}$ is related to the solid-angle integration of the two-point vacuum electric-field correlation function summed over polarization states, given in Eq. (5). After rotation in the complex plane ($\omega \rightarrow i\omega$), we can re-write this expression as

$$W_{ee}(R) = -\frac{\hbar}{2\pi} \int_0^\infty du \alpha_e^A(iu)\alpha_e^B(iu)u^4 \text{Tr} [G(r_A, r_B, iu) \cdot G(r_A, r_B, iu)].$$  

The purely magnetic part, similarly, comes from the $m \cdot H$ interaction:

$$W_{mm}(R) = -\frac{1}{2} \int_0^\infty d\omega d\omega' \alpha_m^A(\omega)e^{i(\omega-\omega')t}\langle \hat{H}_0(r_A, \omega) \cdot \hat{H}_1^\dagger(r_A, \omega') \rangle + h.c.$$  

$$= -\frac{\hbar}{2\pi} \int_0^\infty d\omega \alpha_m^A(\omega)\alpha_m^B(\omega)\omega^2 \text{Re} [\kappa(\omega)k^2(\omega)G(r_A, r_B, \omega)]_{ij} \text{Im} [\nabla_{r_A} \times \nabla_{r_B} \times G(r_A, r_B, \omega)]_{ij}$$  

$$= -\frac{\hbar}{2\pi} \int_0^\infty du \alpha_m^A(iu)\alpha_m^B(iu)u^2 \text{Tr} [G(r_A, r_B, iu) \cdot G(r_A, r_B, iu)].$$  

where Eq. (21) and a rotation in the complex plane were used in obtaining the last equality. Finally, there are two electric-magnetic terms, one arising from the $p \cdot E$ interaction, and one from the $m \cdot H$ interaction. They result in the mixed interaction

$$W_{em}(R) = -\frac{1}{2} \int_0^\infty d\omega d\omega' e^{i(\omega-\omega')t}[\alpha_e^A(\omega)\langle \hat{E}_0(r_A, \omega) \cdot \hat{E}_1^\dagger(r_A, \omega') \rangle + \alpha_m^A(\omega)\langle \hat{H}_0(r_A, \omega) \cdot \hat{H}_1^\dagger(r_A, \omega') \rangle + h.c.]$$  

$$= \frac{\hbar}{\pi} \int_0^\infty d\omega [\alpha_e^A(\omega)\alpha_m^B(\omega) + \alpha_m^A(\omega)\alpha_e^B(\omega)]\omega^2 \text{Re} [\kappa(\omega)G(r_A, r_B, \omega)]_{ij} \text{Im} [\nabla_{r_A} \times \nabla_{r_B} \times G(r_A, r_B, \omega)]_{ij}$$  

$$= \frac{\hbar}{2\pi} \int_0^\infty du [\alpha_e^A(iu)\alpha_m^B(iu) + \alpha_m^A(iu)\alpha_e^B(iu)]u^2 \text{Tr} [\nabla_{r_A} \times G(r_A, r_B, iu) \cdot \nabla_{r_B} \times G(r_A, r_B, iu)].$$  

The traces appearing in the integrands of Eqs. (37, 38, 39) can be explicitly computed given the form of the dyadic Green function evaluated at the imaginary frequency $w = i\omega$. The final result for the complete van der Waals interaction energy between two ground state atoms embedded in an absorbing and dispersive medium is then

$$W(R) = -\frac{\hbar}{16\pi R^5} \int_0^\infty du e^{-2n(iu)aR/c} F \left[ \frac{2n(iu)R}{c} \left[ \frac{\alpha_e^A(iu)\alpha_e^B(iu)}{e^2(iu)} + \frac{\alpha_m^A(iu)\alpha_m^B(iu)}{\mu^2(iu)} \right] \right]$$  

$$+ \frac{\hbar}{4\pi R^4} \int_0^\infty du e^{-2n(iu)R/c} G \left[ \frac{2n(iu)R}{c} \left[ \frac{\alpha_e^A(iu)\alpha_e^B(iu)}{e^2(iu)} + \frac{\alpha_m^A(iu)\alpha_m^B(iu)}{\mu^2(iu)} \right] \right],$$  

where again $F(z) = z^4 + 4z^3 + 20z^2 + 48z + 48$ and $G(z) = (z + 2)^2$. This is identical to the complete van der Waals interaction obtained in Section III

IV. DISCUSSION

Since the van der Waals interaction between electrically polarizable particles is the most important, a rough model for the modification of the vacuum interaction by the medium might be of interest. Let us consider a two-level model in which the polarizabilities are

$$\alpha_e^A(\omega) = \alpha_e^B(\omega) \equiv \alpha(\omega) = \frac{2\omega_0 d^2}{\omega_0^2 - \omega^2}.$$  

(41)
where $d$ and $\omega_0$ are respectively the (real) transition electric dipole moment and the transition angular frequency, and similarly

$$n(\omega) = [1 + 4\pi N\alpha(\omega)]^{1/2},$$

where $N$ is the atomic density of the host medium and we take $\mu = 1$. It is useful to normalize (11) to the familiar, nonretarded London form of the interaction:

$$W_L(R) = -\frac{3h\omega_0\alpha^2(0)}{4R^6} = -\frac{3h\omega_0}{4R^6} \left( \frac{2d^2}{3h\omega_0} \right)^2.$$

We define

$$D(R) \equiv \frac{W_{ee}(R)}{W_L(R)} = \frac{4}{3\pi} \int_0^\infty dy \left( \frac{1}{y^2 + 1} \right)^2 \frac{1}{\epsilon^2} \left[ n^4 r^4 y^4 + 2n^3 r^3 y^3 + 5n^2 r^2 y^2 + 6nry + 3 \right] e^{-2nry},$$

where $r \equiv \omega_0R/c$ and $\epsilon$ and $n$ are evaluated at $i\omega_0y$:

$$\epsilon(i\omega_0y) = n^2(i\omega_0y) = 1 + \frac{C}{y^2 + 1},$$

where $C = 8\pi N\alpha^2/3h\omega_0$. In the limit $C = 0$ and $r \to 0$, $D \to 1$; for $r \to \infty$, $D \to 23/3\pi r$, or $W_{ee} = -23h\alpha\alpha^2(0)/4\pi R^7$, the famous Casimir-Polder result. Figure 1 plots $D(R)$ for $C = 0$ (vacuum) and $C = 3$.

This simple model is not in any sense meant to be a realistic, quantitative description of the van der Waals interaction between two atoms embedded in a dielectric. It does, however, suggest that the predominant effect of the host medium on the van der Waals interaction is to weaken it, without substantially changing the distance dependence in either the nonretarded or retarded regimes. More realistic models of the van der Waals interaction in a liquid, for example, must take into account local field corrections, as has been discussed, for instance, by Abrikosov, et al. [8] and McLachlan [14].

Our results for the electric-electric, magnetic-magnetic, and electric-magnetic van der Waals interactions between two atoms embedded in a dispersive magneto-dielectric medium are in agreement with those obtained previously [3]-[7]. Unlike previous derivations, however, we obtained the interactions directly rather than inferring them from the dilute-medium limit of van der Waals interactions between macroscopic bodies, where the interactions can be obtained by pairwise summations of interatomic interactions [1]-[8].

What is perhaps more interesting, however, concerns the role of absorption in the host medium. In the preceding section we accounted for dissipation (absorption) in the host magneto-dielectric, as in previous work. In Section [1] however, we ignored any possibility that the medium could be absorbing; we worked with expressions for the electric and magnetic fields that derive directly from the assumption that the medium is non-absorbing [11].

The same situation holds, for instance, in the case of the Lifshitz formula for the van der Waals interaction between infinite, plane-parallel dielectric media. Lifshitz’s original derivation, and various derivations that followed, include the imaginary (dissipative) part of the dielectric function $\epsilon(\omega)$. Derivations of the Lifshitz formula based on changes...
in zero-point field energy arising from the dielectric media, however, make no reference to dissipation. In other words, the Lifshitz formula can be derived without explicit accounting for absorption.

Ginzburg has also noted that (macroscopic) van der Waals interactions can be correctly derived based on changes in zero-point field energy, without accounting for absorption. He remarks, in connection with such derivations, that “oddly enough there is no mention that they consider directly only transparent media,” and then gives reasons why the van der Waals (electric-electric) interaction for real media can be obtained by presuming non-absorbing media: “Firstly, the permittivities are functions. Secondly, the function is always real on the imaginary axis.”

The results all change sign at frequencies for which the dipole-dipole interaction in a negative-index medium does not change sign or undergo any other significant change in a magneto-dielectric medium. The same conclusion applies to the refractive index is negative. To see that this is not the case, note that Eq. (8) is unchanged whenever the permeability and permeability on the imaginary axis.

A simple and more physical explanation can be given for why (ground-state) van der Waals interactions calculated for non-absorbing media apply directly to real (absorbing) media, as we have found for the electric-electric, magnetic-magnetic, and electric-magnetic van der Waals interactions between two atoms embedded in a magnetic-dielectric medium. At zero-temperature, for instance, any atom of the host medium is in its ground state and can absorb radiation that is resonant with one of its transitions to an excited state. It cannot, of course, absorb from the vacuum field: in this case the fluctuations in the field that might induce absorption are exactly cancelled by fluctuations in the atom itself. The same is true when the atom is part of a magneto-dielectric in which is embedded, as in examples considered in this paper, two guest atoms. Any atom of the host medium still finds itself in a vacuum field state, regardless of the nature or the number of guest atoms. The host and guest atoms modify the modes of the field from the simple plane waves of a pure vacuum, but the field remains in a vacuum state for any mode. Just as in free space, therefore, there is no absorption unless it is possible to populate one or more field modes; this would be the case only if there were an applied external field or if one or more atoms is excited, leading to the possibility that a different atom could absorb its emitted (real) photon. In other words, absorption by the host medium would play a role if we were to consider a van der Waals interaction involving excited atoms. Otherwise one can expect to obtain correct van der Waals interactions without having to account for the absorption that is always present in a real medium. This expectation applies, of course, regardless of how many atoms are involved and regardless of the shape of any macroscopic bodies for which the van der Waals forces are to be calculated.

In light of recent interest in negative-index media, let us reconsider specifically the most important of the van der Waals interactions we have calculated in this paper, namely that between two electrically polarizable atoms. First we note that the expressions for the quantized fields in Section are directly applicable to negative-index media. In a negative index medium , , and are all negative at some frequency or range of frequencies. This would at first glance suggest that the dipole-dipole interaction changes sign at frequencies for which the refractive index is negative. To see that this is not the case, note that Eq. is unchanged whenever , , and all change sign within any frequency range. This means that the dipole-dipole interaction in a negative-index medium does not change sign, and in particular that is directly applicable in general provided we just replace , , and by their absolute values. It follows similarly that the van der Waals interaction does not change sign or undergo any other significant change in a magnet-dielectric medium. The same conclusion applies to the magnetic-magnetic and mixed van der Waals interactions, and is in agreement with the conclusions of Buhmann et al.

The pairwise electric-electric and magnetic-magnetic van der Waals interactions are always attractive, whereas the pairwise electric-magnetic interaction is always repulsive. These results apply also in the case of negative-index media, at least to the extent that such media can be modeled as continua. It is well known, however, that non-pairwise van der Waals interactions can be repulsive. Evidently repulsive Casimir effects such as those recently suggested by Henkel and Joulain and by Leonhardt and Philbin in the case of negative-index media must in some way involve either non-pairwise interactions or electric-magnetic van der Waals interactions.
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