Near-field induction heating of metallic nanoparticles due to infrared magnetic dipole contribution

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We revisit the electromagnetic heat transfer between a metallic nanoparticle and a metallic semi-infinite substrate, commonly studied using the electric dipole approximation. For infrared and microwave frequencies, we find that the magnetic polarizability of the particle is larger than the electric one. We also find that the local density of states in the near field is dominated by the magnetic contribution. As a consequence, the power absorbed by the particle in the near field is due to dissipation by fluctuating eddy currents. These results show that a number of near-field effects involving metallic particles should be affected by the fluctuating magnetic fields.

I. INTRODUCTION

A lot of attention has recently been devoted to the interaction between nano-objects like atoms, nanoparticles or AFM tips, and surfaces, which is mediated by fluctuating thermal fields. A great variety of phenomena like Casimir-Polder forces, friction forces, or near-field heat transfer is governed by the associated stochastic thermal currents. A common assumption is that the electric dipole approximation can be used to model the nano-object because it varies like \((R/\lambda)^3\) whereas the next term in the Mie expansion varies as \((R/\lambda)^5\) (\(\lambda\) is the wavelength in vacuum). In this work, we will show that the interaction between the magnetic dipole and the large magnetic fields in the near field may give the dominant contribution to the heat transfer.

II. ABSORPTION BY A METALLIC NANOPARTICLE

Let us first compute the power absorbed by a small metallic particle. In what follows, we will use an isotropic, homogeneous and local form of the complex dielectric constant \(\varepsilon_{r}\) is generally assumed to give the leading contribution because it varies like \((R/\lambda)^3\) whereas the next term in the Mie expansion varies as \((R/\lambda)^5\) (\(\lambda\) is the wavelength in vacuum). In this work, we will show that the interaction between the magnetic dipole and the large magnetic fields in the near field may give the dominant contribution to the heat transfer.

In the next section, we compare the absorption cross section of the electric and magnetic dipole moment. The third section is devoted to the analysis of the electric and magnetic energy density in the near field of a metal-vacuum interface. The final section analyses the heat transfer and discusses the physical mechanism.
where $\epsilon_0$ is the dielectric permittivity in vacuum and $\vec{E}$ is the external electric field. Another contribution is given by the magnetic dipolar moment $\vec{m}$ characterized by its magnetic polarizability $\alpha_H$

$$\vec{m} = \alpha_H \vec{H},$$

(2)

where $\vec{H}$ is the external magnetic field. Higher multipoles can be neglected if $|\epsilon_r| \gg 1$ and $R/\lambda \ll 1$. The contributions of the electric and magnetic dipoles to the power dissipated in the particle at a positive frequency $\omega$ are given by [35, 36]:

$$P_{abs}^E(\omega) = \omega 2 Im(\alpha_E) \epsilon_0 \frac{\langle |\vec{E}|^2 \rangle}{2},$$

(3)

$$P_{abs}^M(\omega) = \omega 2 Im(\alpha_H) \mu_0 \frac{\langle |\vec{H}|^2 \rangle}{2},$$

(4)

where $\mu_0$ is the magnetic permeability in vacuum. $\alpha_E$ and $\alpha_H$ can be found in ref. [37]:

$$\alpha_E = 4\pi R^3 \frac{\epsilon_r - 1}{\epsilon_r + 2},$$

(5)

$$\alpha_H = \frac{2\pi}{15} R^3 \left(\frac{2\pi R}{\lambda}\right)^2 (\epsilon_r - 1),$$

(6)

where $\epsilon_r$ is the relative dielectric permittivity.

Here, we do not take into account the diamagnetism of the material. Instead, the magnetic dipole moment is due to eddy currents in the particle. The polarizabilities are calculated assuming that $R$ is much smaller than the skin depth $\delta$. A different form [35] can be derived when dealing with particles such that $\delta \ll R \ll \lambda$. In what follows, we should keep in mind that the dipole model is a fair approximation provided that the distance $d$ between the center of the particle and a surface is much larger than $R$. Note that we have used for simplicity the extinction cross section of the electric dipole. The exact form of the absorption cross section is discussed in ref. [38]. The difference for a metallic nanoparticle is negligible.

As seen from Eqs (3,4), the absorption is the product of two terms, the imaginary part of the polarizability and the local density of energy. We shall show that for metallic nanoparticles at low frequencies, both terms are larger for the magnetic contribution. Let us first analyse the role of the polarizability. It appears from Eqs (3-6) that for values of the dielectric constant on the order of unity, the electric dipole contribution to losses is much larger than the magnetic one because $R/\lambda \ll 1$. Yet, for values $\epsilon_r > 1$ such that $|\epsilon_r| > 1$, as it is the case for metals at low frequencies, the magnetic dipole may provide the leading contribution. The physical reason is that the magnetic fields are continuous at an interface so that they can penetrate in the material. By contrast, the electric field in a spherical particle $E_{int}$ is related to the external electric field by $E_{int} = [3/(\epsilon_r + 2)]E_{ext}$. Surface charges induced at the interface prevent the electric field to penetrate efficiently in the metallic particle. This screening effect takes place on a length scale given by the Thomas-Fermi length. It does not depend on the skin depth.

We consider a non-magnetic metallic particle characterized by a Drude model $\epsilon_r = 1 - \omega_p^2/(\omega^2 + i \omega \nu)$ where $\omega_p$ is the plasma frequency and $\nu$ is the damping coefficient. To account for the confinement effects, the bulk dielectric constant $\epsilon_r$ [35] is corrected by modifying the damping constant $\nu = \nu_0 + A v_F/R$, where $\nu_0$ is the bulk damping coefficient, $v_F$ the Fermi velocity and $A$ a sample-dependent coefficient. Figure 1 shows $Im(\alpha_E)$ and $Im(\alpha_H)$ as a function of circular frequency for two gold spheres with radii $R = 5 \text{ nm}$ and $R = 10 \text{ nm}$. It is seen that the electric polarizability is larger than the magnetic polarizability at optical frequencies. As explained before, this is no longer the case at low frequencies (typically smaller than $\nu$), where $Im(\alpha_H)$ is larger than $Im(\alpha_E)$.
III. LOCAL DENSITY OF ENERGY NEAR A METALLIC SURFACE

To derive the energy absorbed by a particle in the vicinity of an interface, we need to consider the local densities of energy \( \epsilon_0 \frac{\langle |E|^{2} \rangle}{2} \) and \( \mu_0 \frac{\langle |H|^{2} \rangle}{2} \). In a vacuum, both contributions are equal. The energy per unit volume \( U(z, \omega) \) at a distance \( z \) from the interface increases dramatically in the near field due to the presence of evanescent waves as discussed in ref. [36, 40]. \( U(z, \omega) \) is the product of the local density of states (LDOS) \( \rho(z, \omega) \) by the mean energy of a mode \( \Theta(\omega, T) = \hbar \omega / [\exp(\hbar \omega / k_B T) - 1] \), where \( 2\pi \hbar \) is Planck constant, \( k_B \) is Boltzmann constant and \( T \) the temperature of the substrate. The final expression for the evanescent part of the LDOS [36, 40] is the sum of the four following contributions:

\[
\rho_{s}^{E}(z, \omega) = \rho_{o} \int_{\omega/c}^{\infty} \frac{dK}{2|\gamma_{0}|} \frac{cK}{\omega} \text{Im}(r_{s}) e^{-2\gamma_{s}z} \quad (7)
\]

\[
\rho_{s}^{M}(z, \omega) = \rho_{o} \int_{\omega/c}^{\infty} \frac{dK}{2|\gamma_{0}|} \frac{cK}{\omega} f(K, \omega) \text{Im}(r_{s}) e^{-2\gamma_{s}z} \quad (8)
\]

\[
\rho_{p}^{E}(z, \omega) = \rho_{o} \int_{\omega/c}^{\infty} \frac{dK}{2|\gamma_{0}|} \frac{cK}{\omega} f(K, \omega) \text{Im}(r_{p}) e^{-2\gamma_{p}z} \quad (9)
\]

\[
\rho_{p}^{M}(z, \omega) = \rho_{o} \int_{\omega/c}^{\infty} \frac{dK}{2|\gamma_{0}|} \frac{cK}{\omega} \text{Im}(r_{p}) e^{-2\gamma_{p}z} \quad (10)
\]

where the subscripts \( E \) and \( M \) denote the electric and magnetic evanescent components, \( c \) is the light velocity in vacuum, \( \rho_{o}(\omega) = \omega^{2}/\pi^{2} c^{3} \) is the vacuum density of states, \( f(K, \omega) = 2|\gamma_{0}| \omega / \left( \omega^{2} - \omega_{s}^{2} \right)^{2} - 1, \right) \) \( \gamma_{s} = \frac{\gamma_{0}}{\omega^{2} - \omega_{s}^{2}} \) and \( \gamma_{M} = \frac{-\gamma_{0}}{\omega_{s}^{2}} \) are the Fresnel TE and TM reflection factors, and the complex number \( \gamma_{s} = \gamma_{s}^{pr} + i \gamma_{s}^{ev} \) is defined as the perpendicular part of the wave vector at a frequency \( \omega : K^{2} + \gamma_{0}^{2} = \epsilon_{s} \omega^{2} \) where \( i = 0 \) in vacuum (\( \epsilon_{0} = 1 \)) and \( i = 1 \) in the metal. We have neglected non-local effects as we consider distances larger than the Thomas-Fermi screening length.

On Figure 2, we plot the LDOS versus the frequency \( \omega \) for \( d = 30 \) nm. The first conclusion is that the contribution due to the evanescent waves dominates. The second conclusion is that the s-polarized magnetic contribution is dominant for frequencies below \( \omega_{M} = 2.4 \times 10^{14} \text{s}^{-1} \) which are relevant for heat transfer at 300 K. Indeed, in the expression of the energy density \( U(z, \omega), \Theta(\omega, T) \) acts as a temperature-dependent frequency filter. At a given temperature, we define a cut-off frequency \( \omega_{M} \) by \( f_{\omega_{M}}^{\infty} \Theta(\omega, T) d\omega / \int_{0}^{\infty} \Theta(\omega, T) d\omega = 99 / 100 \). Frequencies much higher than \( \omega_{M} \) are not relevant for heat transfer. We note that the p-polarized contribution associated with the surface plasmon polariton dominates at optical frequencies but does not contribute significantly in the infrared.

The dominant contribution of magnetic energy can be understood by considering the analytical expressions of \( \rho_{s}^{M} \) (Eq. 8) and \( \rho_{p}^{E} \) (Eq. 9). Both expressions are exactly symmetric, involving the same factor \( f(K, \omega) \), and the imaginary part of the reflection factor, respectively \( \text{Im}(r_{s}) \) and \( \text{Im}(r_{p}) \). The physical origin of the factor \( f(K, \omega) \) lies in a fundamental difference of structure between propagating and evanescent waves. It is interesting to see why the magnetic energy dominates the electric energy. Indeed, in a vacuum, the electric and the magnetic energy are equal. This is no longer true for an s-polarized evanescent wave close to an interface.

Let us denote the wave vector as follows

\[
\vec{K}_{0} = K \, \vec{e}_{x} - \gamma \, \vec{e}_{z}
\]

where \( K \) and \( \gamma \) are the interface parallel and perpendicular wave vectors and \( \vec{e}_{x} \) and \( \vec{e}_{z} \) unit vectors. The Helmholtz equation in a vacuum yields \( K^{2} + \gamma^{2} = \omega^{2} / c^{2} \). For an s-polarized field, the electric field is given by

\[
\vec{E} = (0, E, 0), \quad (12)
\]

and the magnetic field follows from the Maxwell-Faraday equation in vacuum (\( \text{curl} \, \vec{E} = -\frac{\partial}{\partial t} \vec{B} \))

\[
\vec{B} = \frac{E}{\omega}(\vec{e}_{x} \times \gamma, 0, K). \quad (13)
\]
It follows that \(|\vec{B}|^2 = \vec{B} \cdot \vec{B}^* = \frac{|E|^2}{c^2} (|\gamma|^2 + K^2)\). For a propagating wave, this yields the well-known result
\[
|\vec{B}| = c |\vec{E}|, \tag{14}
\]
whereas for an evanescent wave, \(\gamma\) is purely imaginary so that \(|\gamma|^2 = -\gamma^2\). We get:
\[
\frac{|\vec{B}|}{c |\vec{E}|} = \sqrt{\frac{2K^2}{k_0^2} - 1}. \tag{15}
\]
For evanescent waves, \(K \gg k_0\) so we find that the magnetic energy stored in an s-polarized evanescent wave is much larger than the electric energy. We have plotted in Figure 3 the function \(\frac{B}{cE} = f(K, \omega) = \sqrt{\frac{|\gamma|^2 + K^2}{k_0^2}}\) for different frequencies. It is seen that the density of energy, which is proportional to the density of states, is driven by the magnetic contribution.

\[\text{FIG. 3: Ratio } \frac{|\vec{B}|}{c|\vec{E}|} \text{ for s-polarized evanescent waves.}\]

Now, a similar reasoning can be done for p-polarized waves. In this case, the inverse ratio \(\frac{|\vec{E}|}{c|\vec{B}|}\) is also equal to \(\sqrt{f(K, \omega)} \approx \sqrt{2} \frac{K}{k_0}\) which shows that the electric field dominates in this case. Near field is thus always dominated by an s-polarized evanescent magnetic field and a p-polarized evanescent electric field. The relative weight of both contributions is then given by the values of the imaginary part of the reflection factors. Since \(\text{Im}(r_s)\) is larger than \(\text{Im}(r_p)\) for a metal at low frequencies, the s-polarized contribution to the LDOS dominates as seen in Figure 2. As a consequence, the energy density is dominated by its s-polarized magnetic contribution.

\[\text{FIG. 4: Radiative power dissipated in the gold particle (radius } R = 5 \text{ nm)} \text{ by the semi-infinite planar gold substrate at 300 K, and the asymptotic behaviours}\]
The above analysis can be summarized by the following scenario. Random currents flowing parallel to the interface can excite the s-polarized evanescent electromagnetic fields at infrared frequencies. As explained above, the associated magnetic fields take large values in the near-field. They are continuous across a vacuum-metal interface so that they penetrate efficiently in the nanoparticle and can generate large eddy currents. These currents are dissipated through Joule effect. Thus, thermal heat transfer appears to be due to near-field induction heating. Radiative heat transfer between two parallel metallic surfaces can also be revisited with a similar scenario [41].

V. CONCLUSION

In summary, we have shown that the heat transfer between a metallic nanoparticle and a metallic surface is dominated by the magnetic contribution. Heat is mainly dissipated by fluctuating eddy currents. The widely used electric dipole approximation is valid for dielectrics but breaks down for metals. As a consequence, the $1/z^3$ dependence of the flux between dielectrics is not valid for metals. A number of other effects due to thermal radiation (e.g. forces, friction) between metallic bodies are expected to be driven by their magnetic contribution, even if the media are non-magnetic. We note that the heat exchanged by two metallic nanoparticles separated by a submicronic distance should be driven as well by the interaction between their magnetic dipoles.

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