Scattering of evanescent wave by two cylinders near a flat boundary

O. V. Belai¹, L. L. Frumin¹,³, S. V. Perminov² and D. A. Shapiro¹,³

¹ Institute of Automation and Electrometry, Siberian Branch, Russian Academy of Sciences
² A.V. Rzhanov Institute of Semiconductor Physics, Siberian Branch, Russian Academy of Sciences
³ Novosibirsk State University - 2 Pirogov Street, Novosibirsk 630090, Russia

received 7 October 2011; accepted in final form 23 November 2011
published online 3 January 2012

PACS 02.70.Pt – Boundary-integral methods
PACS 42.25.Fx – Diffraction and scattering
PACS 78.67.Qa – Nanorods

Abstract – The two-dimensional problem of the evanescent wave scattering by dielectric or metallic cylinders near the interface between two dielectric media is solved. A semianalytical method involving a special Green function and a numerical solution of the boundary integral equations is proposed. A configuration with a circular and a prolate elliptic cylinders is suggested to simulate the sample and the probe in near-field optical microscopy. The far-field energy flux through the probe is calculated as a function of its position. The oscillations of the signal are interpreted as a result of the interference between evanescent and cylindrical waves.

The diffraction limit in optics, known to originate from the wave nature of light, gives a striking example of a physical restriction being a target of increasing efforts to overcome. More than a century ago it was realized that the wavelength limits the smallest spot the electromagnetic energy can be localized within, as well as the smallest details one can optically resolve are comparable to the wavelength. However, further study showed that these obstacles, in fact, concern a traveling (homogeneous) electromagnetic wave. Unlike the inhomogeneous (also referred to as evanescent) waves, which cannot propagate far away from their source, open the way to clearly see the limits due to diffraction and go towards the optics of tiny objects. For instance, the nanosized highly polarizable (i.e. metal) particles well manage to concentrate the light energy within a few-nanometer range [1].

A near-field scanning optical microscopy (NSOM) was suggested to obtain optical signals from the objects at nanoscale (see [2,3] and references therein) using sharp tips; the latter serve much like an optical antenna [4], which receives the energy of the local field and then transmits it to a detector. Thus, nanophotonics is, basically, an optics of evanescent waves, and, consequently, the fundamental optical processes (such as diffraction, interference, scattering) are to be reconsidered.

In the past two decades a substantial progress has been achieved in nano-optics [5–7]. However, a significant methodological deficiency persists even for plain, basic problems, like the scattering of evanescent wave by a body. The trouble is that the evanescent wave cannot be considered in isolation from its source (for instance, the interface where the total internal reflection takes place), therefore the source is certainly affected by the scatterer as being located within a few wavelengths. In paper [8] a general analytical approach is suggested that makes it possible to do very effective calculations of the evanescent wave scattering on a 2D particle (a cylinder) near a flat boundary. In the present work we take the next step and consider the problem of two optically coupled objects placed into the inhomogeneous wave. The illumination by evanescent waves is the basic scheme in near-field optics, including the total internal reflection microscopy and tomography [9–12]. Thus, our main goal is to get a physical insight into the near-field scanning optical microscopy, which minimally involves two small bodies—the studied object and the probe. We believe our work is a promising starting point for the analysis of particular NSOM schemes which will allow for the correct extraction of the near field and structural information from NSOM data. Keeping in mind this application, we
should focus our attention on first-principles approaches, avoiding restricting assumptions and approximations.

We start from the homogeneous scalar two-dimensional Helmholtz equation

\[(\Delta + k^2)\mathcal{H} = 0,\tag{1}\]

where \(\Delta\) is the Laplace operator, \(k\) is the wave vector. The scalar approach is fully adequate to our case of isotropic susceptibilities under an appropriate polarization choice. Otherwise, the vector equation should be studied instead of (1). Consider a domain \(D\) with permittivity \(\varepsilon_{\text{in}}\) and its boundary \(\Gamma = \partial D\). We assume the boundary being twice continuously differentiable. Let us denote as \(\varepsilon_{\text{out}}\) the permittivity of the exterior of \(D\). The Green theorem can be written inside the domain \(D\):

\[\mathcal{H}(\mathbf{r}) = -\oint_{\Gamma} \left( \frac{\partial g}{\partial n'} \mathcal{H}' - g \frac{\partial \mathcal{H}'}{\partial n} \right) \, ds'.\tag{2}\]

Here \(\mathbf{r} \in \mathbb{R} \setminus \Gamma\), \(\mathcal{H}' \equiv \mathcal{H}(\mathbf{r}')\), \(\partial / \partial n'\) is the derivative along the internal normal, \(g(\mathbf{r}, \mathbf{r}')\) is a fundamental solution to the inhomogeneous Helmholtz equation

\[(\Delta + k^2)g = \delta(\mathbf{x} - \mathbf{x}')\delta(y - y').\tag{3}\]

Let us assume that \(\mathcal{H}\) is the magnetic field. It has to satisfy the boundary conditions for the field and its normal derivative

\[[\mathcal{H}]_{\Gamma} = \left[ \frac{1}{\varepsilon} \frac{\partial \mathcal{H}}{\partial n} \right]_{\Gamma} = 0,\tag{4}\]

where the square brackets denote the jump, \(\varepsilon\) corresponds to either \(\varepsilon_{\text{in}}\) or \(\varepsilon_{\text{out}}\). Conditions (4) mean that the magnetic field is always continuous at \(\Gamma\), whereas its normal derivative has a jump depending on \(\varepsilon_{\text{in}}, \varepsilon_{\text{out}}\).

To find the field with the help of the Green theorem (2) we need to know \(\mathcal{H}\) and \(\partial \mathcal{H} / \partial n\) along boundary \(\Gamma\). Let there be \(N\) boundaries \(\Gamma_i\) and

\[\Gamma = \bigcup_{i=1}^{N} \Gamma_i.\]

There are \(2N\) independent equations at \(\mathbf{r} \in \Gamma_i\):

\[\frac{1}{2} \mathcal{H}(\mathbf{r}) = -\int_{\Gamma_i} \left( \frac{\partial g_{\text{in}}}{\partial n'} \mathcal{H}' - g_{\text{in}} \frac{\varepsilon_{\text{in}} \partial \mathcal{H}'}{\varepsilon_{\text{in}} \partial n'} \right) \, ds',\tag{5}\]

\[\frac{1}{2} \mathcal{H}(\mathbf{r}) = \sum_{j=1}^{N} \int_{\Gamma_j} \left( \frac{\partial g_{\text{out}}}{\partial n'} \mathcal{H}' - g_{\text{out}} \frac{\varepsilon_{\text{out}} \partial \mathcal{H}'}{\varepsilon_{\text{out}} \partial n'} \right) \, ds' + \mathcal{H}_0(\mathbf{r}),\tag{6}\]

which are obtained by approaching \(\Gamma_i\) from either inner or outer domains. Here the fundamental solution inside and outside \(D\) is denoted as \(g_{\text{in}}\) and \(g_{\text{out}}\), respectively. One may add any solution of the homogeneous Helmholtz equation (1) to the fundamental solutions. We do not require the radiation condition for internal fundamental solutions, whereas function \(g_{\text{out}}\) should satisfy the Sommerfeld radiation condition, i.e. be the diverging spherical or cylindrical wave at infinity. Hence, the implicit integral over an arbitrary contour enveloping all the scatterers in eq. (6) reduces to the field in the absence of scatterer (or scatterers) \(D\), denoted as \(\mathcal{H}_0\).

After solving eqs. (5), (6), the field in the arbitrary point can be calculated using the Green theorem. This approach is the basis for the boundary element method (BEM) [13,14]. Its advantage consists in diminishing the problem dimension. For instance, in two-dimensional geometry the method deals with one-dimensional contour \(\Gamma\), and then appears to be very fast and accurate.

Hereafter we consider the two-dimensional problem in the cross-section perpendicular to the cylinder axis, fig. 1. The plane running wave \(\mathcal{H}(\mathbf{r}, t) = \mathcal{H}_{\text{inc}} \exp(-i\omega t + ik_1 \cdot \mathbf{r})\) goes from the dielectric medium \(\varepsilon_1\) to the medium with permittivity \(\varepsilon_2\); \(\theta_1\) is the incident angle between the wave vector \(k_1\) and the normal to the boundary. While it is greater than \(\theta_0 = \arcsin \sqrt{\varepsilon_2 / \varepsilon_1}\), the angle of the total internal reflection, only the evanescent wave with coordinate dependence \(\exp(-k y + ik_2 x)\) penetrates into the medium 2, where

\[\kappa = \frac{\omega}{c} \sqrt{\varepsilon_1 \sin^2 \theta_1 - \varepsilon_2}, \quad k_{2x} = k_{1x} = \frac{\omega}{c} \sqrt{\varepsilon_1 \sin \theta_1},\tag{7}\]

\(\omega, c\) are the frequency and the speed of light.

In an isotropic medium two states of polarization are possible, depending on the incident wave polarization. We consider the TM-wave with magnetic-field vector perpendicular to the plane of incidence. This case is more interesting in view of the plasmon resonances study, since the electric-field vector lies in the \(xy\)-plane where the cylinder has a finite size. As the magnetic field of the wave obeys the Helmholtz equation (1), then the BEM is applicable. The solution for the TE-wave can be considered in the same way.
We look for the specific Green function \( G(x, y; x', y') \) satisfying the inhomogeneous equation (3) in media 1 and 2. The function \( G \) depends on the difference \( x - x' \) only, due to translational symmetry. The boundary condition at \( y = 0 \) is

\[
[G(x, y; x', y')]_{y=0} = \left[ \frac{1}{\varepsilon} \frac{\partial G(x, y; x', y')}{\partial y} \right]_{y=0} = 0. \tag{8}
\]

After the Fourier transformation

\[
G(x, y; 0, y') = \frac{1}{2\pi} \int_{-\infty}^{\infty} G_q(y, y') e^{i\varphi} dq \tag{9}
\]
eq (3) is reduced to an ordinary differential equation having exponential solutions. Using conditions (8) we obtain the function at \( y' > 0 \) in \( q \)-domain

\[
G_q = -\frac{1}{2\mu_2} \begin{cases} 
(1+r(q))e^{\mu_2 y-y'+q}, & y < 0, \\
(1-q) e^{-\mu_2 |y-y'|+q}, & y > 0.
\end{cases} \tag{10}
\]

Here \( r(q) = (\varepsilon_1\mu_2 - \varepsilon_2\mu_1)/(\varepsilon_1\mu_2 + \varepsilon_2\mu_1) \) is the Fresnel reflection coefficient of the \( p \)-wave at normal incidence [15], \( \mu_{1,2}^2 = q^2 - k_{1,2}^2 \). Carrying out the Fourier transformation (9) of function (10) at \( y > 0 \) we have two terms:

\[
G_1 = -\int_{-\infty}^{-i0} e^{-\mu_2 |y-y'|+iq(x-x')} dq \frac{4\pi\mu_2}{4}, \tag{11}
\]

\[
G_2 = -\int_{-i0}^{\infty} e^{-\mu_2 |y-y'|+iq(x-x')} r(q) dq \frac{4\pi\mu_2}{4}, \tag{12}
\]

and \( G = G_1 + G_2 \), where the sign of the square root is given by the rule \( \sqrt{q^2 - k_2^2} \rightarrow -i\sqrt{k_1^2 - q^2}, q^2 < k_3^2 \).

The first term (11) can be calculated analytically and reduces to the Green function in the homogeneous space,

\[
G_1(r, r') = \frac{1}{4\pi} H_0^{(1)}(k_2|\rho_-), \tag{13}
\]

where \( H_0^{(1)} \) denotes the Hankel function of the first kind [16], \( \rho_2^2 = (x-x')^2 + (y \pm y')^2 \). The second term \( G_2 \) gives the effect of the reflected image source. The amplitude of the source at each \( q \) is equal to the reflectivity coefficient \( r(q) \). Thus, along the point source at \( (x', y') \) we have to consider the mirror-image source \( r(q) \) at \( (x', -y') \).

The total field at the upper half-plane is the sum of the fields generated by the source and its image at each \( q \). This approach naturally takes into account the multiple scattering. The Green function of this type was studied for homogeneous waves: spherical acoustic, see [17,18], or cylindrical electromagnetic waves [19]. A general approach has been built, based on integral equations, to study the field scattered by an object near a two-media interface, also for excitation by homogeneous light wave [20].

The asymptotic behavior of the Green function in the far field can be found by the steepest-descent method [21]. The stationary point is \( q_0 = k_2 |x-x'|/\rho \), where \( \rho = \rho_- \) for \( G_1 \) and \( \rho = \rho_+ \) for \( G_2 \). The result is the sum of cylindrical waves \( G \propto \rho^{-1/2} e^{ik_2 \rho} + r_0 \rho_+^{-1/2} e^{ik_2 \rho_-} \), where

\[
r_0 = \frac{\varepsilon_1 \sin \varphi - \sqrt{\varepsilon_2(\varepsilon_1 - \varepsilon_2 \cos^2 \varphi)}}{\varepsilon_1 \sin \varphi + \sqrt{\varepsilon_2(\varepsilon_1 - \varepsilon_2 \cos^2 \varphi)}} \tag{14}
\]
is the reflection coefficient at \( q = q_0 \). \( \varphi \) is the polar angle of observation counted out from the \( x \)-direction. The reflection coefficient is \( r_0 = -1 \) at \( \varphi = 0 \) and \( (\varepsilon_2 - \sqrt{\varepsilon_1^2 - \varepsilon_2^2})/(\varepsilon_1 + \sqrt{\varepsilon_1^2 - \varepsilon_2^2}) \) at \( \varphi = \pi/2 \). It turns to zero at \( \varphi = \arctan \sqrt{\varepsilon_1/\varepsilon_2} \), i.e., at the Brewster angle.

The Green functions (11), (12) are exploited as the external fundamental solution in eqs. (5), (6), i.e., we set \( g_{out} = G \). The BEM algorithm has been tested in the case of one contour and homogeneous wave when \( \varepsilon_1 = \varepsilon_2 \). There are analytical formulas for a circle [22] and numerical calculations for an ellipse [23]. The comparison demonstrates the relative consistency within \( 10^{-4} \) for \( N = 360 \) panels approximating contour \( \Gamma \).

During calculations we get two sorts of integrals:

\[
I_1 = \int_{\Gamma} g(r; r') \frac{\partial H_{0}^{'(1)}}{\partial n'} ds' \tag{15}
\]

\[
I_2 = \int_{\Gamma} \frac{\partial g(r; r')}{\partial n'} H_{0}^{(1)} ds'. \tag{16}
\]

Here \( r, r' \) correspond to the points of observation and integration, respectively. The integration goes along curvilinear contours. In the stage of discretization the singular integrals appear, if these points belong to the same discrete segment of the curve \( \Gamma \), i.e., the diagonal elements in the matrix of system have to be estimated. For simplicity we present the corresponding estimation only for circle contour. For the elliptic boundary the expressions become much more lengthy.

Note that the expressions for the elements of the main matrix are slightly different for internal and external regions. As mentioned above, within the external domain the Green function \( g_{out} \) has to satisfy the Maxwell boundary condition and the Sommerfeld radiation condition and consists of two parts: the Hankel function \( G_1 \) (13) that gives the cylindrical wave scattered from the point scatterer, and the integral part \( G_2 \) describing the image reflected with respect to the plane boundary. The latter part gives a non-singular contribution into the integrals (15), (16). It is of the same order, as non-singular elements, except for the case where the considered segment of the contour touches the boundary between the half-spaces. This type of singularity was studied analytically [24].

For the calculation of the singular integrals we use the known asymptotic expression of the Hankel function:

\[
G_1 \approx \frac{1}{4\varepsilon} + \frac{C + \ln(k_2 r/2)}{2\pi}, \tag{17}
\]

where \( r = |r - r'| \), and \( C \approx 0.5772 \) is the Euler constant. Suppose that both vectors \( r \) and \( r' \) point to the same arc.
of the circle with radius $R$ and corner angle $\Delta \phi$ with $r$ placed in the center of the arc. The first singular integral is written in the form

$$\frac{\partial \mathcal{G}}{\partial n} \int_{-\Delta \phi/2}^{\Delta \phi/2} \left( \frac{1}{4i} + \frac{C + \ln(k_2 r/2)}{2 \pi R} \right) R d\phi = \frac{\partial^2 \mathcal{G}}{\partial n \partial R} R \Delta \phi \left( \frac{1}{4i} + \frac{C - 1 + \ln(k_2 R \Delta \phi/4)}{2 \pi} \right).$$

(18)

The estimation has the second order of approximation.

For the second singular integral first write the normal derivative of the Green function:

$$\frac{\partial G(r; r')}{\partial n'} \approx \frac{\partial \ln(k_2 r/2)}{2 \pi \partial n'} = -\frac{1}{2 \pi r^2} \frac{\partial r}{\partial r'} = -\frac{1}{4 \pi R}.$$  

(19)

The singularity vanishes, and the second singular integral is equal to

$$\int_{\Gamma} \frac{\partial G(r; r')}{\partial n'} 3l' ds' \approx -3l \Delta \phi/4.$$  

(20)

For off-diagonal elements we do not differentiate the asymptotic expression (17), but use the Hankel function $H_1^{(1)}$ that appears when we calculate the derivative analytically.

Without the scattering body the electric field vector $\mathbf{E}$ in medium 2 has only the $y$-component. The scatterer produces a small component $\mathbf{E}_y$ and the evanescent wave is partially converted into a diverging one. The corresponding Pointing vector

$$\mathbf{S} = \frac{c}{8 \pi} \text{Re} (\mathbf{E} \times \mathbf{H})$$

(21)

acquires a nonzero $y$-component, $S_y$, so the energy flux leaving the plane arises.

First, we consider the scattering on two identical cylinders and calculate the Pointing vector at a distance nearly $2\lambda$, i.e. in the wave zone, and normalized by the average flux of the incoming wave in the first medium $S_{inc} = c \mathcal{G}_{inc} / 8 \pi \sqrt{\nu}$. The indicatrix of the scattering into the upper half-space is shown in fig. 2. The insets in this and the next figures depict the arrangement of scatterers. The scattering is minimal in the normal or longitudinal direction and maximal at some medium angles. It is the quadrupole contribution due to the field of image source (14) and the result of the decay of the evanescent wave amplitude with $y$. The angle $\varphi$ of the first maximum increases with $\varepsilon_1$. For $\varepsilon_1 = 2$ the first maximum is $23^\circ$ and for $\varepsilon_1 = 3$ it is $41^\circ$. The forward-backward asymmetry of the indicatrix is a consequence of the violation of the dipole approximation owing to finite sizes ($2kR \approx 0.8$).

The method can be applied to different shapes of contours. We choose the first contour ($\varepsilon_{in} = \varepsilon_3$) to be a circle whereas the second ($\varepsilon_{in} = \varepsilon_4$) is a prolate ellipse with minor semiaxis $a = 0.04 \mu m$ and the axis ratio $b/a = 10$.

The major axis is directed along $y$. Figure 3 shows the field near the bottom tip of the ellipse as a function of coordinate $x$ at $y = 0.15 \mu m$, when the ellipse is moving along the $x$-axis. The pattern in the near-field region is much more complicated than in the wave zone. We see the slight low-frequency oscillations on the right side and the deep high-frequency ones on the left. The physical nature of the oscillations is the interference between the incident evanescent wave and the diverging wave scattered by the circle. They are counterpropagating on the left and copropagating on the right. Their spatial frequencies are different: $k_\perp = k_z \pm k_2$. However, in the considered case $k_\perp = k_2 \sin \theta_1 \approx k_2$, then the interference oscillation at the right side has a low frequency. The slight interference pattern on the right in this case is caused by the scattering by the tip and these oscillations vanish without the tip, as the lower curve demonstrates.

Figure 4 shows the energy going along the major axis through the ellipse middle cross-section (in the wave zone) as a function of coordinate $x$, namely, the component $S_y$ of vector (21) averaged over the horizontal cross-section. As follows from eq. (10), the higher spatial harmonics with $q^2 > k_2^2$ decay exponentially with distance as $\exp(-\mu_2 y)$. 

Fig. 2: (Colour on-line) Energy flux through a distant semicircle (arbitrary units) from two equal circles with distance between centers $L = 0.21 \mu m$ as a function of polar angle $0 < \varphi < \pi$ at $\lambda = 1.51 \mu m$, $R = 0.1 \mu m$, $\delta = 0.01 \mu m$, $\theta_1 = \pi/4$, $\varepsilon_2 = 1$, $\varepsilon_3 = \varepsilon_4 = -91.5 + 10.3i$, and different $\varepsilon_1 = 2$ (solid line), 2.25 (dotted line), 3 (dashed line).

Fig. 3: (Colour on-line) The magnetic field at the bottom tip of the ellipse as a function of coordinate $x$ at distance $y = 0.15$ from the plane (solid line). Parameters are $\lambda = 1.512 \mu m$, $R = 0.06 \mu m$, $\delta = 0.01 \mu m$, $\varepsilon_1 = \varepsilon_3 = \varepsilon_4 = 2.25$, $\varepsilon_2 = 1$, $\theta_1 = \pi/4$. The same without the ellipse ($\varepsilon_4 = 1$, dashed line).
Scattering of evanescent wave

Fig. 4: (Colour on-line) Energy going through the central cross-section of the ellipse in the $y$-direction as a function of its horizontal position $x$ at distance from the plane $y = 0.15$ (solid line), $0.35$ (dashed line), $0.50 \mu m$ (dot-dashed line). Parameters are $\lambda = 1.512$, $R = 0.06$, $\delta = 0.01 \mu m$, $\varepsilon_1 = \varepsilon_3 = \varepsilon_4 = 2.25$, $\varepsilon_2 = 1$, $\theta_1 = \pi/4$.

Therefore, the small details are not visible in the far-field pattern. However, at $k y \ll 1$ the exponent is not negligible, and then the fine details become apparent. The near field can be observed if one extracts the signal and transfers it to the far zone. In our calculation the stretched ellipse plays the role of such a transmitter.

The height of the tip above the interface is $d = 0.15–0.50 \mu m$, then the minimal distance between the tip of the ellipse and the circle starts from $0.02 \mu m$. We see interference oscillations in the coordinate dependence, their amplitude decreases with the tip distance from the boundary. The oscillations are caused by the interference between the falling evanescent wave and the cylindrical wave scattered by the circle, by analogy to those in fig. 3. The oscillations at $x > 0$ have the lower frequency $k_1 \sin \theta_1 - k_2 \cos \alpha$, where $\alpha$ is the angle of incidence for the cylindrical wave.

Moreover, the oscillations on the right part decreases with distance $x$ between the circle and the tip. The effect of frequency decreasing with distance becomes especially apparent if we plot the middle curve from fig. 4 throughout a wider interval of coordinate. The corresponding curve is shown in fig. 5(a), the interval is wider by an order. The reason is the decreasing angle $\alpha$ with distance $x$, as shown in the inset in fig. 5(a).

Then the oscillations at $x > 0$ are a spurious effect caused by the lateral penetration of the electromagnetic field into the ellipse, that is a source of artifacts in NSOM. At the same time the oscillations at $x < 0$ are transferred to the midplane from the near field. The wanted signal is the dip in the center $x = 0$ and several nearest peaks.

Figure 5(b) illustrates the interference between evanescent and cylindrical waves. It displays the function

$$I(x) = \left| e^{ik_1 x \sin \theta_1} - e^{ik_2 \sqrt{x^2 + y^2}} \right|^2$$

(22)

with fixed $y$. We see a similar pattern as in fig. 5(a). The oscillations at $x > 0$ have also decreasing frequency. At the same time the dip in the center and several nearest peaks by its sides are absent.

The energy flux through the far central plane is the simplest (two-dimensional in our case) model of NSOM [2,3,5]. The curves in fig. 4 correspond to the instrumental function of the microscope. Although, this statement should not be taken literally. The multiple scattering leads to back influence of the probe to the object then it is not a usual linear function of response [25,26]. The slight oscillations in the upper curve in fig. 3 at $x > 0$ are an example of the back influence.

Thus, the coupled boundary equations describing the scattering of the evanescent wave are solved for two cylinders. Asymmetries of the indicatrix and oscillations in the coordinate dependence are observed. The forward-backward scattering asymmetry has quite the same origin as in the case of homogeneous waves (see, for instance, [27]). In fact, we account for higher multipoles here. The oscillating dependence is a result of the interference between evanescent and cylindrical waves. The BEM with the proposed Green function is rather general and applicable for any contour $\Gamma$ or several contours. It can be extended also to 3D geometry. The Green
function could find applications in other calculations, e.g. the volume integral equations including the Born series and the discrete dipole approximation.

***

Authors are grateful to E. V. Podivilov for helpful discussions. This work is supported by the Government program NSh-4339.2010.2, program No. 21 of the Russian Academy of Sciences Presidium, and interdisciplinary grant No. 42 from the Siberian Branch of RAS.

REFERENCES

[1] Stockman M. I., Phys. Today, 64, issue No. 2 (2011) 39.
[2] Greffet J.-J. and Carminati R., Prog. Surf. Sci., 56 (1997) 133.
[3] Hecht B., Sick B., Wild U. P., Deckert V., Zenobi R., Martin O. J. F. and Pohl D. W., J. Chem. Phys., 112 (2000) 7761.
[4] Novotny L. and van Hulst N., Nat. Photon., 5 (2011) 83.
[5] Novotny L. and Hecht B., Principles of Nano-Optics (Cambridge University Press, Cambridge, New York) 2006.
[6] Girard C., Rep. Prog. Phys., 68 (2005) 1883.
[7] Kawata S. and Shalaev V. M. (Editors), Nanophotonics with Surface Plasmons (Elsevier, Oxford) 2007.
[8] Belai O. V., Frumin L. L., Perminov S. V. and Shapiro D. A., Opt. Lett., 36 (2011) 954.
[9] Cragg G. E. and So P. T. C., Opt. Lett., 25 (2000) 46.
[10] Carney P. S. and Schotland J. C., Opt. Lett., 26 (2001) 1072.
[11] Fischer D. G., Opt. Lett., 25 (2000) 1529.
[12] Belkebir K., Chaumet P. C. and Sentenac A., J. Opt. Soc. Am. A, 22 (2005) 1889.
[13] Colton D. L. and R. K., Integral Equation Methods in Scattering Theory (Wiley, New York) 1983.
[14] Brebbia C. A., Telles J. C. F. and Wrobel L. C., Boundary Element Techniques (Springer, Berlin) 1984.
[15] Raftian S. G., Introduction to Physical Optics (URSS, Moscow) 2009 (in Russian).
[16] Ovler F. W. J., Lozier D. W., Boisvert R. F. and Clark C. W., NIST Handbook of Mathematical Functions (Cambridge University Press, Cambridge, New York) 2010.
[17] Brekhovskikh L. M., Waves in Layered Media (Academic Press, New York) 1980.
[18] Landau L. D. and Lifshitz E. M., Fluid Mechanics, Course of Theoretical Physics, Vol. 6 (Pergamon, New York) 1959.
[19] Pincemin F., Sentenac A. and Greffet J.-J., J. Opt. Soc. Am. A, 11 (1994) 1117.
[20] Greffet J.-J., Opt. Commun., 72 (1989) 274.
[21] Bleistein N. and Handelsman R. H., Asymptotic Expansion of Integrals (Dover, New York) 1986.
[22] Harrington R. F., Time-Harmonic Electromagnetic Fields (Wiley, New York) 2001.
[23] Zymovetz S. V. and Geshev P. I., Tech. Phys., 51 (2006) 291.
[24] Saillard M. and Toso G., Radio Sci., 32 (1997) 1347.
[25] Girard C. and Bouju X., J. Opt. Soc. Am. B, 9 (1992) 298.
[26] Sun J., Carney P. S. and Schotland J. C., J. Appl. Phys., 102 (2007) 103103.
[27] Landau L. D. and Lifshitz E. M., Electrodynamics of Continuous Media, Course of Theoretical Physics, Vol. 8 (Pergamon, New York) 1960.