Anomalous light absorption by small particles

M. I. Tribelsky (a)

A.N. Nesmeyanov Institute of Organoelement Compounds, Russian Academy of Sciences
28 Vavilova St., Moscow 119991, Russia,
Moscow State Institute of Radioengineering, Electronics and Automation (Technical University)
78 Vernadskiy Ave., Moscow 119454, Russia and
Max-Planck-Institut für Physik komplexer Systeme - Nöthnitzer Str. 38, Dresden 01187, Germany, EU

received 14 December 2010; accepted in final form 8 March 2011
published online 4 April 2011

PACS 42.25.Bs – Wave propagation, transmission and absorption
PACS 42.25.Fx – Diffraction and scattering
PACS 78.67.Bf – Nanocrystals, nanoparticles, and nanoclusters

Abstract – Light absorption by a spatially uniform spherical nanoparticle in the vicinity of surface plasmon (polariton) resonances is studied in detail based on the exact Mie solution. It is shown that the maximal absorption is achieved for a particle from weakly dissipating materials and may have very unusual properties. A simple universal formula describing the resonant absorption lineshape as a function of the particle size and its complex dielectric permittivity is obtained. Possible comparison with experiment is discussed.

Since the first quantitative study by Lord Rayleigh [1], the problem of light scattering by small particles has remained one of the most important and appealing issues of electrodynamics. There are thousands of articles and numerous monographs devoted to this subject, see, e.g., [2–5] and references therein. Being important subfields of the general problem, plasmon (polariton) resonances and their role in the light scattering as well as a related issue of interplay between radiative and dissipative damping are not new too [6], but they still remain topics of intense study [7–9]. Meanwhile certain important properties of the fundamental problem of light scattering by a small spatially uniform symmetric particle remain either unrevealed, or underappreciated.

Recently it has been pointed out [10,11] that peculiarities of the light scattering by small spatially uniform particles made of weakly dissipating materials actually provide grounds to talk about a new type of the scattering, the so-called anomalous light scattering (ALS), which occurs when the dissipative losses are small enough to become negligible relative to the radiative damping. ALS has very little in common with the Rayleigh scattering, other than both occur without changes in the frequency of an incident light $\omega$. It is characterized by giant optical resonances with narrow linewidths and inverted hierarchy, i.e., the partial extinction cross-section for the quadrupole resonance is larger than that for the dipole one, etc.\(^1\) [12,13]; very unusual and complicated field structure in the vicinity of the particle (nearfield); singular dependence of the scattering amplitudes on the particle size, etc.

Maximization of light absorption by medium (perfect absorption) also attracts attention of researchers, see, e.g., ref. [14]. In addition to purely academic interest, the problem is crucial for some applications, for example, to stimulate chemical reactions, or to enhance laser sintering, etc. Recent achievements in therapy, when laser heating of nanoparticles absorbed by cancer and/or pathogenic bacteria cells kills these cells selectively [15,16] deserve to be mentioned separately.

It motivates extensive study of light absorption by nanoparticles. Among new publications a very interesting paper [17] should be mentioned. In this paper the absorption and scattering cross-sections of a cylindrical subwavelength structure are studied based upon the so-called temporal coupled-mode model [18,19].

In the present paper the simplest case of light absorption by a small spherical spatially uniform nanoparticle is analyzed based upon the exact Mie (also known as Lorenz-Mie [5]) solution of the Maxwell equations. The absorption cross-section $\sigma_{abs}(R, \epsilon'_p, \epsilon''_p)$ of a particle with

\(^{1}\)It does not affect the convergence of the multipole expansion because i) for a spherical particle each resonance occurs at its own frequency; ii) the inverted hierarchy lasts up to a certain finite order of the resonances, then the normal hierarchy is restored [12,13].
radius $R$, magnetic permeability $\mu = 1$ and complex dielectric permittivity $\epsilon_p = \epsilon_p^\prime + i\epsilon_p^\prime\prime$ is inspected. It is shown that there is a range of parameters, where the behavior of the Mie solution is quite unusual and despite the particle smallness, the corresponding $\sigma_{abs}$ differs from the one for the Rayleigh scattering dramatically. The phenomenon occurs in the vicinity of the plasmon (polariton) resonances, so that the absorption is also resonant and its magnitude increases greatly. The phenomenon is observed for small values of $\epsilon_p^\prime\prime$. The values of the maxima are functions of the resonance order and the wave number of the incident light solely and do not depend on the particle size and its optical properties. The maxima should be regarded as the upper theoretical limits for the partial absorption cross-section, which cannot be exceeded for a spatially uniform particle, cf. ref. [17]. A simple universal expression describing the lineshape of the resonant absorption cross-section as a function of $R, \epsilon_p^\prime\prime$ is obtained. The resonant absorption in the vicinity of the plasmon (polariton) resonances inherits many features of ALS, which allows naming it the anomalous light absorption (ALA). Detailed calculations for an aluminum particle indicate that the anomalous absorption may be observed experimentally, and the absorption cross-section may be quite close to the upper theoretical limit. Generalizations of the results obtained to the case of a shell (layered) structure is possible, as well as for magnetic particles with $\mu \neq 1$ are just a matter of more cumbersome algebra.

To understand the physical nature of the Mie solution, let us estimate the power $P$ dissipated in a unit volume of the particle. Neglecting the magnetic part of the dissipative loss owing to the particle smallness, one obtains $P \propto \epsilon^\prime\prime |\mathbf{E}|^2$, where $\mathbf{E}(r)$ is a complex amplitude of the monochromatic electric field in a given point inside the particle [20]. On the other hand, in the vicinity of each plasmon (polariton) resonance $\mathbf{E}(r)$ undergoes a sharp increase. At these resonances in the case of the Rayleigh scattering, $\mathbf{E}(r)$ is limited by dissipative losses, so that $\mathbf{E}(r) \propto 1/\epsilon_p^\prime\prime$. Thus at the resonances, $P \propto 1/\epsilon_p^\prime\prime \to \infty$ at $\epsilon_p^\prime\prime \to 0$, which gives rise to the corresponding divergence of $\sigma_{abs}$.

The divergence occurs because the approximation used is invalid. In fact, $\sigma_{abs}$ does not diverge. At small enough $\epsilon_p^\prime\prime$, the radiative damping prevails over the dissipative losses and the Rayleigh scattering is replaced by ALS, which provides a different cutoff for $\mathbf{E}(r)$ [10,11,13].

To inspect the problem accurately, one should employ the exact Mie solution. It is convenient to introduce dimensionless cross-sections $Q_{ext, sca, abs} = q^{1/2}Q_{ext, sca, abs}/\pi R^2$, where $\sigma_{ext}$, $\sigma_{abs}$ and $\sigma_{sca}$, respectively, stand for the extinction, scattering and absorption cross-sections of the particle. Then, the Mie solution yields the following well-known expressions:

$$Q_{ext} = \sum_{l=1}^{\infty} Q^{(l)}_{ext}, Q_{sca} = \sum_{l=1}^{\infty} Q^{(l)}_{sca}, Q_{abs} = Q_{ext} - Q_{sca},$$

$$Q^{(l)}_{ext} = \frac{2}{q^2} (2l + 1) \operatorname{Re}(a_l + b_l),$$

$$Q^{(l)}_{sca} = \frac{2}{q^2} (2l + 1) \left| a_l \right|^2 + \left| b_l \right|^2. \quad (3)$$

The summation over the corresponding partial multipole cross-sections, the so-called size parameter $q$ equals $n_m R k_0$, $n_m = \sqrt{\epsilon_m}$ is a purely real refractive index of the environmental medium, and $k_0$ stands for the wave number of the incident light in vacuum ($k_0 = \omega/c$, where $c$ is the speed of light).

Amplitudes $a_l$ and $b_l$ may be presented as follows [10,12]:

$$a_l = \frac{F_l^{(a)}(q, \epsilon)}{F_l^{(a)}(q, \epsilon) + iG_l^{(a)}(q, \epsilon)}. \quad (4)$$

Here $\epsilon = \epsilon_p / \epsilon_m(\omega)$ and $F_l^{(a)} \to F_l^{(b)}$; $G_l^{(a)} \to G_l^{(b)}$ for $b_l$. $F_l^{(a,b)}$, $G_l^{(a,b)}$ are expressed in terms of the Bessel $J_{l+1/2}(z)$ and Neumann $N_{l+1/2}(z)$ functions, whose expansions in power series give rise to the following formulae, valid for small $q$:

$$F_l^{(a)}(q, \epsilon) \approx q^{2l+1} \frac{l + 1}{(2l + 1)!!} (\epsilon - 1) + \ldots, \quad (5)$$

$$G_l^{(a)}(q, \epsilon) \approx \frac{l}{2l + 1} \left\{ \epsilon + \frac{l + 1}{l} \right. - q^2 \frac{\epsilon - 1}{2} \left[ \frac{\epsilon}{2l + 3} + \frac{l + 1}{l(2l - 1)} \right] \right. + \ldots \right\}, \quad (6)$$

(ellipses denote omitted higher order terms in $q$). The explicit expressions for $F_l^{(b)}$ and $G_l^{(b)}$ are not required because of the estimate $|b_l| \sim q^{2l+3} \ll |a_l|$, which allows to neglect $b_l$ compared to $a_l$ at $q \ll 1$.

At small dissipations the surface plasmon (polariton) resonances are defined by the condition $G_l^{(a)}(q, \epsilon) = 0$. Here and in what follows prime and double prime denote the corresponding real and imaginary parts, respectively.

The condition $G_l^{(a)}(q, \epsilon) = 0$ determines the resonant values of the real part of the dielectric permittivity $\epsilon_l(q)$ and through the dependence $\epsilon_m(\omega)$, the resonant frequencies $\omega_l(q)$. For the problem in question the maxima of $\sigma_{abs}$ should be situated in the vicinity of the resonances. Let us inspect the vicinity of a certain $l$-th resonance (from now on $\ell$ designates a given order of the resonance, while $l$ remains dummy and may have any value), presenting the dielectric permittivity in the form

$$\epsilon = -\frac{l + 1}{\ell} + \delta \epsilon' + i\epsilon'',$$

where $\delta \epsilon'$ and $\delta \epsilon''$ are small, see eq. (6). Then, the leading approximations for $F_l^{(a)}$, $G_l^{(a)}$ and $Q_{ext}^{(a)} = Q_{ext} - Q_{sca}$ read

$$F_l^{(a)} \approx q^{2l+1} \left[ \frac{l + 1}{(2l + 1)!!} \left( -\frac{2l + 1}{\ell} + i\epsilon'' \right) \right], \quad (8)$$

$$G_l^{(a)} \approx 2q^2 \frac{\ell}{(2l + 1)!!} \left[ \left( \delta \epsilon' + q^2 \frac{2(2l + 1)(\ell + 1)}{\ell^2(2l - 1)(2l + 3)} \right) + i\epsilon'' \right], \quad (9)$$

$$Q_{sca}^{(a)} \approx \frac{2}{q^2} (2l + 1) \left( \left| a_l \right|^2 + \left| b_l \right|^2 \right). \quad (3)$$
According to eq. (8), \( F_{\ell}^{(a)''} \) is negative and does not depend on \( \delta \epsilon', \epsilon'' \). It allows to consider \( G_{\ell}^{(a)''}, G_{\ell}^{(a)'''} \) as new independent variables instead of \( \delta \epsilon', \epsilon'' \), see eq. (9). The maximum of \( Q_{\text{abs}}^{(f)} \) corresponds to \( G_{\ell}^{(a)''} = 0 \) (i.e., it is achieved exactly at the resonant frequency \( \omega_r \)) and \( G_{\ell}^{(a)'''}, -F_{\ell}^{(a)''} \) (it should be remembered that \( G_{\ell}^{(a)'''}, -F_{\ell}^{(a)''} \) cannot be negative, see eq. (9)). Note, the two terms \( G_{\ell}^{(a)'''}, -F_{\ell}^{(a)''} \) in the denominator of eq. (10) correspond to the dissipative and radiative damping, respectively. So, \( Q_{\text{abs}}^{(f)} \) is maximized when the former equals the latter, cf. [17].

Thus on the \((\epsilon', \epsilon'')\)-plane, in the vicinity of each surface plasmon (polariton) resonance, \( Q_{\text{abs}}^{(f)} \) reaches the local maximum

\[
Q_{\text{abs max}}^{(f)} = \frac{2(2\ell + 1)}{q^2} \frac{-F_{\ell}^{(a)'}, G_{\ell}^{(a)''}}{(G_{\ell}^{(a)''} - F_{\ell}^{(a)'})^2 + (G_{\ell}^{(a)''})^2},
\]

which equals 1/4 of the maximal ALS resonant partial extinction cross-section, \( Q_{\text{ext max}}^{(f),\text{ALS}} \) [12,13], and also linearly increases with an increase in the order of resonance \( \ell \) (inverted hierarchy). However, in contrast to \( Q_{\text{ext max}}^{(f),\text{ALS}} \) (achieved at \( \epsilon'' = 0 \)), the maximum of \( Q_{\text{abs}}^{(f)} \) is realized at

\[
\epsilon' = \epsilon'_t \simeq -\frac{\ell + 1}{\ell} + q^2 \frac{2(2\ell + 1)(\ell + 1)}{(2\ell - 1)(2\ell + 3)},
\]

\[
\epsilon'' = \epsilon''_t \simeq q^{2\ell + 1} \frac{\ell + 1}{[\ell(2\ell - 1)]^2} \ll 1.
\]

Note, that eq. (11) provides the theoretical upper limit, for the partial absorption cross-section for a small, spatially uniform, non-magnetic particle, which cannot be exceeded.

Conditions eqs. (12), (13) correspond to \( \sigma_t \simeq 1/2 \). According to the usual expressions [2,3] it yields the following singular dependence of the amplitudes of the scattered electric \( E^{(s)} \) and magnetic \( H^{(s)} \) fields on \( q \) in the nearfield, identical to those for ALS [10]:

\[
E^{(s),p,r} \propto q^{-(\ell+2)}, H^{(s),p,r} \propto q^{-(\ell+1)},
\]

dependence \( H_{r}^{(s)}(q) \) remains the same as that for the Rayleigh scattering: \( H_{r}^{(s)} \propto q^{(\ell+1)} \), where \( H_{r}^{(s)} \) stands for the radial component. We remind that for the Rayleigh scattering one would have \( E_{p}^{(s),p,r} \propto q^{(\ell+1)}, H_{p}^{(s),p,r} \propto q^{\ell} \).

The linewidths of the resonant absorption measured at the level \( Q_{\text{abs}}^{(f)} = Q_{\text{abs max}}^{(f),2} \) are as follows:

\[
\Delta \epsilon' = q^2 \frac{4(\ell + 1)}{(2\ell - 1)[\ell^2]}, \quad \epsilon'' = \epsilon''_t,
\]

\[
\Delta \epsilon'' = q^2 \frac{4(\ell + 1)}{(2\ell - 1)[\ell^2]}, \quad \epsilon' = \epsilon'_t.
\]

Though the values of \( \Delta \epsilon', \Delta \epsilon'' \) are close to each other, the lineshapes are quite different. To study this issue in detail let us introduce new variables:

\[
k = \frac{q^2Q}{2(2\ell + 1)}, \quad \xi = -\frac{G_{\ell}^{(a)'}}{F_{\ell}^{(a)'}}; \quad \zeta = -\frac{G_{\ell}^{(a)''}}{F_{\ell}^{(a)''}}.
\]

Then, (in the given approximation) the resonant partial absorption cross-section is reduced to the following universal, \( q- \) and \( \ell\)-independent form:

\[
k = \frac{\zeta}{(1 + \zeta)^2 + \zeta^2}, \quad \zeta \geq 0.
\]

Thus, the line along the \( \xi \) (i.e., \( \epsilon' \)) axis has a typical symmetric Lorentzian shape, while the one along the \( \zeta \) (i.e., \( \epsilon'' \)) axis is strongly asymmetric, see also fig. 1.

It is interesting to mention that at the resonant points of the absorption cross-sections given by eqs. (12), (13), the scattering cross-section is \( Q_{\text{abs}}^{(f)}(\epsilon'_t, \epsilon''_t) \simeq Q_{\text{abs max}}^{(f)} = Q_{\text{ext max}}^{(f),\text{ALS}} \). Then, at this point \( Q_{\text{ext}}^{(f)}(\epsilon'_t, \epsilon''_t) \simeq 2Q_{\text{abs max}}^{(f),\text{ALS}}/2 \).

The dimensional absorption cross-section which corresponds to eq. (11) \( \sigma_{\text{abs max}}^{(f)} = \pi(2\ell + 1)/(2\ell n_m k_0)^2 \) does not depend on the optical properties of the particle \( (i.e., \epsilon_p) \) and its size and hence does not vanish at \( R \to 0 \). This looks confusing. The confusion is resolved if one notes that for the function \( \sigma_{\text{abs}}^{(f)}(R, \epsilon', \epsilon'') \), the point \( R = 0, \epsilon' = \epsilon'_t, \epsilon'' = \epsilon''_t \).
$e'' = e''_\ell$ is singular; the function does not have a definite limit at this point and may take any value varying from 0 to $\sigma_{abs\, max}$ depending on the way to approach this point in $(R, e', e'')$ space. In particular, if one first fixes values of $e'' \neq 0$ and then tends $R$ to zero (which is the closest case to a possible experimental situation), $\sigma^{(\ell)}_{abs}(R, e', e'')$ vanishes, as it should be for a particle with zero radius.

Regarding off-resonant partial cross-sections, their values are given by the expression

$$Q^{(l)}_{abs} \sim q^{2\ell-1} e'^l 2(2l+1)(l+1)e'^2_{\ell}(l-l)(2l-1)|v|^2, \quad l \neq \ell. \quad (18)$$

Comparison with eq. (11) reveals that at small $q$ in the vicinity of the resonances, the contribution of the partial resonant cross-section to the net cross-section basically is determined by the plasma cross-section of an aluminium nanoparticle in vacuum is

$$\sigma_{abs} \approx \lambda, \quad \text{case to a possible experimentalsituation},$$

depending on the way to approach this point in $q(\ell, \ell', \ell'')$ space. In particular, if one first fixes values of $q'' \neq 0$ and then tends $\ell$ to zero (which is the closest case to a possible experimental situation), $\sigma^{(\ell)}_{abs}(R, \ell, \ell', \ell'')$ vanishes, as it should be for a particle with zero radius.

Let us discuss a possible experimental observation of ALA and related issues. In the preceding theoretical analysis three quantities $R$, $\ell$, and $\ell''$ have been regarded as independent. In reality it is difficult (if possible) to tune $\ell'$ and $\ell''$ independently. Actual independent control parameters are $R$ and $\omega$. Then to maintain the resonance conditions, these quantities should satisfy the following set of equations:

$$e'(\omega, R) = e'_\ell(\omega, R),$$

$$e''(\omega, R, \ell) = e''_\ell(\omega, R),$$

where the left-hand sides follow from the dispersion properties of the particle material, and the right-hand sides are given by eqs. (12), (13), respectively. Solutions of the equations (if any) yield a unique discrete set of pairs $(\omega, R)$, where $R_\ell$ should satisfy the additional constraint following from the restriction $q < 1$. These conditions are strict and the entire set of them is extremely difficult to fulfill in any real experiment. Thus, it seems the phenomenon discussed is just a “virtual” effect.

Fortunately, the situation is not so dramatic. First, the resonance lines are narrow in the $(e', e'')$-plane, but this is not necessarily the case for the linewidths along the $\omega$-axis. The latter is determined by the dispersion properties of the particle materials, so often the resonance conditions imposed on $\omega$ are not as strict as those for $e'$. Second, while the decay of the absorption cross-section at $e'' < e''_\ell$ is very sharp, it is rather slow at $e'' > e''_\ell$, see fig. 1b. Finally, one has to take into account sharpness of the $R$-dependence of the right-hand side of the resonance condition, see eq. (13). Thus, to observe the anomalous absorption it suffices to select material with small values of $e''_\ell$ for $e'$ lying in the region of possible resonances (actually, for $-2.5 \leq e' \leq -1.5$). Then, tuning $\omega$ (and/or $R$), one inevitably passes through the vicinity of the resonance(s).

To illustrate this general reasoning, the absorption cross-section of an aluminium nanoparticle in vacuum is studied in detail. In case of metals the position of the resonant region basically is determined by the plasma frequency. For different metals it may lie from UV to IR ranges of the spectrum. Though qualitatively manifestation of the resonances always remains the same, aluminium is selected owing to its high plasma frequency, which allows the most accurate comparison of the resonances with the developed approximate theory. It should be stressed that the plasma frequency may be reduced by doping the metal with impurities bounding free electrons. This red-shifts the resonant frequencies. The proper choice of the impurities and their concentration allows tuning of the resonances within rather broad limits, which may be important for possible applications.

To describe the light scattering by the aluminium particle the following model is employed [21]: the empirical dependence $\epsilon$ on $\omega$ (known as a table [22]) is approximated by the Drude formula:

$$\epsilon = 1 - \frac{\omega_p^2}{\omega^2 + \gamma^2} + \frac{\gamma\omega_p^2}{\omega(\omega^2 + \gamma^2)}. \quad (19)$$

To enhance the accuracy, quantities $\omega_p$ and $\gamma$ are regarded as functions of $\omega$, calculated at every point in the table with polynomial interpolation between the points. To take into account collisions of free electrons with the particle surface, $\gamma(\omega)$ is given by the formula [22] $\gamma(\omega) = \gamma(\omega_p + v_F/R)$, where $v_F = 10^8$ cm/s stands for the Fermi velocity of the free electrons. The obtained $\epsilon(\omega)$ is used to calculate the absorption cross-section derived from the exact Mie solution.

Results of these calculations are presented in fig. 2. To understand how the absorption cross-section changes when $q$ exceeds the small particle limit the upper bound for $R$ is set to 50 nm, which gives rise to the upper bound for $q$ varying from $\pi/2$ (for $\lambda = 200$ nm) to $\pi$ (for $\lambda = 100$ nm). Qualitative comparison of the results with the theoretically calculated analytical theory is shown in table 1. The resonant values $\lambda_R$ and $R_R$ correspond to the local maxima of $Q_{abs}$ given by eq. (11), while $Q_{abs}$, $Q_{abs max}$ are taken from the described numerical calculations at the values $\lambda_R$ and $R_R$ specified in table 1.

Note that the obtained $Q^{(\ell)}_{abs}$ for the aluminium particle are considerably larger than those given by eq. (13), and that the resonant value of $q$ for $\ell = 2$ lies at the very margin of the small particle approximation. Nevertheless fig. 2 and table 1 exhibit good qualitative and reasonable quantitative agreement with the developed theory. Well-pronounced maxima of the dipole, quadrupole and even octupole resonances are seen in fig. 1 straightforwardly. The absolute maximum of the relative absorption cross-section, $\max(Q_{abs}) \approx 11.516$, is achieved at $\lambda \approx 143$ nm; $R \approx 6.45$ nm, $q \approx 0.283$ and corresponds to the dipole resonance ($\ell = 1$), see fig. 2a. Meanwhile fig. 2b exhibits the pronounced inverted hierarchy of resonances: at $R \approx 16.85$ nm the absolute maximum of $\sigma_{abs}$ shifts from
Table 1: Approximate theory vs. exact Mie solution for Al.

| ℓ | λ_ℓ (nm) | R_ℓ (nm) | q | ε'_f | ε''_f | T | T_f |
|---|---|---|---|---|---|---|---|
| 1 | 143 | 6.45 | 0.283 | -2.16 | -2.19 | 0.046 | 0.216 | 0.617 | 0.617 |
| 2 | 136 | 19.27 | 0.890 | -1.67 | -1.88 | 0.047 | 0.159 | 0.901 | 0.765 |

Fig. 2: (Color online) Contour plots of the dimensionless (a) and dimensional (b) absorption cross-sections for an aluminium particle in the vacuum calculated based upon the exact Mie solution and the actual empirical dependence ϵ(ω) for aluminium [22]. Note that the localized maxima corresponding to different orders of the resonances for the cross-section normalized over πR² (a) become more extended and merge for the dimensional cross-section (b).

The developed theory does not consider effects of spatial dispersion. These effects bring about new sharp resonances related to excitation of longitudinal electromagnetic modes (volume plasmones). However, these resonances lie beyond the plasma frequency while the discussed ones lie below it. Being well-separated the two groups of the resonances practically do not interfere. As long as the surface plasmon resonances are concerned, the spatial dispersion effects give rise just to a small shift of the resonant frequencies ω_ℓ, while the lineshapes in fact remain unchanged [23]. Regarding the new resonances lying beyond the plasma frequency, extension of the developed approach to the case is straightforward.

The presented analysis has thusly revealed a small particle made of weakly dissipating materials may exhibit the anomalous absorption, which has much in common with the anomalous scattering, though they correspond to the opposite limits (the maximal and vanishing absorption, respectively). At the vicinity of the resonances the absorption cross-section as a function of R and complex ϵ is described by a simple universal formula eq. (17), which does not depend on either the size parameter q, or the order of the resonance ℓ. While at a fixed ℓ" the corresponding lineshape as a function of ℓ" is typical Lorentzian, it exhibits the pronounced deviations from the Lorentzian lineshape in transversal (on the (ℓ", ϵ")-plane) directions. The calculations for an aluminium particle indicate that the anomalous absorption should be an experimentally observable phenomenon and the absorption cross-section for such a particle may approach the upper theoretical bound \( σ_{abs max} = \frac{π(2ℓ + 1)}{2(n_m k_0)^2} \) quite closely. A future endeavor will be to carry out the corresponding experimental study, which hopefully may be motivated by this paper.

REFERENCES

[1] LORD RAYLEIGH, Philos. Mag., 41 (1871) 107; 274; 447.
[2] BOHREN C. F. and HUFFMAN D. R., Absorption and Scattering of Light by Small Particles (Wiley) 1998.
[3] VAN DE HULST H. C., Light Scattering by Small Particles (Dover) 2000.
[4] MISCHENKO M. I., HOVENIER J. W. and TRAVIS L. D. (Editors), Light Scattering by Nonspherical Particles: Theory Measurements, and Applications (Academic Press) 2000.
[5] MISCHENKO M. I., TRAVIS L. D. and LACIS A. A., Scattering, Absorption, and Emission of Light by Small Particles (Cambridge University Press, Cambridge) 2002.
[6] FUCHS R. and KLIEWER K. L., J. Opt. Soc. Am., 58 (1968) 319.
[7] FAN J. A. et al., Science, 328 (2010) 1135.
[8] LASSETER J. B. et al., Nano Lett., 10 (2010) 2694.
[9] HAO F. et al., Nano Lett., 8 (2008) 3983.
[10] TRIBEISKY M. I. and LUK'YANCHUK B. S., Phys. Rev. Lett., 97 (2006) 263902.
[11] Tribelsky M. I. et al., Phys. Rev. Lett., 100 (2008) 043903.

[12] Gil'denburg V. B. and Kondrat'ev I. G., Sov. Phys. Tech. Phys., 8 (1963) 221; Radio Eng. Electron. Phys., 10 (1965) 560.

[13] Tribelski˘i M. I., Sov. Phys. JETP, 59 (1984) 534.

[14] Chong Y. D. et al., Phys. Rev. Lett., 105 (2010) 053901.

[15] Zharov V. P. et al., Biophys. J., 90 (2006) 1.

[16] Kogan B. et al., Laser microexplosions of nanoparticles in tumor, in 25th Annual Meeting of the European Society for Hyperthermic Oncology, 2009, Verona, Italy, Abstract book, May 2009, p. 58.

[17] Ruan Z. and Fan S., Phys. Rev. Lett., 105 (2010) 013901.

[18] Hamam R. E. et al., Phys. Rev. A, 75 (2007) 053801.

[19] Ruan Z. and Fan S., J. Phys. Chem. C, 114 (2010) 7324.

[20] Landau L. D. and Lifshitz E. M., Electrodynamics of Continuous Media (Pergamon Press, Oxford, New York, Toronto, Sydney, Frankfurt) 1984.

[21] Luk’yanchuk B. S. et al., J. Phys.: Conf. Ser., 59 (2007) 234.

[22] Palik E. D., Handbook of Optical Constants of Solids (Academic Press, Orlando) 1985–1998.

[23] Ruppin R., Phys. Rev. B, 11 (1975) 2871.