Perfect Absorption of a Focused Light Beam by a Single Nanoparticle

Alexey Proskurin, Andrey Bogdanov, and Denis G. Baranov*

Absorption of electromagnetic energy by a dissipative material is one of the most fundamental electromagnetic processes that underlies a plethora of applied problems, including sensing, radar detection, wireless power transfer, and photovoltaics. Common wisdom is that a finite dissipative object illuminated by a plane wave removes only a finite amount of the wave’s energy flux, which is determined by the object’s absorption cross-section. Thus, it is of fundamental interest to see if any far-field waveform can be perfectly absorbed by a finite object. Here, it is theoretically demonstrated that a precisely tailored light beam containing only far-field components can be perfectly absorbed by a finite scatterer on a substrate. The self-consistent scattering problem in the dipole approximation is analytically solved and finds a closed-form expression for the spatial spectrum of the incident field and the required complex polarizability of the particle. All analytical predictions are confirmed with full-waves simulations. The results introduce a qualitatively novel class of perfect absorption phenomena in electromagnetics and other wave processes.

1. Introduction

The absorption of electromagnetic energy into a material is a phenomenon that underlies many applied problems, including molecular sensing, photovoltaics, and photodetection. The efficiency of absorption is the key parameter for those. Commonly, the incident energy is delivered to the system through a trivial single channel, such as a plane wave incident on one side of an absorber.[1] A classic example of an electromagnetic absorber is the Salisbury screen,[2] consisting of a thin resistive sheet placed one-quarter of the wavelength above a flat reflector. Perfect absorption can be realized in many other planar systems by the virtue of critical coupling, which requires equal radiative and dissipative decay rates of the system’s eigenmode.[3,4]

When a dissipative obstacle (e.g., a metallic nanoparticle) is illuminated by a plane wave, it absorbs a fraction of its energy flux determined by the object’s absorption cross-section. For an object of limited dimensions, this cross-section is always finite meaning the object removes only a finite amount of energy from the plane wave with an infinitely wide front.[5] Moreover, a dissipative object illuminated with a plane wave must also scatter some of the energy in order to comply with the optical theorem.[5] This makes impossible perfect absorption of a plane wave by a single finite-size object; recently the ultimate upper limit for absorption cross-section was revealed for scatterers of arbitrary shapes.[6]

Exploiting the interference of multiple incident waves allows making the electromagnetic absorption more efficient and controllable.[7–10] This is the case of the so-called coherent perfect absorbers (CPA)—devices in which complete absorption of electromagnetic radiation is achieved by the interference of several incident waves. It is possible to realize coherent perfect absorption in confined geometries, for example with transversely localized waveguide modes,[11–13] or compact scatterers in a given structured environment,[14] in a disordered environment,[15] and in free space.[16] In the latter scenario, the main difficulty is that the incident waveforms are converging cylindrical or spherical waves, which contain a lot of evanescent components in their angular spectrum.[17] The listed examples of perfect absorption mostly consider highly symmetric scatterers in homogeneous environments and neglect the substrate effect almost unavoidable in photonic systems. Thus, there is still an open question of whether a single particle of finite size on a substrate can completely absorb the incident far-field radiation.

Here, we demonstrate theoretically that a dipolar particle placed on a substrate can perfectly absorb a focused vectorial light beam, which does not have any near-field components in its spatial spectrum. Therefore, such a beam can be created in the far-field by conventional optical tools, such as phase plates and spatial light modulators, and focused onto the nanoparticle, where it will be absorbed. We reveal an analytical solution to the scattering problem, suggest a feasible system supporting this effect, and verify it with numerical full-wave simulations.
2. Results

The basic intuition behind our idea is the following: a homogeneous plane wave has an infinitely wide front and, thus, it carries an infinite energy flux \( P = \frac{1}{2} \text{Re}(E \times H) \text{d}x \). A nanoparticle having only a finite absorption cross-section approximately bounded by \( \lambda^2 \) (or a finite absorption length for a nanowire scatterer)\(^{[6,18]}\) cannot absorb all the energy of a plane wave. However, if the incident field is focused down to the diffraction limit of \( \frac{\lambda}{4} \), its cross-section becomes comparable to the maximal absorption cross-section of the nanoparticle, thus suggesting that the focused incident field can, in principle, be perfectly absorbed. Similar ideas have been discussed previously in the context of perfect reflection of a focused beam by a dipolar particle\(^{[19]}\); partial extinction of a beam by a single molecule has also been observed.\(^{[20]}\) A similar problem has been studied in Ref. \(^{[21]}\) for a cylinder, which showed that absorption cross-section can be controlled by tailoring the excitation field profile.

The system under study is illustrated in Figure 1a. We start by considering an absorbing subwavelength scatterer placed at a distance \( h \) above a perfectly conducting substrate. We will assume that the response of the scatterer is limited to an electric dipole resonance, which is a good assumption for a metallic or dielectric subwavelength particle, or an atom.

Before we proceed to treat the scattering problem, we note that the perfect absorption phenomenon is often demonstrated in the basis of spherical harmonics converging at the scatterer. That the response of the scatterer is limited to an electric dipole combination of plane waves with different \( k_p \) sharing the common \( k_z \). The red arrow illustrates the direction of the electric field of a particular \( k_p \) plane wave component of the radially polarized Bessel beam.

Thanks to the cylindrical symmetry of the problem, the solution can be sought for as a combination of radially and azimuthally polarized cylindrical beams.\(^{[21]}\) The azimuthally polarized beam has zero electric field on its axis, thus, it does not interact with a point scatterer. Therefore, we consider only the radially polarized incident beams. The electric field of an incident beam propagating in the negative \( z \)-direction can be written in cylindrical coordinates (with the radial direction \( \rho \) parallel to the interface and the scatterer at \( \rho = 0, z = 0 \)) as \( e^{-i\omega t} \text{timedependence is assumed} \)

\[
E_{\text{inc}}(r) = \begin{pmatrix} E_r \mathcal{E}_r(k) \mathcal{E}_r(k) \mathcal{E}_r(k) & E_i \mathcal{E}_i(k) \mathcal{E}_i(k) \mathcal{E}_i(k) \\ \end{pmatrix} e^{-ikz \text{dk}_p}
\]

where \( k_0 = \omega/c, k_p = \sqrt{k_0^2 - k_z^2} > 0 \), and \( E_{\text{inc}}(k) \) is the spectral amplitude of the incident field. This field represents a linear combination of Bessel beams with various \( k_0 < k_p \), each being a combination of \( p \)-polarized plane waves with a fixed \( k_z \) and all possible \( (k_0, k_p) \) satisfying \( k_0^2 + k_p^2 = k_z^2 \). The integration limit from 0 to \( k_p \) is imposed to make sure the incident field contains only far-field components.

Total scattered field can be written as the sum of three components [Figure 1a]: \( E_{\text{scat}} = E_{\text{scat,ref}} + E_{\text{scat,ref}} + E_{\text{scat,ref}} \), where \( E_{\text{ref}} \) is the electric field of the beam reflected directly from the substrate without any interaction with the scatterer, \( E_{\text{scat}} \) is the electric field scattered by the particle as if it was in a homogeneous dielectric environment, and \( E_{\text{scat,ref}} \) is the scattered electric field additionally reflected back by the substrate. Thanks to the symmetry of the problem, all three components share the common vector factor \( \mathcal{V}(k, \rho) = (-\frac{1}{\sqrt{2}} J_0(k, \rho), 0, J_0(k, \rho))^T \) allowing us to describe the scattered field with a scalar amplitude (see Note S1, Supporting
Information):

\[ E_{\text{scat}}(k_j) = \int_0^1 \left[ \tilde{E}_{\text{ref}}(k_j) + \tilde{E}_{\text{scat}}(k_j) + \tilde{E}_{\text{scat-ref}}(k_j) \right] \nabla (k_j \rho_j) e^{ik_jz} \, dk_j \]  

(2)

With the use of the multiple-scattering method\(^{23-25}\) we find the spectral amplitudes of the reflected, scattered, and scattered-reflected fields \( \tilde{E}_{\text{ref}}(k_j) \), \( \tilde{E}_{\text{scat}}(k_j) \), and \( \tilde{E}_{\text{scat-ref}}(k_j) \) (see Note S1, Supporting Information).

Perfect absorption arises when the total scattered field vanishes for all propagating channels \( k_j < k_0 \): \( \tilde{E}_{\text{ref}}(k_j) + \tilde{E}_{\text{scat}}(k_j) + \tilde{E}_{\text{scat-ref}}(k_j) = 0 \). Note that we do not require vanishing of scattered components with \( k_j > k_0 \); these evanescent components do not carry energy along the \( z \)-direction, although do carry energy along \( x \) in the general case. However, if the substrate does not support propagating guided modes, such as in the case of a perfect electric conductor (PEC) substrate, these spectral components of the scattered field do not contribute to the energy transfer in the horizontal plane. Thus, we are looking for a specific geometry of the system and incident spectral amplitude \( E_0(k_j) \) yielding the perfect absorption condition. Incorporating the expressions for the total scattered field (see Note S1, Supporting Information), we arrive at an integral Fredholm equation of the second kind:

\[ \tilde{E}_0(k_j) + \frac{i}{4\pi} \int_{k_0}^{k_j} \tilde{K}(k_j, k_j') \tilde{E}_0(k_j') \, dk_j' = 0 \]  

(3)

where \( \tilde{K}(k_j, k_j') = \frac{k_j k_j'}{k_0 k_j'} (1 + R' e^{i\delta'}) \), \( R(k_j) \) is the Fresnel reflection coefficient of the substrate (relating the \( z \)-components of the electric field), \( \delta(k_j) = 2k_j h \) is the double phase delay between the scatterer and the substrate, \( \tilde{E}_0(k_j) \) is the dressed polarizability, and \( \alpha_0 \) is the bare dipole polarizability of the scatterer. Variables with prime in \( \tilde{K}(k_j, k_j') \) indicate that they are calculated as a function of \( k_j' \): \( R' = R(k_j') \), \( \delta' = \delta(k_j') \). Thanks to the degenerate (separable) kernel, the problem admits an analytical solution. Indeed, according to Fredholm alternative,\(^{26}\) Equation (3) has a nontrivial solution if and only if

\[ \frac{1}{\alpha} = -\frac{i}{4\pi} \int_{k_0}^{k_0} \frac{1}{k_0} \left( 1 + R e^{i\delta} \right)^2 \, dk_0 \]  

(4)

The corresponding solution satisfying the integral equation with this \( \tilde{E}_0(k_j) \) is

\[ \tilde{E}_0(k_j) = A \frac{k_j^3}{Re^{i\delta}} H(k_0 - k_j) \]  

(5)

where \( A \) is a normalization constant, and \( H(k_0 - k_j) \) is the Heaviside step-function. The corresponding bare polarizability of the scatterer supporting perfect absorption for the given incident field is

\[ \alpha_0 = 4\pi i \left( \frac{k_j^3}{Re^{i\delta}} \int_{k_0}^{k_0} \frac{1}{k_0} \left( 1 + R e^{i\delta} \right)^2 \, dk_0 - \int_{k_0}^{k_0} \frac{k_j^3}{Re^{i\delta}} \, dk_0 \right)^{-1} \]  

(6)

Equations (5) and (6) are the main result of the paper. Particularly, the spectral amplitude of the incident beam \( E_0(k_j) \) marks the key difference with earlier proposals of perfect absorption by compact objects, such as Refs.\(^{14,16}\). While the incident field given by Equation (5) contains far-field propagating components only, proposals by Sentenac et al.\(^{14}\) and Noh et al.\(^{16}\) relied on the incident field containing evanescent components.

Figure 1b shows the resulting spectral amplitude \( E_0(k_j) \) of the incident field given by Equation (5) for a PEC substrate (\( R(k_j) = 1 \); note that \( R \) relates \( z \)-components of the electric field) and \( h = \lambda/4 \). Interestingly, it shows that for a perfect absorption, most of the energy should be delivered to the sphere by harmonics with \( k \approx k_0 \), propagating at large oblique angles. This is in agreement with the fact that a vertical dipole radiates mostly in the horizontal plane. Realizing such a beam experimentally might be challenging, although the current progress in high-NA metalenses may offer a solution.\(^{27}\) In addition, we expect that by engineering the shape of the absorbing nanoparticle and the substrate reflectivity spectrum, one can achieve this effect for incident beams with more Gaussian-like spectra; however, this will be the subject of future study.

Corresponding spatial distributions of the incident and total scattered electric fields (their \( z \)-components) are plotted in Figure 2. The argument of the incident field clearly reveals a phase gradient indicating the energy transfer from infinity towards the sphere [Figure 2b]. The phase of the scattered field, however, exhibits no gradual variation with only abrupt \( \pi \)-jumps [Figure 2d], reminiscent of the phase of a standing wave and visualizing lack of the energy transfer away from the scatterer.

To confirm the absence of the scattered power, we calculate the dissipation rate in the particle \( W_{\text{inc}} = \frac{1}{2} \varepsilon_0 \int \text{Im} \left| E(\eta) \right|^2 \, dV \) and compare it to the total energy flux of the incident beam \( P_{\text{inc}} = \int Re(E_{\text{inc}} \times H_{\text{inc}}^*) \, ds \) (see Note S2, Supporting Information). The analytical calculations verify that the dissipated power exactly equals the total incident flux of the perfectly absorbing beam (see Note S2, Supporting Information). Calculating the normalized absorption rate \( W_{\text{abs}} / P_{\text{inc}} \) for the ideal geometry as a function of the NA of the incident beam reveals that the amount of absorbed light drops rapidly as soon as the NA goes below unity, Figure 2f, thus indicating the critical role of the large-momentum components of the incident field in the perfect absorption phenomenon.

Equations (5) and (6) dictate not only the angular spectrum of the incident field, but also the sphere’s polarizability \( \alpha_0 \) required for perfect absorption. Figure 3a shows the real and imaginary parts of \( \alpha_0 \) as a function of the distance \( h \) between the sphere and PEC substrate. At moderate to large distances, the real part of the polarizability oscillates near zero, while the imaginary part approaches a positive constant. Interestingly, the value of this constant is exactly \( 3\pi i \frac{k_0^2}{Re^{i\delta}} \) (as one can verify by integrating Equation (6) in the limit \( h \rightarrow \infty \)), that is, the polarizability of a critically coupled dipolar scatterer having equal scattering and absorption free-space cross-sections.\(^{28,29}\) This is a remarkable result, since this is exactly the condition for a perfect absorption of a spherical harmonic by a resonant dipolar scatterer.\(^{30}\)

That and nearly zero (compared to \( \alpha_0 \)) real part of the polarizability indicate that a resonant scatterer is required to perfectly absorb the impinging beam. The required electric dipole polarizability \( \alpha_0 \) can be realized either with a subwavelength negative permittivity nanosphere, or with a Mie-resonant dielectric
Figure 2. Electric field distribution of a perfectly absorbing solution. a,b) The absolute value and phase of the $E_z$ electric field component of the incident field satisfying the perfect absorption condition for $h = \lambda/4$ and a PEC substrate. The substrate is located at $z = -h$, the absorbing sphere is at $x = 0, z = 0$. c,d) The same as (a,b) for the total scattered field, Equation (2). e) The absolute value of the $E_z$ component of the total field. f) Normalized absorption amplitude $W_{abs}/P_{inc}$ for the ideal scatterer and geometry satisfying the perfect absorption as a function of the NA of the incident field.
We choose a negative permittivity nanosphere as it is better approximated by a pure electric dipole scatterer. We can consider a few different options for the sphere’s material: either a metal giving rise to plasmonic resonance, or a polar crystal giving rise to a phonon-polariton resonance.\[^{[31]}\] To study the effect of the material’s permittivity on the perfect absorption regime, we show in Figure 3b the variation of $|\alpha_0 - \alpha_1|$ with real and imaginary parts of the sphere’s permittivity, where $\alpha = \alpha_{C-M}(1 - \frac{1}{5} e^{2} - \frac{2}{\pi} a \alpha_{C-M})^{-1}$ is the corrected electric dipole polarizability with $\alpha_{C-M} = 4\pi a^{3}/\lambda_{0}^{3}$ being the Clausius–Mossotti polarizability of a sphere of radius $a$.\[^{[30]}\] The map was calculated with a fixed $h = \lambda_{0}/4$ and $a/\lambda_{0} = 0.05$. When $a_0$ hits exactly the analytical value $\alpha_0$ at some value of the complex permittivity $\Re \varepsilon + i \Im \varepsilon$, the perfect absorption condition becomes satisfied, which can be seen as a dip in Figure 3b.

In order to evaluate if this regime can be reached with existing materials, we show parametric trajectories of complex permittivities yielding the perfect absorption condition parameterized with $a/\lambda_0$ for a series of values of $h/\lambda_0$. Thick lines denote complex permittivities of silver, aluminum, and silicon carbide, crossing the analytical solution in specific points.

Figure 3. a) Bare electric dipole polarizability $\alpha_0$ required for perfect absorption (Equation 11) as a function of the sphere-to-substrate distance $h$ for a PEC substrate; dashed line—polarizability of a critically coupled dipolar scatterer $3\pi(\frac{a}{\lambda_0})^3$. b) Logarithmic density plot of $|\alpha_0 - \alpha_1|$ as a function of the sphere’s permittivity for $h = \lambda_0/4$ and $a = 0.05\lambda_0$, where $\alpha_0$ is the polarizability of a perfectly absorbing sphere given by the analytical solution, and $\alpha_1$ is the corrected Clausius–Mossotti polarizability. c) Trajectories of complex permittivities yielding the perfect absorption condition parametrized with $a/\lambda_0$ for a series of values of $h/\lambda_0$.

The dependence of the absorption rate $A$ on the incident wavelength comes in via a number of factors. First, a deviation of the wavelength of the incident field may affect the wavefront of the incident field; it also changes the $h/\lambda$ ratio, as well as the permittivity (and dressed polarizability) of the particle. To investigate the bandwidth of the perfect absorption, we study the effect of these variations independently.

Figures S5 and S6, Supporting Information, demonstrate how the absorption rate by a particle satisfying the perfect absorption condition at a certain wavelength $\lambda_0$ changes when other parameters are varied. Figure S5, Supporting Information, shows the dependence of the absorption rate on the wavelength for a fixed particle illuminated by a field corresponding to perfect absorption at a different wavelength $\lambda$, while Figure S6, Supporting Information, shows the absorption rate calculated versus the $h/\lambda$ ratio for a fixed wavelength. They indicate that variation of these parameters does not affect strongly the perfect absorption, as the absorption rate remains above the 0.98 level when the $\lambda/\lambda_0$ or $h/\lambda$ ratio deviates by 5%.

Figure 4a demonstrates how the variation of the real and imaginary parts of the permittivity affects the absorption efficiency. Indeed, accounting for the wavelength dispersion of the nanoparticle’s permittivity near the perfect absorption point $\delta \varepsilon / \delta \lambda (\lambda_0)$, one can estimate the permittivity at the shifted wavelength $\varepsilon (\lambda_0 + \delta \lambda)$ is very close to the analytical one with some deviations that account for higher-order multipole moments. Figure S3, Supporting Information. The absorption rate calculated for the refined field reveals a maximum as high as 0.994, indicating the near-perfect absorption, Figure 4a. Its position in the permittivity space ($\varepsilon = -2.24 + 0.092i$) agrees with the one predicted by the analytical dipole approximation for the same geometry; Figure 3b. Figure 4b shows the spatial distribution of the total field obtained at the point of maximal absorption, showing agreement with the analytical total field, Figure 2e. We also compared the trajectory of complex permittivities yielding the near-perfect absorption with the analytical trajectory (see Figure S4, Supporting Information).

The dependence of the absorption rate $A$ on the incident wavelength comes in via a number of factors. First, a deviation of the wavelength of the incident field may affect the wavefront of the incident field; it also changes the $h/\lambda$ ratio, as well as the permittivity (and dressed polarizability) of the particle. To investigate the bandwidth of the perfect absorption, we study the effect of these variations independently.

Figures S5 and S6, Supporting Information, demonstrate how the absorption rate by a particle satisfying the perfect absorption condition at a certain wavelength $\lambda_0$ changes when other parameters are varied. Figure S5, Supporting Information, shows the dependence of the absorption rate on the wavelength for a fixed particle illuminated by a field corresponding to perfect absorption at a different wavelength $\lambda$, while Figure S6, Supporting Information, shows the absorption rate calculated versus the $h/\lambda$ ratio for a fixed wavelength. They indicate that variation of these parameters does not affect strongly the perfect absorption, as the absorption rate remains above the 0.98 level when the $\lambda/\lambda_0$ or $h/\lambda$ ratio deviates by 5%.

Figure 4a demonstrates how the variation of the real and imaginary parts of the permittivity affects the absorption efficiency. Indeed, accounting for the wavelength dispersion of the nanoparticle’s permittivity near the perfect absorption point $\delta \varepsilon / \delta \lambda (\lambda_0)$, one can estimate the permittivity at the shifted wavelength $\varepsilon (\lambda_0 + \delta \lambda)$ is very close to the analytical one with some deviations that account for higher-order multipole moments. Figure S3, Supporting Information.
and find the resulting absorption amplitude at that permittivity. For example, considering a particle whose permittivity is described by the Drude model that passes the perfect absorption point \( \varepsilon = -2.24 + 0.092i \), one can calculate how a small wavelength deviation changes the permittivity and what it yields in terms of the absorption rate. In the case of results presented in Figure 4a, a 0.5% wavelength deviation maintains the absorption rate above the 0.95 level, thus indicating the crucial role of the absorbing material’s dispersion.

Next, we analyzed how replacing the PEC substrate with a realistic metallic one affects the perfect absorption. We replaced PEC with aluminum and simulated the scattering of the incident beam at a fixed wavelength of 10.8 \( \mu m \), approximately corresponding to the point where the perfect absorption trajectory crosses the SiC permittivity [Figure 3c]. The incident field is composed according to Equation (5) with \( h = \lambda_0/4 \) and \( a = 0.05 \lambda_0 \). The dashed line indicates a variation of the Drude permittivity passing the perfect absorption point upon a wavelength sweep of 0.5% from the ideal wavelength. Other parameters are the same as in panel (a).

We finally note that an identical regime can be reached with a cylindrical scatterer (having an infinitely long axis) on a PEC substrate. Instead of illumination with radially polarized Bessel beams, the cylinder should be illuminated with a TE or TM linearly polarized beam. For TM polarization and an electric dipole response, the problem becomes scalar and can be solved easily (see Note S3, Supporting Information).

### 3. Discussion and Conclusion

The demonstrated effect of perfect absorption bears a close connection to the problem of dipole emission and time-reversal symmetry. It is well known that perfect absorption is a time-reversed process of lasing, which is the emission of coherent radiation.\(^{[14]}\) In this context, time-reversal has been used for subwavelength focusing in the far-field.\(^{[16,17]}\) In the problem addressed here, the obtained angular spectrum of the incident beam and the particle’s polarizability represent a stationary monochromatic solution. Applying the time-reversal operator yields another stationary solution. The stationarity of the time-reversal process is ensured by the proper polarizability of the particle precisely balancing the total incoming power carried by the fields and work performed by the field on the induced currents. Therefore, the incident beam required for perfect absorption can be found by reversing the far-field of a dipole placed above the metallic substrate, whereas the required polarizability of the dipole ensures the stationarity of the solution.

An oscillating dipole, however, produces both near and far-fields. It might appear that in order to realize perfect absorption both the near- and far-fields should be reversed. However, it is easy to see that time-reversal of the near field does not modify the field. Indeed, the near field of a vertical dipole above a substrate is a cylindrically symmetric combination of evanescent waves with all possible \( k_z = (k_x, k_y) > k_r \) and imaginary \( k_r \). Reversing this field in time corresponds to complex conjugation, flipping the sign of \( k_z \) of each spectral component, but maintaining the same \( k_r \). Therefore, this operation yields the near field identical to the initial one. The same argument applies to the near field of a linear dipole considered in Supporting Information. In other words, converting (absorbing) and diverging (radiating) counterparts of the near-field part of a dipole emission are equivalent and do not require time-reversal.

To conclude, we have demonstrated that a focused incident beam containing only far-field propagating components in its spatial spectrum can be ideally absorbed by a compact point scatterer located above a reflective substrate. We have found an analytical solution of the scattering problem in the dipole approximation, which provides the spectrum of the incident beam and the required polarizability of the absorbing particle. We have also verified the effect with full-wave simulations. Our findings significantly expand the class of the perfect absorption phenomena and offer a new tool for electromagnetic energy harvesting. An interpretation of the effect in terms of time-reversal operation also provides a simple way to generalize the perfect absorption to cases of arbitrary multipole excitations and substrates.
Note added: after acceptance of the manuscript we became aware of another paper reporting a similar effect.[39]

Supporting Information

Supporting Information is available from the Wiley Online Library or from the author.

Acknowledgements

Analytical treatment of the scattering problem was supported by the Russian Science Foundation grant 19-79-00362. D.G.B. acknowledges the Grant of the President of the Russian Federation (MK-1211.2021.1.2). A.A.B. acknowledges the RFBR (18-29-20063), the Grant of the President of the Russian Federation (MK-2224.2020.2), and the BASIS foundation.

Conflict of Interest

The authors declare no conflict of interest.

Data Availability Statement

Data sharing is not applicable to this article as no new data were created or analyzed in this study.

Keywords

focused beams, nanophotonics, perfect absorption

Received: September 30, 2020
Revised: March 19, 2021
Published online: June 13, 2021

[1] Y. Ra’di, C. R. Simovski, S. A. Tretyakov, Phys. Rev. Appl. 2015, 3, 037001.
[2] R. L. Fante, M. T. Mccormack, IEEE Trans. Antennas Propag. 1988, 36, 1443.
[3] L. Zhu, F. Liu, H. Lin, J. Hu, Z. Yu, X. Wang, S. Fan, Light: Sci. Appl. 2016, 5, e16062.
[4] K. X. Wang, Z. Yu, S. Sandhu, V. Liu, S. Fan, Optica 2014, 1, 388.
[5] H. C. van de Hulst, Light Scattering by Small Particles, Courier Corporation, New York 1981.
[6] A. E. Miroshnichenko, M. I. Tribelsky, Phys. Rev. Lett. 2018, 120, 033902.
[7] D. G. Baranov, A. Krasnok, T. Shegai, A. Alù, Y. Chong, Nat. Rev. Mater. 2017, 2, 17064.
[8] W. Wan, Y. Chong, L. Ge, H. Noh, A. D. Stone, H. Cao, Science 2011, 337, 889.
[9] Y. D. Chong, L. Ge, H. Cao, A. D. Stone, Phys. Rev. Lett. 2010, 105, 053901.
[10] J. Zhang, K. F. MacDonald, N. I. Zheleudev, Light: Sci. Appl. 2012, 1, e18.
[11] R. R. Grote, J. B. Driscoll, R. M. Osgood Jr., Opt. Lett. 2013, 38, 3001.
[12] J. M. Rothenberg, C. P. Chen, J. J. Ackert, J. I. Dadap, A. P. Knights, K. Bergman, R. M. Osgood, R. R. Grote, Opt. Lett. 2016, 41, 2537.
[13] A. I. Ignatov, I. A. Nechepurenko, D. G. Baranov, Ann. Phys. 2016, 528, 537.
[14] A. Sentenac, P. C. Chaumet, G. Leuchs, Opt. Lett. 2013, 38, 818.
[15] K. Pichler, M. Kühmayer, J. Böhm, A. Brandstötter, P. Ambichl, U. Kuhl, S. Rotter, Nature 2019, 567, 351.
[16] H. Noh, Y. Chong, A. D. Stone, H. Cao, Phys. Rev. Lett. 2012, 108, 186805.
[17] L. Novotny, B. Hecht, Principles of Nano-Optics, Cambridge University Press, Cambridge 2006.
[18] O. D. Miller, C. W. Hsu, M. T. H. Reid, W. Qiu, B. G. DeLacy, J. D. Joannopoulos, M. Soljačić, S. G. Johnson, Phys. Rev. Lett. 2014, 112, 123903.
[19] G. Zumofen, N. M. Mojarad, V. Sandoghdar, M. Agio, Phys. Rev. Lett. 2008, 101, 180404.
[20] I. Gerhardt, G. Wrigge, P. Bushev, G. Zumofen, M. Agio, R. Pfab, V. Sandoghdar, Phys. Rev. Lett. 2007, 98, 033601.
[21] J. Y. Lee, Y.-H. Chung, A. E. Miroshnichenko, R.-K. Lee, Opt. Lett. 2019, 44, 5310.
[22] Z. Bouchal, M. Olivik, J. Mod. Opt. 1995, 42, 1555.
[23] P. C. Waterman, R. Truell, J. Math. Phys. 1961, 2, 512.
[24] M. Laroche, J. J. Sáenz, R. Gómez-Medina, Opt. Express 2006, 14, 3730.
[25] F. J. G. De Abajo, Rev. Mod. Phys. 2007, 79, 1267.
[26] H. Hochstadt, Integral Equations, John Wiley & Sons, New York 2011.
[27] R. Paniagua-Dominguez, Y. F. Yu, E. Khaidarov, S. Choi, V. Leong, R. M. Bakker, X. Liang, Y. H. Fu, V. Valuckas, L. A. Krivitsky, A. I. Kuznetsov, Nano Lett. 2018, 18, 2124.
[28] S. Trey yakov, Plasmonics 2014, 9, 935.
[29] V. Grigoriev, N. Bonod, J. Wenger, B. Stout, ACS Photonics 2015, 2, 263.
[30] A. I. Kuznetsov, A. E. Miroshnichenko, M. L. Brongersma, Y. S. Kivshar, B. Luk’yanchuk, Science 2016, 354, aag2472.
[31] J. B. Khurgin, Nanophotonics 2018, 7, 305.
[32] E. D. Palik, Handbook of Optical Constants of Solids, Academic Press, Orlando 1985.
[33] A. D. Rakí, Appl. Opt. 1995, 34, 4755.
[34] T. E. Tiwald, J. A. Woollam, S. Zollner, J. Christiansen, R. B. Gregory, T. Wetteroth, S. R. Wilson, A. R. Powell, Phys. Rev. B 1999, 60, 11464.
[35] COMSOL Multiphysics v. 5.4, www.Comsol.com (accessed: February 2021).
[36] J. De Rosny, G. Lerosey, A. Tourin, Lect. Notes Comput. Sci. Eng. 2008, 59, 187.
[37] R. Carminati, R. Pierrat, J. de Rosny, M. Fink, Opt. Lett. 2016, 41, 1555.
[38] R. Carminati, R. Pierrat, J. de Rosny, M. Fink, ACS Photonics 2015, 2, 263.
[39] A. I. Kuznetsov, A. E. Miroshnichenko, M. L. Brongersma, Y. S. Kivshar, B. Luk’yanchuk, Science 2016, 354, aag2472.