Photonic forces in the near field of statistically homogeneous fluctuating sources

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Electromagnetic sources, e.g., lasers, antennas, diffusers, or thermal sources, produce a wave field that interacts with objects to transfer to them its momentum. We show that the photonic force exerted on a dipolar particle in the near field of a planar statistically homogeneous fluctuating source uniquely depends on and acts along the coordinate perpendicular to its surface. The gradient part of this force is contributed by only the evanescent components of the emitted field, its sign being opposite that of the real part of the particle polarizability. The nonconservative force is due to the propagating components, which are repulsive and constant. In addition, the source coherence length adds a degree of freedom since it largely affects these forces. The excitation of plasmons in the source surface drastically enhances the gradient force while slightly weakening the nonconservative scattering plus curl force. Hence these consequences obtained for random and partially coherent wave fields emitted by the sources addressed here should be relevant for particle manipulation at the subwavelength scale.

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I. INTRODUCTION

Photonic forces have been increasingly studied due to their potential in many disciplines ranging from physics and chemistry to biology [1–3]. Of special importance is the manipulation of dipolar particles, which are understood as those for which the incident wave excites their first electric and/or magnetic Mie coefficients [4,5]. Extensive research done on light from quasicoherent sources shows that it exerts a mechanical action on these particles through both their conservative (gradient) and nonconservative components, thus allowing the design of optical tweezers that rely on the former component [1,2,6] and their recent extensions to the subwavelength, particularly nanometric, scale [6,7]. In contrast, the scattering or radiation pressure force, which until recently was believed to push objects [3,4,7–9], has recently been designed to exert a pulling action toward the coherent source, as recently shown by exciting the induced magnetic dipole or multipoles of the particle [10,11] as well as by an appropriate choice of the illuminating wave-field angular spectrum [12].

We report here a study of the mechanical action on a dipolar particle due to a fluctuating wave field emitted by a statistical source. The subject of random fields has been extensively addressed in different contexts, ranging from macroscopic physics (e.g., in coherence theory [13,14], atmospheric turbulence [15,16], wave propagation in random and dense media [17–19], speckle formation from random phase screens [20,21], and reflection from rough surfaces [22–25]) to the microscopic and nanoscopic scales (such as systems of randomly distributed nanoparticles [14,26], quantum dots [27], and disordered photonic crystals [28], including dispersion forces between fluctuating atoms or molecules of separated objects as thermal sources and blackbodies at the nanoscale [29–37]).

Studying the optical force from a fluctuating field is of interest for optical manipulation at the subwavelength scale by random and partially coherent fields emanating from statistical sources. The nature and behavior of these forces strongly depend on the source statistics. Interestingly, we find that planar sources, of a general class such as those that are statistically stationary and homogeneous, produce gradient forces that may be either attractive or repulsive. In turn, we demonstrate that these force components are dramatically enhanced as the coherence length of the source decreases as well as when surface plasmon polaritons are excited on its surface. In contrast, the nonconservative part of the force, composed of the radiation pressure plus the spin density of the angular momentum of the electric wave vector, is pushing and constant throughout the emission half space. In this way, one can control either the tractor or the pushing effect of the resulting force on the particle according to the sign of the real part of its polarizability [12,38–49].

II. FLUCTUATING OPTICAL FORCES

Let us consider a fluctuating source emitting into $z \geqslant 0$ from the plane $z = 0$, with its volume being in the region $z < 0$ (see Fig. 1). We shall assume that the radiated random field is described by an statistical ensemble that is stationary; then we may work in the space-frequency domain [14] so that its electric vector is expressed at frequency ω as an angular spectrum of plane waves propagating throughout the half space $z > 0$ [14,50]:

$$\mathbf{E}(\mathbf{r},\omega) = \int_{-\infty}^{\infty} \mathbf{e}(k\mathbf{s}_\perp,\omega) e^{ik\mathbf{s}\cdot\mathbf{r}} d^2\mathbf{s}_\perp, \quad (1)$$

where $k = \omega/c$, with c being the speed of light in a vacuum. The propagation vector $\mathbf{k} = k\mathbf{s}$ is expressed as $\mathbf{k} = k(\mathbf{s}_\perp, s_z)$, so that $\mathbf{s}_\perp = (s_x, s_y)$ are the transversal components of $\mathbf{s}$ and $s_z = \sqrt{1 - |\mathbf{s}_\perp|^2}$ ($|\mathbf{s}_\perp|^2 \leqslant 1$) for homogeneous or propagating waves and $s_z = i\sqrt{|\mathbf{s}_\perp|^2 - 1}$ ($|\mathbf{s}_\perp|^2 > 1$) for evanescent components.

We shall describe the source as planar [14] on characterizing it by the limiting value at $z = 0$: $\mathbf{E}^{(0)}(\boldsymbol{\rho},\omega)$ [$\mathbf{r} = (\boldsymbol{\rho}, z)$] of the random field $\mathbf{E}(\mathbf{r},\omega)$ emitted into free space. It is known [14,50]

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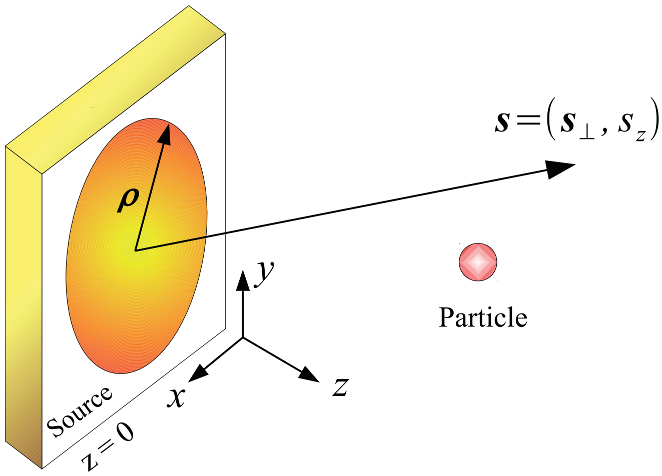

FIG. 1. (Color online) Illustration of the notation.

that

$$\mathbf{e}(k\mathbf{s}_\perp,\omega) = \left(\frac{k}{2\pi}\right)^2 \int_\Sigma \mathbf{E}^{(0)}(\boldsymbol{\rho},\omega)e^{-ik\mathbf{s}_\perp\cdot\boldsymbol{\rho}}d^2\rho, \qquad (2)$$

where Σ denotes the source domain of integration at $z = 0$.

In the Appendix we present a discussion of the relationship between the correlation of the planar source limiting value in $z = 0$, $\mathbf{E}^{(0)}(\boldsymbol{\rho},\omega)$, and the correlations of the volumetric fluctuating currents and/or polarizations employed in the study of fluctuating wave fields. There we state that in $z \simeq 0$, the subwavelength correlation length of $\mathbf{E}^{(0)}$ is similar to that of those currents and/or polarizations. Even when these latter are δ correlated, so is $\mathbf{E}^{(0)}$.

Let a dipolar particle with dynamic electric polarizability α_e be placed in the source vicinity. Since $\mathbf{p} = \alpha_e\mathbf{E}$ is the dipole moment induced in the particle by the $\mathbf{E}$ field, the ith Cartesian component ($i = 1,2,3$) of the mean force that the emitted wave field exerts on it at frequency ω is [3,4,51]

$$\begin{aligned} F_i(\mathbf{r},\omega) &= \tfrac{1}{2}\mathrm{Re}\{\alpha_e\langle E_j^*\partial_i E_j\rangle\} \\ &= \tfrac{1}{4}\mathrm{Re}\alpha_e\partial_i\langle E_j^* E_j\rangle + \tfrac{1}{2}\mathrm{Im}\alpha_e\mathrm{Im}\{\langle E_j^*\partial_i E_j\rangle\} \\ &= F_i^{\mathrm{grad}}(\mathbf{r},\omega) + F_i^{\mathrm{nc}}(\mathbf{r},\omega) \quad (i,j = 1,2,3), \qquad (3) \end{aligned}$$

expressed as the sum of a conservative, or gradient, force F_i^{grad} proportional to $\mathrm{Re}\alpha_e$ and a nonconservative term F_i^{nc} proportional to $\mathrm{Im}\alpha_e$, where Re and Im stand for the real and imaginary parts, respectively, the asterisk denotes a complex conjugate, the angular brackets mean an ensemble average, and Einstein's convention of omitting the sum symbol $\sum_{j=1}^3$ on the repeated index j has been used.

It should be noted that in writing the force as in Eq. (3) we have assumed that the particle dipole does not fluctuate itself. Otherwise, one should add a term similar to that of Eq. (3) containing both the fluctuating dipole moment and the electric field that it emits. This occurs, for example, in studies of van der Waals and Casimir forces between bodies of fluctuating atoms or molecules at a given temperature, whose current and polarization correlations are expressed by the fluctuation-dissipation theorem [29,30,32–37,41,52].

On introducing Eq. (1) into Eq. (3) one obtains

$$\begin{aligned} F_i^{\mathrm{grad}}(\mathbf{r},\omega) &= -i\frac{k}{4}\mathrm{Re}\alpha_e \iint_{-\infty}^{\infty} \mathrm{Tr}\mathcal{A}_{jk}^{(e)}(k\mathbf{s}_\perp,k\mathbf{s}'_\perp\omega) \\ &\quad \times (s_i^* - s_i')e^{-ik(\mathbf{s}^*-\mathbf{s}')\cdot\mathbf{r}}d^2\mathbf{s}_\perp d^2\mathbf{s}'_\perp, \qquad (4) \end{aligned}$$

$$\begin{aligned} F_i^{\mathrm{nc}}(\mathbf{r},\omega) &= \frac{1}{2}\mathrm{Im}\alpha_e\mathrm{Im}\left\{ik \iint_{-\infty}^{\infty} \mathrm{Tr}\mathcal{A}_{jk}^{(e)}(k\mathbf{s}_\perp,k\mathbf{s}'_\perp\omega)\right. \\ &\quad \left.\times s_i' e^{-ik(\mathbf{s}^*-\mathbf{s}')\cdot\mathbf{r}}d^2\mathbf{s}_\perp d^2\mathbf{s}'_\perp\right\}, \qquad (5) \end{aligned}$$

(with $i,j,k = 1,2,3$), where Tr denotes the trace of the electric angular correlation tensor $\mathcal{A}_{jk}^{(e)}(k\mathbf{s}_\perp,k\mathbf{s}'_\perp,\omega) = \langle e_j^*(k\mathbf{s}_\perp,\omega)e_k(k\mathbf{s}'_\perp,\omega)\rangle$. Notice that since $\langle E_j^* E_j\rangle$ is real and non-negative, F_i^{grad}, given by Eq. (4), which equals $\frac{1}{4}\mathrm{Re}\alpha_e\partial_i\langle E_j^* E_j\rangle$ according to Eq. (3), is a real quantity. Equations (4) and (5) reveal that whereas the gradient force depends on a weighted sum of the difference vectors $\mathbf{s}^* - \mathbf{s}'$ and, as we shall see, it has a negative sign if $\mathrm{Re}\alpha_e$ is positive, thus pulling the particle towards the source, the nonconservative force associated with $\mathrm{Im}\alpha_e$, which is always non-negative, depends only on the weighted sum of vectors $\mathbf{s}$ and pushes the particle forward.

A. Statistically homogeneous sources: Gradient and nonconservative forces

Let us address the wide variety of statistically homogeneous sources [38,39]. Their electric cross-spectral density tensor [14] $\mathcal{E}_{ij}(\mathbf{r}_1,\mathbf{r}_2,\omega) = \langle E_i^*(\mathbf{r}_1)E_j(\mathbf{r}_2)\rangle$ in the source plane $z = 0$ is [53] $\mathcal{E}_{ij}^{(0)}(\boldsymbol{\rho}_1,\boldsymbol{\rho}_2,\omega) = \langle E_i^{(0)*}(\boldsymbol{\rho}_1)E_j^{(0)}(\boldsymbol{\rho}_2)\rangle = \mathcal{E}_{ij}^{(0)}(\boldsymbol{\rho},\omega)$, with $\boldsymbol{\rho} = \boldsymbol{\rho}_2 - \boldsymbol{\rho}_1$, and $\mathbf{r}_\alpha = (\boldsymbol{\rho}_\alpha,z_\alpha)$, with $\alpha = 1,2$.

It is well known [14] that $\mathcal{A}_{jk}^{(e)}(k\mathbf{s}_\perp,k\mathbf{s}'_\perp\omega) = k^2\tilde{\mathcal{E}}_{jk}(k\mathbf{s}_\perp,k\mathbf{s}'_\perp\omega)$, where $\tilde{\mathcal{E}}_{jk}(k\mathbf{s}_\perp,k\mathbf{s}'_\perp\omega)$ is the four-dimensional inverse Fourier transform of $\mathcal{E}_{ij}^{(0)}(\boldsymbol{\rho}_1,\boldsymbol{\rho}_2,\omega)$. In addition, it was proven [40] that for a homogeneous source the components of the electric angular correlation tensor are δ correlated as

$$\begin{aligned} \mathcal{A}_{jk}^{(e)}(k\mathbf{s}_\perp,k\mathbf{s}'_\perp,\omega) &= k^4\delta^{(2)}[k(\mathbf{s}_\perp - \mathbf{s}'_\perp),\omega] \\ &\quad \times \tilde{\mathcal{E}}_{jk}^{(0)}\left[\frac{k}{2}(\mathbf{s}_\perp + \mathbf{s}'_\perp),\omega\right], \qquad (6) \end{aligned}$$

where $\delta^{(2)}$ represents the two-dimensional Dirac delta function.

On introducing the above δ-function expression for $\mathcal{A}_{jk}^{(e)}(k\mathbf{s}_\perp,k\mathbf{s}'_\perp,\omega)$ into Eqs. (4) and (5) one straightforwardly obtains for the gradient force

$$\begin{aligned} F_i^{\mathrm{grad}}(z,\omega) &= F_{z,\mathrm{ev}}^{\mathrm{grad}}(z,\omega) \\ &= -i\frac{k^3}{4}\mathrm{Re}\alpha_e \int_{|\mathbf{s}_\perp|^2>1} \mathrm{Tr}\tilde{\mathcal{E}}_{jk}^{(0)}(k\mathbf{s}_\perp,\omega) \\ &\quad \times (s_i^* - s_i)e^{-2k\sqrt{|\mathbf{s}_\perp|^2-1}z}d^2\mathbf{s}_\perp. \qquad (7) \end{aligned}$$

The subindex in the integral of Eq. (7) means that the integration extends only to the nonradiative region because the difference vector $\mathbf{s}^* - \mathbf{s}$ in Eq. (4) is clearly zero for propagating waves ($|\mathbf{s}_\perp|^2 \leqslant 1$). Therefore the radiative components of the field emitted by statistically homogeneous sources do not contribute to the gradient force, which depends only on the evanescent components ($|\mathbf{s}_\perp|^2 > 1$) for which $\mathbf{s}^* - \mathbf{s} = (0,0,s_z^* - s_z) = (0,0,-2i\sqrt{|\mathbf{s}_\perp|^2 - 1})$. Hence this force exists only in the near field and depends on the distance z of the particle to the source, having solely a z component normal to its surface. In addition, this force is attractive or repulsive depending on the sign of $\mathrm{Re}\alpha_e$. Small particles with relative permittivity $\varepsilon > 1$ have $\mathrm{Re}\alpha_e > 0$ out of resonance and thus $F_z^{\mathrm{grad}}(z,\omega)$ will drag them toward the source. Conversely, near a resonance $\mathrm{Re}\alpha_e$ may be negative [3]; thus this force is repulsive. However, further study is required in this latter case since then the particle strongly scatterers the field emitted by the source and therefore the analysis developed here should not be exact due to multiple scattering of the radiation between the source and the particle. Hence it is shown that the gradient force near a statistically homogeneous source is entirely of nonradiative nature and may work as a tractor force [10–12,46].

Analogously, from Eq. (5) one also derives for the nonconservative force F_i^{nc} a dependence on z only:

$$\begin{aligned} F_i^{\mathrm{nc}}(z,\omega) &= F_{i,h}^{\mathrm{nc}}(z,\omega) + F_{i,\mathrm{ev}}^{\mathrm{nc}}(z,\omega) \\ &= \frac{k^3}{2}\mathrm{Im}\alpha_e \mathrm{Im}\left\{ i \int_{|\mathbf{s}_\perp|^2 \leqslant 1} \mathrm{Tr}\tilde{\mathcal{E}}_{jk}^{(0)}(k\mathbf{s}_\perp,\omega) s_i d^2\mathbf{s}_\perp \right\} \\ &\quad + \frac{k^3}{2}\mathrm{Im}\alpha_e \mathrm{Im}\left\{ i \int_{|\mathbf{s}_\perp|^2 > 1} \mathrm{Tr}\tilde{\mathcal{E}}_{jk}^{(0)}(k\mathbf{s}_\perp,\omega) \right. \\ &\quad \left. \times\, s_i e^{-2k\sqrt{|\mathbf{s}_\perp|^2-1}z} d^2\mathbf{s}_\perp \right\}, \end{aligned} \tag{8}$$

where $F_{i,h}^{\mathrm{nc}}$ and $F_{i,\mathrm{ev}}^{\mathrm{nc}}$, denote propagating and evanescent wave contributions, which correspond to the first and second integral terms of Eq. (8), respectively. Notice that $F_{i,h}^{\mathrm{nc}} > 0$ is constant throughout $z > 0$.

Let the source also be statistically isotropic [14] so that $\mathcal{E}_{ij}^{(0)}(\boldsymbol{\rho}_1,\boldsymbol{\rho}_2,\omega) = \mathcal{E}_{ij}^{(0)}(\rho,\omega)$, where $\rho = |\boldsymbol{\rho}_1 - \boldsymbol{\rho}_2|$. The spatial coherence function of the field in $z = 0$ is [53–55] $\mathrm{Tr}\mathcal{E}_{ij}^{(0)}(\rho,\omega)$ and the spectral degree of spatial coherence $\mu^{(0)}(\rho,\omega) = \mathrm{Tr}\mathcal{E}_{ij}^{(0)}(\rho,\omega)/S^{(0)}(\omega)$, where the wave-field spectrum on the source is $S^{(0)}(\omega) = \mathrm{Tr}\mathcal{E}_{ij}^{(0)}(0,\omega)$.

To illustrate these results, we shall consider a Gaussian spectral degree of coherence $\mu^{(0)}(\rho,\omega) = \exp[-\rho^2/2\sigma^2]$, so that taking the Fourier inverse one obtains

$$\begin{aligned} \mathrm{Tr}\tilde{\mathcal{E}}_{jk}^{(0)}(k\mathbf{s}_\perp,\omega) &= S^{(0)}(\omega)\tilde{\mu}^{(0)}(k\mathbf{s}_\perp,\omega) \\ &= S^{(0)}(\omega)(\sigma^2/2\pi)\exp[-(k\sigma|\mathbf{s}_\perp|)^2/2], \end{aligned} \tag{9}$$

where σ is the correlation, or coherence, length of the source. We shall write [45]

$$S^{(0)}(\omega) = \mathcal{S}(\omega)/2\pi\sigma^2 \tag{10}$$

to express that when $\sigma \to 0$, the source is δ correlated so that $\mathrm{Tr}\mathcal{E}_{ij}^{(0)}(\rho,\omega) \to \mathcal{S}(\omega)\delta^{(2)}(\rho)$. Here $\mathcal{S}(\omega)$ is a positive quantity such that $\mathcal{S}(\omega) = \int \mathrm{Tr}\mathcal{E}_{ij}^{(0)}(\rho,\omega)d^2\rho$ for any σ. Equations (9) and (10) convey a normalization of $\mathrm{Tr}\mathcal{E}_{jk}^{(0)}(\rho,\omega)$ in the $\boldsymbol{\rho}$ space that is analogous to the ω normalization of the source spectrum in the theory of partial coherence (see Refs. [56,57]).

On introducing Eq. (9) for $\mathrm{Tr}\tilde{\mathcal{E}}_{jk}^{(0)}(k\mathbf{s}_\perp,\omega)$ into the force equations (7) and (8), employing cylindrical coordinates $s_x = s_\perp\cos\phi$ and $s_y = s_\perp\sin\phi$, and making use of the rotational symmetry of the source, we obtain that the transversal components of the nonconservative force are zero, viz., $F_{x,y}^{\mathrm{nc}}(z,\omega) = 0$, as are the corresponding integrals of Eq. (8) when one performs the azimuthal angle ϕ integration. Also, since $s_z = i\sqrt{|\mathbf{s}_\perp|^2 - 1}$ for $|\mathbf{s}_\perp|^2 > 1$, the second integral in Eq. (8) is purely imaginary, which implies that $F_{z,\mathrm{ev}}^{\mathrm{nc}} = 0$. Hence

$$\begin{aligned} F_i^{\mathrm{nc}}(z,\omega) &= F_{i,h}^{\mathrm{nc}}(z,\omega) \\ &= \frac{k^3}{2}\mathrm{Im}\alpha_e \mathrm{Im}\left\{ i \int_{|\mathbf{s}_\perp|^2 \leqslant 1} \mathrm{Tr}\tilde{\mathcal{E}}_{jk}^{(0)}(k\mathbf{s}_\perp,\omega) s_i d^2\mathbf{s}_\perp \right\}. \end{aligned} \tag{11}$$

Thus, while $F_{z,h}^{\mathrm{nc}}(z,\omega) > 0$ is constant throughout $z > 0$, as is the spectrum $S^{(0)}(\omega)$ propagating into $z > 0$ [43], the evanescent waves do not contribute to the nonconservative force $F_z^{\mathrm{nc}}(\mathbf{r},\omega)$.

In summary, there are two force components acting on the particle, $F_{z,\mathrm{ev}}^{\mathrm{grad}}(z,\omega)$ and $F_{z,h}^{\mathrm{nc}}(z,\omega)$, which are perfectly distinguishable from each other since the former is due to the nonradiative plane-wave components of the emitted field, whereas to the latter only the radiative components contribute. As the distance from the particle to the source plane grows to values $z > \lambda$, $F_{z,\mathrm{ev}}^{\mathrm{grad}}(z,\omega)$ tends to zero due to its evanescent wave composition. Nevertheless, as we shall see, the source coherence length σ plays an important role in these contributions.

The integration of Eqs. (7) and (11) using the Gaussian spectral degree of coherence, quoted before, $\mu^{(0)}(\rho,\omega) = \exp[-\rho^2/2\sigma^2]$, leads to an analytical expression for the gradient and for the nonconservative force. For the latter, Eq. (19) yields the proportion of radiation pressure and curl components for unpolarized emission. This calculation is straightforwardly done again on setting $s_x = s_\perp\cos\phi$ and $s_y = s_\perp\sin\phi$ and leads to

$$\begin{aligned} F_z^{\mathrm{grad}}(z,\omega) &= \mathrm{Re}\alpha_e \mathcal{S}(\omega) e^{-k^2\sigma^2/2}\frac{1}{\sigma^2}\left[\frac{z}{\sigma^2} \right. \\ &\quad \left. - \sqrt{\frac{\pi}{2}}\left(\frac{2z^2}{\sigma^3} + \frac{1}{2\sigma}\right) e^{2z^2/\sigma^2}\mathrm{erfc}(\sqrt{2}z/\sigma)\right], \end{aligned} \tag{12}$$

$$\begin{aligned} F_z^{\mathrm{nc}}(z,\omega) &= \mathrm{Im}\alpha_e \mathcal{S}(\omega)\frac{1}{2\sigma^2} \\ &\quad \times \left[k - \frac{1}{\sigma}\sqrt{\frac{\pi}{2}} e^{-k^2\sigma^2/2}\mathrm{erfi}(k\sigma/\sqrt{2})\right], \end{aligned} \tag{13}$$

where $\mathrm{erfc}(x) = 1 - \mathrm{erf}(x)$, with $\mathrm{erf}(x)$ being the error function $\mathrm{erf}(x) = 2/\sqrt{\pi}\int_0^x e^{-t^2}dt$, and $\mathrm{erfi}(x)$ is a positive real function defined as $\mathrm{erfi}(x) = \mathrm{erf}(ix)/i$.

B. The curl force

It is well known [9,51] that the nonconservative part of the force F_i^{nc} is the sum of a scattering force, or radiation pressure

$$\begin{aligned} F_i^{\mathrm{nc}} &= (k/2)\mathrm{Im}\alpha_e \mathrm{Re}\langle \mathbf{E}\times\mathbf{B}^*\rangle_i \\ &= (1/2)\mathrm{Im}\alpha_e \mathrm{Im}\{\langle E_j^*\partial_i E_j\rangle - \langle E_j^*\partial_j E_i\rangle\}, \end{aligned} \tag{14}$$

given by the averaged field Poynting vector, plus the curl of an electric spin density

$$F_i^{\rm nc,curl} = (1/2){\rm Im}\alpha_e {\rm Im}\langle(\mathbf{E}^* \cdot \nabla)\mathbf{E}\rangle_i = (1/2){\rm Im}\alpha_e {\rm Im}\langle E_j^* \partial_j E_i\rangle. \quad (15)$$

If the field emitted by the source is unpolarized, $\mathcal{E}_{jk}^{(0)}(\rho,\omega) = \mathcal{F}^{(0)}(\rho,\omega)\delta_{jk}$, with $\mathcal{F}^{(0)}(\rho,\omega)$ being a scalar spatial correlation function whose two-dimensional Fourier transform will be denoted as $\tilde{\mathcal{F}}^{(0)}(k\mathbf{s}_\perp,\omega)$, then

$$\mathrm{Tr}\tilde{\mathcal{E}}_{jk}^{(0)}(k\mathbf{s}_\perp,\omega) = 3\tilde{\mathcal{F}}^{(0)}(k\mathbf{s}_\perp,\omega) \quad (16)$$

and the radiation pressure contribution $F_i^{\rm nc,pr}$ to the nonconservative force is

$$\begin{aligned} F_i^{\rm nc,pr} &= \frac{k^3}{2}{\rm Im}\alpha_e \int_{|\mathbf{s}_\perp|\leqslant 1} \left[\mathrm{Tr}\tilde{\mathcal{E}}_{jk}^{(0)}(k\mathbf{s}_\perp,\omega)s_i - \tilde{\mathcal{E}}_{ji}^{(0)}(k\mathbf{s}_\perp,\omega)s_j\right] d^2\mathbf{s}_\perp \\ &= \frac{k^3}{2}{\rm Im}\alpha_e \int_{|\mathbf{s}_\perp|\leqslant 1} [3\tilde{\mathcal{F}}^{(0)}(k\mathbf{s}_\perp) - \tilde{\mathcal{F}}^{(0)}(k\mathbf{s}_\perp)] s_i d^2\mathbf{s}_\perp \\ &= k^3{\rm Im}\alpha_e \int_{|\mathbf{s}_\perp|\leqslant 1} \tilde{\mathcal{F}}^{(0)}(k\mathbf{s}_\perp) s_z d^2\mathbf{s}_\perp = F_z^{\rm nc,pr} \end{aligned} \quad (17)$$

since the azimuthal angle integration when s_i is either s_x or s_y is zero. In a similar manner, the curl force contribution $F_i^{\rm nc,curl}$ to $F_i^{\rm nc}$ is

$$\begin{aligned} F_i^{\rm nc,curl} &= \frac{k^3}{2}{\rm Im}\alpha_e \int_{|\mathbf{s}_\perp|\leqslant 1} \tilde{\mathcal{E}}_{ji}^{(0)}(k\mathbf{s}_\perp,\omega)s_j d^2\mathbf{s}_\perp \\ &= \frac{k^3}{2}{\rm Im}\alpha_e \int_{|\mathbf{s}_\perp|\leqslant 1} \tilde{\mathcal{F}}^{(0)}(k\mathbf{s}_\perp)s_z d^2\mathbf{s}_\perp = F_z^{\rm nc,curl}. \end{aligned} \quad (18)$$

Namely, for unpolarizad radiation

$$F_z^{\rm nc,pr} = 2F_z^{\rm nc,curl} = \tfrac{2}{3}F_z^{\rm nc}. \quad (19)$$

III. EXCITATION OF SURFACE PLASMON POLARITONS: NUMERICAL RESULTS

A. Normalized force

Without loss of generality, we shall also address surface plasmon polaritons (SPPs) that are excited on the source plane $z = 0$. Let this be gold, for example, choosing, for instance, $\lambda = 459.9$ nm; its permittivity is $\varepsilon = -2.546 + i3.37$ [58]. The SPP wave vector $k\mathbf{s}_\perp = k\mathbf{s}_\perp^{\rm SPP} = \pm k[\varepsilon/(\varepsilon+1)]^{1/2}$ corresponds to a pole of the Fresnel transmission coefficient (in the assumed transmission setup configuration) $T^{(p)}(k\mathbf{s}_\perp,\omega)$ [50,59] (see also the Appendix). Then it is easy to obtain that Eqs. (7) and (8) are valid upon replacing $\tilde{\mathcal{E}}_{jk}^{(0)}(k\mathbf{s}_\perp,\omega)$ by $\tilde{\mathcal{E}}_{jk}^{(0)}(k\mathbf{s}_\perp,\omega)|T^{(p)}(k\mathbf{s}_\perp,\omega)|^2$ [45].

Figure 2 shows the normalized value $F_z^{\rm grad}(z,\omega)/[k^3\mathcal{S}(\omega){\rm Re}\alpha_e/(2\pi)^2]$ of the attractive gradient optical force due to evanescent components for two random sources: one without and one with excited SPPs [see Figs. 2(a) and 2(b), respectively]. As predicted by Eq. (7), the normalized gradient force drags the particle toward the source plane (notice that since this normalization does not include ${\rm Re}\alpha_e$, it does not

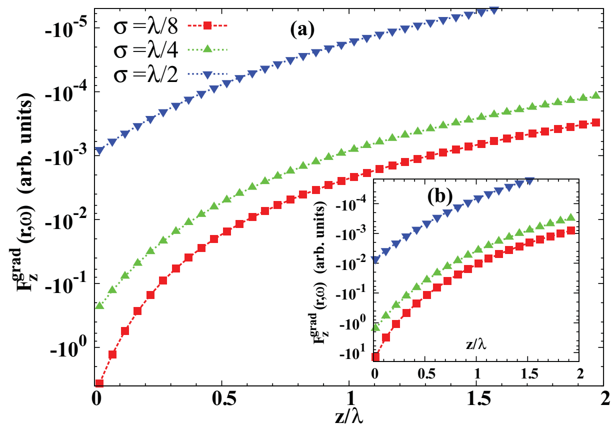

FIG. 2. (Color online) Pulling gradient optical force due to evanescent components. (a) Gradient force versus distance to the source z/λ for different values of the source coherence length σ. (b) Same force as in (a) when SPPs are excited in the source. A significant decrease of the magnitude of this force is clearly seen as σ increases to $\lambda/2$ and beyond. The normalized value $F_z^{\rm grad}(z,\omega)/[k^3\mathcal{S}(\omega){\rm Re}\alpha_e/(2\pi)^2]$ is represented in arbitrary units.

contain an eventual negative value of this quantity). In both figures we observe its exponential increase as the distance z of the particle to the source decreases. Nevertheless, this force is mainly governed by the coherence length σ. For $\sigma = \lambda/8$ (red line with squares), the magnitude of this force is maximum, but we observe that it presents an important decrease, with values between 10^{-3} and 10^{-4}, at approximately $\sigma = \lambda/2$ (blue line with down triangles) and beyond, even at subwavelength distances z, which are practically zero ($F_z^{\rm grad} \simeq 10^{-8}$) for $\sigma = \lambda$ and $z = 0$ (this latter curve is not shown).

Hence we demonstrate that the decrease of the source coherence length gives rise to an increase of the gradient force and its effect is larger than that of the distance z of the particle to the source plane. Eventually, a δ-correlated source (thus $\sigma \to 0$), e.g., a thermal source, will maximize this force. In addition, we show in Fig. 2(b) that the excitation of SPPs in the source increases the strength of this near-field force by approximately one order of magnitude. This is due to the then larger values of $\tilde{\mathcal{E}}_{jk}^{(0)}(k\mathbf{s}_\perp,\omega)|T^{(p)}(k\mathbf{s}_\perp,\omega)|^2$ stemming from the pole of $|T^{(p)}(k\mathbf{s}_\perp,\omega)|^2$ at $k\mathbf{s}_\perp^{\rm SPP}$.

Correspondingly, Figs. 3(a) and 3(b) show the normalized total force $F_z^{\rm tot}(z,\omega) = [(2\pi)^2/k^3\mathcal{S}(\omega)][F_z^{\rm grad}(z,\omega)/{\rm Re}\alpha_e + F_z^{\rm nc}(z,\omega)/{\rm Im}\alpha_e]$, in arbitrary units, without and with SPP excitation, respectively. At large distances ($z > \lambda$), the total force is a constant of the distance z and repulsive according to the behavior of the nonconservative component $F_z^{\rm nc}$, which dominates in this region of z, regardless of the value of σ. In addition, this nonconservative force increases as σ decreases.

One might think that, due to its evanescent wave composition, the magnitude of the normalized gradient force at subwavelength distances would be higher than that of the corresponding nonconservative force; however, this is not completely true due the larger effect of the source coherence length on $F_z^{\rm grad}$ rather than on $F_z^{\rm nc}$. In the near field $F^{\rm tot} \simeq F^{\rm grad}$ for $\sigma \leqslant \lambda/4$; however, as σ increases, $F^{\rm grad}$ becomes negligible, being for $\sigma > \lambda/4$ $F^{\rm tot} \simeq F^{\rm nc}$. These effects appear

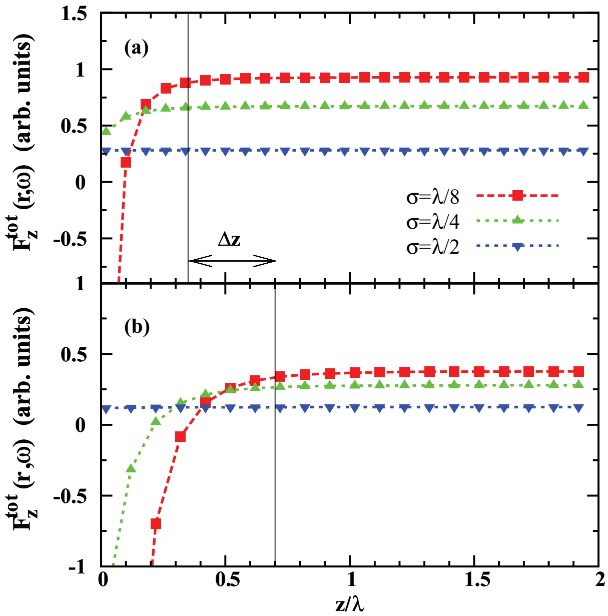

FIG. 3. (Color online) (a) Normalized total optical force (see the text) in arbitrary units versus the distance z/λ to the source for different values of the coherence length σ. (b) Same as in (a) when SPPs are excited. In this latter case we observe an increment Δz in the value z/λ at which the magnitude of the gradient force starts to exponentially increase.

in Figs. 3(a) and 3(b). In particular we see in Fig. 3(b) that if SPPs are excited, an increment Δz appears in the distance z/λ where the magnitude of the attractive gradient component is noticeable [compare Figs. 3(a) and 3(b)]. The enhancement of the near-field intensity due to SPP resonances then conveys a larger range of the gradient force.

IV. A SOURCE SPECTRUM MODEL TO ILLUSTRATE THE OPTICAL FORCE ON A DIPOLAR PARTICLE

So far we have calculated the relative weights of the gradient and nonconservative forces without taking into account the strength of the particle-induced dipole. This was done by normalizing those forces to the corresponding real and imaginary values of the particle polarizability. However, it is well known that this latter quantity largely influences the values of these forces [4,5]; hence its presence should be relevant in estimating the actual mechanical action of the emitted light. We shall now address this with a certain source spectrum model.

A small particle of radius r_0 with relative permittivity ε_p in the Rayleigh limit ($kr_0 \ll 1$) is assumed. We adopt the expression for the dynamic electric polarizability that conserves energy on scattering [4], namely,

$$\alpha_e = \alpha^{(0)} \left(1 - i\tfrac{2}{3}k^3\alpha^{(0)}\right)^{-1}, \tag{20}$$

with $\alpha^{(0)}$ being the static polarizability

$$\alpha^{(0)} = r_0^3 \frac{\varepsilon_p - 1}{\varepsilon_p + 2}. \tag{21}$$

For a nonmonochromatic source, the total mean force exerted by the random field on the dipolar particle is determined by ω integration of each frequency component, i.e.,

$$F_i(\mathbf{r}) = \int F_i(\mathbf{r},\omega)d\omega \quad (i = 1,2,3). \tag{22}$$

Of course, as discussed above, in our study the source contributes with the Cartesian component $F_z(\mathbf{r},\omega)$ only.

We assume a Gaussian spectrum model [14] [cf. Eq. (10)] so that

$$\mathcal{S}(\omega) = \frac{A}{\sigma_\omega\sqrt{2\pi}} e^{-(\omega-\omega_0)^2/2\sigma_\omega^2}, \tag{23}$$

where A, σ_ω, and ω_0 are positive constants. Incidentally, it is straightforward to see that in the monochromatic case at frequency ω_0 ($\sigma_\omega \to 0$), we recover Eqs. (12) and (13) with $\mathcal{S}(\omega) = A\delta(\omega - \omega_0)$.

As an example we perform the integral (22) for a small dielectric particle of radius $r_0 = 25$ nm with a constant value of $\varepsilon_p = 2{,}25$ in the range of studied frequencies. In the case of SPP excitation, an Au source surface is considered like in Sec. III A, with a frequency variation of its permittivity approximated by [60,61]

$$\varepsilon(\omega) = \varepsilon_\infty - \frac{\omega_d^2}{\omega^2 + i\gamma\omega} + \sum_{p=1}^{2} A_p\Omega_p \left(\frac{e^{i\phi_p}}{\Omega_p - \omega - i\Gamma_p} + \frac{e^{-i\phi_p}}{\Omega_p + \omega + i\Gamma_p}\right), \tag{24}$$

where the values of the parameters are considered to be the same as in Ref. [61].

An optical power of 300 mW impinging the particle is assumed at the central wavelength $\lambda_0 = 2\pi c/\omega_0 = 579$ nm. Then $\mathrm{Re}\alpha_e = 4593\ \mathrm{nm}^3 \gg \mathrm{Im}\alpha_e = 17\ \mathrm{nm}^3$. The spectral width is taken as $\sigma_\omega = 0.01\omega_0$. The constant A is then adjusted to these values.

The total force calculated on introducing Eqs. (12) and (13) into Eq. (22) (that is to say, this time without introducing any normalization) is shown in Figs. 4(a) and 4(b), without and with excitation of surface plasmon polariton resonances on the source surface, respectively.

As predicted in Sec. III, the total force is governed in the near field by its gradient component for $\sigma \leqslant \lambda/4$, with its magnitude increasing as the coherence length σ decreases. In addition, its exponential growth as the particle approaches the source is remarkable. In particular, at $z/\lambda = 0.5$ and for $\sigma = \lambda/8$, the magnitude of the gradient force when SPPs are excited is practically double (3×10^{-16} N) than when they are absent.

For $z/\lambda > 1$ the total force is due to its nonconservative part; however, the distance at which this force begins to dominate is larger than that shown in Fig. 3 for the normalized force. This is due to the fact that now we have introduced in the calculations $\mathrm{Re}\alpha$, which is much larger than $\mathrm{Im}\alpha$.

In contrast, Fig. 4 manifests a behavior of both the gradient and scattering plus curl forces similar to that of their nonintegrated normalized counterparts, shown in Fig. 3, both without and with SPP excitation. However, the action of the gradient force reaches larger distances from the source than

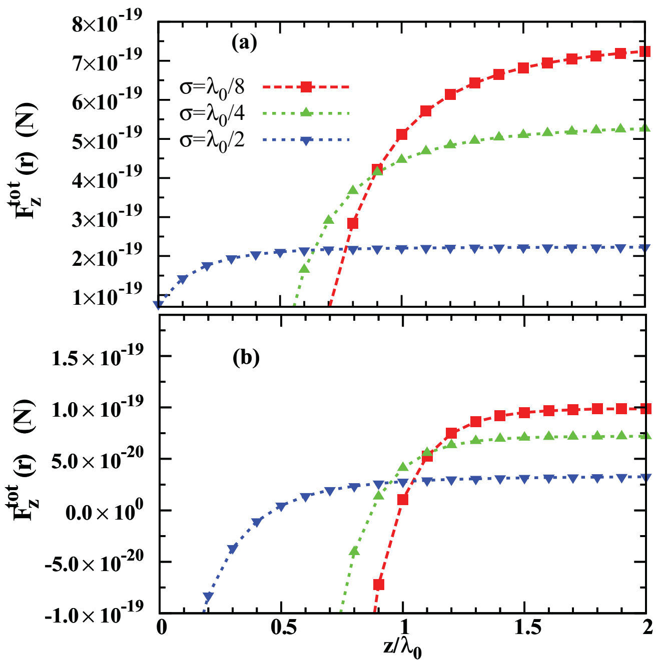

FIG. 4. (Color online) Total optical force in Newtons versus distance z/λ_0, obtained by integration over a Gaussian spectrum of the source. Results for different values of the coherence length σ_0 are shown. Surface plasmons polaritons (a) are absent and (b) are excited on the metallic surface of the source.

its normalized counterpart as Fig. 4 shows in comparison with Fig. 3, at least within the scale of values shown here.

V. CONCLUSION

We have reported a study of photonic forces exerted on dipolar particles by discussing near-field effects due to fluctuating sources. The behavior of gradient and nonconservative forces opens the possibility of new theoretical and experimental observations on subwavelength phenomena, particularly at the nanoscale, concerning optical manipulation in physics cases such as those ranging from emission by partially correlated primary sources (i.e., beyond δ-correlated thermal sources and blackbodies), e.g., fluctuations in nanoantennas, to secondary sources resulting from light propagation through the turbulent atmosphere [29], e.g., speckle patterns from a large variety of statistical structures including scatterers, random rough surfaces, phase screens, and optical diffusers [62–64].

We have seen that in the large variety of stationary statistically homogeneous and isotropic sources, only the evanescent components contribute to the gradient forces, while the nonconservative part that contains radiation pressure and curl forces is due solely to emitted propagating components. Hence the subwavelength information is encoded in the gradient forces. Same numerical examples were given for statistically isotropic unpolarized emitted wave fields, showing the important effect that the source coherence length has on these forces, especially on the gradient component. In addition, due to the higher concentration of energy in the near field when there is excitation of surface waves in the source, this largely enhances the magnitude of the gradient part of these forces while slightly diminishing the strength of their nonconservative part. We expect that these findings will stimulate further experiments and applications in this particle manipulation scenario, which may also have implications when fluctuations in the small particles, which are analogous to those leading to van der Waals and Casimir interactions between bodies, are also addressed.

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APPENDIX

In this appendix we relate the limiting value of the emitted field at $z = 0$, $\mathbf{E}^{(0)}(\boldsymbol{\rho},\omega)$ $[\mathbf{r} = (\boldsymbol{\rho},z)]$, to the fluctuating currents $\mathbf{j}(\mathbf{r},\omega)$ and polarizations $\mathbf{P}(\mathbf{r},\omega)$ of the source (assumed to be nonmagnetic) that characterize the emission. If we let the source occupy a volume V in $z < 0$, the electric field emitted into the free space $z \geqslant 0$ by these distributions is

$$\mathbf{E}(\mathbf{r},\omega) = \int_V \left[\frac{ik}{c}\mathbf{j}(\mathbf{r}',\omega) + k^2\mathbf{P}(\mathbf{r}',\omega)\right]\mathcal{G}(\mathbf{r},\mathbf{r}')d^3\mathbf{r}', \tag{A1}$$

where $\mathcal{G}(\mathbf{r},\mathbf{r}')$ is the dyadic Green's function whose elements $\mathcal{G}_{ij}$ $(i,j = 1,2,3)$ are [44,65]

$$\mathcal{G}_{ij}(\mathbf{r},\mathbf{r}',\omega) = \frac{i}{8\pi^2}\int_{-\infty}^{\infty}\frac{\hat{s}_i T^{(s)}(k\mathbf{s}_\perp,\omega)\hat{s}_j + \hat{p}_i T^{(p)}(k\mathbf{s}_\perp,\omega)\hat{p}_j^{(1)}}{ks_z} \times e^{ik(\mathbf{s}\cdot\mathbf{r}-\mathbf{s}'\cdot\mathbf{r}')}d^2\mathbf{s}_\perp. \tag{A2}$$

Here the caret denotes a unit vector and $\hat{\mathbf{s}} = \hat{\mathbf{s}}_\perp \times \hat{z}$, $\hat{\mathbf{p}} = |\mathbf{s}_\perp|\hat{z} + s_z\hat{\mathbf{s}}_\perp$, $\hat{\mathbf{p}}^{(1)} = |\mathbf{s}_\perp|\hat{z} + s_z'\hat{\mathbf{s}}_\perp$, $\mathbf{s} = (\mathbf{s}_\perp,s_z)$, $\mathbf{s}' = (\mathbf{s}_\perp,s_z')$, and $s_z' = \sqrt{\varepsilon - |\mathbf{s}_\perp|^2}$ (for $|\mathbf{s}_\perp|^2 \leqslant \varepsilon$) and $i\sqrt{|\mathbf{s}_\perp|^2 - \varepsilon}$ (for $|\mathbf{s}_\perp|^2 > \varepsilon$). The superindices (s) and (p) denote the Fresnel transmission coefficient under s and p polarization, respectively. The source medium permittivity is denoted by ε, which we shall assume is isotropic.

From Eqs. (A1) and (A2) one obtains the correlation function of the emitted field

$$\begin{aligned}\langle E_i^*(\mathbf{r},\omega)E_j(\mathbf{r}',\omega')\rangle &= \iint_V \mathcal{G}_{ik}^*(\mathbf{r},\mathbf{r}'')\mathcal{G}_{jl}(\mathbf{r}',\mathbf{r}''')\left[\left(\frac{k}{c}\right)^2\langle j_k^*(\mathbf{r}'',\omega)j_l(\mathbf{r}''',\omega')\rangle + k^4\langle P_k^*(\mathbf{r}'',\omega)P_l(\mathbf{r}''',\omega')\rangle\right]d^3\mathbf{r}''d^3\mathbf{r}''' \\ &= \int_V \mathrm{Im}\mathcal{G}_{ij}(\mathbf{r}-\mathbf{r}'-\mathbf{R})\frac{1}{\mathrm{Im}\varepsilon}\left(\frac{1}{c}\right)^2[\langle j_k^*(\mathbf{r}'',\omega)j_k(\mathbf{r}''+\mathbf{R},\omega')\rangle + k^2\langle P_k^*(\mathbf{r}'',\omega)P_k(\mathbf{r}''+\mathbf{R},\omega')\rangle]d^3\mathbf{R},\end{aligned} \tag{A3}$$

where $\langle j_k^*(\mathbf{r},\omega)j_l(\mathbf{r}',\omega')\rangle = \delta_{kl}\langle j_k^*(\mathbf{r},\omega)j_k(\mathbf{r}',\omega')\rangle\delta(\omega-\omega')$ (with an analogous expression for the polarization correlations) due to the isotropy of the source medium.

In addition, we have taken into account that $\mathcal{G}(\mathbf{r}'',\mathbf{r}''') = \mathcal{G}(\mathbf{r}'' - \mathbf{r}''')$, having written $\mathbf{R} = \mathbf{r}''' - \mathbf{r}''$ and using the equality [44,66]

$$k^2 \int_V \mathrm{Im}\varepsilon \mathcal{G}^*_{ik}(\mathbf{r} - \mathbf{r}'')\mathcal{G}_{jk}(\mathbf{r}' - \mathbf{r}'')d^3\mathbf{r}'' = \mathrm{Im}\mathcal{G}_{ij}(\mathbf{r} - \mathbf{r}'). \tag{A4}$$

In particular, in the Lifshitz-Rytov theory in which the emitted wave is a thermal field at temperature T, the sources are characterized by the currents $\mathbf{j}(\mathbf{r},\omega)$, obeying the fluctuation-dissipation theorem [26,29,52,67]

$$\langle j_i^*(\mathbf{r},\omega) j_j(\mathbf{r}',\omega')\rangle = \frac{\omega\Theta(\omega,T)}{4\pi^2}\mathrm{Im}\varepsilon\delta_{ij}\delta(\mathbf{r} - \mathbf{r}')\delta(\omega - \omega'), \tag{A5}$$

with $\Theta(\omega,T)$ the well-known Planck energy of the quantum oscillator. If the source is magnetic, then the magnetization $\mathbf{M}(\mathbf{r},\omega)$ and/or magnetic currents $\mathbf{j}_m(\mathbf{r},\omega)$ are necessary, and similarly introduced, to describe the radiated field.

In contrast, when the source is secondary, namely, the emitted field is due to scattering by, e.g., a random medium or rough surface with inhomogeneities in its dielectric and/or magnetic susceptibility $\chi(\mathbf{r},\omega)$ and/or $\eta(\mathbf{r},\omega)$, respectively, then the source is characterized by these latter constitutive parameters that act as scattering potentials. For example, if the medium is nonmagnetic, the integrand in Eq. (A1) should be replaced by [50,68] $\mathbf{P}(\mathbf{r},\omega) = \chi(\mathbf{r},\omega)\mathbf{E}(\mathbf{r},\omega)$. Then, based on Eqs. (A1)–(A3), calculations of the emitted field are very similar in either case, whether the fluctuating source is primary or secondary.

It is known [44,69,70] that if $\mathbf{P}(\mathbf{r},\omega)$ and $\mathbf{j}(\mathbf{r},\omega)$ are statistically homogeneous and isotropic and each of them has a correlation length $\mathcal{L}$, when the distance z to the source surface $z = 0$ holds, $z \ll \mathcal{L} \ll \lambda$, then the electric cross-spectral density tensor elements in the source plane $z = 0$ given by the correlations $\mathcal{E}^{(0)}_{ij}(\boldsymbol{\rho}_1,\boldsymbol{\rho}_2,\omega) = \langle E^{(0)*}_i(\boldsymbol{\rho}_1)E^{(0)}_j(\boldsymbol{\rho}_2)\rangle = \mathcal{E}^{(0)}_{ij}(\boldsymbol{\rho},\omega)$ ($\boldsymbol{\rho} = \boldsymbol{\rho}_2 - \boldsymbol{\rho}_1$) (see Sec. II A) have a correlation length $\sigma \simeq \mathcal{L}$.

In contrast, when $\mathcal{L} \ll z \ll \lambda$, one has $\sigma \simeq z$ and the amplitude of the correlation functions at such distances z, $\mathcal{E}_{ij}(\boldsymbol{\rho},z \geqslant 0;\omega)$, is of the order of $1/z$ [44,71]. Hence, when $\mathbf{P}(\mathbf{r},\omega)$ and $\mathbf{j}(\mathbf{r},\omega)$ are both δ correlated so that $\mathcal{L} \to 0$, then so is the emitted field at $z = 0$, $\mathbf{E}^{(0)}(\boldsymbol{r},\omega)$, and therefore $\sigma \to 0$. We base on these facts the use made in Sec. II A of this limiting value of the wave field as the quantity that characterizes the fluctuating source at its exit plane $z = 0$, even in the case of δ-correlation currents and/or polarizabilities.

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