Relation between spatial confinement of light and optical tunneling

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With the aim of obtaining a better insight in the optical tunneling process the near-field electrodynamics of a current-density (equivalently polarization) sheet is investigated, taking as a starting point the near-field optics of a single atom, and afterwards the tunneling field of a macroscopic medium is determined integrating over a distribution of sheets. The total electric field hitherto used to study tunneling times and effective tunneling velocities is divided into a nonretarded (matter attached) longitudinal part of standing-wave character, and a retarded (detached) transverse part propagating away from the matter-vacuum interface with the vacuum speed of light. For a current-density distribution phase shifted with a wave number $q_\parallel$ along the interface, the transverse part is nonzero in the vacuum and decays exponentially with a decay constant $q_\parallel^{-1}$ as a function of the distance from the interface. Since the source domain of photons is precisely the domain of the transverse current density, the optical tunneling process attains an important contribution associated with the lack of spatial localizability of a photon in the evanescent regime. It is shown that in an observationally equivalent electromagnetic propagator description of the space-time dynamics, where the source domain of the photons is identified with the domain of the total electron current density, the retarded transverse dynamics necessarily must include spacelike couplings in the evanescent regime. Since these are destroyed with the vacuum speed of light as the light-cone coupling moves away from the matter-vacuum interface, the Einstein causality is always obeyed. The link to previous studies of the optical tunneling process is established by investigating the transverse and longitudinal dynamics in the frequency domain. Finally, it is shown that surface currents may play an important role in the optical tunneling process, in particular in cases where the incident electromagnetic field generates divergence-free currents in the bulk of the source medium. [S1050-2947(99)08408-5]

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I. INTRODUCTION

A paradigm of optical tunneling appears in the case of frustrated total internal reflection (FTIR); see, e.g., Refs. [1,2]. Thus, if a plane monochromatic electromagnetic wave is incident on a planar glass-vacuum interface at an angle larger than the critical angle, an electric field decaying exponentially with the distance from the interface is generated in the vacuum. Alhough, traveling along the interface, the field in the vacuum has standing-wave-like character in the direction perpendicular to the interface. This standing-wave form at a first glance seems to indicate that no energy is transported away from the boundary, and yet, if another planar glass medium is placed in the evanescent tail a traveling electromagnetic wave appears in this medium. In the glass-vacuum-glass system, the field now consists of a superposition of incident and reflected traveling waves in the first glass region, a sum of two evanescent modes decaying in opposite directions in the vacuum gap, and a single wave traveling away from the interface in the second glass medium. The fact that this electric mode pattern along the direction perpendicular to the interfaces mathematically is exactly the same as the wave-function mode pattern obtained when solving the stationary-state Schrödinger-equation problem for an electron incident upon a one-dimensional square potential barrier, a formal analogy between electron (massive particle) and photon tunneling emerges [1]. If instead of a vacuum gap, one has an air gap between the glass media, a small traveling-wave component is introduced in the direction perpendicular to the interface and the formal analogy with electron tunneling becomes less obvious. Optical ''tunneling'' in a dielectric-air-dielectric system (excited at an angle of incidence larger than the critical angle) resembles optical ''tunneling'' in a dielectric-metal-dielectric system (or a system consisting of a thin metallic film suspended in vacuum) when the monochromatic optical field has a frequency below the (bulk) plasma frequency of the metal [3]. Only if relaxation mechanisms for the electron system are neglected is the electromagnetic field in the metal purely evanescent. An optical tunneling-barrier problem also arises if one considers the ''evanescent-wave propagation'' of electromagnetic fields inside a one-dimensional (1D) photonic band-gap material, since here a mathematical analogy to the evanescent propagation of electrons in a Kronig-Penney periodic potential appears [4,5]. In microwave studies of the electromagnetic tunneling process one often makes use of a (rectangular) undersized waveguide [6,7] or a waveguide filled with a dielectric material and interrupted by an air gap that serves as the barrier [8]. If the frequency of the monochromatic wave progressing along the waveguide is chosen to be above the cutoff frequency in the dielectric-filled sections but below cutoff in the air gap, a mathematical analogy to the stationary-state electron tunneling through a one-dimensional rectangular barrier again emerges.

Although the two-dimensional optical tunneling in a dielectric-vacuum-dielectric system in the stationary-state situation resembles that of one-dimensional electron tunneling across a rectangular barrier (of the same width as the optical gap), in particular for *s*-polarized light where the (textbook) matching conditions for the electric field are the same as for the electron wave function (viz., continuity of the field and its first-order spatial derivative in the direction perpendicular to the interfaces), the state of things is more complicated in the dynamical situation where the transient behavior comes into play. To achieve a better understanding of the physics of the optical tunneling process and the relation of this process to tunneling of massive particles, it is necessary to deal with the process not only in the frequency domain but also, and primarily, in the time domain. To address *the two important issues in (optical) tunneling, viz., the tunneling-time question, and the ''effective'' tunneling-velocity problem*, it is necessary to pay full attention to the dynamics in both space and time. To investigate the question of the tunneling velocity, one traditionally studies the reflection and transmission of (short) electromagnetic pulses incident on the barrier, and in the theoretical treatment one must distinguish carefully between different kinds of velocities, such as the phase velocity, the group velocity, the energy velocity, the signal velocity, and the front velocity [9]. Though it has been claimed from time to time that Einstein causality can be violated in optical tunneling processes, no rigorous theoretical analyses exist that contradict the Einstein causality in relation to the velocity with which information can be transferred; see, e.g. [10], and references therein. From a perspective that is somewhat different from the ones hitherto adopted, the analysis of this paper also ends up with the conclusion that the speed of information transfer never exceeds the vacuum speed of light, even in a tunneling process. Studies of so-called superluminality in relation to propagation of laser pulses with group velocities greater than the vacuum velocity of light is of interest in their own right, and superluminality has for instance been investigated in the context of off-resonance pulse propagation through a medium with inverted atomic populations [11–13]. Using a collection of (classical) Lorentz oscillators as a model for two-level atom population inversion seems to make it possible to obtain tachyonlike propagation [13]. On a tachyonlike branch of a dispersion relation, the group velocity is always larger than the vacuum velocity of light, and a tachyonlike excitation possesses an *effective* mass that is imaginary. A tachyon dispersion-relation branch also is present for surface electromagnetic waves (surface polaritons) propagating on a BCS-paired superconductor surface if the frequencies lie slightly above the superconducting gap frequency [14,15]. Readers interested in a detailed overview of the entire field of optical tunneling and its relation to massive-particle tunneling may consult the recent review articles by Chiao and Steinberg [10] and Nimtz and Heitman [16].

In the present paper I shall seek to demonstrate for the reader that there is an important link between our ability to confine electromagnetic fields in space and the optical tunneling process, and that additional insight in the physics underlying the tunneling phenomenon may appear from such an understanding. In theoretical treatments of optical tunneling, evanescent waves necessarily always appear, and since numerous investigations in near-field optics have emphasized that a close formal relation exists between near fields prevailing around mesoscopic particlelike objects and evanescent fields [17–20], I have found it intriguing to investigate whether near-field electrodynamics already might contain the ingredients needed to obtain a better physical understanding of the optical tunneling process.

To set the scene, in Sec. II, we briefly review a recently established electromagnetic propagator theory dealing with the spatial confinement of light emitted from a single atom [21], paying particular attention to those aspects that appear to be of most relevance to the tunneling problem. Although the above-mentioned theory is formulated on the basis of rigorous (nonrelativistic) quantum electrodynamics (QED), and thus gives a field-quantized description of the near-field electrodynamics of the atom, it is sufficient here to consider the electromagnetic field as a c-number quantity, because, once the classical problem of optical tunneling is formulated correctly, the quantization in a gauge that is particularly adequate for a propagator description is not so difficult. Much of the classical near-field electrodynamics of an atom, but not all of it, of course, may be understood starting from a rigorous description of the attached and radiated near fields of an electric point dipole, as I have shown not long ago [22]. ''Seen,'' so to speak, with the eyes of the photon, the source domain of light quanta is to be identified with the region occupied by the transverse part of the current density induced in the atom by the prevailing local field. Since the transverse part of the induced atomic current density in general spreads over a region of spatial extension much larger than that of the electron orbitals, in the quantum-statistical sense photons emitted in an atomic decay process are generated (born) in the entire domain of the transverse current density. In the domain of the longitudinal part of the atomic current density (which has the same extension as the transverse part) an attached longitudinal electric self-field evolves, which in the QED description can be eliminated in favor of the particle-position variables. If one wants to consider, with what one may call an electron eye, the region in space where the total atomic current density is different from zero, i.e., roughly speaking the region occupied by the electrons of the atom, as the source region for the emitted field, the price one has to pay is the acceptance of the occurrence of a transverse self-field in the transverse current-density domain, and a retarded transverse field containing a spacelike part in the near-field zone of the atom. In the far field only couplings on the light cone remain. In passing I stress that the photon- and electron-eye views though picturally different, lead in every respect to the same physical *observations*. This is so because a change in the standard Lagrangian of the Coulomb gauge (or a unitary transformation) to a physically equivalent new Lagrangian that is adequate for a transformation from the photon- to the electron-eye picture can be found [21]. In the propagator QED description only the retarded part of the transverse field is quantized and hence only this part leads to photons. Furthermore, the related photon propagator explicitly demonstrates that this massless particle necessarily moves with the vacuum speed of light. In the description of photon propagation and localization in real space-time, the retarded part of the transverse field multiplied by the square root of the vacuum permittivity is precisely the Riemann-Silberstein wave function [23,24], which describes the so-called energy wave function of the photon [24,25] and which gives rise to the ''clicks of a detector'' when single photons are registered [26,27].

Equipped with the understanding of the atomic near-field electrodynamics, in Sec. IV and the remaining part of the paper, we turn our attention towards the optical tunneling problem, and we begin with an analysis of the near-field electrodynamics of a current-density sheet. *Although the electric field outside the sheet is divergence free it cannot be identified entirely with the photon field, because it also contains a part that is rotational free.* Theoretically, a vector field can be classified as transverse (longitudinal) only if it is divergence free (rotational-free) in its entire domain of definition. To classify the electromagnetic field correctly in the vacuum outside the sheet, the proper split must also be done in the plane of the sheet, and though the field diverges here (for an infinitesimally thin sheet), a rigorous division of the field into its transverse and longitudinal components can be made. The same kind of problem arises in models where a particle current-density distribution is confined to a point, the rest of space being vacuum [22]. Though the electric field certainly is divergence free in the entire vacuum domain, in this case it is not a transverse vector field because one point, namely the point where the particle is located and where the charge density thus is not zero, is missing. A rotational-free component also is present in the near-field zone of the point particle, and only if a proper split of the total current density into its transverse and longitudinal parts is made can the retarded part of the vacuum field (photon field) be determined. If the sheet current density is phase shifted along a given direction, as characterized by a wave number $q_{\parallel}$, the attached part of the electric field, which is different from zero also outside the sheet, decays exponentially with the distance from the sheet and the decay length is $q_{\parallel}^{-1}$. The transverse part of the field has a spacelike part with the same exponential decay length. Altogether this means that photons emitted from the sheet are only exponentially localizable (with a localization length $q_{\parallel}^{-1}$). After a propagator analysis of the retarded field in the space-time domain, we briefly study the field obtained in the case where the sheet current density is monochromatic.

In Sec. V, the optical tunneling from macroscopic current density distributions is investigated by composing them from sheets, and we are led to the conclusion that for a glass-vacuum-glass system a substantial part of the tunneling phenomenon must be related to the lack of photon localizability as long as the vacuum gap has a width not substantially larger than $q_{\parallel}^{-1}$. In the phenomenological approach where no distinction is made between transverse and longitudinal fields an exponential decay length equal to $[q_{\parallel}^2-(\omega/c_0)^2]^{-1/2}$ enters for monochromatic waves (of frequency ω) and only for them. In the expression above, c_0 denotes the vacuum speed of light. Though the total field *in the monochromatic situation* is of standing-wave character, a retarded field *propagating* always with c_0 *is* present and on top of this a significant photon delocalizability effect plays an important role.

The lack of photon localizability in an interesting fashion allows one to introduce a ‘‘tunneling time’’ with no need for a related tunneling velocity. From a quantum-statistical point of view this photon tunneling time with no associated velocity is related to the spacelike part of the transverse electromagnetic field, since this part basically gives the probability that a given photon is created at a certain distance from the sheet. The tunneling time in an essential (inevitable) manner thus is linked to the sensitivity of the detector. As in the atomic case, the transfer of information is linked to the discontinuous trailing edge of the spacelike coupling, and this edge always propagates with the speed of light.

In Sec. VI, we investigate the particular role of surface currents within the framework of a macroscopic sharp-boundary refractive-index model for the medium under study, and in Sec. VII tasks lying immediately a head of us following the present line of reasoning are outlined.

II. NEAR-FIELD ELECTRODYNAMICS OF A SINGLE ATOM: LIGHT-CONE AND SPACELIKE COUPLINGS

We begin our study of the relation between optical tunneling and spatial localization of light by reviewing certain aspects of the near-field electrodynamics of a single atom. This is done because electromagnetic couplings over *finite* distances in the near-field zone of the atom can be pictured as consisting of two parts, viz., a part *propagating* through space with the vacuum speed of light, and a part *evolving in time* in a standing-wave-like fashion. When an immense number of atoms are joined together so as to form a condensed-matter medium, the first part, related to the photon concept, tells us that Einstein causality is not broken in optical tunneling processes, though the photon field has a spacelike electric-field component, and the second part allows us to relate to a time scale for tunneling that does not involve time-space propagation effects, but appears because of fundamental quantum electrodynamic limitations on the possible degree of spatial confinement of matter-attached electromagnetic fields. From the above-mentioned point of view ‘‘superluminality’’ always stems from the interference of these two parts. Outside the near-field zone only couplings on the light cone remain. In macroscopic descriptions of optical tunneling evanescent fields play a prominent role, and since near-field and evanescent-mode electrodynamics are closely related, even in microscopic approaches, only by examining the near-field zone of the atom may physical ingredients needed for a better understanding of the tunneling process be found.

Let us now assume that the electrodynamics of the atom is driven by a prescribed external (ext) electric field $\vec{E}_T^{ext}(\vec{r},t)$, which is transverse (strictly speaking, divergence free), as indicated by the subscript T. The impressed field induces a current-density distribution $\vec{J}(\vec{r}\,',t')$ in the atom, which in turn is the source for a transverse and retarded (R) electromagnetic field $\vec{E}_T^R(\vec{r},t)$ that is emitted from the atom. To describe the propagation characteristics of the electromagnetic coupling between a space-time source point located at $(\vec{r}\,',t')$ and a field observation point $(\vec{r},t)$, it is *eo ipso* necessary to use an electromagnetic propagator formalism (picture). Explicitly stated, the retarded electric field is linked to the time derivative of the atomic current density via an integral relation of the form

$$\vec{E}_T^R(\vec{r},t)=\mu_0\int_{-\infty}^{\infty}\overleftrightarrow{D}_0^R(\vec{r}-\vec{r}\,',t-t')\cdot\frac{\partial\vec{J}(\vec{r}\,',t')}{\partial t'}d^3r'dt', \tag{1}$$

where, with $\vec{r}-\vec{r}\,'=\vec{R}$ and $t-t'=\tau$,

$$\overleftrightarrow{D}_0^R(\vec{R},\tau) = -\frac{1}{4\pi R}\delta\left(\frac{R}{c_0}-\tau\right)(\overleftrightarrow{U}-\vec{e}_R\vec{e}_R) + \frac{c_0^2\tau}{4\pi R^3}\Theta(\tau)\Theta\left(\frac{R}{c_0}-\tau\right)(\overleftrightarrow{U}-3\vec{e}_R\vec{e}_R) \quad (2)$$

is the transverse and retarded (photon) propagator. Before commenting on Eqs. (1) and (2), let me emphasize that a relation of the form appearing in Eq. (1) can be derived using rigorous quantum electrodynamics (QED) [21] or starting within the framework of semiclassical electrodynamics (SCED) [22]. In the QED description, Eq. (1) is a relation between the retarded part of the transverse field operator and the atomic current-density operator (in second quantization). Within the framework of the propagator picture the photon concept may be established from a Hamiltonian description where precisely (only) $\vec{E}_T^R$ is subjected to a canonical quantization [21]. In a nonrelativistic SCED theory, where the electromagnetic field is a *c*-number quantity, $\vec{J}(\vec{r}\,',t')$ is calculated via the Schrödinger equation. Before proceeding let me stress that the induced atomic current density in general need to be calculated self-consistently [29]. Apart from a few remarks given below, I shall only return to this problem when discussing the optical tunneling in condensed matter media; see Secs. V and VI.

The photon propagator given in Eq. (2) consists of a far-field part (proportional to R^{-1}) and a near-field part ($\sim R^{-3}$). The far-field part is different from zero (and singular) only on the retarded light cone $|\vec{r}-\vec{r}\,'| = c_0(t-t')$, c_0 being the speed of light in vacuum, as it is evident from the appearance of the Dirac delta function $\delta(R/c_0-\tau)$, and polarized perpendicular to the local direction of propagation, as one readily realizes from the tensor $\overleftrightarrow{U}-\vec{e}_R\vec{e}_R$ ($\overleftrightarrow{U}$: unit tensor, $\vec{e}_R=\vec{R}/R$). The Heaviside unit step functions $\Theta(\tau)$ and $\Theta(R/c_0-\tau)$ appearing in the near-field part of $\overleftrightarrow{D}_0^R(\vec{R},\tau)$ show that the near-field coupling is causal ($t>t'$) and different from zero only for spacelike events, i.e., those for which $|\vec{r}-\vec{r}\,'|>c_0(t-t')$. The spacelike form of the near-field coupling ensures that the unpleasant R^{-3}-singularity, which in this case would make the space integral in Eq. (1) conditionally convergent, does not appear, a physically satisfactory feature. Although the time delay (τ) and source-observation distance (R) do not appear in the form $R/c_0-\tau$ in the factor in front of the step functions, the near-field coupling is destroyed with the vacuum speed of light as the far-field light-cone coupling sweeps the near-field region. As we shall realize when discussing optical tunneling, the electromagnetic energy flows with the vacuum speed of light. I shall demonstrate below that the photon field that appears in the observation point prior to the arrival of the light-cone pulse can be considered as stemming from a fundamental inability to localize photons precisely in space, but before doing so, we shall briefly consider the matter-attached part of the electromagnetic (near) field.

From the vector field $\vec{J}(\vec{r},t)$, giving the induced atomic current density, we now project out the longitudinal (properly speaking, rotational-free) part $\vec{J}_L(\vec{r},t)$ by means of the tensorial longitudinal delta function $\overleftrightarrow{\delta}_L(\vec{r}-\vec{r}\,')$, i.e.,

$$\vec{J}_L(\vec{r},t) = \int_{-\infty}^{\infty} \overleftrightarrow{\delta}_L(\vec{r}-\vec{r}\,')\cdot\vec{J}(\vec{r}\,',t)d^3r'. \quad (3)$$

The relation between $\vec{J}$ and $\vec{J}_L$ is *nonlocal in space* but *local in time*, since the projection operator ($\overleftrightarrow{\delta}_L$) works on the spatial current-density distribution at a fixed time. Upon a split of the (operator) Maxwell equations into sets describing the longitudinal and transverse electrodynamics, respectively, one finds from the longitudinal set that the induced atom-attached longitudinal electric field $\vec{E}_L(\vec{r},t)$ is given by

$$\vec{E}_L(\vec{r},t) = -\frac{1}{\epsilon_0}\int_{-\infty}^{t}\vec{J}_L(\vec{r},t')dt', \quad (4)$$

provided that the induced longitudinal field vanishes in the remote past. As we shall realize in Sec. IV B, the last condition is (weakly) linked to the principle of causality. Since the longitudinal delta function is different from zero in (and only in) the near-field zone of the atom the *nonpropagating* attached field extends over the entire near-field region. In a QED description formulated in the Coulomb gauge, the longitudinal part of the electric-field operator is not a dynamical variable, since it can be eliminated in favor of the dynamical particle-position variables of the atom (redundancy).

In the propagator picture of the near-field electrodynamics of the atom a transversely polarized attached field component is also present [21]. This field, $\vec{E}_T^{\mathrm{SF}}(\vec{r},t)$, named the transverse atomic self-field (SF) is linked to the transverse part

$$\vec{J}_T(\vec{r},t) = \int_{-\infty}^{\infty} \overleftrightarrow{\delta}_T(\vec{r}-\vec{r}\,')\cdot\vec{J}(\vec{r}\,',t)d^3r' \quad (5)$$

of the atomic current density by means of the relation

$$\vec{E}_T^{SF}(\vec{r},t) = -\frac{1}{3\epsilon_0}\int_{-\infty}^{t}\vec{J}_T(\vec{r},t')dt'. \quad (6)$$

Since the transverse delta function $\overleftrightarrow{\delta}_T(\vec{r}-\vec{r}\,')$ added to the longitudinal delta function equals the usual Dirac delta function $\delta(\vec{r}-\vec{r}\,')$ times the unit tensor $\overleftrightarrow{U}$, i.e., $\overleftrightarrow{\delta}_T(\vec{r}-\vec{r}\,')+\overleftrightarrow{\delta}_L(\vec{r}-\vec{r}\,') = \overleftrightarrow{U}\delta(\vec{r}-\vec{r}\,')$, the transverse self-field extends over the same spatial domain as the longitudinal self-field (attached field). The transverse self-field dynamics is nonlocal in space but local in time, and the factor 3^{-1} appearing in Eq. (6) shows that the transverse and longitudinal self-fields do not cancel outside the spatial region where the atomic current density itself is different from zero. In the QED description [21] a change of the Coulomb Lagrangian into a new one closely related (but not identical) to the Power-Zinau-Woolley Lagrangian [30,31] allows one to transfer the transverse self-field to the particle Hamiltonian. The retarded part of the transverse-field operator can now be subjected to the canonical quantization procedure leading to the (spherical) photon concept. The length scale for the abovementioned self-field phenomena is readily obtained from the spherical representation of the transverse and longitudinal δ functions, viz.,

$$\overleftrightarrow{\delta}_T(\vec{R}) = \overleftrightarrow{U}\delta(\vec{R}) - \overleftrightarrow{\delta}_L(\vec{R}) = \frac{2}{3}\delta(\vec{R})\overleftrightarrow{U} - \frac{1}{4\pi R^3}(\overleftrightarrow{U} - 3\vec{e}_R\vec{e}_R). \tag{7}$$

Now it is seen that the self-field effects are present precisely in the near-field zone of the atom, and in the same zone (and only here) also the spacelike retarded coupling is present; cf. Eq. (2).

Altogether, we have thus realized that around a single (spinless) atom, excited by an externally impressed and prescribed transverse field $\vec{E}_T^{ext}(\vec{r},t)$, a total field

$$\vec{E}(\vec{r},t) = \vec{E}_T^{ext}(\vec{r},t) + \vec{E}_L(\vec{r},t) + \vec{E}_T^{SF}(\vec{r},t) + \vec{E}_T^R(\vec{r},t) \tag{8}$$

emerges with space and time properties given by Eqs. (1), (2), (4), and (6).

At this stage it is instructive to look at the total transverse field from a different point of view, named the photon-eye view [21]. In Eqs. (6) [combined with Eq. (5)] and (1), the transverse field is looked upon as being driven by the total induced current-density distribution of the atom, a view one may call the electron-eye view [21]. If one eliminates the current density itself, $\vec{J}(\vec{r}\,',t')$, and its first-order time derivative $\partial\vec{J}(\vec{r}\,',t')/\partial t'$ in favor of the (time-derivative of) the transverse current density, one obtains the integral relation [21]

$$\vec{E}_T^{SF}(\vec{r},t) + \vec{E}_T^R(\vec{r},t) = \mu_0 \int_{-\infty}^{\infty} \overleftrightarrow{d}^R(\vec{r}-\vec{r}\,',t-t') \cdot \frac{\partial \vec{J}_T(\vec{r}\,',t')}{\partial t'} d^3r'\,dt', \tag{9}$$

where now only an (the) isotropic electromagnetic propagator $\overleftrightarrow{d}^R(\vec{R},\tau)$ with a far-field-like distance dependence appears, i.e.,

$$\overleftrightarrow{d}^R(\vec{R},\tau) = -\frac{1}{4\pi R}\,\delta\!\left(\frac{R}{c_0} - \tau\right)\overleftrightarrow{U}. \tag{10}$$

The propagation characteristics of the transverse field in the photon-eye view indeed is simple; only retarded couplings on the light cone $R = c_0\tau$ are present, the spreading of the field is isotropic, and the coupling exhibits an R^{-1}-dependence only. The price we have paid to achieve such a simple picture is that it has become necessary to consider the $\vec{J}_T$ domain as the source domain, instead of the (much better) localized $\vec{J}$ domain. I have named the aforementioned view, the photon-eye view, because photons (introduced via the transverse field operator dynamics) apparently in the statistical sense can be no better localized in space than given by the extension of the $\vec{J}_T$ domain in the initial stage of emission ''from'' the atom. In terminating this section I emphasize that the photon-eye view and the electron-eye view lead to equivalent (in every respect) observable effects in all situations, but the manner in which we in a propagator description *picture* the near-field electrodynamics is different.

III. DEFINITION OF THE CONCEPT OPTICAL (PHOTON) TUNNELING

Based on the considerations put forth in the preceding section I now *define* optical (photon) tunneling phenomena as those phenomena that occur only in the presence of longitudinal and transverse self-field effects and spacelike retarded couplings. Seen with the ''eyes of the photon,'' optical tunneling thus is equivalent to the presence of two fundamental aspects of QED, namely, (i) spatial localization of transverse photons in the quantum-statistical sense, and (ii) longitudinal self-field interactions mediated in a relativistically invariant (Lorenz gauge) description by longitudinal and scalar photons. The definition above relates at first sight to single-atom electrodynamics, because it appeared from a single-atom analysis, but, and this may be the important aspect for traditional studies of optical tunneling, the definition when carried over (applied) to many-particle systems (condensed matter, *par excellence*) seems to allow one to obtain better insight in the physics of tunneling in these systems. This is what I shall try to argue in the main body of this paper. To my understanding, photon tunneling thus is a ''naturally'' occurring process already in the near-field electrodynamics of a single atom. By extrapolation from this view it is tempting for me to claim that optical tunneling also is an indispensible phenomenon in near-field diffraction from small holes, slits, etc.

IV. NEAR-FIELD ELECTRODYNAMICS OF A CURRENT-DENSITY SHEET

The principles governing the near-field electrodynamics of single atoms we now extend to a study of the spatial localization and radiation of light from a current-density sheet. In the course of the analysis, information is obtained on aspects of the electrodynamic tunneling process, which later on (in Secs. V and VI) allow us to describe the physics underlying the optical tunneling across a vacuum gap separating macroscopic media.

A. Retarded and nonretarded electromagnetic propagators

Let us consider a current-density distribution of the form

$$\vec{J}(\vec{r}\,',t') = \vec{J}_0(t')e^{iq_\parallel x'}\delta(z'-z_0), \tag{11}$$

in a Cartesian xyz-coordinate system. The current density in Eq. (11) is translationally invariant in the y direction, and, as indicated by the Dirac delta function $\delta(z'-z_0)$, concentrated on a sheet located at the plane $z'=z_0$. Along the x direction the current density exhibits a simple phase shift given by the wave number $q_\parallel$. The ansatz in Eq. (11) implies that the total electric field in the surroundings of the sheet takes the generic form

$$\vec{E}(\vec{r},t) = \vec{E}(z,t;q_\parallel\vec{e}_x)e^{iq_\parallel x}, \tag{12}$$

emphasizing that the amplitude $\vec{E}(z,t;q_\parallel\vec{e}_x)$ depends in a parametric fashion on the wave vector $q_\parallel\vec{e}_x$, $\vec{e}_x$ being a unit vector along the direction of the x axis. Unit vectors along the two other Cartesian axes appear below and will be denoted by $\vec{e}_y$ and $\vec{e}_z$. In the frequency (ω) domain the amplitudes $\vec{E}(z;q_\parallel\vec{e}_x,\omega)$ and $\vec{J}_0(\omega)$ are related via

$$\vec{E}(z;q_\parallel\vec{e}_x,\omega)=-i\mu_0\omega\overleftrightarrow{D}_0(z-z_0;q_\parallel\vec{e}_x,\omega)\cdot\vec{J}_0(\omega) \tag{13}$$

in a propagator description. In explicit form, the vacuum Green function $\overleftrightarrow{D}_0(z-z_0;q_\parallel\vec{e}_x,\omega)$ is given by the dyadic formula

$$\overleftrightarrow{D}_0(z-z_0;q_\parallel\vec{e}_x,\omega)=\frac{e^{i\kappa_\perp^0|z-z_0|}}{2i\kappa_\perp^0 q_0^2}[(\kappa_\perp^0)^2\vec{e}_x\vec{e}_x+q_\parallel^2\vec{e}_z\vec{e}_z+q_0^2\vec{e}_y\vec{e}_y-\kappa_\perp^0 q_\parallel\,\mathrm{sgn}(z-z_0)\times(\vec{e}_x\vec{e}_z+\vec{e}_z\vec{e}_x)], \tag{14}$$

where $\kappa_\perp^0=(q_0^2-q_\parallel^2)^{1/2}$, $q_0=\omega/c_0$ being the vacuum wave number of light, and $\mathrm{sgn}(z-z_0)=+1$ for $z>z_0$ and -1 for $z<z_0$. The singular nature of the current-density ansatz in Eq. (11) implies that a contact term giving the electric field in the source plane ($z=z_0$) has to be omitted from the Green function. In Sec. V, where nonsingular current-density distributions are considered, the role of the contact term will be discussed. In the single-particle case described in Sec. II the contact term would also have to be omitted if one from the outset had assumed that the electron confinement were complete (pointlike atom). For the atom electrodynamics the contact interaction enters via the singular contributions $\frac{2}{3}\delta(\vec{R})\overleftrightarrow{U}$ and $\frac{1}{3}\delta(\vec{R})\overleftrightarrow{U}$ to the transverse and longitudinal δ functions; see Eq. (7).

The Green function in Eq. (14) contains a nonradiative (NR) part, denoted by $\overleftrightarrow{D}_0^{NR}(z-z_0;q_\parallel\vec{e}_x,\omega)$. The explicit expression for this part is readily obtained by letting c_0 approach infinity, and remembering that $\mu_0c_0^2=\epsilon_0^{-1}$. Since $\kappa_\perp^0\to+iq_\parallel$, necessarily, for $c_0\to\infty$ one finds

$$\overleftrightarrow{D}_0^{NR}(z-z_0;q_\parallel\vec{e}_x,\omega)=c_0^2\lim_{c_0\to\infty}\{c_0^{-2}\overleftrightarrow{D}_0(z-z_0;q_\parallel\vec{e}_x,\omega)\}=\frac{q_\parallel}{2q_0^2}e^{-q_\parallel|z-z_0|}[\vec{e}_x\vec{e}_x-\vec{e}_z\vec{e}_z+i(\vec{e}_x\vec{e}_z+\vec{e}_z\vec{e}_x)\mathrm{sgn}(z-z_0)]. \tag{15}$$

The difference between the two Green functions appearing in Eqs. (14) and (15), denoted by $\overleftrightarrow{D}_0^{T}(z-z_0;q_\parallel\vec{e}_x,\omega)$, i.e.,

$$\overleftrightarrow{D}_0^{T}(z-z_0;q_\parallel\vec{e}_x,\omega)=\overleftrightarrow{D}_0(z-z_0;q_\parallel\vec{e}_x,\omega)-\overleftrightarrow{D}_0^{NR}(z-z_0;q_\parallel\vec{e}_x,\omega), \tag{16}$$

is the transverse (T) propagator responsible for retarded (R) (with speed c_0) interactions between the source plane at z_0 and a plane of observation located at z.

In the frequency domain, the transverse part of the field amplitude in Eq. (13) thus is given by

$$\vec{E}_T^R(z;q_\parallel\vec{e}_x,\omega)=-i\mu_0\omega\overleftrightarrow{D}_0^{T}(z-z_0;q_\parallel\vec{e}_x,\omega)\cdot\vec{J}_0(\omega). \tag{17}$$

The amplitude of the nonradiative electric field attached to the sheet, which is denoted by $\vec{E}_L^{NR}(z;q_\parallel\vec{e}_x,\omega)$, may be obtained from

$$\vec{E}_L^{NR}(z;q_\parallel\vec{e}_x,\omega)=-i\mu_0\omega\overleftrightarrow{D}_0^{NR}(z-z_0;q_\parallel\vec{e}_x,\omega)\cdot\vec{J}_0(\omega). \tag{18}$$

If the current-density distribution is spread over a finite interval in the z' direction so that it is no longer singular, a contact contribution is added to the field. This contribution plus $\vec{E}_L^{NR}$ gives the total longitudinal (nonradiative) field, *the curl of which is identically zero in the entire space*. The subscript L on $\vec{E}_L^{NR}$ in Eq. (18) is meant to indicate that this field plus the contact field is longitudinal for *all* z, including $z=z_0$. The contact field does not change the transverse dynamics (see Sec. V) and therefore $\vec{E}_T^R$ correctly describes the transverse dynamics even if the current-density distribution in Eq. (11) is smeared in the z' direction. The propagator $\overleftrightarrow{D}_0^T$ hence is the correct photon propagator for excitations that are invariant on the y axis and contain only one wave-number component ($q_\parallel$) in the x direction. In the subsequent two subsections we shall investigate the two parts of the entire field

$$\vec{E}(z,t;q_\parallel\vec{e}_x)=\vec{E}_T^R(z,t;q_\parallel\vec{e}_x)+\vec{E}_L^{NR}(z,t;q_\parallel\vec{e}_x) \tag{19}$$

in the space-time (z,t) domain.

B. Attached field and its spatial confinement

Starting from the microscopic Maxwell-Lorentz equations in the space-frequency domain it readily appears that the longitudinal parts of the electric field and current density are related via [28]

$$\vec{E}_L(\vec{r};\omega)=\frac{1}{i\epsilon_0\omega}\vec{J}_L(\vec{r};\omega). \tag{20}$$

Making use of the result

$$\frac{1}{2\pi}\int_{-\infty}^{\infty}e^{-i\omega\tau}\frac{d\omega}{\omega}=\frac{1}{2i}\mathrm{sgn}\tau, \tag{21}$$

Eq. (20) takes the form

$$\vec{E}_L(\vec{r},t)=-\frac{1}{2\epsilon_0}\int_{-\infty}^{\infty}\mathrm{sgn}(t-t')\vec{J}_L(\vec{r},t')dt' \tag{22}$$

in the space-time domain, or equivalently

$$\vec{E}_L(\vec{r},t)=-\frac{1}{2\epsilon_0}\int_{-\infty}^{t}\vec{J}_L(\vec{r},t')dt'+\frac{1}{2\epsilon_0}\int_{t}^{\infty}\vec{J}_L(\vec{r},t')dt'. \tag{23}$$

If one assumes that the induced longitudinal field vanishes in every space point in the remote past, i.e., for $t\to-\infty$, one has

$$\vec{E}_L(\vec{r},-\infty)=\frac{1}{2\epsilon_0}\int_{-\infty}^{\infty}\vec{J}_L(\vec{r},t')dt'=\vec{0}. \tag{24}$$

By means of Eq. (24), the last term on the right-hand side of Eq. (23) may now be eliminated, giving the result cited in Eq. (4). Under the assumption that $\vec{J}_L(\vec{r},t')$ is a *prescribed* source field, the *noncausal* term $(2\epsilon_0)^{-1}\int_t^\infty \vec{J}_L(\vec{r},t')dt'$ in Eq. (23) must vanish, and this is ensured by the condition $\vec{E}_L(\vec{r},-\infty)=\vec{0}$.

To determine the electric field attached to the current-density sheet, one transforms the quantity $-i\mu_0\omega\overleftrightarrow{D}_0^{NR}(z-z_0;q_\parallel\vec{e}_x,\omega)$ to the time domain. A glance at Eq. (15) shows that this essentially amounts to a calculation of the integral in Eq. (21). Hence, we obtain

$$\vec{E}_L^{NR}(z,t;q_\parallel\vec{e}_x)=-\frac{q_\parallel}{2\epsilon_0}e^{-q_\parallel|z-z_0|}\times\begin{pmatrix}1 & 0 & i\,\mathrm{sgn}(z-z_0)\\ 0 & 0 & 0\\ i\,\mathrm{sgn}(z-z_0) & 0 & -1\end{pmatrix}\cdot\int_{-\infty}^{t}\vec{J}_0(t')dt'. \tag{25}$$

It appears from Eq. (25) that the attached field has standing-wave character, as it must have, and that it decays exponentially with the distance from the plane of the sheet (placed at z_0) with a decay constant $q_\parallel$. The nonradiative field disappears for $q_\parallel=0$ and also in the limit $q_\parallel\to\infty$. If the current density is linearly polarized in the y direction the attached field is zero.

The result in Eq. (25) can be understood from a (slightly) different point of view, realizing that the longitudinal delta function $\overleftrightarrow{\delta}_L(z-z';q_\parallel\vec{e}_x)$ of relevance for the present vector-field problem is given by [28]

$$\overleftrightarrow{\delta}_L(z-z';q_\parallel\vec{e}_x)=\vec{e}_z\vec{e}_z\delta(z-z')+\frac{q_\parallel}{2}e^{-q_\parallel|z-z'|}[\vec{e}_x\vec{e}_x-\vec{e}_z\vec{e}_z+i(\vec{e}_x\vec{e}_z+\vec{e}_z\vec{e}_x)\mathrm{sgn}(z-z')]. \tag{26}$$

For the sheet problem the singular term $\vec{e}_z\vec{e}_z\delta(z-z')$ is of no relevance, but it plays a role for spatially extended current-density distributions, as we shall see in Sec. V. A combination of Eqs. (25) and (26) allows us to write

$$\vec{E}_L^{NR}(z,t;q_\parallel\vec{e}_x)=-\frac{1}{\epsilon_0}\overleftrightarrow{\delta}_L(z-z_0)\cdot\int_{-\infty}^{t}\vec{J}_0(t')dt',\quad z\neq z_0 \tag{27}$$

and this tells us that the $|z-z_0|$ range of the attached field is identical to the range of the integral of the longitudinal current density $\vec{J}_L(z-z_0)=\overleftrightarrow{\delta}_L(z-z_0)\cdot\vec{J}_0$. In a quantum-electrodynamic approach the attached field is eliminated as a dynamical variable in favor of the particle-position variables to remove redundancy. Since the total current density of the sheet *per* definition is different from zero only for $z=z_0$, the transverse part of the current density, characterizing the spatial localization range for the photon field, extends in the z direction in a manner given by $\exp(-q_\parallel|z-z_0|)$; see Fig. 1.

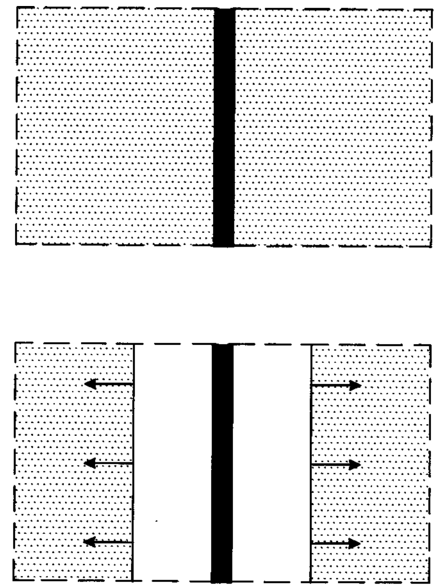

FIG. 1. Schematic illustration of the optical tunneling response of a current-density sheet (black strip). The response consists of an attached longitudinal part (top figure), different from zero only in the evanescent (shaded) region characterized by the decay length $q_\parallel^{-1}$, plus a detached retarded and transverse part (bottom figure) nonvanishing only for spacelike events belonging to the evanescent zone, as indicated by the shading. In the wake of the two light cones (shown with small arrows attached) moving away from the source plane with the vacuum speed of light, a Bessel-function-like coupling persists.

The tunneling regime for phenomena generated by a current-density sheet source thus is characterized by the length $q_\parallel^{-1}$.

C. Detached field: Space- and timelike couplings

It appears from Eq. (17) that the retarded and transverse electromagnetic field emitted from the current-density sheet is given by

$$\vec{E}_T^R(z,t;q_\parallel\vec{e}_x)=\mu_0\int_{-\infty}^{\infty}\overleftrightarrow{D}_0^T(z-z_0,t-t';q_\parallel\vec{e}_x)\cdot\frac{\partial\vec{J}_0(t')}{\partial t'}dt' \tag{28}$$

in the space-time representation. The propagation characteristics of the detached field are hidden in the photon propagator $\overleftrightarrow{D}_0^T(Z,\tau;q_\parallel\vec{e}_x)$, with $Z=z-z_0$ and $\tau=t-t'$, and to establish the explicit expression for this, it is convenient to start from the form the plane-wave expansion of the $\overleftrightarrow{D}_0^R(\vec{R},\tau)$-propagator in Eq. (2) takes in the space-frequency domain, viz. (see, e.g., [22,32])

$$\overleftrightarrow{D}_0^R(\vec{R};\omega)=(2\pi)^{-3}\int_{-\infty}^{\infty}\frac{\overleftrightarrow{U}-\vec{e}_{\vec{q}}\vec{e}_{\vec{q}}}{q_0^2-q^2}e^{i\vec{q}\cdot\vec{R}}d^3q, \tag{29}$$

where $\vec{e}_{\vec{q}}=\vec{q}/q$ is a unit vector in the q direction, and $q_0=\omega/c_0$ is the vacuum wave number for light of angular frequency ω. From Eq. (29) one readily obtains the integral expression

$$\overleftrightarrow{D}_0^T(Z;q_\parallel\vec{e}_x,\omega)=\frac{1}{2\pi}\int_{-\infty}^{\infty}\frac{\overleftrightarrow{U}-\vec{e}_{\vec{q}}\vec{e}_{\vec{q}}}{q_0^2-q^2}e^{iq_\perp Z}dq_\perp, \tag{30}$$

where now

$$\vec{e}_{\vec{q}}=\frac{1}{q}(q_\parallel\vec{e}_x+q_\perp\vec{e}_z), \tag{31}$$

with $q=(q_\parallel^2+q_\perp^2)^{1/2}$. In the (Z,τ) domain, the photon propagator thus may be represented by the integral form

$$\overleftrightarrow{D}_0^T(Z,\tau;q_\parallel\vec{e}_x)=(2\pi)^{-2}\int_{-\infty}^{\infty}\frac{\overleftrightarrow{U}-\vec{e}_{\vec{q}}\vec{e}_{\vec{q}}}{q_0^2-q^2}e^{i(q_\perp Z-\omega\tau)}d\omega\; dq_\perp. \tag{32}$$

To obey the principle of causality in its most general form one must have $\tau=t-t'>0$. The ω integration in Eq. (32) is readily carried out, giving

$$\frac{1}{2\pi}\int_{-\infty}^{\infty}\frac{e^{-i\omega\tau}d\omega}{\left(\frac{\omega}{c_0}\right)^2-q^2}=-\frac{c_0}{2q}\sin(qc_0\tau),\quad \tau>0. \tag{33}$$

By combining Eqs. (32) and (33) we therefore get

$$\overleftrightarrow{D}_0^T(Z,\tau;q_\parallel\vec{e}_x)=-\frac{c_0}{4\pi}\int_{-\infty}^{\infty}(\overleftrightarrow{U}-\vec{e}_{\vec{q}}\vec{e}_{\vec{q}})\frac{\sin(qc_0\tau)}{q}e^{iq_\perp Z}dq_\perp. \tag{34}$$

It is to some extent possible to carry out the integration in Eq. (34) [for details of particular interest the reader is referred to Appendix A], and the final result is as follows:

$$\begin{aligned}\overleftrightarrow{D}_0^T(Z,\tau;q_\parallel\vec{e}_x)=&\frac{1}{4}c_0^2\tau q_\parallel e^{-q_\parallel|Z|}\begin{pmatrix}1&0&i\,\mathrm{sgn}\,Z\\0&0&0\\i\,\mathrm{sgn}\,Z&0&-1\end{pmatrix}\\&\times\theta(\tau)\theta(|Z|-c_0\tau)\\&+\left[\begin{pmatrix}\pi_1&0&i\pi_2\,\mathrm{sgn}\,Z\\0&0&0\\i\pi_2\,\mathrm{sgn}\,Z&0&-\pi_1\end{pmatrix}\right.\\&-\frac{1}{4}c_0J_0[q_\parallel\sqrt{(c_0\tau)^2-Z^2}]\\&\left.\times\begin{pmatrix}1&0&-iq_\parallel Z\\0&1&0\\-iq_\parallel Z&0&0\end{pmatrix}\right]\theta(c_0\tau-|Z|),\end{aligned} \tag{35}$$

where

$$\pi_1=\frac{c_0q_\parallel^2}{2\pi}\int_0^{\infty}\frac{1}{(q_\parallel^2+q_\perp^2)^{3/2}}\sin(c_0\tau\sqrt{q_\parallel^2+q_\perp^2})\cos q_\perp Z dq_\perp, \tag{36}$$

and

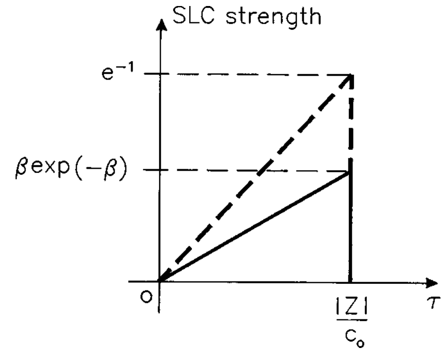

FIG. 2. The amplitude strength of the spacelike coupling (SLC), as given by the associated transverse propagator part (divided by $c_0/4$), as a function of the delay time (τ). At a given distance ($|Z|$) from the sheet plane, the coupling strength grows linearly in time until the light-cone coupling and its Bessel-function-like wake signal take over (full-line graph). At maximum the coupling strength is equal to $\beta\exp(-\beta)$ where $\beta=q_\parallel|Z|$. At a given distance from the sheet source, the largest coupling is obtained for a wave number $q_\parallel=|Z|^{-1}$, and equals e^{-1} (broken line).

$$\pi_2=\frac{c_0^2\tau q_\parallel}{2\pi}\int_0^{\infty}\frac{q_\perp}{q_\parallel^2+q_\perp^2}\cos(c_0\tau\sqrt{q_\parallel^2+q_\perp^2})\sin q_\perp Z\,dq_\perp. \tag{37}$$

It was mentioned in Sec. II that the retarded field emitted from a single atom in the near-field zone contains a spacelike contribution, the propagation characteristics of which are given by the second part of the propagator in Eq. (2). It appears from Eq. (35) that the propagation characteristics of the retarded field emerging from a current-density sheet also contain a spacelike part [the term with the Heaviside unit step function $\theta(|Z|-c_0\tau)$]. The spacelike coupling decays exponentially with the distance ($|Z|$) from the sheet plane with a decay constant $q_\parallel$, and, as expected, the nonretarded coupling and the retarded spacelike coupling hence have the same spatial range, as illustrated schematically in Fig. 1. As long as $|Z|>c_0\tau$, the coupling increases linearly in time (τ) at a fixed distance from the sheet (see Fig. 2), but although the trailing edge of the spacelike interaction moves outwards from the sheet with the vacuum speed of light, the Z and τ dependences enter in product form, viz., as $\tau\exp(-q_\parallel|Z|)$. The linear rise of the coupling in time together with a given detector sensitivity allow us to introduce a velocity-independent tunneling time in a concrete experiment. If the time-derivative of the sheet current density points in the y direction, i.e., perpendicular to the plane of field propagation, the spacelike coupling vanishes. The coupling also depends on $q_\parallel$, is zero for $q_\parallel=0$ and $q_\parallel\to\infty$, and for fixed $|Z|$ is largest for $q_\parallel=|Z|^{-1}$; cf. Fig. 2. The particular tensor form of the spacelike term in Eq. (35), which is the same as the one appearing in the relation between the nonretarded field $\vec{E}_L^{NR}$ and the source current-density amplitude $\vec{J}_0$ in Eq. (25), has an important physical meaning, as I shall show in Sec. V.

The timelike part of the transverse propagator, given by the factor in front of the step function $\theta(c_0\tau-|Z|)$ in Eq. (35), I have divided into two pieces. In the case where the time derivative of the current density is perpendicular to the plane of field propagation, i.e., in the y direction, only the last piece, proportional to the zeroth-order Bessel function

$J_0[q_\parallel\sqrt{(c_0\tau)^2-Z^2}]$, contributes. In passing we note the Lorentz invariance of the quantity $(c_0\tau)^2-Z^2$, when the field emerging from the sheet is observed in inertial systems moving relative to each other with (uniform) velocity $\vec{v}=v\vec{e}_z$ in the direction perpendicular to the plane of the sheet. For a time-varying current-density source distribution pointing in the y direction this kind of relativistic invariance might have been anticipated. In the single-atom case far-field couplings are present only on the retarded light cone, i.e., for $|\vec{r}-\vec{r}\,'|=c_0(t-t')$; cf. the presence of the Dirac delta function $\delta(R/c_0-\tau)$ in Eq. (2). In the sheet case the couplings exist also when $c_0(t-t')>|z-z'|$, and for a given $Z=z-z'$, these timelike couplings die out in the slow fashion dictated by the τ dependence of the zeroth-order Bessel function $J_0[q_\parallel\sqrt{(c_0\tau)^2-Z^2}]$. For time-varying source current densities confined to the plane of field propagation (the xz plane), off-diagonal elements (xz and zx) appear in the timelike part of the retarded propagator. Apart from a term proportional to $iq_\parallel Z$ times the zeroth-order Bessel function, these (identical) elements contain a contribution π_2 [Eq. (37)] given only in the integral form. Also a zz component $-\pi_1$ and an additional xx component π_1 now appear, the explicit integral expression for π_1 being given in Eq. (36). If $q_\parallel=0$, the retarded propagator takes the particularly simple form

$$\overleftrightarrow{D}_0^T(Z,\tau;\vec{0})=\frac{c_0}{4}(\vec{e}_z\vec{e}_z-\overleftrightarrow{U})\,\theta(c_0\tau-|Z|), \tag{38}$$

and no tunneling phenomena exist.

D. Monochromatic sheet current density

In Secs. IV B and IV C, the attached and detached fields were studied in the space-time domain, and the results obtained were valid for sheet current densities with arbitrary time dependence. Optical tunneling is, however, often discussed assuming the source dynamics to be monochromatic, and many types of experiments are carried out with (quasi-)monochromatic excitation. Also, on a more formal basis, the analogy between electron and (so-called) photon tunneling is investigated starting from the form the wave equation for the electromagnetic field takes for monochromatic waves, i.e., the Helmholtz equation [10,16]. Though a direct fingerprint of the physics hidden in the optical tunneling process is obtained only in the space-time domain, and in the framework of a propagator description, reminiscences of the optical tunneling phenomenon do appear also in the monochromatic case, as we shall realize below.

Let us now assume that the time dependence of the sheet current density is given by

$$\vec{J}_0(t')=\vec{J}_0(\omega)e^{-i\omega t'}e^{\epsilon t'}, \tag{39}$$

where $\epsilon=0^+$ is an infinitesimal but positive number needed to ensure that the excitation disappears in the remote past.

By inserting Eq. (39) into Eq. (25) it follows that the attached field oscillates monochromatically (angular frequency ω) with an amplitude $\vec{E}_L^{NR}(z;q_\parallel\vec{e}_x,\omega)$ given by

$$\vec{E}_L^{NR}(z;q_\parallel\vec{e}_x,\omega)=\frac{q_\parallel}{2i\epsilon_0\omega}e^{-q_\parallel|z-z_0|}\times\begin{pmatrix}1 & 0 & i\,\mathrm{sgn}(z-z_0)\\ 0 & 0 & 0\\ i\,\mathrm{sgn}(z-z_0) & 0 & -1\end{pmatrix}\cdot\vec{J}_0(\omega). \tag{40}$$

Expressed in terms of the nonretarded propagator $\overleftrightarrow{D}_0^{NR}(z-z_0;q_\parallel\vec{e}_x,\omega)$ given in Eq. (15), Eq. (40) takes the form displayed in Eq. (18).

For a harmonically oscillating sheet current density the amplitude of the retarded electric field, $\vec{E}_T^R(z;q_\parallel\vec{e}_x,\omega)$, is given by Eq. (17), and to obtain the explicit expression for the propagator $\overleftrightarrow{D}_0^T(z-z_0;q_\parallel\vec{e}_x,\omega)$ we just need the Fourier transform of Eq. (35), namely,

$$\overleftrightarrow{D}_0^T(Z;q_\parallel\vec{e}_x,\omega)=\overleftrightarrow{D}_{space}^T(Z;q_\parallel\vec{e}_x,\omega)+\overleftrightarrow{D}_{time}^T(Z;q_\parallel\vec{e}_x,\omega). \tag{41}$$

The spacelike part of the propagator, $\overleftrightarrow{D}_{space}^T(Z;q_\parallel\vec{e}_x,\omega)$, is readily found using the result

$$\int_{-\infty}^{\infty}\tau\theta(\tau)\,\theta(|Z|-c_0\tau)e^{i\omega\tau}d\tau=\frac{1}{\omega^2}[(1-iq_0|Z|)e^{iq_0|Z|}-1], \tag{42}$$

where $q_0=\omega/c_0$ is the vacuum wave number of light. Hence,

$$\overleftrightarrow{D}_{space}^T(Z;q_\parallel\vec{e}_x,\omega)=\frac{q_\parallel}{4q_0^2}[(1-iq_0|Z|)e^{iq_0|Z|}-1]e^{-q_\parallel|Z|}\times\begin{pmatrix}1 & 0 & i\,\mathrm{sgn}\,Z\\ 0 & 0 & 0\\ i\,\mathrm{sgn}\,Z & 0 & -1\end{pmatrix}. \tag{43}$$

Though necessarily an oscillating factor with a spatial period $2\pi/q_0$ is also present in Eq. (43), the spatial part of the retarded response vanishes exponentially with the distance from the sheet plane, the decay constant being $q_\parallel$.

To determine the timelike part of the propagator, $\overleftrightarrow{D}_{time}^T(Z;q_\parallel\vec{e}_x,\omega)$, one just needs to combine Eqs. (16) and (41). This gives

$$\overleftrightarrow{D}_{time}^T(Z;q_\parallel\vec{e}_x,\omega)=\overleftrightarrow{D}_0(Z;q_\parallel\vec{e}_x,\omega)-\overleftrightarrow{D}_0^{NR}(Z;q_\parallel\vec{e}_x,\omega)-\overleftrightarrow{D}_{space}^T(Z;q_\parallel\vec{e}_x,\omega), \tag{44}$$

and since the three terms on the right-hand side of this equation have already been found, see Eqs. (14), (15), and (43), $\overleftrightarrow{D}_{time}^T(Z;q_\parallel\vec{e}_x,\omega)$ is obtained. The structures of $\overleftrightarrow{D}_0^{NR}$ and $\overleftrightarrow{D}_{space}^T$ are closely related, and the sum of these propagators is

$$\overleftrightarrow{D}_0^{NR}(Z;q_\parallel\vec{e}_x,\omega)+\overleftrightarrow{D}_{space}^T(Z;q_\parallel\vec{e}_x,\omega)$$
$$=\frac{q_\parallel}{4q_0^2}[(1+iq_0|Z|)e^{iq_0|Z|}+1]e^{-q_\parallel|Z|}$$
$$\times\begin{pmatrix} 1 & 0 & i\,\mathrm{sgn}\,Z \\ 0 & 0 & 0 \\ i\,\mathrm{sgn}\,Z & 0 & -1 \end{pmatrix}. \qquad (45)$$

Before finishing this section on the near-field electrodynamics of a sheet carrying a monochromatic current density, let us briefly reflect on the distance dependence of the electromagnetic field and its various parts. Thus, if $q_\parallel < q_0$, it appears from Eq. (14) that the total electric field oscillates as a function of $|z-z_0|$ with a period $2\pi/(q_0^2-q_\parallel^2)^{1/2}$. Since both $\overleftrightarrow{D}_0^{NR}$ and $\overleftrightarrow{D}_{space}^T$ vanish exponentially (with a decay constant $q_\parallel$), it follows that $\overleftrightarrow{D}_{time}^T(Z\to\infty;q_\parallel\vec{e}_x,\omega)=\overleftrightarrow{D}_0(Z\to\infty;q_\parallel\vec{e}_x,\omega)$. Far from the sheet the field is therefore purely transverse and only timelike events are coupled. In the special case where the current density is y polarized, the relation $(\overleftrightarrow{D}_{time}^T)_{yy}=(\overleftrightarrow{D}_0)_{yy}$ holds for all $|z-z_0|$. If $q_\parallel > q_0$, the situation becomes particularly interesting. The total field now decays exponentially as a function of $|z-z_0|$ with a (real) decay constant $\alpha_\perp^0=(q_\parallel^2-q_0^2)^{1/2}$. Despite the fact that the total field is proportional to $\exp(-\alpha_\perp^0|z-z_0|)\exp(-i\omega t)$ and hence at a first glance seems to have ''standing-wave character,'' we know that it contains a retarded transverse component carrying information away from the sheet with the vacuum speed of light. Since $\overleftrightarrow{D}_0^T=\overleftrightarrow{D}_0-\overleftrightarrow{D}_0^{NR}$ [see Eq. (16)], the retarded response in general is described via a Green function containing a combination of two exponential decay lengths, namely, $2\pi/(q_\parallel^2-q_0^2)^{1/2}$ and $2\pi/q_\parallel$. For $q_\parallel$ values only slightly larger than q_0, the propagating transverse field, although exponentially decaying, effectively reaches much farther away from the sheet plane than does the attached longitudinal field. Beyond the decay length $2\pi/q_\parallel$ the timelike part of the retarded response dominates; cf. Eq. (45). For y-polarized current densities only the retarded response with its decay length $2\pi/(q_\parallel^2-q_0^2)^{1/2}$ is present.

In a field-quantized description not yet developed, the retarded response is the one to be linked to the photon concept, and, *provided* one has a detector sensitive *only* to the photon part of the sheet field [26,27], the detected field should exhibit a distance dependence that is a linear combination of the two exponential forms $\exp[-(q_\parallel^2-q_0^2)^{1/2}|z-z_0|]$ and $\exp(-q_\parallel|z-z_0|)$.

E. Photon- and electron-eye views

We have seen in Sec. IV C that the transverse electromagnetic field emerging from a sheet current-density distribution, though retarded, contains both space- and timelike contributions. Based on the description of the near-field electrodynamics of a single atom (Sec. II) we know that the spacelike part of the detached field is present only if one identifies the source region with the spatial domain occupied by the induced current-density distribution of the atom. This so-called electron-eye view follows in a rigorous manner from a quantum electrodynamical description and is not in conflict with the principle of causality. Another view of the transverse electrodynamics called the photon-eye view appears if one identifies the source region with that of the transverse part of the induced current density. Though the electron- and photon-eye views lead to exactly the same predictions for all measurable quantities, the intuitive pictures they offer look quite different, and, if compared, it appears to me that a better insight into the near-field electrodynamics is achieved. The electron-eye view is convenient because it can be related in a direct and simple manner to the energy wave function [23–25] describing single-photon dynamics in space and time via the relativistically invariant photon propagator [21]. The photon-eye view on the other hand is particularly useful for discussing the spatial localizability of a photon emitted from a given current-density source.

Holding the point of view that the photon localizability plays an important role for our understanding of the optical tunneling process, it is fruitful to study the photon-eye view for the sheet electrodynamics. In this view only events on the light cone are coupled, and in the space-frequency domain the transverse electric field $\vec{E}_T(\vec{r};\omega)$ is given by [cf. Eq. (9)]

$$\vec{E}_T(\vec{r};\omega)=-i\mu_0\omega\int_{-\infty}^{\infty}\overleftrightarrow{d}^R(\vec{r}-\vec{r}\,';\omega)\cdot\vec{J}_T(\vec{r}\,';\omega)d^3r', \qquad (46)$$

where

$$\overleftrightarrow{d}^R(\vec{R};\omega)=-\frac{e^{iq_0R}}{4\pi R}\overleftrightarrow{U}. \qquad (47)$$

To apply Eq. (46) in the sheet case we make use of the Weyl expansion for a spherical scalar wave [27]

$$\frac{e^{iq_0R}}{R}=\frac{i}{2\pi}\int_{-\infty}^{\infty}\frac{1}{\kappa_\perp^0}e^{i\kappa_\perp^0|Z|}e^{i\vec{q}_\parallel\cdot\vec{R}_\parallel}d^2q_\parallel, \qquad (48)$$

where $R=|\vec{R}|=|\vec{R}_\parallel+Z\vec{e}_z|$; and upon a comparison to the general Weyl expansion

$$\overleftrightarrow{d}^R(\vec{R};\omega)=(2\pi)^{-2}\int_{-\infty}^{\infty}\overleftrightarrow{d}(Z;\vec{q}_\parallel,\omega)e^{i\vec{q}_\parallel\cdot\vec{R}_\parallel}d^2q_\parallel \qquad (49)$$

for the isotropic propagator, one obtains

$$\overleftrightarrow{d}^R(Z;q_\parallel,\omega)=\frac{1}{2i\kappa_\perp^0}e^{i\kappa_\perp^0|Z|}\overleftrightarrow{U}, \qquad (50)$$

with again $\kappa_\perp^0=(q_0^2-q_\parallel^2)^{1/2}$. By replacing $\vec{q}_\parallel$ by $q_\parallel$ in the argument of the propagator we have stressed that this depends only on the magnitude of $\vec{q}_\parallel$. For the sheet electrodynamics ''seen with the eyes of the photon,'' the relevant integral relation between the retarded field $\vec{E}_T(z;\vec{q}_\parallel,\omega)=\vec{E}_T^R(z;\vec{q}_\parallel,\omega)$ and the transverse current density $\vec{J}_T(z';\vec{q}_\parallel,\omega)$ therefore is

$$\vec{E}_T(z;\vec{q}_\parallel,\omega) = -i\mu_0\omega\int_{-\infty}^{\infty}\overleftrightarrow{d}^R(z-z';q_\parallel,\omega)\cdot\vec{J}_T(z';\vec{q}_\parallel,\omega)dz', \quad (51)$$

with $\overleftrightarrow{d}^R(z-z';q_\parallel,\omega)$ given by Eq. (50).

In the electron-eye view the relation equivalent to Eq. (50) is Eq. (17), and a comparison of the two shows that the one gives an algebraic relation between field and current density, and the other an integral relation. The reason for this stems from the fact that the transverse (or longitudinal) part of a vector field $\vec{V}(\vec{r}\,')=\vec{V}_0\delta(z'-z_0)\exp(i\vec{q}_\parallel\cdot\vec{r}\,')$, which is different from zero only in the plane $z'=z_0$, is nonzero outside this plane also. To quantify this statement we consider a current density of the form $\vec{J}(\vec{r}\,',t)=\vec{J}(z',t;q_\parallel\vec{e}_x)\exp(iq_\parallel x')$. The associated transverse current density necessarily has the generic form $\vec{J}_T(\vec{r},t)=\vec{J}_T(z,t;q_\parallel\vec{e}_x)\exp(iq_\parallel x)$, and the relation between the two amplitudes is given by the nonlocal equation

$$\vec{J}_T(z,t;q_\parallel\vec{e}_x)=\int_{-\infty}^{\infty}\overleftrightarrow{\delta}_T(z-z';q_\parallel\vec{e}_x)\cdot\vec{J}(z',t;q_\parallel\vec{e}_x)dz', \quad (52)$$

where [28]

$$\overleftrightarrow{\delta}_T(z-z';q_\parallel\vec{e}_x)=(\overleftrightarrow{U}-\vec{e}_z\vec{e}_z)\,\delta(z-z') - \frac{q_\parallel}{2}e^{-q_\parallel|z-z'|}[\vec{e}_x\vec{e}_x-\vec{e}_z\vec{e}_z + i(\vec{e}_x\vec{e}_z+\vec{e}_z\vec{e}_x)\mathrm{sgn}(z-z')] \quad (53)$$

is the relevant transverse δ function. The corresponding longitudinal δ function, $\overleftrightarrow{\delta}_L(z-z';q_\parallel\vec{e}_x)$, was given in Eq. (26), and the sum of the two is equal to the Dirac δ function of $z-z'$ times the unit tensor, i.e., $\overleftrightarrow{\delta}_T(z-z';q_\parallel\vec{e}_x)+\overleftrightarrow{\delta}_L(z-z';q_\parallel\vec{e}_x)=\overleftrightarrow{U}\delta(z-z')$. If the current density itself is confined to the plane $z'=z_0$, so that $\vec{J}(z',t;q_\parallel\vec{e}_x)=\vec{J}_0(t)\delta(z'-z_0)$, the transverse current density, given by

$$\vec{J}_T(z,t;q_\parallel\vec{e}_x)=\Big\{\delta(z-z_0)(\overleftrightarrow{U}-\vec{e}_z\vec{e}_z)-\frac{q_\parallel}{2}e^{-iq_\parallel|z-z_0|}\times[\vec{e}_x\vec{e}_x-\vec{e}_z\vec{e}_z+i(\vec{e}_x\vec{e}_z+\vec{e}_z\vec{e}_x)\times\mathrm{sgn}(z-z_0)]\Big\}\cdot\vec{J}_0(t), \quad (54)$$

extends over a $|z-z_0|$ strip characterized by the exponential decay length $q_\parallel^{-1}$, and is singular at $z=z_0$ [as $\vec{J}(z',t;q_\parallel\vec{e}_x)$ of course]. In the special case where the current density is polarized in the y direction, $\vec{J}_0(t)=J_0(t)\vec{e}_y$, the transverse current density is confined to the sheet plane, and in fact equal to the total current density, i.e., $\vec{J}_T(z,t;q_\parallel\vec{e}_x)=\vec{J}(z,t;q_\parallel\vec{e}_x)=\vec{J}_0(t)\delta(z-z_0)\vec{e}_y$. As expected, the decay length $(q_\parallel^{-1})$ for the transverse current density coincides with that of the spacelike part of the coupling in the electron-eye view.

By inserting the frequency transform of Eq. (54) into Eq. (51), a subsequent comparison with Eq. (17) reveals that the transverse propagator can be represented by the significant integral formula

$$\overleftrightarrow{D}_0^T(z-z_0;q_\parallel\vec{e}_x,\omega) = \int_{-\infty}^{\infty}\overleftrightarrow{d}^R(z-z';q_\parallel,\omega)\cdot\overleftrightarrow{\delta}_T(z'-z_0;q_\parallel\vec{e}_x)dz'. \quad (55)$$

A direct proof that Eq. (55) is correct can be established by carrying out the z' integration; see Appendix B.

V. OPTICAL TUNNELING AND MACROSCOPIC CURRENT-DENSITY DISTRIBUTIONS

In the preceding sections we have studied the spatial confinement of light emerging from an atom (or a pointlike particle) and from a sheet, and I have argued that a relation exists between the near-field electrodynamics of charged-particle distributions and optical tunneling. Hitherto, optical tunneling effects have always seemed to have been investigated in the context of macroscopic media in the literature; cf., e.g., tunneling across a vacuum gap between dielectric prisms, tunneling in thin metal films suspended in vacuum or placed between dielectric media, tunneling in photonic band-gap materials, tunneling across air gaps in waveguides, etc. Recently, it has also been discussed among scientists in the optical near-field community whether optical tunneling may be observed in their field. To demonstrate that the microscopic considerations put forth in the first parts of this paper are closely related to the conventional macroscopic approach to tunneling, we shall now embark on an extension to macroscopic media.

A. General considerations

Though we shall aim at a rather general description of the optical tunneling process, we nevertheless assume that the medium under study exhibits translational invariance against arbitrary displacements parallel to the xy plane of our Cartesian coordinate system. This assumption is not crucial for the analysis nor for a basic understanding of the underlying physics, and may easily be lifted. For simplicity we also assume that the induced current density is independent of the y coordinate. Despite invoking the two aforementioned assumptions, we are still able to make contact with the key experiment: optical tunneling across a vacuum gap between dielectric prisms.

Beginning thus with a current-density distribution

$$\vec{J}(\vec{r}\,',t')=\vec{J}(z',t';q_\parallel\vec{e}_x)e^{iq_\parallel x'}, \quad (56)$$

the total electric field necessarily takes the form

$$\vec{E}(\vec{r},t)=\vec{E}(z,t;q_\parallel\vec{e}_x)e^{iq_\parallel x}. \quad (57)$$

Although the current density we start with (and thus the field) has plane-wave character along the x axis, a Fourier superposition of the obtained results for different $q_\parallel$ values readily allows one to generalize the considerations to arbitrary x distributions of $\vec{J}$. Such a generalization is needed if one wants to examine the possible link between near-field diffraction from slits and line sources and optical tunneling. In integral form the field $[\vec{E}(z,t;q_\parallel\vec{e}_x)]$ and current-density $[\vec{J}(z',t';q_\parallel\vec{e}_x)]$ amplitudes are related via the expression

$$\vec{E}(z,t;q_\parallel\vec{e}_x)=\mu_0\int_{-\infty}^{\infty}\overleftrightarrow{G}_0(z-z',t-t';q_\parallel\vec{e}_x)\cdot\frac{\partial\vec{J}(z',t';q_\parallel\vec{e}_x)}{\partial t'}dz'dt' \tag{58}$$

or, equivalently, in the space-frequency representation

$$\vec{E}(z;q_\parallel\vec{e}_x,\omega)=-i\mu_0\omega\int_{-\infty}^{\infty}\overleftrightarrow{G}_0(z-z';q_\parallel\vec{e}_x,\omega)\cdot\vec{J}(z';q_\parallel\vec{e}_x,\omega)dz'. \tag{59}$$

Apart from a contact term, the Green function $\overleftrightarrow{G}_0(z-z';q_\parallel\vec{e}_x,\omega)$ is identical to $\overleftrightarrow{D}_0(z-z';q_\parallel\vec{e}_x,\omega)$ given in Eq. (14), i.e., [28,32],

$$\overleftrightarrow{G}_0(z-z';q_\parallel\vec{e}_x,\omega)=\overleftrightarrow{D}_0(z-z';q_\parallel\vec{e}_x,\omega)+q_0^{-2}\delta(z-z')\vec{e}_z\vec{e}_z. \tag{60}$$

If the contact term is added to the nonretarded part of $\overleftrightarrow{D}_0$, one obtains the longitudinal δ function $\overleftrightarrow{\delta}_L(z-z';q_\parallel\vec{e}_x)$ divided by q_0^2, i.e.,

$$q_0^{-2}\overleftrightarrow{\delta}_L(z-z';q_\parallel\vec{e}_x)=\overleftrightarrow{D}_0^{NR}(z-z';q_\parallel\vec{e}_x,\omega)+q_0^{-2}\delta(z-z')\vec{e}_z\vec{e}_z, \tag{61}$$

as one may readily verify looking at Eqs. (15) and (26). By dividing $\overleftrightarrow{G}_0(z-z';q_\parallel\vec{e}_x,\omega)$ into two pieces as follows:

$$\overleftrightarrow{G}_0(z-z';q_\parallel\vec{e}_x,\omega)=\overleftrightarrow{D}_0^T(z-z';q_\parallel\vec{e}_x,\omega)+q_0^{-2}\overleftrightarrow{\delta}_L(z-z';q_\parallel\vec{e}_x), \tag{62}$$

it is realized that the retarded and transverse electric field is given by

$$\vec{E}_T^R(z;q_\parallel\vec{e}_x,\omega)=-i\mu_0\omega\int_{-\infty}^{\infty}\overleftrightarrow{D}_0^T(z-z';q_\parallel\vec{e}_x,\omega)\cdot\vec{J}(z';q_\parallel\vec{e}_x,\omega)dz' \tag{63}$$

and the longitudinal (attached) and nonretarded field by

$$\vec{E}_L^{NR}(z;q_\parallel\vec{e}_x,\omega)=\frac{1}{i\epsilon_0\omega}\int_{-\infty}^{\infty}\overleftrightarrow{\delta}_L(z-z';q_\parallel\vec{e}_x)\cdot\vec{J}(z';q_\parallel\vec{e}_x,\omega)dz'=\frac{1}{i\epsilon_0\omega}\vec{J}_L(z;q_\parallel\vec{e}_x,\omega). \tag{64}$$

Since the transverse propagator appearing in Eq. (63) is identical to the one used in the sheet case, the transverse dynamics of macroscopic media both in the frequency domain and in the time domain, where Eq. (63) reads

$$\vec{E}_T^R(z,t;q_\parallel\vec{e}_x)=\mu_0\int_{-\infty}^{\infty}\overleftrightarrow{D}_0^T(z-z',t-t';q_\parallel\vec{e}_x)\cdot\frac{\partial\vec{J}(z',t';q_\parallel\vec{e}_x)}{\partial t'}dz'dt', \tag{65}$$

can be discussed along the same lines as for the sheet source. Furthermore, because Eq. (64) in the space-time domain takes the form

$$\vec{E}_L^{NR}(z,t;q_\parallel\vec{e}_x)=-\epsilon_0^{-1}\vec{e}_z\vec{e}_z\cdot\int_{-\infty}^{t}\vec{J}(z,t';q_\parallel\vec{e}_x)dt'-\frac{q_\parallel}{2\epsilon_0}\int_{-\infty}^{\infty}\left[e^{-q_\parallel|z-z'|}\begin{pmatrix}1 & 0 & i\,\mathrm{sgn}(z-z')\\ 0 & 0 & 0\\ i\,\mathrm{sgn}(z-z') & 0 & -1\end{pmatrix}\cdot\int_{-\infty}^{t}\vec{J}(z',t';q_\parallel\vec{e}_x)dt'\right]dz', \tag{66}$$

the qualitative analysis of the attached-field dynamics can be carried out in a fashion similar to the one used in the sheet case; one just has to remember the contact term [first term on the right-hand side of Eq. (66)] and to add the effects from the current densities $\vec{J}(z',t';q_\parallel\vec{e}_x)dz'$ of the various infinitesimally thin strips $(z',z'+dz')$. For observation planes in the vacuum, the contact term does not contribute, of course, but it is nevertheless needed to ensure that the nonretarded (attached) field is rotational free (longitudinal) for all z values [28].

B. Many-body linear response theory

So far, we have only investigated the propagator relation between the local electric field and the prevailing current density, but this current density itself is in fact determined by the sum of the prescribed external field impressed on the medium and the yet unknown induced electric field. To close the loop problem an extra relation is needed between the field and the current density. This relation is provided by the Schrödinger equation in the nonrelativistic regime, and as is most often done we shall assume that the relation is linear, but in contrast to conventional (macroscopic) studies of optical tunneling, we shall allow the relation to be spatially nonlocal.

In linearized many-body (MB) response theory the constitutive equation in our case takes the general form [32]

$$\vec{J}(z;q_\parallel\vec{e}_x,\omega)=\int_{-\infty}^{\infty}\overleftrightarrow{\sigma}^{MB}(z,z';q_\parallel\vec{e}_x,\omega)\cdot[\vec{E}_T(z';q_\parallel\vec{e}_x,\omega)+\vec{E}_L^{ext}(z';q_\parallel\vec{e}_x,\omega)]dz', \tag{67}$$

where $\vec{E}_L^{ext}(z';q_\parallel\vec{e}_x,\omega)$ is the longitudinal part of the external (ext) field, and $\overleftrightarrow{\sigma}^{MB}(z,z';q_\parallel\vec{e}_x,\omega)$ is the microscopic many-body conductivity tensor, a nonlocal object in general. Usually the external light source (laser, etc.) is placed so far from the medium under study that $\vec{E}_L^{ext}=\vec{0}$ inside the medium, and we also know that the transverse field does not contain a self-field part for current densities of the form given in Eq. (56), so that $\vec{E}_T=\vec{E}_T^R$ in Eq. (67).

Taking into account the transverse external field $\vec{E}_T^{ext}(z;q_\parallel\vec{e}_x,\omega)$ acting on the medium, we are thus led to a loop equation

$$\vec{E}_T^R(z;q_\parallel\vec{e}_x,\omega)=\vec{E}_T^{ext}(z;q_\parallel\vec{e}_x,\omega)-i\mu_0\omega\int_{-\infty}^{\infty}\left[\int_{-\infty}^{\infty}\overleftrightarrow{D}_0^T(z-z'';q_\parallel\vec{e}_x,\omega)\cdot\overleftrightarrow{\sigma}^{MB}(z'',z';q_\parallel\vec{e}_x,\omega)dz''\right]\cdot\vec{E}_T^R(z';q_\parallel\vec{e}_x,\omega)dz' \tag{68}$$

for the retarded transverse field. Loop equations of the form given in Eq. (68) can be solved (approximately) using different schemes [32]. Once $\vec{E}_T^R(z;q_\parallel\vec{e}_x,\omega)$ has been obtained, the current density can be determined from Eq. (67) (leaving out $\vec{E}_L^{ext}$), and a knowledge of $\vec{J}(z;q_\parallel\vec{e}_x,\omega)$ allows one to calculate the attached field from Eq. (64).

In the present context of optical tunneling we need not have any explicit solution for the local field; it is sufficient just to realize how the induced current density $\vec{J}(z';q_\parallel\vec{e}_x,\omega)$ [and its transverse (longitudinal) part] emerges microscopically. Hence, if one denotes the various many-body energy eigenstates by the quantum labels $M,N,\ldots$, the many-body conductivity tensor has a structure [32]

$$\overleftrightarrow{\sigma}^{MB}(z,z')=\sum_{M,N}\mathcal{A}_{M,N}\vec{\mathcal{J}}_{M\to N}(z)\vec{\mathcal{J}}_{N\to M}(z'), \tag{69}$$

leaving out for simplicity the reference to $q_\parallel\vec{e}_x$ and ω. The quantity $\vec{\mathcal{J}}_{M\to N}(\vec{\mathcal{J}}_{N\to M})$ denotes the many-body transition current density involved in an electronic excitation from state M to state N (or opposite). From a knowledge of the stationary-state wave functions of these [many-body] states, $\vec{\mathcal{J}}_{M\to N}(z)$ $[\vec{\mathcal{J}}_{N\to M}(z')]$ can be obtained. The transverse [longitudinal] part of this current density, $\vec{\mathcal{J}}^T_{M\to N}(z)$ $[\vec{\mathcal{J}}^L_{M\to N}(z)]$, and *in particular its spill-out in vacuum* is the crucial one for the tunneling process, as one readily realizes by combining Eqs. (67) and (69). The final spill-out is determined by superimposing the weighted spill-outs belonging to the participating transitions. The weight factor of a given $M\to N$ transition is $\mathcal{A}_{M,N}\int_{-\infty}^{\infty}\vec{\mathcal{J}}_{N\to M}(z')\cdot\vec{E}_T(z')dz'$, and its explicit value may be determined once the loop equation for $\vec{E}_T(z)$ has been solved.

VI. PHENOMENOLOGICAL DESCRIPTION OF OPTICAL TUNNELING; IN PARTICULAR, THE ROLE OF SURFACE CURRENTS

In theoretical studies of optical tunneling across a vacuum gap separating two (nonmagnetic) macroscopic media, it is usually assumed that the tangential components of the electric and magnetic fields are continuous across the sharp medium/vacuum boundaries. For plane electromagnetic waves propagating perpendicular to the boundaries, i.e., in the z direction, or for s-polarized waves as such, this choice implies that the electric field and its first derivative with respect to z are continuous at the medium/vacuum surfaces, conditions which make the stationary-state problems for electron and optical tunneling mathematically equivalent [1]. As long as one can justify the assumption that no surface currents are induced at the boundaries, other matching choices for the electromagnetic field can be chosen without altering the physical result. If induced surface currents cannot be neglected one must be more careful. Hence, if one relies on the standard (textbook) jump (boundary) conditions for the field, inconsistencies are likely to appear, because the possible presence of induced surface currents perpendicular to the medium/vacuum boundary is neglected. This omission usually leads to results for the amplitude and reflection and transmission coefficients that depend on the choice of the set of jump conditions [33], an unacceptable situation, not least for optical tunneling studies, as we shall realize below. For electromagnetic transients and finite-frequency responses nothing prevents surface current oscillations from being induced with a component perpendicular to the medium/vacuum surface. In linear and nonlinear surface optics it is often crucial to keep the z component in the induced surface (or interface) current-density distribution. Although it has been claimed occasionally that the optical tunneling (time) is independent of the boundary conditions [34], the analysis below do not support such a point of view.

A. Heuristic sharp-boundary model

Let us now consider a model consisting of a semi-infinite medium occupying the half-space $z<0$ and separated from a vacuum half-space $(z>0)$ by a sharp boundary (at $z=0$). The sharp-boundary assumption is of course an abstraction, realized as early as 1860 by Lorenz [35] when analyzing the results of the reflection experiments carried out by Jamin 12 years earlier, but not understood in the intervening years [36]. The spill-out of the electron distribution will in general be somewhat less for dielectric than for metallic and semiconducting media. For the following qualitative discussion the sharp-boundary model is sufficient. In our treatment we shall consider monochromatic waves with a single wave-vector component parallel to the surface, and to simplify the notation we therefore omit the reference to $q_\parallel \vec{e}_x$ and ω from the notation, i.e., $\vec{J}(z;q_\parallel \vec{e}_x,\omega)=\vec{J}(z)$, etc. The induced current density hence is given by

$$\vec{J}(x,z)=\vec{J}_B(z)e^{iq_\parallel x}\theta(-z), \tag{70}$$

where one may consider $\vec{J}_B(z)$ [multiplied by $\exp(iq_\parallel x)$] as the bulk (B) contribution to the total current density. The presence of the Heaviside unit step function $\theta(-z)$ allows us to estimate the role of an induced surface current, albeit in a heuristic manner. To elaborate on this let us look at the divergence of the current density, i.e.,

$$\vec{\nabla}\cdot\vec{J}(x,z)=\theta(-z)\vec{\nabla}\cdot\vec{J}_B(x,z)-\delta(z)\vec{e}_z\cdot\vec{J}_B(x,0). \tag{71}$$

If the induced bulk current-density distribution is divergence free, $\vec{\nabla}\cdot\vec{J}_B(x,z)=0$, so that $\vec{J}_B(x,z)=\vec{J}_B^T(x,z)$, the part of the optical tunneling process that is associated with the attached longitudinal field originates solely in the induced surface current density and is present only if this has a component perpendicular to the medium/vacuum boundary; cf. the form of the second term on the right-hand side of Eq. (71). To underscore the importance of surface currents let us therefore analyze the case where the current-density distribution of the bulk is transverse. If the induced current density is s polarized in the bulk and surface regions the attached field and the spacelike part of the retarded field will vanish, as we have realized in Sec. IV, and it is therefore sufficient to restrict ourselves to studies of p-polarized distributions. In turn this means that the y component of all vector fields is zero. Two-component vectors and related 2×2-component tensors may hence to be used to simplify the notation.

The attached field can be calculated everywhere in space inserting Eq. (70) into Eq. (64) and utilizing the explicit expression given in Eq. (26) for the longitudinal δ function. By making use of the assumption that the bulk current density is transverse, i.e.,

$$iq_\parallel J_{B,x}(z)+\frac{dJ_{B,z}(z)}{dz}=0, \tag{72}$$

one obtains, as shown in Appendix C, the following expression for the nonretarded longitudinal field:

$$\vec{E}_L^{NR}(z)=\frac{e^{-q_\parallel|z|}}{2\epsilon_0\omega}J_{B,z}(0)\begin{pmatrix}1\\ i\,\mathrm{sgn}\,z\end{pmatrix}. \tag{73}$$

The simple result in Eq. (73) illustrates main principles of the optical tunneling process in a fine manner. Hence, when the induced current density is divergence free in the bulk, the nonretarded longitudinal field generated originates solely in the surface current density induced perpendicular to the boundary [$J_{B,z}(0)\vec{e}_z$], and, as we know, the attached field decays exponentially to both sides of the surface plane with a characteristic decay constant $q_\parallel$. In the vacuum ($z>0$) the field is right-hand circularly polarized and inside the medium ($z<0$) it is left-hand circularly polarized, in both domains, with the polarization unit vectors $(1,i\,\mathrm{sgn}\,z)/\sqrt{2}$ lying in the xz plane, of course. The particular form

$$\vec{E}_L^{NR}(x,z)=\frac{J_{B,z}(0)}{2\epsilon_0\omega}e^{iq_\parallel x}e^{-q_\parallel|z|}\begin{pmatrix}1\\ i\,\mathrm{sgn}\,z\end{pmatrix} \tag{74}$$

shows that *the nonretarded longitudinal field is not only rotational free* [$\partial E_{L,x}^{NR}(x,z)/\partial z-iq_\parallel E_{NR,z}^{L}(x,z)=0$] *as it by definition must be but also divergence free* [$iq_\parallel E_{L,x}^{NR}(x,z)+\partial E_{L,z}^{NR}(x,z)/\partial z=0$]. In the vacuum, where there is no charge density, the nonretarded field must of course also be divergence free, as the Maxwell equation $\vec{\nabla}\cdot(\vec{E}_T^R+\vec{E}_L^{NR})=\vec{0}$ or equivalently $\vec{\nabla}\cdot\vec{E}_L^{NR}=0$ immediately shows.

The spacelike part of the retarded field is determined from the integral relation

$$\vec{E}_{space}^T(z)=-i\mu_0\omega\int_{-\infty}^{\infty}\overleftrightarrow{D}_{space}^T(z-z')\cdot\vec{J}(z')dz' \tag{75}$$

upon insertion of the expressions given in Eqs. (43) and (70) for the propagator and current density. By assuming as before that the bulk current density is transverse, a tedious but straightforward calculation (see Appendix C) leads to the following result for the spacelike part of the retarded field:

$$\begin{aligned}\vec{E}_{space}^T(z)=\frac{1}{4\epsilon_0\omega}\Bigg\{&q_0^2\int_{-\infty}^{0}\begin{pmatrix}1\\ i\,\mathrm{sgn}(z-z')\end{pmatrix}(z-z')\\ &\times e^{(iq_0-q_\parallel)|z-z'|}J_{B,z}(z')dz'+\begin{pmatrix}1\\ i\,\mathrm{sgn}\,z\end{pmatrix}\\ &\times[(1-iq_0|z|)e^{iq_0|z|}-1]e^{-q_\parallel|z|}J_{B,z}(0)\Bigg\}.\end{aligned} \tag{76}$$

The term proportional to $J_{B,z}(0)$ in Eq. (76) gives the contribution to the spacelike field from the z component of the surface current density. This contribution is right- and left-hand circularly polarized in the vacuum and medium regions, respectively. It also vanishes when the observation plane approaches the surface, i.e., for $z\to0$, as one would expect for a retarded spacelike contribution; cf. Eqs. (2) and (35). The integral term in Eq. (76) represents the contribution to the spacelike field from the induced bulk current-density distribution. The infinitesimal contribution from the z component of the current density in the strip located between z' and $z'+dz'[J_{B,z}(z')dz']$ is right- or left-hand circularly polarized, depending on whether the plane of observation is located to the right ($z>z'$) or left ($z<z'$) of this plane, and again this contribution vanishes as the plane of observation approaches the source plane ($z\to z'$).

Up to this point we have only considered the electric field generated by a *given* induced current-density distribution. As discussed in Sec. V B, a self-consistent theory is obtained by relating the induced current density in the medium under study to the prevailing field. In the present context it is sufficient to limit the considerations to the linear regime and assume that the response is isotropic, linear, and local in space. The relevant single-body conductivity tensor hence takes the form

$$\overleftrightarrow{\sigma}(z,z';q_\parallel\vec{e}_x,\omega)=\sigma(\omega)\delta(z-z')\overleftrightarrow{U} \tag{77}$$

in the frequency domain. Within the framework of the random-phase-approximation approach the bulk current density appearing in Eq. (70) is therefore given by

$$\vec{J}_0(z)=\sigma(\omega)\vec{E}(z), \tag{78}$$

and contact with previous studies of tunneling across a vacuum gap [1] is established, assuming the local electric field to consist of a spatially single-mode incident (inc) field ($\vec{E}^{inc}$) plus the associated reflected (refl) field ($\vec{E}^{refl}$). In most cases the source of the incident field is located so far from the medium under investigation that the transverse current-density domains of the source and medium do not overlap. In such situations the incident field is transverse ($\vec{E}^{inc}=\vec{E}_T^{inc}$), and since the medium is assumed to be isotropic, the reflected field must also be transverse, i.e., $\vec{E}^{refl}=\vec{E}_T^{refl}$. Altogether, the bulk current density thus becomes

$$\vec{J}_B(z)=\sigma(\omega)(\overleftrightarrow{U}e^{iq_\perp z}+\overleftrightarrow{r}e^{-iq_\perp z})\cdot\vec{E}_T^{inc}(0), \tag{79}$$

where $\vec{E}_T^{inc}(0)=\vec{E}_T^{inc}(q_\parallel\vec{e}_x,\omega)$ is the amplitude of the incident field, and

$$\overleftrightarrow{r}=r_p(\vec{e}_z\vec{e}_z-\vec{e}_x\vec{e}_x)+r_s\vec{e}_y\vec{e}_y \tag{80}$$

the reflection matrix. Fresnel's amplitude reflection coefficients r_p and r_s for *p*- and *s*-polarized fields, respectively ,are those belonging to reflection from the medium side of the boundary. In terms of the complex relative dielectric constant $\epsilon(\omega)=1+i\sigma(\omega)/(\epsilon_0\omega)$, the wave-vector component of the incident field perpendicular to the surface is given by

$$q_\perp=[q_0^2\epsilon(\omega)-q_\parallel^2]^{1/2}. \tag{81}$$

In the manner we have introduced $\sigma(\omega)$ here, the associated $\epsilon(\omega)$ is able to describe the optical response of dielectric as well as semiconducting and metallic media. For the standard situation where the incident field is divergence free, it appears from Eq. (79) that the bulk current density is transverse, and therefore Eqs. (74) and (76) can be used to calculate the attached and spacelike field parts. Since $J_{B,z}(0)=(1+r_p)\sigma(\omega)E_{T,z}^{inc}(0)$, the field contributions from the surface current density are readily expressed in terms of the amplitude of the incident electric field. The entire $\vec{E}_{space}^T(z)$ field by is obtained inserting Eq. (79) into Eq. (76). The explicit result is not needed here. A schematic illustration of the optical tunneling across a vacuum gap separating two dielectric prisms is presented in Fig. 3, in a manner that is meant to underscore the physical picture established in this work.

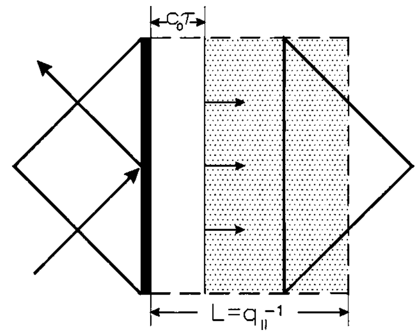

FIG. 3. Schematic illustration of the optical tunneling across a vacuum gap between two dielectric prisms as it is pictured in this work. The superposition of incident and reflected electromagnetic fields (indicated by the two big arrows) gives rise to a current density in and at the surface (black strip) of the prism to the left. If the other prism is placed within the evanescent zone (characterized by the exponential decay length $q_\parallel^{-1}$) of the first prism, optical tunneling occurs. In the picture suggested in this paper the tunneling field has two components, namely, a spacelike retarded (and necessarily transverse) component and a standing-wave-like longitudinal component. The back edge of the spacelike component (indicated by small arrows attached) moves away from the source prism with the vacuum speed of light, and is thus only nonvanishing in the shaded part of the evanescent region. From a quantum electrodynamic point of view the photons emitted by the surface current density induced in the source prism cannot be better localized spatially than what is dictated by the exponential decay length ($q_\parallel^{-1}$) of the transverse part of the surface current density.

B. Enhanced surface-generated tunneling

When the current density induced in the bulk of a macroscopic medium is divergence free the part of the optical tunneling process one may associate with the attached field originates solely in the currents generated in the surface region. In Sec. VI A, where a naive sharp-boundary model was adopted, the amplitude strength of the nonretarded field was proportional to the normal component of the bulk current density at the edge. To go beyond the heuristic approach it is necessary to take into account the fact that the electron density changes from its bulk value to zero over a finite distance in the *z* direction. Once a (self-consistent) surface potential has been determined, the bound energy eigenstates of the electrons may be found, and from a knowledge of these the light-induced surface current density can be calculated. For metallic and semiconducting media where highly delocalized Bloch states play a particular role, the field-induced longitudinal currents terminating at the surface can extend many Fermi wavelengths into the solid, and the overall contribution from surface states localized to within a few atomic monolayers may be rather weak.

To investigate the role of surface currents in the optical tunneling process it may therefore be fruitful to seek to enhance the currents induced in the surface region relative to those generated in the bulk. One possibility for doing this might be to deposit an ultrathin metallic or semiconducting film on top of a homogeneous dielectric substrate. If the film is sufficiently thin, the electron motion would be subjected to an essential spatial quantization perpendicular to the plane of the film, and resonance excitation between selected pairs of these so-called quantum-well states may lead to strong oscillating surface currents perpendicular to the well plane [32]. A few-monolayers thick metallic film may thus behave like a two-level system and strong currents can be induced in the z direction with p-polarized light incident at an oblique angle. If the electromagnetic field exciting the quantum well (from the medium side) is transverse, the attached tunneling field stems from the quantum well, and with a dielectric substrate even the spacelike part of the transverse tunneling field may be dominated by the quantum-well current source. A semiconducting $GaAs/Ga_{1-x}Al_xAs$ film/substrate combination seems to be adequate for optical tunneling studies in the (near)-infrared regime, and a film thickness of the order of ~ 100 Å can serve as a two-level system [37].

Surface currents may also be enhanced by nonlinear optical methods. Hence, in centrosymmetric media, the optical second-harmonic generation primarily is a surface effect, and even though the first harmonic (fundamental) field is transverse strong longitudinal field components can be induced in the surface region [38].

In certain wavelength regions so-called electromagnetic surface waves can be excited on an interface between two macroscopic media, or at a medium/vacuum surface. Such waves, which constitute part of the electromagnetic eigenmode spectrum of these systems, in certain frequency regions, contain a substantial fraction of nonretarded longitudinal fields so important for optical tunneling [39]. At a metal/vacuum surface the electromagnetic surface waves are even dominated by the longitudinal field contribution for frequencies close to the surface plasmon frequency. At this frequency the waves are circularly polarized in the plane of propagation both inside the metal and in the vacuum domain, and, in fact, the *total* field takes precisely the form given in Eq. (74) [40].

VII. OUTLOOK

Taking as a starting point for optical tunneling studies the framework suggested in this paper, a number of important issues should be addressed. Thus, instead of forming wave packets of the total electromagnetic field, it would be interesting to build these from the transverse part of the electromagnetic field, and again investigate the role of the various velocities introduced in the literature, the pulse reshaping, etc. In the transverse photon propagator description used in this work the inherent role of the spatial photon delocalizability for the wave packet analysis might show up clearly. Since the retarded field studied here essentially is identical to the Riemann-Silberstein energy wave function for photons in real space, a rigorous single-photon tunneling theory can be constructed by adjusting the prevailing current density in such a manner that the eigenvalue of the number operator equals unity. The approach presented in this paper for improving our understanding of optical tunneling should be of relevance also for a proper interpretation of the two-photon coincidence experiments of Chiao and co-workers [4,10], but a detailed account can only be given after an extension of the present theory along the lines indicated above has been worked out. Work on a quantum electrodynamic theory for the spatial localization (and birth) of polychromatic single-photon wave packets and, once generated, their Einstein-causal propagation is in progress. To make quantitative contact with the experimental microwave studies of optical tunneling of many-photon pulses [10,16] and thus with the issue of superluminality, the semiclassical theory presented here should be sufficient, but it is necessary to insert a specific current-density pulse form in Eqs. (65) and (66), and carry out a calculation of the transverse and longitudinal fields numerically. The choice of the pulse form should be in accordance with the form given by the incident microwave pulse; cf. Eq. (79). By decomposing the energy wave function in a basis set where the photon eigenstates have definite energy, momentum, and helicity, the role of the photon spin in optical tunneling possibly may be addressed. Building up by numerical methods the tunneling field originating in surface and bulk currents from the individual near fields of two- and three-dimensional regular distributions of atoms, respectively, it would be interesting to see how well the continuum theory developed here describes the time and space behavior of the tunneling field in media with strongly localized atomic (molecular) orbitals. Since there seems to exist a relation between optical tunneling and near-field diffraction, it would be interesting to investigate, for instance, the near-field diffraction from small holes in a scheme where a clear distinction between the transverse (or longitudinal) field and the total field is made. A proper identification of the attached field in the vicinity of the hole might provide us with a better insight into the selfconsistently induced transverse current density in the wall surrounding the hole, a quantity that appears to be of utmost importance in the vector theory for near-field diffraction. A correct identification of the retarded transverse field may also allow us to study the near-field diffraction of individual photons in space and time.

APPENDIX A: PHOTON PROPAGATOR FOR SHEET RADIATION

To verify that the expression for the retarded transverse propagator given in Eq. (35) follows from Eq. (34), let us first consider the yy component, i.e.,

$$D^T_{0,yy}(Z,\tau)=-\frac{c_0}{4\pi}\int_{-\infty}^{\infty}\frac{\sin(qc_0\tau)}{q}e^{iq_\perp Z}dq_\perp$$
$$=-\frac{c_0}{2\pi}\int_0^{\infty}\frac{\sin(c_0\tau\sqrt{q_\parallel^2+q_\perp^2})}{\sqrt{q_\parallel^2+q_\perp^2}}\cos q_\perp|Z|dq_\perp\,, \tag{A1}$$

since $\cos q_\perp Z=\cos q_\perp|Z|$. The last integral is different from zero only for $c_0\tau>|Z|$, and one finds [41]

$$D^T_{0,yy}(Z,\tau)=-\frac{c_0}{4}\theta(c_0\tau-|Z|)J_0[q_\parallel\sqrt{(c_0\tau)^2-Z^2}]. \tag{A2}$$

The zz component of the photon propagator, given in integral form by

$$D^T_{0,zz}(Z,\tau)=-\frac{c_0q_\parallel^2}{4\pi}\int_{-\infty}^{\infty}\frac{\sin(qc_0\tau)}{q^3}e^{iq_\perp Z}dq_\perp$$
$$=-\frac{c_0q_\parallel^2}{2\pi}\int_0^{\infty}\frac{\sin(c_0\tau\sqrt{q_\parallel^2+q_\perp^2})}{(q_\parallel^2+q_\perp^2)^{3/2}}\cos q_\perp|Z|dq_\perp, \quad \text{(A3)}$$

can be calculated explicitly for $|Z|>c_0\tau>0$ (causality implies that $\tau>0$ always), and we have [41]

$$D^T_{0,zz}(Z,\tau)=-\pi_1\theta(c_0\tau-|Z|)$$
$$-\frac{1}{4}c_0^2\tau q_\parallel e^{-q_\parallel|Z|}\theta(\tau)\theta(|Z|-c_0\tau), \quad \text{(A4)}$$

where π_1 is given by Eq. (36). Since $D^T_{0,xx}=D^T_{0,yy}-D^T_{0,zz}$, one readily obtains

$$D^T_{0,xx}(Z,\tau)=\left(\pi_1-\frac{c_0}{4}J_0[q_\parallel\sqrt{(c_0\tau)^2-Z^2}]\right)\theta(c_0\tau-|Z|)$$
$$+\frac{1}{4}c_0^2\tau q_\parallel e^{-q_\parallel|Z|}\theta(\tau)\theta(|Z|-c_0\tau). \quad \text{(A5)}$$

By now only the calculation of the identical off-diagonal elements remains,

$$D^T_{0,xz}(Z,\tau)=D^T_{0,zx}(Z,\tau)=\frac{c_0q_\parallel}{4\pi}\int_{-\infty}^{\infty}\frac{q_\perp}{q^3}\sin(qc_0\tau)e^{iq_\perp Z}dq_\perp$$
$$=\frac{ic_0q_\parallel}{2\pi}\int_0^{\infty}\frac{q_\perp}{q^3}\sin(qc_0\tau)\sin q_\perp Z dq_\perp$$
$$=-\frac{ic_0q_\parallel}{2\pi}\int_0^{\infty}\left[\frac{d}{dq_\perp}(q_\parallel^2+q_\perp^2)^{-1/2}\right]$$
$$\times\sin(qc_0\tau)\sin q_\perp Z dq_\perp. \quad \text{(A6)}$$

Since $(q_\perp^2+q_\parallel^2)^{-1/2}\sin(qc_0\tau)\sin q_\perp Z=0$ for $q_\perp=0$ and ∞, an integration by parts now gives

$$D^T_{0,xz}(Z,\tau)=\frac{ic_0q_\parallel}{2\pi}\left[Z\int_0^{\infty}\frac{\sin(qc_0\tau)}{q}\cos q_\perp Z dq_\perp\right.$$
$$\left.+c_0\tau\int_0^{\infty}\frac{q_\perp}{q^2}\cos(qc_0\tau)\sin q_\perp Z dq_\perp\right]. \quad \text{(A7)}$$

The first integral in Eq. (A7) has already been determined [see Eqs. (A1) and (A2)], and the second one can be calculated in explicit form for $0<c_0\tau<|Z|$ (see Ref. [41]). Hence

$$\int_0^{\infty}\frac{q_\perp}{q^2}\cos(qc_0\tau)\sin q_\perp Z dq_\perp=\text{sgn } Z\int_0^{\infty}\frac{q_\perp}{q_\parallel^2+q^2_\perp}\cos(c_0\tau\sqrt{q_\parallel^2+q_\perp^2})\sin q_\perp|Z|dq_\perp$$
$$=\left[\frac{\pi}{2}e^{-q_\parallel|Z|}\theta(\tau)\theta(|Z|-c_0\tau)+\frac{2\pi}{c_0^2\tau q_\parallel}\pi_2\theta(c_0\tau-|Z|)\right]\text{sgn } Z. \quad \text{(A8)}$$

Altogether we therefore obtain

$$D^T_{0,xz}(Z,\tau)=D^T_{0,zx}(Z,\tau)$$
$$=\frac{i}{4}c_0^2\tau q_\parallel\text{sgn } Z e^{-q_\parallel|Z|}\theta(\tau)\theta(|Z|-c_0\tau)$$
$$+\theta(c_0\tau-|Z|)\left\{i\pi_2 \text{ sgn } Z\right.$$
$$\left.+\frac{i}{4}c_0q_\parallel ZJ_0[q_\parallel\sqrt{(c_0\tau)^2-Z^2}]\right\}. \quad \text{(A9)}$$

Gathered in matrix form, the results in Eqs. (A2), (A4), (A5), and (A9) give $\overleftrightarrow{D}^T_0(Z,\tau)$ of Eq. (35).

APPENDIX B: INTEGRAL RELATION BETWEEN THE TRANSVERSE AND ISOTROPIC PROPAGATORS

To prove the integral relation in Eq. (55) one inserts the explicit expressions for $\overleftrightarrow{d}^R(z-z')$ [Eq. (50)] and $\overleftrightarrow{\delta}_T(z'-z_0)$ [Eq. (53)] on the right-hand side of the equation. This gives

$$\int_{-\infty}^{\infty}\overleftrightarrow{d}^R(z-z')\cdot\overleftrightarrow{\delta}_T(z'-z_0)dz'$$
$$=\frac{1}{2i\kappa_\perp^0}\left[e^{i\kappa_\perp^0|z-z_0|}\begin{pmatrix}1&0&0\\0&1&0\\0&0&0\end{pmatrix}\right.$$
$$\left.-\frac{q_\parallel}{2}\begin{pmatrix}I_1(z-z_0)&0&iI_2(z-z_0)\\0&0&0\\iI_2(z-z_0)&0&-I_1(z-z_0)\end{pmatrix}\right], \quad \text{(B1)}$$

where

$$I_1(z-z_0)=\int_{-\infty}^{\infty}e^{i\kappa_\perp^0|z-z'|}e^{-q_\parallel|z'-z_0|}dz' \quad \text{(B2)}$$

and

$$I_2(z-z_0)=\int_{-\infty}^{\infty}e^{i\kappa_\perp^0|z-z'|}e^{-q_\parallel|z'-z_0|}\,\mathrm{sgn}(z'-z_0)dz'. \tag{B3}$$

Since the integrals in Eqs. (B2) and (B3) are given by

$$I_1(z-z_0)=\frac{2}{q_0^2}(i\kappa_\perp^0 e^{-q_\parallel|z-z_0|}+q_\parallel e^{i\kappa_\perp^0|z-z_0|}), \tag{B4}$$

and

$$I_2(z-z_0)=\frac{2i\kappa_\perp^0}{q_0^2}(e^{-q_\parallel|z-z_0|}-e^{i\kappa_\perp^0|z-z_0|})\mathrm{sgn}(z-z_0), \tag{B5}$$

it appears that the integral relation in Eq. (B1) consists of terms proportional to $\exp(-q_\parallel|z-z_0|)$ and $\exp(i\kappa_\perp^0|z-z_0|)$, respectively. By gathering the two sets of terms in each of their tensors one obtains

$$\int_{-\infty}^{\infty}\overleftrightarrow{d}^R(z-z')\cdot\overleftrightarrow{\delta}_T(z'-z_0)dz'=\overleftrightarrow{D}_0(z-z_0)-\overleftrightarrow{D}_0^{NR}(z-z_0), \tag{B6}$$

where $\overleftrightarrow{D}_0(z-z_0)$ and $\overleftrightarrow{D}_0^{NR}(z-z_0)$ are given by Eqs. (14) and (15). Since the difference between $\overleftrightarrow{D}_0(z-z_0)$ and $\overleftrightarrow{D}_0^{NR}(z-z_0)$ is just the transverse propagator $\overleftrightarrow{D}_0^T(z-z_0)$ [see Eq. (16)], the claim in Eq. (55) has been proven.

APPENDIX C: TUNNELING FIELDS IN THE CASE OF TRANSVERSE BULK CURRENTS

1. Attached field

To determine within the framework of the sharp-boundary model the nonretarded longitudinal field in the vacuum half-space ($z>0$), we start from the expression

$$\vec{E}_L^{NR}(z)=\frac{q_\parallel}{2i\epsilon_0\omega}e^{-q_\parallel z}\times\int_{-\infty}^{0}e^{q_\parallel z'}[J_{B,x}(z')+iJ_{B,z}(z')]dz'\begin{pmatrix}1\\ i\end{pmatrix}, \tag{C1}$$

readily obtained by combining Eqs. (26), (64), and (70); and for simplicity all the p-polarized fields are written in two-component notation. An integration by parts of the term containing $J_{B,z}(z')$ in Eq. (C1) now gives

$$\vec{E}_L^{NR}(z)=\frac{1}{2\epsilon_0\omega}e^{-q_\parallel z}\times\left[J_{B,z}(0)-\int_{-\infty}^{0}e^{q_\parallel z'}\times\left(iq_\parallel J_{B,x}(z')+\frac{dJ_{B,z}(z')}{dz'}\right)dz'\right]\begin{pmatrix}1\\ i\end{pmatrix}, \tag{C2}$$

since $\exp(q_\parallel z')J_{B,z}(z')\to 0$ for $z'\to-\infty$. Under the assumption that the bulk current density is divergence free [see Eq. (72)] one immediately obtains the result in Eq. (73) for $z>0$.

For $z<0$, the contact term in the longitudinal δ functions must be included, and the integral over z' has to be divided into two parts over $-\infty<z'\leqslant z$ and $z<z'\leqslant 0$, respectively. Hence

$$\vec{E}_L^{NR}(z)=\frac{1}{i\epsilon_0\omega}\Bigg[J_{B,z}(z)\begin{pmatrix}0\\ 1\end{pmatrix}+\frac{q_\parallel}{2}e^{-q_\parallel z}\times\int_{-\infty}^{z}e^{q_\parallel z'}[J_{B,x}(z')+iJ_{B,z}(z')]dz'\begin{pmatrix}1\\ i\end{pmatrix}+\frac{q_\parallel}{2}e^{q_\parallel z}\times\int_{z}^{0}e^{-q_\parallel z'}[J_{B,x}(z')-iJ_{B,z}(z')]dz'\begin{pmatrix}1\\ -i\end{pmatrix}\Bigg]. \tag{C3}$$

A partial integration of the terms containing $J_{B,z}(z')$, followed by a use of the transversality condition in Eq. (72), leads to

$$\vec{E}_L^{NR}(z)=\frac{1}{i\epsilon_0\omega}\left\{J_{B,z}(z)\begin{pmatrix}0\\ 1\end{pmatrix}+\frac{i}{2}J_{B,z}(z)\left[\begin{pmatrix}1\\ i\end{pmatrix}-\begin{pmatrix}1\\ -i\end{pmatrix}\right]+\frac{i}{2}e^{q_\parallel z}J_{B,z}(0)\begin{pmatrix}1\\ -i\end{pmatrix}\right\}, \tag{C4}$$

a result one readily verifies as being identical to the one cited in Eq. (73) for $z<0$. We have thus shown that the attached field is given by Eq. (73) for all z.

2. Spacelike part of the detached field

To determine the above-mentioned part of the electric field, Eqs. (43), (70), and (75) have to be combined. Doing this, and using afterwards the transversality condition in Eq. (72) to eliminate $J_{B,x}(z')$ in favor of $(i/q_\parallel)dJ_{B,z}(z')/dz'$, one obtains for $z<0$

$$\vec{E}_{space}^T(z)=\frac{q_\parallel}{4\epsilon_0\omega}\left\{e^{-q_\parallel z}\int_{-\infty}^{z}F_+(z,z')\left[\frac{1}{q_\parallel}\frac{dJ_{B,z}(z')}{dz'}+J_{B,z}(z')\right]dz'\begin{pmatrix}1\\ i\end{pmatrix}+e^{q_\parallel z}\int_{z}^{0}F_-(z,z')\times\left[\frac{1}{q_\parallel}\frac{dJ_{B,z}(z')}{dz'}-J_{B,z}(z')\right]dz'\begin{pmatrix}1\\ -i\end{pmatrix}\right\}, \tag{C5}$$

where

$$F_\pm(z,z')=\{[1\pm iq_0(z'-z)]e^{\pm iq_0(z-z')}-1\}e^{\pm q_\parallel z'}. \tag{C6}$$

Utilizing the fact that

$$F_+(z,z')|_{z'\to-\infty}=F_\pm(z,z')|_{z'=z}=0 \tag{C7}$$

and

$$F_-(z,z')\big|_{z'=0} = (1+iq_0 z)e^{-iq_0 z} - 1, \tag{C8}$$

upon partial integrations of the terms with $dJ_{B,z}(z')/dz'$, we next get

$$\begin{aligned}\vec{E}^{T}_{space}(z) = \frac{1}{4\epsilon_0\omega}\Bigg\{&[(1+iq_0 z)e^{-iq_0 z}-1]e^{q_\| z}J_{B,z}(0)\begin{pmatrix}1\\-i\end{pmatrix}\\&-e^{-q_\| z}\int_{-\infty}^{z}\left(\frac{dF_+(z,z')}{dz'}-q_\| F_+(z,z')\right)\\&\times J_{B,z}(z')dz'\begin{pmatrix}1\\i\end{pmatrix}-e^{q_\| z}\int_{z}^{0}\left(\frac{dF_-(z,z')}{dz'}\right.\\&\left.+q_\| F_-(z,z')\right)J_{B,z}(z')dz'\begin{pmatrix}1\\-i\end{pmatrix}\Bigg\}.\end{aligned} \tag{C9}$$

By inserting the formulas

$$\frac{dF_\pm(z,z')}{dz'} \mp q_\| F_\pm(z,z') = q_0^2(z'-z)e^{\pm iq_0 z}e^{\pm(q_\| - iq_0)z'} \tag{C10}$$

into Eq. (C9), it is a straightforward matter to show that the resulting Eq. (C9) equals Eq. (76) for $z<0$.

For $z>0$, one begins from

$$\begin{aligned}\vec{E}^{T}_{space}(z) = \frac{1}{4\epsilon_0\omega}e^{-q_\| z}\int_{-\infty}^{0}&F_+(z,z')\\&\times\left(q_\| J_{B,z}(z')+\frac{dJ_{B,z}(z')}{dz'}\right)dz'\begin{pmatrix}1\\i\end{pmatrix},\end{aligned} \tag{C11}$$

cf. Eq. (C5); and a partial integration of the $dJ_{B,z}(z')/dz'$-term, followed by a use of Eq. (C10) for the plus sign, gives the result in Eq. (76) for $z>0$.

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