Dyadic Green’s function analysis of the non-stationary nanoelectrodynamic polaritonic response of a two-dimensional excitonic layer

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ABSTRACT

The non-stationary nanoelectrodynamics of polariton mode excitation and radiation in a two-dimensional excitonic layer due to a finite current pulse in the layer are analyzed in this paper. Using dyadic Green’s functions we obtain an exact result for the radiated field in closed form. This exact solution exhibits physical features of striking interest: in particular, spatial growth of the radiated inhomogeneous exciton-polariton electromagnetic plane wave field away from the nanolayer, while decaying in time. This analysis also provides insight into the resolution of an old controversy concerning differing approaches to the description of radiative modes of a two-dimensional excitonic system.

Keywords: Nanoelectrodynamics, polaritonic-radiation, excitonic-layer, dyadic-Green’s-function

I. INTRODUCTION

The dispersion properties of radiatively decaying polariton modes of a two-dimensional (2D) excitonic layer have been analyzed theoretically in a number of papers [1-9]. Three optically active exciton-polariton modes were identified: T, L and Z modes [2]. These modes are distinguished by their electric field polarization. The T mode has s (or TE) polarization, whereas the L and Z modes have p (or TM) polarization. In a symmetrical dielectric environment, L and Z modes exhibit even and odd distributions of in-plane components of the electric field across the 2D excitonic layer, respectively. Radiative exciton-polaritons have been experimentally observed in single and multiple GaAs/AlGaAs quantum-well (QW) structures [10-12]. Furthermore, an enhanced (superradiant) spontaneous radiative decay rate of free excitons in QWs was predicted [13,14] and observed in time-resolved photoluminescence (PL) experiments [15,16] in the picosecond domain.

In most of the earlier theoretical papers [1-7], a perturbative approach was employed, representing a photon quasi-continuum by a spectrum of electromagnetic (EM) eigen-modes for a cavity of infinite size. Mathematically speaking, this model is conditioned upon finiteness of the EM field at infinity, at large distance from the 2D excitonic layer. This involves the presence of both incoming and outgoing homogeneous plane EM waves at the plane of the 2D excitonic system: Moreover, resonant exciton-polaritons acquire a finite lifetime due to coupling to the quasi-continuum of photons in the infinite cavity. Within the framework of this model, it was concluded [2,6,7] that resonant polariton branches, which are essentially exciton-like having almost no dispersion, terminate at the points of their crossing with the free-photon dispersion line of the host dielectric medium. In other words, the conclusions of the earlier theories [2,6,7] can be
summarized as follows: (i) only polaritons within the light cone (i.e. those having in-plane phase velocity exceeding the speed of light in the host medium, and can be thus described as “fast polaritons”) can be optically active and (ii) the frequency of luminescence from the radiative exciton-polariton modes should be essentially independent of the angle of emission.

Naturally, a perturbative approach is fairly reasonable for treating systems with relatively weak exciton-photon coupling, but it may fail in application to systems with strong exciton-photon coupling. In addition, the above theoretical model involving both incoming and outgoing EM waves is quite appropriate for analyzing light transmission/reflection, but it seems inconsistent to apply it to the problem of decaying spontaneous emission, which, for example, occurs in time-resolved PL experiments [15,16] involving only outgoing EM waves.

In other theoretical papers [8,9] a self-consistent elecrodynamical approach was employed, representing the radiation field in the host medium in terms of a single outgoing inhomogeneous plane EM wave that satisfies the scattering condition, requiring that only outgoing waves exist at infinity far from the 2D excitonic system. It was found that in this approach [8] the field of every radiative exciton-polariton mode grows in space away from the excitonic layer, while decaying in time. It was shown in [8] that the radiative-polariton dispersion branches enter the slow-wave region of the dispersion plane (i.e., escape the light cone) with the radiative decay rate of the slow-polariton mode vanishing at the terminal point of the dispersion branch. This approach was applied in the study of exciton-polaritons in PbI-based self-organized QW structures in [8,9], where it was shown that the radiative polaritons may exhibit substantial frequency dependence on in-plane wavevector (leading to angular dependence of the emission frequency). The dispersion of the radiative exciton-polaritons in the vicinity of the light cone appeared to be much stronger than might be expected from perturbative considerations.

In fact, both above mentioned approaches are well justified from a mathematical point of view. They differ by the condition chosen at infinity, far from the 2D excitonic layer. However, in real physical situations the exact behavior of the EM field at infinity is not dominantly important because the field cannot reach infinitely distant points during finite observation time. A more revealing approach is needed to study exciton-polariton mode excitation in a 2D excitonic system, which embeds consideration of initial conditions. In this paper, we analyze a non-stationary process of excitation of the exciton-polariton modes in a 2D excitonic layer due to in-layer current pulses. We employ EM dyadic Green’s functions for this purpose.

In Sec. II, the dyadic Green’s function of the structure with a 2D excitonic layer embedded in a host medium is constructed. In Sec. III, we use the Green’s function to obtain the EM response of the structure to current pulses of two different forms, obtaining exact solutions. The main conclusions are summarized in Sec. IV.

II. NANOELECTRODYNAMIC ANALYSIS BY DYADIC GREEN’S FUNCTION FOR A 2D EXCITONIC LAYER

We consider a 2D excitonic layer located at the plane $z = 0$, embedded in a three-dimensional (3D) host medium with dielectric constant $\varepsilon_h$. The response of the excitonic material is described by a local dielectric function of the Lorentzian type [8]

$$\varepsilon(\omega) = \varepsilon_0 \left( 1 + \frac{\omega_p^2}{\omega_p^2 - \omega^2 - i \gamma \omega} \right),$$

(1)
where $\varepsilon_b$ is the background dielectric constant, $\omega_{ex}$ is the exciton frequency, $\omega_{LT}$ is the longitudinal-transverse splitting, and $\gamma_{ex} > 0$ describes exciton damping due to all relaxation processes except radiative decay. Strictly speaking, the dielectric response of excitons in QW’s may be both anisotropic and non-local, described by a non-local dielectric response dyad (tensor) quite generally [3,4,7]. However, the phenomena analyzed here are prominent only in 2D excitonic layers with sufficiently large LT splitting, which leads to a very short radiative lifetime for excitons-polaritons [8]. For example, the LT splitting of excitons in PbI-based self-organized QW structures [17,18] is two orders of magnitude as great as in GaAs/AlGaAs QW structures. Since the excitons in PbI-based QW structures are largely bound excitations, the local scalar dielectric function (1) is an adequate approximation for excitonic response in such structures [8,9,18].

The Helmholtz equation for the electric field $E$, derived from the Maxwell equations, is driven by the current density, $J$, which is the sum of induced, $J_{ind}(r,t) = \int dt' \int drr' \hat{\sigma}(r,r',t,t')E(r',t')$, and external, $J_{ext}$, parts. Here, $\hat{\sigma}$ is the conductivity tensor (nonlocal and spatially inhomogeneous in general case). For the problem at hand, $\hat{\sigma}$ can be represented as the sum of the bulk conductivity tensor of the host medium, which we take simply as $\hat{\sigma}_{3D} = -i\omega \chi_b \mathbf{I}$ ($\chi_b$ is the dielectric susceptibility of the host medium with dielectric constant $\varepsilon_h = 1 + 4\pi\chi_b$) and the conductivity tensor of the 2D excitonic layer, $\hat{\sigma}_{2D}$. In this case, the dyadic Helmholtz equation can be written in temporal Fourier representation ($t \to \omega$) as

$$
\left[ \hat{\mathbf{I}} \left( \nabla^2 + \frac{\omega^2\varepsilon_h}{c^2} \right) - \nabla \nabla \right] E(r,\omega) + \frac{4\pi\omega}{c^2} \int drr' \hat{\sigma}_{2D}(r,r',\omega)E(r',\omega) = -\frac{4\pi\omega}{c^2} J_{ext}(r,\omega),
$$

(2)

where $\hat{\mathbf{I}}$ is the unit dyad, and $c$ is the speed of light in vacuum.

The description of electrodynamics in terms of dyadic (tensor) Green’s functions has a long history in physics and in electrical engineering [19,20]. The present interest in the electrodynamics of low dimensional semiconductor nanostructures [21, for example], which involve currents that are geometrically confined in narrow regions, offers a fertile ground for the application of dyadic Green’s functions, $\hat{\mathbf{G}}(r,r',t,t')$, as defined by

$$
\left[ \hat{\mathbf{I}} \left( \nabla^2 + \frac{\omega^2\varepsilon_h}{c^2} \right) - \nabla \nabla \right] \hat{\mathbf{G}}(r,r',t,t') - \frac{4\pi\omega}{c^2} \int drr' \hat{\sigma}_{2D}(r,r',\omega)\hat{\mathbf{G}}(r',r',t,t') = \hat{\mathbf{I}} \delta(r-r';t,t').
$$

(3)

Employing a position-space matrix notation with frequency suppressed, $\hat{\mathbf{G}}(r,r';t,t') \equiv \langle r|\hat{\mathbf{G}}|r' \rangle$, we define

$$
(\hat{\mathbf{G}}_0)^{-1} = \hat{\mathbf{I}} \left( \nabla^2 + \frac{\omega^2}{c^2} \right) - \nabla \nabla,
$$

and further define

$$
\hat{\mathbf{G}}_{3D}^{-1} = (\hat{\mathbf{G}}_0)^{-1} + \frac{4\pi\omega}{c^2} \hat{\sigma}_{3D},
$$

(4)
with the resulting integral equation for $\hat{G}$ given by [19,20]

$$
\hat{G}(r,r',\omega) = \hat{G}_{3D}(r,r';\omega) - \frac{4\pi i \omega}{c^2} \int dr'' \int dr'''' \hat{G}_{3D}(\omega, r, r'') \hat{\sigma}_{2D}(r'', r'''; \omega) \hat{G}(r'', r'''; \omega),
$$

(5)

where $\hat{G}_{3D}(r,r';\omega)$ is the dyadic Green’s function of bulk 3D homogeneous space with dielectric constant $\epsilon_h$. In spatial Fourier representation along the layer, $r \rightarrow \mathbf{k}$, Eq. (5) becomes

$$
\hat{G}(z,z',\mathbf{k}_z,\omega) = \hat{G}_{3D}(z,z',\mathbf{k}_z,\omega)
- \frac{4\pi i \omega}{c^2} \int dz'' \int dz''' \hat{G}_{3D}(z,z'',\mathbf{k}_z,\omega) \hat{\sigma}_{2D}(z'', z''', \mathbf{k}_z, \omega) \hat{G}(z'', z''', \mathbf{k}_z, \omega).
$$

(6)

In this analysis we describe the response of the narrow, uniform 2D excitonic layer by the local conductivity tensor $\hat{\sigma}_{2D} = \hat{I} \sigma_{2D}(\omega) \delta(z) \delta(z')$ confined to the 2D sheet, where $\sigma_{2D} = i \omega \epsilon_h - \epsilon(\omega)d/4\pi$, and $d$ is the thickness of the excitonic layer. Then Eq. (6) readily reduces to

$$
\hat{G}(z,z',\mathbf{k}_z,\omega) = \hat{G}_{3D}(z,z',\mathbf{k}_z,\omega) - \frac{4\pi i \omega}{c^2} \hat{G}_{3D}(z,0,\mathbf{k}_z,\omega) \hat{\sigma}_{2D}(\mathbf{k}_z, \omega) \hat{G}(0, z', \mathbf{k}_z, \omega).
$$

(7)

This reduction illustrates the utility of the dyadic Green’s function integral equation formulation for narrow nanoscale electrodynamic analyses. To determine $\hat{G}(z,z',\mathbf{k}_z,\omega)$ we note that we can express $\hat{G}(0, z', \omega, \mathbf{k}_z)$ in terms of $\hat{G}_{3D}$ by setting $z$ to zero in Eq. (7) and solving algebraically. Substituting the result in the right side of Eq. (7), an exact solution for the dyadic Green’s function of the whole structure is obtained as follows:

$$
\hat{G}(z,z',\mathbf{k}_z,\omega) = \hat{G}_{3D}(z,z',\mathbf{k}_z,\omega) - \frac{4\pi i \omega}{c^2} \hat{G}_{3D}(z,0,\mathbf{k}_z,\omega) \hat{\sigma}_{2D}(\mathbf{k}_z, \omega) \times
\left[ \hat{I} + \frac{4\pi i \omega}{c^2} \hat{G}_{3D}(0,0,\mathbf{k}_z,\omega) \hat{\sigma}_{2D}(\mathbf{k}_z, \omega) \right]^{-1} \hat{G}_{3D}(0, z', \mathbf{k}_z, \omega).
$$

(8)

Dyad $\hat{G}_{3D}(z,z',\mathbf{k}_z,\omega)$ is obtained from the Fourier transform of the well-known bulk dyadic Green’s function for 3D homogeneous space with dielectric constant $\epsilon_h$ [19,20],

$$
\hat{G}_{3D}(\mathbf{k}, \omega) = \frac{\hat{I} - \mathbf{k}\mathbf{k}/k^2}{\omega^2 \epsilon_k/c^2 - k^2} + \frac{\mathbf{k}\mathbf{k}/k^2}{\omega^2 \epsilon_h/c^2},
$$
by integrating it over \( k_z \), whence

\[
\hat{G}_{3D}(z, z', k_1, \omega) = 
\frac{1}{2i k_\perp} \left\{ \hat{I} - \frac{c^2}{\omega^2 \epsilon_h} \left[ k_1 k_\parallel + (k_1 e_\epsilon + e_\epsilon k_\parallel) \frac{1}{i} \frac{\partial}{\partial z} + e_\epsilon e_\epsilon \left( \frac{1}{i} \frac{\partial}{\partial z} \right)^2 \right] \right\} \exp(ik_\perp |z - z'|),
\]

(9)

where \( k_\perp = \sqrt{(\omega^2/c^2)\epsilon_h - k_1^2} \). The sign before the radical is chosen in such a way that the field in the host medium satisfies the radiation condition for \( k_1 < \omega/\sqrt{\epsilon_h/c} \) and the evanescent field condition for \( k_1 > \omega/\sqrt{\epsilon_h/c} \).

**III. RESPONSE OF THE SYSTEM TO A CURRENT PULSE**

Addressing the linear electromagnetic response of the system to an applied current pulse, the excited electric field can be written in position-time representation using Eqns. (2) and (3) (with linear superposition) as

\[
E(\mathbf{r}, t) = \frac{4\pi}{c^2} \int_{-\infty}^{\infty} dt' \int d\mathbf{r}' \hat{G}(\mathbf{r}, \mathbf{r}'; t, t') \frac{\partial \mathbf{J}_{ext}(r', t')}{\partial t'},
\]

where \( \mathbf{J}_{ext}(r', t') \) is the external current density applied in the layer and \( \hat{G}(\mathbf{r}, \mathbf{r}'; t, t') \) is the dyadic Green’s function of the structure. Alternatively, in Fourier representation we have \( (t - t' \to \omega, \mathbf{r} - \mathbf{r}' \to \mathbf{k}_1) \)

\[
E(z, k_1, \omega) = -\frac{4\pi \omega}{c^2} \int dz' \hat{G}(z, z', k_1, \omega) \mathbf{J}_{ext}(z', k_1, \omega),
\]

(10)

where \( \hat{G}(z, z', k_1, \omega) \) is given by Eq. (8) for the problem at hand. We will employ this to obtain the electric field induced by external current pulses of two different forms. In both cases we consider that the pulses, as well as the dielectric properties on the layer are spatially homogeneous. (We choose \( \epsilon_h = \epsilon_b \) for simplicity.)

**A. Current \( \delta(t) \)-pulse**

To start, we assume a current pulse of the form \( \mathbf{J}_{ext}(\mathbf{r}, t) = \mathbf{J}_0 \delta(t) \delta(z) \) localized on the plane of the excitonic layer as well as in time, so that \( \mathbf{J}_{ext}(z, k_1, \omega) = (2\pi)^3 \mathbf{J}_0 \delta(z) \) in Fourier representation, where \( \mathbf{J}_0 \) is the normalizing amplitude of the sheet current pulse. Performing the integration of Eq. (10), we obtain the induced electric field as

\[
E_\delta(z, k_1, \omega) = -(2\pi) \frac{4\pi \omega}{c^2} \hat{G}(z, 0, k_1, \omega) \mathbf{J}_0 \delta(z)
\]

(11)
where we write $E \to E_d$ for the field responding to the $\delta(t)$-current pulse. This yields

$$E_d(z, r, \omega) = -\frac{4\pi \omega}{c^2} J_0 \int d^2 k \exp(i k \cdot \mathbf{r}) \tilde{G}(z, 0, \mathbf{k}_1, \omega) \delta(k)$$

in position-frequency representation. Due to spatial homogeneity in the layer plane, the $k$-integral of Eq. (12) involves only in-plane diagonal elements of the dyadic Green’s function (which are equal to each other due to the symmetry of the problem in the plane of the 2D excitonic layer). Using Eq. (9), one readily finds the result given by

$$E(z, \omega) = -\frac{4\pi \omega}{c^2} J_0 \left[ -\frac{c}{2i \omega \sqrt{\varepsilon_b}} - \frac{\omega_l \varepsilon_{\omega}^{(L)} d}{4(\omega - \tilde{\omega}^{(L)})} \left( 1 - \frac{1}{\omega - \tilde{\omega}^{(L)}} \right) \right] \exp\left( i \frac{\omega}{c} \sqrt{\varepsilon_b} z \right). \quad (13)$$

The first term in square brackets describes the 3D host medium contribution and the second one describes a contribution of the 2D excitonic layer. The electric field $E_d(z, \omega)$ has a pole at the complex frequency

$$\tilde{\omega}^{(L)} = \omega^{(L)} + i \gamma^{(L)} = \frac{\omega_l - i \gamma_\omega^{(L)}}{1 + i \sqrt{\varepsilon_b} \omega_l^{(L)} d}$$

which exactly satisfies the dispersion relation for the L-polariton mode (the L-mode is degenerate with the T-mode in the case $k_1 = 0$) obtained in [8]. The real part of the complex frequency, $\omega^{(L)}$, is the polariton-mode eigen-frequency and the imaginary part, $\gamma^{(L)} = \gamma_d^{(L)} + \gamma_r^{(L)}$, is the total decay rate of the mode, which involves a dissipative part

$$\gamma_d^{(L)} = -\frac{\gamma^{(L)}_\omega}{1 + \varepsilon_b \left( \omega_l^{(L)} d \right)^2},$$

as well as a radiative part,

$$\gamma_r^{(L)} = -\frac{1}{2c} \frac{\omega_l^{(L)} \sqrt{\varepsilon_b} d}{1 + \varepsilon_b \left( \omega_l^{(L)} d \right)^2}.$$

In position-time representation we obtain the electric field as
\[ E_e(z,t) = \frac{1}{2\pi} \int d\omega \exp(-i\omega t) E(z,\omega) \]

\[ = J_0 \left[ -\frac{1}{c \sqrt{E_b}} + \frac{i\omega_{zL,\omega} d}{2c^2(\omega_{\alpha,\gamma} - i\gamma_{\alpha})} + \frac{i\omega_{xL,\omega} (\tilde{\omega}^{(L)})}{2c^2(\omega_{\alpha,\gamma} - i\gamma_{\alpha})} - \frac{1}{\omega - \tilde{\omega}^{(L)}} \right], \] (15)

where \( t' = t - \sqrt{E_b} |z|/c \). Integration of the first and second terms in braces multiplied by \( \exp(-i\omega t) \) straightforwardly yields \( \delta(t') \). The third term in braces has a pole at \( \omega = \tilde{\omega}^{(L)} \) and we carry out the integration for \( t' < 0 \) closing the contour of integration in Eq. (15) in the upper half \( \omega \)-plane finding a null result as there are no singularities within the closed contour. For \( t' > 0 \), we close the contour of integration in the lower half \( \omega \)-plane and obtain the induced electric field responding to the \( \delta(t') \) current pulse by residues as

\[ E_{\delta}(z,t) = J_0 \left[ -\delta(t') \frac{2\pi}{c \sqrt{E_b}} + \delta(t') i \frac{\pi \tilde{\omega}^{(L)} \omega_{LT,\omega} d}{c^2(\omega_{\alpha,\gamma} - i\gamma_{\alpha})} + \frac{\pi (\tilde{\omega}^{(L)})^2 \omega_{LT,\omega} d}{c^2(\omega_{\alpha,\gamma} - i\gamma_{\alpha})} \exp(-i\tilde{\omega}^{(L)} t') \theta(t') \right]. \] (16)

The first term in the braces describes the instantaneous transient retarded EM field excited in the host medium by the current pulse. The second term in braces describes the contribution of the 2D excitonic layer to the instantaneous transient retarded field, which is shifted in phase (by \( \pi/2 \) at zero dissipation, \( \gamma_{\alpha} = 0 \)) with respect to the current pulse. The third term in braces exhibits the radiative inhomogeneous polariton L-mode emerging from the excitonic layer. The field of this mode decays in time with the total decay rate of the exciton-polariton mode once the current pulse has elapsed, while it grows in space \([22,23]\) away from the excitonic layer. This field vanishes for \( |z| > t c/\sqrt{E} \) in accordance with the causality principle. The peculiarities of such a non-stationary process are illustrated in Fig.1, where the amplitude distribution of the radiatively decaying exciton-polariton mode is shown in the space-time domain.

Note that the contribution of the 2D excitonic layer to the induced electric field (Eq. 16) is proportional to the \( \omega_{LT,\omega}/\omega_{ex,\omega} \)-ratio and, hence, it is much more pronounced in PbI-based self-organized QW structures (\( \omega_{LT} \approx 50 \text{ meV} [8,9,17,18] \)) than in GaAs/AlGaAs QW structures (\( \omega_{LT} \approx 0.5 \text{ meV} \)).

**B. Current pulse of a rectangular form**

Considering a rectangular current pulse of duration \( \tau \) in the plane of the 2D excitonic layer,

\[ J_{ex}(\mathbf{r},t) = J_0 [\theta(t) - \theta(t - \tau)] \delta(z) / \tau, \]

we have

\[ J_{ex}(z,\mathbf{k},\omega) = J_0 (2\pi)^2 \mathcal{I} \delta(z) \delta^\prime(\mathbf{k}) \left[ |\tau| \right] \exp(-i\omega\tau) \] in Fourier representation, with \( J_0 \) as the normalizing amplitude of the sheet current pulse. Employing Eq. (10) and the results above, we have

\[ E(z,t) = \frac{J_0}{\tau} \int_{-\infty}^{\infty} d\omega \left[ \frac{1}{ic \sqrt{E_b}} \frac{\exp(-i\omega t')}{\omega + i\tau} - \frac{1}{ic \sqrt{E_b}} \frac{\exp(-i\omega t')}{\omega + i\tau} \right]. \]
where \( t' = t - \sqrt{\epsilon_b} |z| / c \). All terms in braces in Eq. (17) have poles located in the lower half of the complex \( \omega \)-plane. For \( t' < 0 \) we again close the contour of integration in the upper-half \( \omega \)-plane and integration of the first and third terms in braces yields zero. For \( t' > 0 \) the integration contour is closed in the lower-half \( \omega \)-plane and integration of these two terms by residues yields finite results. Integration of the second and fourth terms in braces similarly vanish for \( t' - \tau < 0 \), and yields finite contributions by residues for \( t' - \tau > 0 \). The final result is given by

\[
E(z, t) = 2\pi i \mathbf{j}_b \left\{ \frac{\theta(t - t') - \theta(-t')}{c \sqrt{\epsilon_b}} - \frac{i \omega^{(l)} \omega^{(l)} d}{2(\omega_{ex} - i \gamma_{ex}) c^2} \left[ \theta(t' - \tau) \exp(i \omega^{(l)} \tau) - \theta(t') \right] \exp(-i \omega^{(l)} t') \right\},
\]

(18)

Here, the first term in braces describes the transient retarded EM field excited by the pulse in the host medium, which is non-zero only within the time duration of the current pulse. The second term in braces describes the radiatively decaying inhomogeneous EM field of the L-radiative-polariton mode excited by the time variations at the beginning and end of the current pulse.

### IV. CONCLUSIONS

We have examined the non-stationary electrodynamics of a 2D exciton layer in which polariton modes are excited by current pulses in the layer. This nanostructure analysis was facilitated by the use of dyadic Green’s functions, which are particularly well suited to the description of low-dimensional systems, and exact, explicit solutions for the radiated electric field due to \( \delta(t) \)- and square-wave current pulses have been obtained. These results exhibit physical features of striking interest, namely spatial growth of the radiated inhomogeneous exciton-polariton electromagnetic plane wave field away from the nanolayer, while decaying in time.

In regard to the old controversy concerning differing approaches to the description of radiative modes of a 2D excitonic system, the present work shows that radiative decay of free polariton modes of 2D exciton layer occurs via emitting the inhomogeneous EM plane waves, which accounts for peculiar dispersion and spatio-temporal properties of these excitations.

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Fig. 1: Spatio-temporal distribution of the amplitude of a radiatively decaying exciton-polariton mode in a 2D excitonic layer. The 2D excitonic layer is at $z=0$. 

$|z| = tc/\sqrt{\varepsilon_b}$