Poynting’s theorem and energy conservation in the propagation of light in bounded media

F. Richter(a), M. Florian and K. Henneberger

Universität Rostock, Institut für Physik - 18051 Rostock, Germany

received 18 December 2007; accepted 2 February 2008
published online 29 February 2008

PACS 78.20.-e – Optical properties of bulk materials and thin films
PACS 03.50.De – Classical electromagnetism, Maxwell equations
PACS 71.35.Cc – Intrinsic properties of excitons; optical absorption spectra

Abstract – Starting from the Maxwell-Lorentz equations, Poynting’s theorem is reconsidered. The electromagnetic energy flux vector is introduced such that it can be related to the kinetic energy of the matter subsystem. Conservation of the total energy follows. In our discussion, the microscopic nature of media is represented exactly by susceptibility functions, which do not necessarily have to be known. On this footing, it can be shown that energy conservation in the propagation of light through bounded media is ensured by Maxwell’s boundary conditions alone, even for some frequently used approximations. This is demonstrated for approaches using additional boundary conditions and the dielectric approximation in detail, the latter of which suspected to violate energy conservation for decades.

Copyright © EPLA, 2008

Introduction. – The effects of a medium on radiation (e.g., in an experiment on the transmission and reflection of light in a dielectric medium) are determined exactly by the susceptibility function \( \hat{\chi}(r, r', t, t'; E) \), which enters Maxwell’s equations through the material equations (also known as constitutive relations). Due to its complexity, it is only in rare cases that \( \hat{\chi} \) can be expressed analytically or valuable relations can be drawn from it. For a long time, the more complex properties of \( \hat{\chi} \) have not been of much interest, and textbooks still introduce the susceptibility as a constant scalar or as a constant tensor, to account at least for anisotropy (e.g., [1,2]). Non-linearity (i.e., dependence on \( E \)) is an own interesting field of research and technology but out of scope here and will be neglected.

But strictly speaking, the presence of a medium boundary alone enforces a non-local and spatially inhomogeneous \( \hat{\chi} \), if spatial dependence in the dispersion relation (\textit{spatial dispersion}) is not explicitly neglected. In crystal and semiconductor optics, some approaches to this problem resort to using the susceptibility of an infinite medium for the description of a bounded real one. Then, multiple polariton modes are predicted to propagate concurrently in the medium, and the precondition of a continuous transition of the electromagnetic fields at the medium boundaries (\textit{Maxwell’s boundary conditions}, MBCs) is no longer sufficient to determine the amplitudes and phases of these modes. As a solution, different phenomenological \textit{additional boundary conditions} (ABCs) were proposed [3,4], the first of which in 1957. In 1973, a related approach was published [5], suggesting to limit the bulk susceptibility to the bounded medium by step functions (\textit{dielectric approximation}, DA). These methods were applied in a vast number of publications in the past decades, and they still are used today and successfully explain experimental data [6].

However, it is known that some of these approximations can conflict with the fundamental law of energy conservation [7,8]. Thus, scientists were prompted to take greatest care in their argumentation or even to reject these approximations despite their successes (see, e.g., [8,9]).

We will show that, in the case of the DA and other important approximations, energy conservation is not a question of the approximation but rather is ensured by the MBCs alone. The DA is addressed in detail: The compatibility of the approximation with energy conservation is proven, and problems in the respective work are pointed out. For all this, we use an approach that one of us recently developed [10], which provides analytically exact relations between radiation and matter in the elusive case of spatial dispersion with the aid of Keldysh’s photon Green’s functions.

Throughout the discussion of bounded media, we regard arbitrary steadily excited media. They are represented by their susceptibility functions, which do not necessarily
have to be known. Besides these latter, the cornerstone of our discussion is *Poynting’s theorem* (PT), the equivalent of the energy conservation law in electrodynamics. It balances the time derivative of electromagnetic and mechanical energy $U_{e,m}$ and sources of respective energy flux $S_{e,m}$:

$$\frac{\partial}{\partial t} (U_e + U_m) + \text{div} (S_e + S_m) = 0. \quad (1)$$

In Poynting’s original work [11] and many established textbooks [1,2,12], the theorem is given —without regarding spatial dispersion— as

$$\text{div} E \mathbf{E} + \text{div} H \mathbf{B} = -j_{\text{ext}} \mathbf{E}, \quad S_e = \text{div} (\mathbf{E} \times \mathbf{H}). \quad (2)$$

A corresponding expression for the electromagnetic energy density $U_e$ can only be given with further assumptions. This is why eq. (2) cannot be generally interpreted as an energy continuity equation and, consequently, $\text{div} (\mathbf{E} \times \mathbf{H})$ cannot be considered as an energy flux vector.

In order to lay solid grounds for our considerations, we will at first give a thorough derivation of the theorem (1) in conjunction with spatial dispersion and mechanical energy, giving hints on why several articles, such as [13–16], or those referenced therein, doubt the validity and form of eq. (2) in more complex situations.

A note on terminology: Historically, the term *Poynting vector* denoted the electromagnetic energy flux $S_e$ only. With regard to eq. (1), we here need to distinguish clearly between the *electromagnetic* and the *mechanical Poynting vector*.

**A clean-room derivation of Poynting’s theorem.**

There are many doubts expressed on under which conditions (static magnetic fields, stationary situations, dispersive media) PT applies and on how Poynting’s energy flux vector has to be interpreted (compare [13–16]). Here, we do not see any reason for these. Instead, we will give a brief derivation of PT, while carefully accounting for any approximations made or conditions implied. This derivation is done in classical physics but can easily be extended to quantum physics by applying the well-known recipes (*i.e.*, operators, symmetrisation, commutation relations).

Surprisingly, we do not need any assumptions about the matter. Avoiding constitutive relations and derived fields, we rather rely on the *microscopic* Maxwell-Lorentz equations [17], which are the most fundamental and universal:

$$\text{rot} \mathbf{E} = -\dot{\mathbf{B}}, \quad \text{rot} \mathbf{B} = \mu_0 \dot{\mathbf{j}}, \quad \mu_0 \frac{1}{c^2} \dot{\mathbf{E}}, \quad (3a)$$

$$j(r,t) = \sum_{\alpha} e_{\alpha} \dot{r}_{\alpha} \delta(r - r_{\alpha}(t)), \quad (3b)$$

with $\mathbf{j}$ as a current of charged particles $\alpha$.

With the vector identity $\text{div}(\mathbf{X} \times \mathbf{Y}) = \mathbf{Y} \cdot \text{rot} \mathbf{X} - \mathbf{X} \cdot \text{rot} \mathbf{Y}$, we immediately have

$$\varepsilon_0 \dot{\mathbf{E}} \mathbf{E} + \frac{1}{\mu_0} \dot{\mathbf{B}} \mathbf{B} + \frac{1}{\mu_0} \text{div}(\mathbf{E} \times \mathbf{B}) = -j \mathbf{E} \quad (4)$$

and may identify the electromagnetic field energy density and the electromagnetic Poynting vector as

$$U_e = \frac{1}{2} \left( \varepsilon_0 \varepsilon_0^2 + \frac{1}{\mu_0} \mathbf{B}^2 \right), \quad (5)$$

$$S_e = \frac{1}{\mu_0} (\mathbf{E} \times \mathbf{B}). \quad (6)$$

Note that in what follows, we will use these new definitions of the two quantities instead of the more conventional given in eq. (2), if not otherwise stated.

Now, a microscopic mechanical expression for $j \mathbf{E}$ shall be developed. We use the density of mechanical (kinetic) energy of the particles $\alpha$, $U_m$, as well as their kinetic energy flux, $S_m$, according to

$$U_m(r,t) = \sum_{\alpha} \frac{m_\alpha}{2} \dot{r}_{\alpha}^2 \delta(r - r_{\alpha}(t)), \quad (7)$$

$$S_m(r,t) = \sum_{\alpha} \frac{m_\alpha}{2} \dot{r}_{\alpha} \dot{r}_{\alpha} \delta(r - r_{\alpha}(t)). \quad (8)$$

The force $m \dot{r}$ appearing in the time derivative of $U_m$ is just the Lorentz force $F = e(\mathbf{E} + \hat{r} \times \mathbf{B})$, yielding

$$\frac{\partial}{\partial t} U_m(r,t) + \text{div} S_m(r,t) = j(r,t) \mathbf{E}(r,t). \quad (9)$$

This can finally be combined with eq. (4) to Poynting’s theorem as shown in eq. (1). No restrictions have been made to the behaviour of the fields or particles or the nature of the latter, except for non-relativistic velocities. This approach will turn out to be inevitable for establishing a relation between electromagnetic and mechanical energy. Mechanical energy here is purely kinetic, any potential energy is attributed to the fields.

The continuity of the energy flux at medium boundaries is an aspect to which much attention is paid in the literature. The electromagnetic Poynting vector $S_e$ derived here is not necessarily continuous, but this does not account for physical meaningfulness. It is the energy continuity eq. (1) that holds. The continuity of one of the quantities involved depends on that of the others, which in turn is subject to the system and model considered.

**Material equations and susceptibility.**

The macroscopic Maxwell equations originally were introduced to obtain a global, effective description of electrodynamics in the presence of a medium. One may consider their form a well-founded historic convention, which can be reproduced simply by splitting $j = j_{\text{ind}} + j_{\text{ext}}$ and the ansatz

$$j_{\text{ind}} = \dot{\mathbf{P}} + \text{rot} \mathbf{M}, \quad (10)$$

---

1. Landau/Lifshitz [2] only give the differential $dU_e = E \text{div} D + H \text{div} B$; Born and Wolf [1] give $U_e = (\varepsilon_0 E^2 + \mu_0 B^2)/2$, valid for material equations (12); Jackson [12] gives $U_e = (E \text{div} D + H \text{div} B)/2$ for the multipole approach.
which is also suggested [18] by the well-known multi-
pole expansion of the electromagnetic properties of the medium [12], and the respective material equations
\[
D(r, t) = \varepsilon_0 E(r, t) + P(r, t),
\]
(11a)
\[
H(r, t) = \frac{1}{\mu_0} B(r, t) - M(r, t),
\]
(11b)
\[
P(r, t) = \varepsilon_0 \int d^3 r' d' t' \chi_{\text{elec}}(r, r', t, t') E(r', t'),
\]
(11c)
\[
M(r, t) = \frac{1}{\mu_0} \int d^3 r' d' t' \chi_{\text{magn}}(r, r', t, t') B(r', t').
\]
(11d)

Note the clear analogy in this notation and the obvious distinction of fields \(E, B\) and derived fields \(D, H\). Usually, textbooks [1,2,18] focus more on the analogy
\[
B = \mu_0 \varepsilon_0 H, \quad D = \varepsilon_0 \varepsilon_0 E,
\]
(12)
i.e., with constant tensorial relations (11c), (11d), and consequently treat \(E\) and \(H\) on the same level. But the physically relevant fields are actually \(E\) and \(B\), because they determine, e.g., the Lorentz force, the energy density, and the electromagnetic energy flux.

Neither the multipole expansion discussion in [12] nor the tensorial descriptions in [1,2] end up with non-local or inhomogeneous material equations and thus cannot be used with spatial dispersion and bounded media. Hence, a more general dependence of the derived fields on \(E\) and \(B\), as shown in eqs. (11), needs to be used. Here, the tensorial susceptibilities \(\chi_{\text{elec}}\), \(\chi_{\text{magn}}\) represent exactly the linear properties of an arbitrarily complex stationary medium. However, PT can still be derived in the same manner without any restrictions to the electromagnetic fields, properties of the medium or mathematical properties of \(\chi_{\text{elec}}\) or \(\chi_{\text{magn}}\), except for the temporal homogeneity of the latter. Also, eq. (2) can still be obtained, but then \(j\) does no longer fulfill eq. (3b) and thus cannot be coupled with the mechanical-energy continuity relation (9), and the field energy density cannot be given.

In optics, the approximation \(\chi_{\text{magn}} = 0\) is widely accepted. Then, \(S_e\) becomes \(E \times H\) and \(j_{\text{ind}} = \hat{P}\). From a higher level of abstraction one can also postulate that \(j_{\text{ind}} = \hat{P}\). This avoids pinpointing magnetisation and other difficulties and results in the following material equations [19]:
\[
D(r, t) = \varepsilon_0 E(r, t) + P(r, t),
\]
(13a)
\[
P(r, t) = \varepsilon_0 \int d^3 r' d' t' \hat{\chi}(r, r', t, t') E(r', t'),
\]
(13b)
\[
H(r, t) = \frac{1}{\mu_0} B(r, t).
\]
(13c)

Now, the \(H\) field is merely a rescaled \(B\) field and any electromagnetic effects are contained in \(P\).

Much of the controversy around the validity of Poynting’s theorem may be due to inconsistent formulation of the material equations or using eq. (2) despite magnetisation. In [13], the author correctly finds that eq. (2) together with the common approximation (12) for the material equations does not apply for wide categories of interactions, such as spatial dispersion. He also bases his further considerations on the Maxwell-Lorentz equations and equally finds eq. (4) but then develops an elaborate expression for \(\varepsilon E\) within the (local) multipole expansion approximation. The result, of course, is an expression for our \(S_m\) in eq. (1) (compare eq. (21) in [13]). This is an honorable effort, even though his results can only hold true in the limits of this approximation. Unfortunately, the author does not identify \((E \times B)/\mu_0\) as a valid electromagnetic Poynting vector and eq. (1) as a more general energy conservation law.

The work of Bishop/Maradudin and its conceptual problems. – In [7], Bishop and Maradudin analyse the energy flow in a semi-infinite spatially dispersive dielectric (i.e., half-space geometry) on which light is incident from the vacuum. The energy flux in the vacuum is given by the electromagnetic Poynting vector \(S_e\) only. In the medium, the mechanical Poynting vector \(S_m\) appears additionally. They state that the DA fails to conserve energy by showing that \(S_m\) does not evolve continuously from zero at the surface, as if there was an energy source. For this purpose, they develop a model of the transport of mechanical energy in a special system, a diatonic cubic crystal, on the basis of fixed, undamped and uncoupled harmonic oscillators only, then derive PT from Maxwell’s equations and adapt it for their model. In [7], eq. (4.7), they postulate that \(S_m\) be represented by the oscillations; the spatial restriction to the medium is forced by an appropriate step function.

This postulate does not necessarily fulfill the continuity condition. \(S_m\) is continuous only if the oscillator amplitude vanishes at the surface. One of the possibilities to ensure this is Pekar’s ABC [3], as the authors confirm in their conclusion. It is no surprise that the step function leads to a delta-shaped term [7], eq. (4.10), when conducting the divergence operation in eq. (1). Of course, such a source term is not compatible with energy conservation. Consequently, the authors generalise their model and derive additional conditions for its free parameters.

This kind of approach is highly doubtful because any predictions might have to be attributed to the model and its deficiencies. Here, this is clearly the case: The mechanical Poynting vector as postulated does neither contain all mechanical-energy transport in the medium (e.g., via phonons, center-of-mass motion of excitons, etc.) nor does it fulfill itself the condition to be analysed.

Compatibility of DA and PT. – We rather suggest an approach free of models. It includes but is not limited to the case of DA; it does not make any assumptions on the microscopic nature of the medium but instead it is based on the susceptibility \(\chi\) as an exact but unknown representation of the latter.
We regard a stationary isotropic medium in the slab geometry. Light is incident perpendicular to the medium surface. Then Poynting’s theorem reduces to:

$$\frac{\partial S_m(x,t)}{\partial x} = -\frac{\partial S_m(x,t)}{\partial x} = -j(x,t)E(x,t) = -W(x,t).$$

By expressing $E$ and $B$ through the vector potential $A(x,t)$, $j$ through $\partial P/\partial t$ and Fourier-transforming $t \rightarrow \omega$, we obtain

$$W(x,\omega) = i\varepsilon_0 \int \frac{d\omega'}{2\pi} \omega'^2(\omega - \omega') \cdot \int dx' \chi(x,x',\omega')A(x',\omega')A(x,\omega - \omega').$$

(15)

There is no need for taking a time average. As $\chi = 0$ in vacuum, it may be replaced by $\chi(x,x',\omega) = \Theta(\frac{L}{2} - |x|)\Theta(\frac{L}{2} - |x'|)\chi(x,x',\omega)$, where $L$ is the length of the slab. Of course, this is also valid for the special case of DA, where $\chi$ just simplifies to $\Theta(\frac{L}{2} - |x|)\Theta(\frac{L}{2} - |x'|)\chi(x-x',\omega)$.

Evaluating the step functions, we have:

$$W(x,\omega) = i\varepsilon_0 \int \frac{d\omega'}{2\pi} \omega'^2(\omega - \omega') \Theta\left(\frac{L}{2} - |x|\right) A(x,\omega - \omega') \cdot \int_{-\frac{L}{2}}^{\frac{L}{2}} dx' \chi(x,x',\omega')A(x',\omega') = : \Theta\left(\frac{L}{2} - |x|\right) W(x,\omega).$$

On the left surface, the step function reduces to $\Theta(\frac{L}{2} + x)$. We now obtain a new, model-free relation for $S_m$ by integrating (14):

$$S_m(x,\omega) = \Theta\left(\frac{L}{2} + x\right) \int_{-\frac{L}{2}}^{x} dx' W(x',\omega).$$

(16)

As can clearly be seen, $S_m$ is zero at the surface and grows continuously for $x > -\frac{L}{2}$, because $W$ is a finite value. It is therefore continuous at the surface. There is no evidence for any energy source or violation of Poynting’s theorem.

Energy conservation with bulk susceptibility and ABCs. – Notwithstanding the first successful microscopic non-local computation of the polarisation [20], which required supercomputers, the pragmatic assumption of a bulk susceptibility $\chi(q,\omega)$ is still widely used to analyse light propagation through bounded media. In this case, the dispersion relation

$$q(\omega)^2 = \frac{\omega^2}{c^2} (1 + \chi(q,\omega))$$

(17)

can have various solutions $q_i(\omega)$, giving rise to different polariton modes. Consequently, the electromagnetic field inside is a superposition of these forward and backward propagating modes (we stick with slab geometry):

$$E(x,\omega) = \sum_i [f_i(\omega)e^{iq_i(\omega)x} + b_i(\omega)e^{-iq_i(\omega)x}].$$

(18)

The extinction theorem,

$$a_1(\omega) = 1 - |r|^2 - |t|^2,$$

(19)

where $a_1$ is the absorption and $r, t$ the coefficients of reflected or transmitted light, respectively, and the MBCs,

$$e^{-iq_0 \frac{\omega}{c} t} + re^{iq_0 \frac{\omega}{c} t} = \sum_i \left( f_i e^{-iq_i \frac{\omega}{c} x} + b_i e^{iq_i \frac{\omega}{c} x} \right),$$

(20a)

$$e^{-iq_0 \frac{\omega}{c} t} - re^{iq_0 \frac{\omega}{c} t} = \sum_i \frac{q_i}{q_0} \left( f_i e^{-iq_i \frac{\omega}{c} x} - b_i e^{iq_i \frac{\omega}{c} x} \right),$$

(20b)

$$te^{iq_0 \frac{\omega}{c} t} = \sum_i \frac{q_i}{q_0} \left( f_i e^{iq_i \frac{\omega}{c} x} + b_i e^{-iq_i \frac{\omega}{c} x} \right),$$

(20c)

$$te^{-iq_0 \frac{\omega}{c} t} = \sum_i \frac{q_i}{q_0} \left( f_i e^{-iq_i \frac{\omega}{c} x} + b_i e^{iq_i \frac{\omega}{c} x} \right),$$

(20d)

where $q_0 = \omega/c$, form a set which could be solved if only a single mode was present but in fact is under-determined due to the polariton modes. The postulation of ABCs helps solving this problem, but some of them were known to conflict with energy conservation. (For a concise summary on ABCs and the controversy around them, see [8] and references therein.) Most of these ABCs, as well as the DA, can be reproduced by generalised ABCs [19].

Using the RHS of the integral representation of eq. (14), an independent expression for the absorption can be derived:

$$a_2(\omega) = \frac{1}{2i} \int dxdx' A^*(x,\omega)\chi(x,x',\omega)A(x',\omega).$$

(21)

with

$$\chi(x,x',\omega) = \chi(x,x',\omega) - \chi^*(x',x,\omega).$$

(22)

It follows by applying a monochromatic wave and some straightforward calculation. Then, static terms ($\propto \delta(\omega)$) and such $\propto \delta(\omega \pm 2\omega_0)$ appear. The former can be regarded separately, elegantly circumventing cycle averaging. The important relations (19) and (21) are presented in [10] in a quantum-mechanical generalisation valid for non-locally dispersive media and oblique incidence. They take almost the same form as shown here.

The polarisation here follows from eqs. (13) as

$$P(x,\omega) = \sum_j \chi_j(\omega) \left[ f_j(\omega)e^{iq_j(\omega)x} + b_j(\omega)e^{-iq_j(\omega)x} \right].$$

It can be used to calculate $a_2(\omega)$ in a variant of eq. (21) based on field strength and polarisation function only,

$$a_2(\omega) = \frac{\omega}{c} \text{Im} \int_{-\frac{L}{2}}^{\frac{L}{2}} dx E^*(x,\omega)P(x,\omega).$$

(23)

Now, constructing eq. (19) as

$$a_1(\omega) = \frac{1}{4} \left| (a(20a) + (20b))^2 - |(20a) - (20b)|^2 \right. \left. - |(20c) + (20d)|^2 + |(20c) - (20d)|^2 \right|,$$

(24)

the last term equaling zero, it can be shown that $a_1 = a_2$. As no ABCs have been used up to now, this means that energy conservation in such a system is ensured by the MBCs alone.
Conclusion. – We hopefully could present the concept of Poynting’s theorem in conjunction with a mechanical-energy flux vector in an illustrative way, concretising and extending textbooks in this sense, offer a best-practice approach, and show that doubts on its validity in certain conditions, such as spatial dispersion, are unsubstantiated. By the given approach, a clear relation between electromagnetic and kinetic energy in a system can be established, and electromagnetic energy density and flux are determined by the fundamental physical fields \( E \) and \( B \) alone.

We then have proven analytically that energy conservation in the propagation of light through bounded media is fulfilled for arbitrary susceptibilities, and thus also for the dielectric approximation.

Often, the use of models and approximations for the microscopic nature of a medium is misleading, and any predictions have to be taken with a grain of salt. This was pointed out as a problem in [7] and in the common perception of Poynting’s theorem in general. Great care must be taken in the application of the material equations, and it seems advisable to always start from (11) and the Maxwell-Lorentz equations (3), especially when spatial dispersion is dealt with. If a model for the mechanical Poynting vector is applied, it must be consistent with the susceptibility. Namely, it can be constructed via eq. (16).

Furthermore, we were able to prove analytically that energy conservation in approximations employing bulk susceptibility is completely assured by Maxwell’s boundary conditions alone, even if multiple polariton modes are predicted. Contrary to what was thought before, additional boundary conditions do not play any role in this respect, whether positive or negative. Consequently, from the energy conservation point of view, it is possible to freely choose the ABCs that match microscopic calculation results best (which would probably be a combination of Pekar’s and a dead layer, see [9]).

All these results save a plethora of recent and older work the shortcoming of possibly violating fundamental principles, and the use of their results in research is finally justifiable, at least in this respect.

**REFERENCES**

[1] Born M. and Wolf E., Principles of Optics, 7th edition (Cambridge University Press) 2003.
[2] Landau L. D. and Lifshitz E. M., Electrodynamics of Continuous Media (Pergamon, Oxford) 1960.
[3] Pekar S., Zh. Eksp. Theor. Fiz., 33 (1957) 1022 (Sov. Phys. JETP, 6 (1985) 785).
[4] Ting C. S., Frankel M. J. and Birman J. L., Solid State Commun., 17 (1975) 1285.
[5] Maradudin A. A. and Mills D. L., Phys. Rev. B, 7 (1973) 2787.
[6] Seemann M., Kieseling F., Stolz H., Manzke G., Henneberger K., Passow T. and Hommel D., Phys. Status Solidi (c), 3 (2006) 2453.
[7] Bishop M. F. and Maradudin A. A., Phys. Rev. B, 14 (1976) 3384.
[8] Venger E. F. and Piskovoi V. N., Phys. Rev. B, 70 (2004) 115107.
[9] Muljarov E. A. and Zimmermann R., Phys. Rev. B, 66 (2002) 235319.
[10] Henneberger K., submitted to EPL, arXiv:0710.5686 [condmat.str-el] (2007).
[11] Poynting J. H., Philos. Trans., 175 (1884) 277.
[12] Jackson J. D., Classical Electrodynamics (Wiley) 1999.
[13] Nelson D. F., Phys. Rev. Lett., 76 (1996) 4713.
[14] Campos I. and Jimenez J., Eur. J. Phys., 13 (1992) 117.
[15] Nelson D. F., Phys. Rev. E, 51 (1995) 6142.
[16] Jiang Y. and Liu M., Phys. Rev. Lett., 77 (1996) 1043.
[17] Cohen-Tannoudji C., Dupont-Roc J. and Grynberg G., Photons and Atoms (Wiley, New York) 1989.
[18] Nolting W., Elektrodynamik, Grundkurs Theoretische Physik, Vol. 3, 8th edition (Springer, Berlin) 2007.
[19] Halevi P., Spatial Dispersions in Solids and Plasma (North-Holland, Amsterdam) 1992.
[20] Tignon J., Hasche T., Chemla D. S., Schneider H. C., Jahneke F. and Koch S. W., Phys. Rev. Lett., 84 (2000) 3382.

The authors would like to thank the Deutsche Forschungsgemeinschaft for support through Sonderforschungsbereich 652.