Vol 18 No 6, June 2009
1674-1056/2009/18(06)/2314-05

Chinese Physics B

© 2009 Chin. Phys. Soc.
and IOP Publishing Ltd

Quantization of the radiation field in an anisotropic dielectric medium*

Li Wei(李 维), Liu Shi-Bing (刘世炳)†, and Yang Wei(杨 巍)

Institute of Laser Engineering, Beijng University of Technology, Beijing 100022, *China*

(Received 22 October 2008; revised manuscript received 25 November 2008)

There are both loss and dispersion characteristics for most dielectric media. In quantum theory the loss in medium is generally described by Langevin force in the Langevin noise (LN) scheme by which the quantization of the radiation field in various homogeneous absorbing dielectrics can be successfully actualized. However, it is invalid for the anisotropic dispersion medium. This paper extends the LN theory to an anisotropic dispersion medium and presented the quantization of the radiation field as well as the transformation relation between the homogeneous and anisotropic dispersion media.

Keywords: dispersion dieletric, anisotropic dielectric, noise current operator
PACC: 4250

1. Introduction

The electromagnetic radiations arising in various dielectric media refer to many applied fields.[1,2] Especially, with the rapid growth of interest in quantum optical experiments which take place inside the material media, the quantization of the radiation field in a nonlinear medium is important. In the early quantization approaches the lossless dielectrics were selected for convenience, in which the dielectric function was a constant[3] or a real function in space and frequency.[4] According to Kramers–Kronig relations, the loss in any dielectric medium with dispersion is inevitable. Subsequently, three approaches have been independently developed to deal with this problem. Firstly, the dispersive and lossy dielectric medium is represented by a collection of interacting matter fields based on microscopy,[5] the radiated electromagnetic field is treated by a fully canonical quantization scheme.[6−10] In the second approach, the quantization of the radiated electromagnetic field which propagates in different dielectric geometries is carried out, where the dielectric function of the homogeneous absorbing medium is treated as a known complex function and the loss is modeled by Langevin forces in terms of noise current operators.[11−16] The realization of quantization, in the last one, is to introduce a set of auxiliary fields to replace noise current based on the classical phenomenological Maxwell equations and in detail introduce a set of constitutive permeabilities to obtain a unitary time evolution canonical formalism for a linear, nonconducting, absorptive, and dispersive medium.[17−19] The recent investigations show that the intrinsic affiliation in these three approaches becomes more and more lucid[20] and attracts the researcher's special interest because the really measured values of the dielectric function can be directly applied to the Langevin noise (LN) approach in the theoretical description of quantized light in media.

The LN approach[12,14] is based on the consideration of Langevin force dictated by fluctuation dissipation theorem and all classical electromagnetic field quantities are directly transferred into operators in order to realize the quantization of classical Maxwell equations. Recently, this approach has been used to quantize the isotropic absorbing medium from one- to three-dimensions[12,15] and applied to the calculation of quantum optical processes in the isotropic dielectric medium with various geometrical shapes. However, it is well known that the most of the media in applied fields are anisotropic dielectrics and the current theory is still limited to the isotropic medium. Therefore, many researchers pay attention to the development of a theory for the anisotropic dielectric medium. Based on LN theory, in this paper, we suggest a general

*Project supported by the National Natural Science Foundation of China (Grant No 10574010).
†Corresponding author. E-mail: sbliu@bjut.edu.cn

approach to realize the quantization of the radiation field in an anisotropic dispersive dielectric medium by defining a modified new coefficient of the Langevin force operator. The transformation relation between the quantization methods for isotropic and anisotropic dispersive media is also presented.[12,15]

2. Basic equations

In classical macroscopic electrodynamics, the medium can be completely described by its dielectric function. According to the Kramers–Kronig relation, a dielectric that shows dispersion must inevitably be lossy, so the dielectric function must be defined as a complex function

$$\varepsilon(\omega) = \mathrm{Re}\,\varepsilon(\omega) + \mathrm{i}\,\mathrm{Im}\,\varepsilon(\omega), \tag{1}$$

where the real part $\mathrm{Re}\,\varepsilon(\omega)$ indicates the dielectric dispersion and the imaginary part $\mathrm{Im}\,\varepsilon(\omega)$ shows the absorption.

It is usually convenient to separate the electromagnetic field operators into the positive and negative frequency components by Fourier transform

$$\hat{\boldsymbol{E}}(\boldsymbol{r},t) = \frac{1}{\sqrt{2\pi}}\int_0^{+\infty} \mathrm{d}\omega \left[\hat{\boldsymbol{E}}^+(\boldsymbol{r},\omega)\mathrm{e}^{-\mathrm{i}\omega t} + \hat{\boldsymbol{E}}^-(\boldsymbol{r},\omega)\mathrm{e}^{\mathrm{i}\omega t}\right], \tag{2}$$

where the positive and negative frequency components involve the annihilation and creation operators and satisfy Maxwell's equations, respectively

$$\boldsymbol{\nabla}\times\hat{\boldsymbol{E}}^+(\boldsymbol{r},\omega) = \mathrm{i}\omega\hat{\boldsymbol{B}}^+(\boldsymbol{r},\omega), \tag{3}$$

$$\boldsymbol{\nabla}\times\hat{\boldsymbol{B}}^+(\boldsymbol{R},\omega) = \mathrm{i}\omega\mu_0\hat{\boldsymbol{D}}^+(\boldsymbol{r},\omega) + \mu_0\hat{\boldsymbol{J}}^+(\boldsymbol{r},\omega). \tag{4}$$

By choosing the gauge, the scalar potential $\varphi = 0$, the electric and magnetic field operators can be obtained

$$\hat{\boldsymbol{E}}^+(\boldsymbol{r},\omega) = \mathrm{i}\,\omega\hat{\boldsymbol{A}}^+(\boldsymbol{r},\omega), \tag{5}$$

$$\hat{\boldsymbol{B}}^+(\boldsymbol{r},\omega) = \boldsymbol{\nabla}\times\hat{\boldsymbol{A}}^+(\boldsymbol{r},\omega). \tag{6}$$

Combining Eqs.(3) and (4), one can present the positive frequency component of the electromagnetic field as

$$\boldsymbol{\nabla}\times[\boldsymbol{\nabla}\times\hat{\boldsymbol{A}}^+(\boldsymbol{r},\omega)] - \mathrm{i}\,\omega\mu_0\hat{\boldsymbol{D}}^+(\boldsymbol{r},\omega) = \mu_0\hat{\boldsymbol{J}}^+(\boldsymbol{r},\omega). \tag{7}$$

By choosing the Coulomb gauge, only the transverse component of the vector potential has a contribution. Taking advantage of the transverse wave condition $\boldsymbol{\nabla}\cdot\hat{\boldsymbol{A}}^+(\boldsymbol{r},\omega) = 0$, we can change Eq.(7) to

$$\boldsymbol{\nabla}^2\hat{\boldsymbol{A}}^+(\boldsymbol{r},\omega) - \mathrm{i}\,\omega\mu_0\hat{\boldsymbol{D}}^+(\boldsymbol{r},\omega) = \mu_0\hat{\boldsymbol{J}}^+(\boldsymbol{r},\omega), \tag{8}$$

where $\varepsilon(\boldsymbol{r},\omega)$ is a scalar function, $\varepsilon_0(\mu_0)$ is the vacuum electric (magnetic) permeability, and

$$\hat{\boldsymbol{D}}^+(\boldsymbol{r},\omega) = \mathrm{i}\,\omega\varepsilon_0\varepsilon(\boldsymbol{r},\omega)\hat{\boldsymbol{A}}^+(\boldsymbol{r},\omega). \tag{9}$$

The negative frequency components can be obtained by the Hermitian conjugates of the positive frequency operators. $\hat{\boldsymbol{J}}^+(\boldsymbol{r},\omega)$ represents the noise current operator associated with the absorptive nature of the medium.

Lagrangian mechanics determines that the generalized momentum operator is conjugate to the vector potential operator in the quantized field theory. In the Coulomb gauge, the conjugate momentum of the vector potential operator $\hat{\boldsymbol{A}}(\boldsymbol{r},\omega)$ is $-\varepsilon_0\hat{E}(\boldsymbol{r},\omega)$, and their canonical commutation relation is

$$[\hat{\boldsymbol{A}}_\alpha(\boldsymbol{r},\omega), -\varepsilon_0\hat{\boldsymbol{E}}_\beta(\boldsymbol{r}',\omega)] = \mathrm{i}\,\hbar\delta_{\alpha\beta}\delta(\boldsymbol{r}-\boldsymbol{r}'), \tag{10}$$

where indexes $\alpha,\beta = 1,2,3$ represent the Cartesian components x, y and z.

The field commutation relations are closely related to the commutation properties of the noise current operator. The noise is assumed to be uncorrelated at different frequencies, its commutation relation has the form

$$[\hat{j}_\alpha^+(\boldsymbol{r},\omega), \hat{j}_\beta^-(\boldsymbol{r}',\omega')] = \alpha(\boldsymbol{r},\omega)\,\delta_{\alpha\beta}\delta(\boldsymbol{r}-\boldsymbol{r}')\,\delta(\omega-\omega'), \tag{11}$$

$$[\hat{j}_\alpha^+(\boldsymbol{r},\omega), \hat{j}_\beta^+(\boldsymbol{r}',\omega')] = \left[\hat{j}_\alpha^-(\boldsymbol{r},\omega), \hat{j}_\beta^-(\boldsymbol{r}',\omega')\right] = 0, \tag{12}$$

where

$$\alpha(\boldsymbol{r},\omega) = 2\varepsilon_0\hbar\omega^2\mathrm{Im}\,\varepsilon(\boldsymbol{r},\omega), \tag{13}$$

in which $\alpha(\boldsymbol{r},\omega)$ is determined by the condition that equation (11) should be consistent with Eq.(10), which is a scalar real function in the isotropic medium.[12] In the next section we will modify the $\alpha(\boldsymbol{r},\omega)$ expression for extending the applicability of the LN approach to the anisotropic medium.

3. Field quantization in an anisotropic medium

In the anisotropic medium, the displacement operator has the more complicated tensor form, which is expressed as

$$\hat{D}_\alpha(\boldsymbol{r},\omega)=\sum\varepsilon_{\alpha\beta}\hat{E}_\beta(\boldsymbol{r},\omega), \tag{14}$$

where the dielectric function $\varepsilon_{\alpha\beta}$ is a tensor function. In this paper all the repeated indexes do not imply a summation over the Cartesian components x, y and z, if necessary the summation will be represented by $\sum$.

Let us start with Eq.(7). By employing Eqs.(2), (3) and (14), one can show that the positive frequency component of the vector potential satisfies an equation in the form

$$\sum_\beta\left[k^2\delta_{\alpha\beta}-\frac{\omega^2}{c^2}\varepsilon_{\alpha\beta}\left(\boldsymbol{r},\omega\right)\right]\hat{A}_\beta^+\left(\boldsymbol{r},\omega\right)=\frac{1}{c^2\varepsilon_0}\hat{J}_\alpha^+\left(\boldsymbol{r},\omega\right), \tag{15}$$

we consider the field quantization in an infinite homogeneous medium, in which the dielectric function has no spatial dependence. We can take advantage of this feature to consider the Fourier transform

$$\hat{A}^+\left(r,\omega\right)=\frac{1}{\left(2\pi\right)^{3/2}}\int\mathrm{d}^3\boldsymbol{k}\hat{A}^+\left(\boldsymbol{k},\omega\right)\mathrm{e}^{\mathrm{i}\boldsymbol{k}\cdot\boldsymbol{r}}. \tag{16}$$

Using Eq.(16) for $\hat{A}^+\left(\boldsymbol{r},\omega\right)$ and $\hat{J}^+\left(\boldsymbol{r},\omega\right)$, we find

$$\sum_\beta\left[k^2\delta_{\alpha\beta}-\frac{\omega^2}{c^2}\varepsilon_{\alpha\beta}\left(\omega\right)\right]\hat{A}_\beta^+\left(\boldsymbol{k},\omega\right)=\frac{1}{c^2\varepsilon_0}\hat{J}_\alpha^+\left(\boldsymbol{k},\omega\right). \tag{17}$$

We assume that non-diagonal elements of the dielectric tensor ($\varepsilon_{\alpha\beta}$, $\alpha\neq\beta$) are very small, only the first-order term is considered. Then three components of $\hat{\boldsymbol{A}}^+\left(\boldsymbol{k},\omega\right)$ can be easily obtained,

$$\begin{aligned}\hat{A}_1^+\left(\boldsymbol{k},\omega\right)=\frac{T^{-1}}{\varepsilon_0}\Big\{&\left[k^2c^2-\omega^2\varepsilon_{22}\left(\omega\right)\right]\left[k^2c^2-\omega^2\varepsilon_{33}\left(\omega\right)\right]\hat{J}_1^+\left(\boldsymbol{k},\omega\right)\\&+\omega^2\varepsilon_{12}\left(\omega\right)\left[k^2c^2-\omega^2\varepsilon_{33}\left(\omega\right)\right]\hat{J}_2^+\left(\boldsymbol{k},\omega\right)\\&+\omega^2\varepsilon_{13}\left(\omega\right)\left[k^2c^2-\omega^2\varepsilon_{22}\left(\omega\right)\right]\hat{J}_3^+\left(\boldsymbol{k},\omega\right)\Big\},\end{aligned} \tag{18}$$

$$\begin{aligned}\hat{A}_2^+\left(\boldsymbol{k},\omega\right)=\frac{T^{-1}}{\varepsilon_0}\Big\{&\omega^2\varepsilon_{21}\left(\omega\right)\left[k^2c^2-\omega^2\varepsilon_{33}\left(\omega\right)\right]\hat{J}_1^+\left(\boldsymbol{k},\omega\right)\\&+\left[k^2c^2-\omega^2\varepsilon_{11}\left(\omega\right)\right]\left[k^2c^2-\omega^2\varepsilon_{33}\left(\omega\right)\right]\hat{J}_2^+\left(\boldsymbol{k},\omega\right)\\&+\omega^2\varepsilon_{23}\left(\omega\right)\left[k^2c^2-\omega^2\varepsilon_{11}\left(\omega\right)\right]\hat{J}_3^+\left(\boldsymbol{k},\omega\right)\Big\},\end{aligned} \tag{19}$$

$$\begin{aligned}\hat{A}_3^+\left(\boldsymbol{k},\omega\right)=\frac{T^{-1}}{\varepsilon_0}\Big\{&\omega^2\varepsilon_{31}\left(\omega\right)\left[k^2c^2-\omega^2\varepsilon_{22}\left(\omega\right)\right]\hat{J}_1^+\left(\boldsymbol{k},\omega\right)\\&+\omega^2\varepsilon_{32}\left(\omega\right)\left[k^2c^2-\omega^2\varepsilon_{11}\left(\omega\right)\right]\hat{J}_2^+\left(\boldsymbol{k},\omega\right)\\&+\left[k^2c^2-\omega^2\varepsilon_{11}\left(\omega\right)\right]\left[k^2c^2-\omega^2\varepsilon_{22}\left(\omega\right)\right]\hat{J}_3^+\left(\boldsymbol{k},\omega\right)\Big\},\end{aligned} \tag{20}$$

where

$$T=\left[k^2c^2-\omega^2\varepsilon_{11}\left(\omega\right)\right]\left[k^2c^2-\omega^2\varepsilon_{22}\left(\omega\right)\right]\left[k^2c^2-\omega^2\varepsilon_{33}\left(\omega\right)\right]. \tag{21}$$

Taking into account the contribution of the negative frequency component, we can express the complete expression for the vector potential operator in the spatial coordinate as

$$\hat{\boldsymbol{A}}\left(\boldsymbol{r},t\right)=\frac{1}{\left(2\pi\right)^2}\int_0^{+\infty}\mathrm{d}\omega\int\mathrm{d}^3\boldsymbol{k}\left\{\hat{\boldsymbol{A}}^+\left(\boldsymbol{k},\omega\right)\exp[-\mathrm{i}(\omega t-\boldsymbol{k}\cdot\boldsymbol{r})]+\mathrm{H.c.}\right\}. \tag{22}$$

The electric field operators are obtained with the help of Eq.(5). We find

$$\hat{\boldsymbol{E}}\left(\boldsymbol{r},t\right)=\frac{\mathrm{i}}{\left(2\pi\right)^{2}}\int_{0}^{+\infty}\mathrm{d}\omega\omega\int\mathrm{d}^{3}\boldsymbol{k}\left\{\hat{\boldsymbol{A}}^{+}\left(\boldsymbol{k},\omega\right)\exp[-\mathrm{i}(\omega t-\boldsymbol{k}\cdot\boldsymbol{r})]-\mathrm{H.c.}\right\}. \tag{23}$$

For the anisotropic medium, the field operators are still required to satisfy the equal-time canonical commutation relation Eq.(10). Substituting Eqs.(18)–(20) into Eqs.(22) and (23), respectively, and using Eq.(10), we find

$$\begin{aligned}&\left[\hat{A}_{1}\left(r,t\right),-\varepsilon_{0}\hat{E}_{1}\left(r',t\right)\right]\\=&\frac{\mathrm{i}}{\left(2\pi\right)^{2}\varepsilon_{0}}\left\{\int_{-\infty}^{+\infty}\mathrm{d}\omega\int\mathrm{d}^{3}k\omega\left|\left[k^{2}c^{2}-\omega^{2}\varepsilon_{22}\left(\omega\right)\right]\left[k^{2}c^{2}-\omega^{2}\varepsilon_{33}\left(\omega\right)\right]\right|^{2}T^{-2}\alpha_{1}\ \mathrm{e}^{\mathrm{i}\boldsymbol{k}\cdot\left(\boldsymbol{r}-\boldsymbol{r}'\right)}\right.\\&+\int_{-\infty}^{+\infty}\mathrm{d}\omega\int\mathrm{d}^{3}k\omega^{5}\left|\varepsilon_{12}\left(\omega\right)\left[k^{2}c^{2}-\omega^{2}\varepsilon_{33}\left(\omega\right)\right]\right|^{2}T^{-2}\alpha_{2}\,\mathrm{e}^{\mathrm{i}\boldsymbol{k}\cdot\left(\boldsymbol{r}-\boldsymbol{r}'\right)}\\&\left.+\int_{-\infty}^{+\infty}\mathrm{d}\omega\int\mathrm{d}^{3}k\omega^{5}\left|\varepsilon_{13}\left(\omega\right)\left[k^{2}c^{2}-\omega^{2}\varepsilon_{22}\left(\omega\right)\right]\right|^{2}T^{-2}\alpha_{3}\,\mathrm{e}^{\mathrm{i}\boldsymbol{k}\cdot\left(\boldsymbol{r}-\boldsymbol{r}'\right)}\right\},\end{aligned} \tag{24}$$

$$\begin{aligned}&\left[\hat{A}_{2}\left(r,t\right),-\varepsilon_{0}\hat{E}_{2}\left(r',t\right)\right]\\=&\frac{\mathrm{i}}{\left(2\pi\right)^{2}\varepsilon_{0}}\left\{\int_{-\infty}^{+\infty}\mathrm{d}\omega\int\mathrm{d}^{3}k\omega^{5}\left|\varepsilon_{21}\left(\omega\right)\left[k^{2}c^{2}-\omega^{2}\varepsilon_{33}\left(\omega\right)\right]\right|^{2}T^{-2}\alpha_{1}\ \mathrm{e}^{\mathrm{i}\boldsymbol{k}\cdot\left(\boldsymbol{r}-\boldsymbol{r}'\right)}\right.\\&+\int_{-\infty}^{+\infty}\mathrm{d}\omega\int\mathrm{d}^{3}k\omega\left|\left[k^{2}c^{2}-\omega^{2}\varepsilon_{11}\left(\omega\right)\right]\left[k^{2}c^{2}-\omega^{2}\varepsilon_{33}\left(\omega\right)\right]\right|^{2}T^{-2}\alpha_{2}\,\mathrm{e}^{\mathrm{i}\boldsymbol{k}\cdot\left(\boldsymbol{r}-\boldsymbol{r}'\right)}\\&\left.+\int_{-\infty}^{+\infty}\mathrm{d}\omega\int\mathrm{d}^{3}k\omega^{5}\left|\varepsilon_{23}\left(\omega\right)\left[k^{2}c^{2}-\omega^{2}\varepsilon_{11}\left(\omega\right)\right]\right|^{2}T^{-2}\alpha_{3}\,\mathrm{e}^{\mathrm{i}\boldsymbol{k}\cdot\left(\boldsymbol{r}-\boldsymbol{r}'\right)},\right.\end{aligned} \tag{25}$$

$$\begin{aligned}&\left[\hat{A}_{3}\left(r,t\right),-\varepsilon_{0}\hat{E}_{3}\left(r',t\right)\right]\\=&\frac{\mathrm{i}}{\left(2\pi\right)^{2}\varepsilon_{0}}\left\{\int_{-\infty}^{+\infty}\mathrm{d}\omega\int\mathrm{d}^{3}k\omega^{5}\left|\varepsilon_{31}\left(\omega\right)\left[k^{2}c^{2}-\omega^{2}\varepsilon_{22}\left(\omega\right)\right]\right|^{2}T^{-2}\alpha_{1}\ \mathrm{e}^{\mathrm{i}\boldsymbol{k}\cdot\left(\boldsymbol{r}-\boldsymbol{r}'\right)}\right.\\&+\int_{-\infty}^{+\infty}\mathrm{d}\omega\int\mathrm{d}^{3}k\omega^{5}\left|\varepsilon_{32}\left(\omega\right)\left[k^{2}c^{2}-\omega^{2}\varepsilon_{11}\left(\omega\right)\right]\right|^{2}T^{-2}\alpha_{2}\,\mathrm{e}^{\mathrm{i}\boldsymbol{k}\cdot\left(\boldsymbol{r}-\boldsymbol{r}'\right)}\\&\left.+\int_{-\infty}^{+\infty}\mathrm{d}\omega\int\mathrm{d}^{3}k\omega\left|\left[k^{2}c^{2}-\omega^{2}\varepsilon_{11}\left(\omega\right)\right]\left[k^{2}c^{2}-\omega^{2}\varepsilon_{22}\left(\omega\right)\right]\right|^{2}T^{-2}\alpha_{3}\,\mathrm{e}^{\mathrm{i}\boldsymbol{k}\cdot\left(\boldsymbol{r}-\boldsymbol{r}'\right)}\right\},\end{aligned} \tag{26}$$

where coefficient α_{α} is different from the one in Eq.(13). Its values are not the same in different directions. According to the equal-time canonical commutation relation Eq.(10), all equations (24)–(26) equal $\mathrm{i}\hbar\delta\left(\boldsymbol{r}-\boldsymbol{r}'\right)$, and taking advantage of

$$\int_{-\infty}^{+\infty}\frac{\omega^{3}\operatorname{Im}\varepsilon\left(\omega\right)}{\left|k^{2}c^{2}-\omega^{2}\varepsilon\left(\omega\right)\right|^{2}}=\pi, \tag{27}$$

finally we find

$$\alpha_{\alpha}=2\varepsilon_{0}\hbar\omega^{2}\operatorname{Im}\varepsilon_{\alpha\alpha}\left(\omega\right)\left(1-\frac{\omega^{4}\sum\limits_{\beta}\left[\left|\varepsilon_{\alpha\beta}\left(\omega\right)\right|^{2}\left(1-\delta_{\alpha\beta}\right)\right]}{\left|\left[k^{2}c^{2}-\omega^{2}\varepsilon_{\alpha\alpha}\left(\omega\right)\right]\right|^{2}}\right), \tag{28}$$

this is the modified coefficient of the Langevin force operator in the anisotropic medium, which is more complicated than that in the isotropic medium. We will discuss its property in the next section.

4. Discussion

We extend the applicability of the LN approach[12,15] to the anisotropic medium by modifying its coefficient. From Eq.(28), it can be easily seen that the modified coefficient is not only related with the dielectric function in its direction, but also influenced by non-diagonal elements of the dielectric tensor. For the isotropic medium, all the non-diagonal elements of the dielectric tensor equal zero and the diagonal elements are the same, then the dielectric tensor can be described as

$$\varepsilon_{\alpha\beta}(\omega) = \varepsilon(\omega)\,\delta_{\alpha\beta}. \tag{29}$$

Substituting Eq.(29) into Eq.(28), we find

$$\alpha_1 = \alpha_2 = \alpha_3 = \alpha = 2\varepsilon_0\hbar\omega^2\,\mathrm{Im}\,\varepsilon(\omega). \tag{30}$$

Equation (30) reverts back to Eq.(13) when the modified coefficient is defined by the isotropic medium. This explains the modified coefficient of quantization for the isotropic medium which is a special case for the anisotropic medium. By employing the new modified coefficient, the LN approach can be directly used to quantize the electromagnetic field in the anisotropic medium.

In addition, it is obvious in Eq.(28), which is different from Eq.(13), that the modified coefficient of the Langevin force operator depends on both real and imaginary parts of the dielectric function. The modified coefficient of the Langevin force operator is different between strong and weak dispersion media, in which the value of the modified coefficient is larger in strong dispersion. Furthermore, the value of the modified coefficient in the anisotropic medium is smaller than that in the isotropic medium.

References

[1] Jiang S E, Yang J M, Zheng Z J and Ding Y K 2007 *Chin. Phys.* **16** 193
[2] Zhang C X, Tian H T, Li M, Jin J and Song N F 2008 *Chin. Phys.* B **17** 573
[3] Jauch J M and Watson K M 1948 *Phys. Rev.* **74** 950
[4] Glauber R J and Lewenstein M 1991 *Phys. Rev.* A **43** 467
[5] Hopfield J J 1958 *Phys. Rev.* **112** 1555
[6] Huttner B, Baumberg J J and Barnett S M 1991 *Europhys. Lett.* **16** 177
[7] Barnett S M, Huttner B and Loudon R 1992 *Phys. Rev. Lett.* **68** 3698
[8] Huttner B and Barnett S M 1992 *Phys. Rev.* A **46** 4306
[9] Yeung M S and Gustafson T K 1996 *Phys. Rev.* A **54** 5227
[10] Dutra S M and Furuya K 1998 *Phys. Rev.* A **57** 3050
[11] Gruner T and Welsch D G 1995 *Phys. Rev.* A **51** 3246
[12] Matloob R, Loudon R, Barnett S M and Jeffers J 1995 *Phys. Rev.* A **52** 4823
[13] Barnett S M, Matloob R and Loudon R 1995 *J. Mod. Opt.* **42** 1165
[14] Matloob R and Loudon R 1996 *Phys. Rev.* A **53** 4567
[15] Matloob R 1999 *Phys. Rev.* A **60** 50
[16] Matloob R 2004 *Phys. Rev.* A **70** 022108
[17] Tip A 1997 *Phys. Rev.* A **56** 5022
[18] Tip A 1998 *Phys. Rev.* A **57** 4818
[19] Lastdrager B, Tip A and Verhoeven J 2000 *Phys. Rev.* E **61** 5767
[20] Tip A, Knöll L, Scheel S and Welsch D G 2001 *Phys. Rev.* A **63** 043806