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Suppression of spontaneous emission and superradiance over macroscopic distances in media with negative refraction

Jürgen Kästel and Michael Fleischhauer
Fachbereich Physik, Technische Universität Kaiserslautern, D-67663 Kaiserslautern, Germany
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The possibility of negative optical path length in left-handed media (LHM) is shown to lead to complete suppression of spontaneous emission of an atom in front of a mirror with a layer of LHM. For the same reason two atoms put at the foci of a perfect lens formed by a parallel LHM slab [J. B. Pendry, Phys. Rev. Lett. **85**, 3966 (2000)] exhibit perfect subradiance and superradiance. It is shown that these effects occur over distances that can be orders of magnitude larger than the transition wavelength and are only limited by the propagation length within the free-space decay time of the atoms. Single- and two-atom decay rates are calculated from the Greens function of the electric field in the presence of a LHM and limitations as well as potential applications are discussed.

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Negative refraction of electromagnetic radiation in so-called left-handed media (LHM) predicted in the late 1960s [1] has recently attracted much attention because of its experimental demonstration in metamaterials in the cm-wave range [2–5]. The most peculiar property of LHM associated with the negative refraction is the possibility of negative optical path length, which, as pointed out by Pendry [6], allows us to construct a perfect lens with a resolution not limited by diffraction. The lens, an infinite parallel slab of LHM, collects all plane waves from a point source on one side of the slab in a focal point on the other side with zero phase difference. Due to the vanishing optical path length between the focal points also evanescent waves emerging from the source are exactly reproduced leading to, in principle, unlimited resolution. This raises the question of what happens to a pair of atoms put in the foci of the lens or to a single atom put in its own focus induced by a nearby mirror in combination with a LHM (Fig. 1).

We here show that the imaginary part of the retarded Greens function $\mathrm{Im}[\mathbf{G}(\mathbf{r}_1,\mathbf{r}_2)]$ between the two focal points $\mathbf{r}_1$,$\mathbf{r}_2$ of the perfect lens [Fig. 1(a)] is identical to the free-space value at the same position $\mathrm{Im}[\mathbf{G}(\mathbf{r},\mathbf{r})]$. As a consequence there occurs perfect subradiance and superradiance [7] of two atoms at $\mathbf{r}_1$ and $\mathbf{r}_2$ as well as mutual dipole-dipole shifts even for distances $|\mathbf{r}_1-\mathbf{r}_2|$ large compared to the transition wavelength. Likewise the imaginary part of the Greens function in the focus before a mirror-LHM combination [Fig. 1(b)] becomes the same as for a combination of mirror and vacuum with $\mathbf{r}_1=\mathbf{r}_2$ being directly on the surface of the mirror leading to a strong radiative back action. In both cases the strong radiative coupling or self-coupling persists as long as the distance between the atoms or the atom and the mirror is smaller than the propagation length during the free-space radiative decay time.

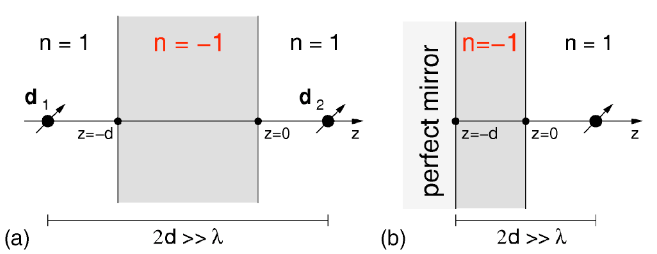

FIG. 1. (a) Two atoms put into the focal points of a Veselago-Pendry lens with $n=-1$. The focal points are all pairs of positions at the two sides of the slab with distance $2d$. The spatial regions $z>0$ (vacuum), $-d\leqslant z\leqslant 0$ (LHM), and $z<-d$ (vacuum) are denoted by the numbers 0,1,2 respectively. (b) One atom in front of a system of a perfect mirror combined with a perfect LHM. The optical length between atom and mirror is zero.

Since the one-atom system [Fig. 1(b)] can be viewed as a special case of the two-atom system [Fig. 1(a)], we here discuss only the second one. Let us therefore consider an infinitely extended slab of homogeneous LHM of thickness d and two atoms put in the focal points as shown in [Fig. 1(a).] The atoms are two-level systems with ground states $|1\rangle$ and excited states $|2\rangle$ and common transition frequency ω_0. The dipole vectors of the atoms are denoted by $\mathbf{d}_1$ and $\mathbf{d}_2$. The coupling of the atoms to the quantized radiation field is described by the interaction Hamiltonian in dipole approximation

$$H_{\mathrm{WW}}=-\hat{\mathbf{d}}_1\cdot\hat{\mathbf{E}}(\mathbf{r}_1)-\hat{\mathbf{d}}_2\cdot\hat{\mathbf{E}}(\mathbf{r}_2), \tag{1}$$

where $\hat{\mathbf{E}}(\mathbf{r})$ is the operator of the electric field in the presence of the LHM. Employing the usual Born-Markov and rotating-wave approximations one derives the standard master equation for the two-atom density matrix in the interaction picture

$$\dot{\rho}=-\sum_{k,l=1}^{2}\frac{\Gamma(\mathbf{r}_k,\mathbf{r}_l)}{2}\left(\hat{\sigma}_l^\dagger\hat{\sigma}_k\rho+\rho\hat{\sigma}_l^\dagger\hat{\sigma}_k-2\hat{\sigma}_k\rho\hat{\sigma}_l^\dagger\right)+i\sum_{k,l=1}^{2}\delta\omega(\mathbf{r}_k,\mathbf{r}_l)\left[\hat{\sigma}_l^\dagger\hat{\sigma}_k,\rho\right]. \tag{2}$$

Here $\hat{\sigma}_k=|1\rangle_{kk}\langle 2|$ are the flip operators of the kth atom. The rates $\Gamma(\mathbf{r}_k,\mathbf{r}_l)$ describe the radiative decay of the two two-level atoms. $\Gamma(\mathbf{r}_k,\mathbf{r}_k)$ corresponds to the single-particle decay rate of an atom at position $\mathbf{r}_k$ and $\Gamma(\mathbf{r}_1,\mathbf{r}_2)$ describes the dissipative cross coupling. Both quantities are determined by the imaginary part of the Greens tensor of the electric field at the atomic transition frequency $\mathbf{G}(\mathbf{r}_k,\mathbf{r}_l,\omega_0) = G_{\mu\nu}(\mathbf{r}_k,\mathbf{r}_l,\omega_0)\hat{\mu}\circ\hat{\nu}$, "$\circ$" denoting a tensorial product with $\hat{\mu}$ and $\hat{\nu}$ being unit vectors [8]

$$\Gamma(\mathbf{r}_k,\mathbf{r}_l) = \frac{2\omega^2 d_\mu d_\nu}{\hbar\varepsilon_0 c^2}\mathrm{Im}[G_{\mu\nu}(\mathbf{r}_k,\mathbf{r}_l,\omega_0)]. \tag{3}$$

$\delta\omega(\mathbf{r}_k,\mathbf{r}_k)$ is the single-atom Lamb shift and $\delta\omega(\mathbf{r}_1,\mathbf{r}_2)$ describes the radiative dipole-dipole shift. It is well known that the single-atom Lamb shift is not correctly described by the two-level model. Its explicit expression as given below diverges and a renormalization is needed. The Lamb shift is, however, of no relevance for the present discussion and will be ignored, i.e., it is assumed to be included in the bare transition frequency. In contrast the dipole-dipole shift can, in principle, correctly be calculated within the present model, when the free-field part is subtracted. This is because the Veselago-Pendry lens can only lead to contributions within a finite frequency window [1]. Subtracting the free-field contribution one finds

$$\delta\omega = \frac{d_\mu d_\nu}{\hbar\pi\varepsilon_0}\mathrm{P}\int_0^\infty d\omega\frac{\omega^2}{c^2}\frac{\mathrm{Im}[\Delta G_{\mu\nu}(\mathbf{r}_1,\mathbf{r}_2,\omega)]}{\omega-\omega_0}, \tag{4}$$

where $\Delta G_{\mu\nu}(\mathbf{r}_1,\mathbf{r}_2,\omega) = G_{\mu\nu}(\mathbf{r}_1,\mathbf{r}_2,\omega) - G^0_{\mu\nu}(\mathbf{r}_1,\mathbf{r}_2,\omega)$, $G^0_{\mu\nu}(\mathbf{r}_1,\mathbf{r}_2,\omega)$ being the components of the free-field retarded Greens tensor.

The master Eq. (2) for the two-atom system can be written in a diagonal form introducing a basis of symmetric and antisymmetric states $|11\rangle$, $|22\rangle$, and $|s\rangle = (|12\rangle + |21\rangle)/\sqrt{2}$ and $|a\rangle = (|12\rangle - |21\rangle)/\sqrt{2}$. This yields for the populations

$$\dot{\rho}_{22} = -2\Gamma_{11}\rho_{22}, \tag{5}$$

$$\dot{\rho}_{\mathrm{ss}} = +(\Gamma_{11}+\Gamma_{12})\rho_{22} - (\Gamma_{11}+\Gamma_{12})\rho_{\mathrm{ss}}, \tag{6}$$

$$\dot{\rho}_{\mathrm{aa}} = +(\Gamma_{11}-\Gamma_{12})\rho_{22} - (\Gamma_{11}-\Gamma_{12})\rho_{\mathrm{aa}}, \tag{7}$$

$$\dot{\rho}_{11} = +(\Gamma_{11}+\Gamma_{12})\rho_{\mathrm{ss}} + (\Gamma_{11}+\Gamma_{12})\rho_{\mathrm{aa}}, \tag{8}$$

where $\Gamma_{11} = \Gamma(\mathbf{r},\mathbf{r})$ and $\Gamma_{12} = \Gamma(\mathbf{r}_1,\mathbf{r}_2)$. In addition there is a level shift of the symmetric and antisymmetric states $|\mathrm{s}\rangle$ and $|\mathrm{a}\rangle$ below or above the single atom energy by the dipole-dipole shift $\delta\omega$, given in Eq. (4).

To obtain the corresponding results for the mirror-LHM system [Fig. 1(b)] it suffices to set the atomic flip operators of the second atom equal to zero and interpret the two-atom density matrix as a single-atom density matrix. The resulting rate equations simply read $\dot{\rho}_{22} = -\Gamma_{11}\rho_{22}$ and $\dot{\rho}_{11} = \Gamma_{11}\rho_{22}$, which describe the spontaneous decay of one two-level atom.

The retarded Greens function corresponding to a slab with a homogeneous and linear magnetodielectric medium can be calculated by a plane wave decomposition. Following [9] one finds for the two positions $\mathbf{r}$ and $\mathbf{r}'$ in vacuum on the same side of the lens

$$\mathbf{G}^{00}(\mathbf{r},\mathbf{r}',\omega) = \frac{i}{8\pi^2}\int d^2k_\perp\frac{1}{k_z} \times\Big\{\big[R^{\mathrm{TE}}\hat{\mathbf{e}}(k_z)e^{i\mathbf{k}\cdot\mathbf{r}} + \hat{\mathbf{e}}(-k_z)e^{i\mathbf{K}\cdot\mathbf{r}}\big]\circ\hat{\mathbf{e}}(-k_z)e^{-i\mathbf{K}\cdot\mathbf{r}'} + \big[R^{\mathrm{TM}}\hat{\mathbf{h}}(k_z)e^{i\mathbf{k}\cdot\mathbf{r}} + \hat{\mathbf{h}}(-k_z)e^{i\mathbf{K}\cdot\mathbf{r}}\big]\circ\hat{\mathbf{h}}(-k_z)e^{-i\mathbf{K}\mathbf{r}'}\Big\}, \tag{9}$$

where $z \leqslant z'$ has been assumed. For $\mathbf{r}$ and $\mathbf{r}'$ being in vacuum on different sides of the lens one finds

$$\mathbf{G}^{20}(\mathbf{r},\mathbf{r}',\omega) = \frac{i}{8\pi^2}\int d^2k_\perp\frac{1}{k_z}e^{i\mathbf{K}\cdot(\mathbf{r}-\mathbf{r}')}\big[T^{\mathrm{TE}}\hat{\mathbf{e}}(-k_z)\circ\hat{\mathbf{e}}(-k_z) + T^{\mathrm{TM}}\hat{\mathbf{h}}(-k_z)\circ\hat{\mathbf{h}}(-k_z)\big]. \tag{10}$$

The superscripts 0,1,2 denote the zones of positions $\mathbf{r}$ and $\mathbf{r}'$: $z>0$ (vacuum), $-d \leqslant z \leqslant 0$ (LHM), and $z < -d$ (vacuum), respectively. We here have used the definitions $k^2 = \omega^2/c^2$, $k_z = \sqrt{(k^2-k_\perp^2)}$ and $d^2k_\perp = dk_x dk_y$. Furthermore, $\mathbf{K} \equiv k_x\hat{\mathbf{x}} + k_y\hat{\mathbf{y}} - k_z\hat{\mathbf{z}}$ and we have introduced the orthogonal unit vectors $\hat{\mathbf{e}} = \mathbf{k}\times\hat{\mathbf{z}}/|\mathbf{k}\times\hat{\mathbf{z}}|$ and $\hat{\mathbf{h}} = p\hat{\mathbf{e}}\times\mathbf{k}/|k|$, where $p=1$ for a normal medium and $p=-1$ for a LHM. $R^{\mathrm{TE}}, R^{\mathrm{TM}}$, and $T^{\mathrm{TE}}, T^{\mathrm{TM}}$ are the reflection and transmission functions of the lens for transverse electric and transverse magnetic modes. They read

$$R^{\mathrm{TE}} = \frac{R_{01}+R_{12}e^{i2k_{1z}d}}{1+R_{01}R_{12}e^{i2k_{1z}d}}, \tag{11}$$

$$R^{\mathrm{TM}} = \frac{S_{01}+S_{12}e^{i2k_{1z}d}}{1+S_{01}S_{12}e^{i2k_{1z}d}}, \tag{12}$$

and correspondingly

$$T^{\mathrm{TE}} = \frac{2\mu k_z}{\mu k_z + k_{1z}}\frac{1+R_{12}}{1+R_{01}R_{12}e^{i2k_{1z}d}}e^{i(k_{1z}-k_z)d}, \tag{13}$$

$$T^{\mathrm{TM}} = \frac{2\varepsilon k_z}{\varepsilon k_z + k_{1z}}\frac{1+S_{12}}{1+S_{01}S_{12}e^{i2k_{1z}d}}e^{i(k_{1z}-k_z)d}. \tag{14}$$

Here $k_{1z} = \sqrt{k_1^2-k_\perp^2}$ and $k_1^2 = \varepsilon(\omega)\mu(\omega)\omega^2/c^2$. R_{ij} and S_{ij} are the reflection coefficients at the boundaries between media i and j for TE and TM modes, respectively,

$$R_{ij} = \frac{\mu_j k_{iz} - \mu_i k_{jz}}{\mu_j k_{iz} + \mu_i k_{jz}}, \quad S_{ij} = \frac{\varepsilon_j k_{iz} - \varepsilon_i k_{jz}}{\varepsilon_j k_{iz} + \varepsilon_i k_{jz}}. \tag{15}$$

The indices $i,j \in \{0,1,2\}$ denote again the spatial region, i.e., $k_0^2 = k_2^2 = k^2 \equiv \omega^2/c^2$ and $k_1^2 = \varepsilon(\omega)\mu(\omega)\omega^2/c^2$.

From expressions (9) and (10) one can calculate $\mathrm{Im}[\mathbf{G}(\mathbf{r}_k,\mathbf{r}_l,\omega_0)]$ for an ideal Veselago-Pendry lens, i.e., for infinite transversal extension and a lossless medium with $n(\omega_0) = -1$. Since in this case $R^{\mathrm{TE}} = R^{\mathrm{TM}} = 0$ one finds $\mathrm{Im}[\mathbf{G}^{00}(\mathbf{r},\mathbf{r},\omega_0)] = (k/6\pi)\hat{1}$, i.e., exactly the free-space value. Most importantly one finds that for all points $\mathbf{r}'$ in region 2 $(z' \leqslant -d)$

$$\mathrm{Im}\big[\mathbf{G}^{20}(\mathbf{r}',\mathbf{r},\omega_0)\big] = \mathrm{Im}\big[\mathbf{G}^{00}(\mathbf{r}'-2d\hat{\mathbf{z}},\mathbf{r},\omega_0)\big], \tag{16}$$

since $T^{\mathrm{TE}} = T^{\mathrm{TM}} = e^{ik_{1z}-k_z d}$ and $k_{1z} = -k_z$. The latter holds because $\mathbf{k}$ points backward in a LHM. Thus with respect to the

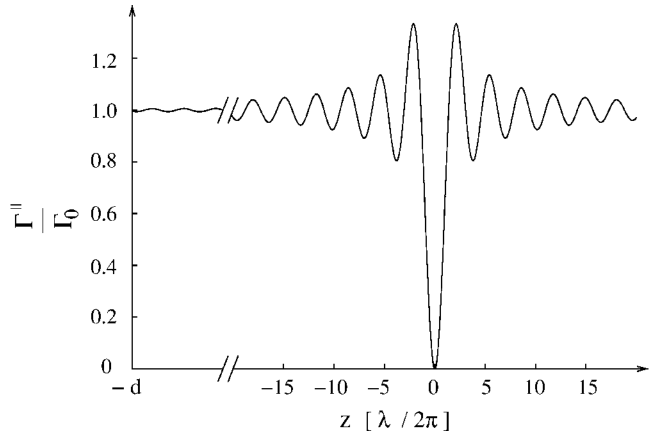

FIG. 2. Normalized rate of spontaneous emission $\Gamma^{\parallel}(z)/\Gamma_0$ as function of displacement from focus z. D is the distance from the focus to the surface of the LHM.

radiative decay, the second atom located in region 2 at $\mathbf{r}'$, i.e., on the other side of the Veselago-Pendry lens, behaves as if it would be located at position $\mathbf{r}'-2d\hat{\mathbf{z}}$, i.e., in region 0. This implies that for an atom pair in the focal points $\mathbf{r}'=\mathbf{r}+2d\hat{\mathbf{z}}$

$$\Gamma_{12} = \Gamma_{11}. \tag{17}$$

Thus the antisymmetric, single excited state $|a\rangle$ has vanishing radiative decay, while the symmetric state $|s\rangle$ decays with twice the free-space rate (6) and (7). That is, the pair of atoms shows perfect subradiance and superradiance.

A similar calculation yields the imaginary part of the Greens tensor for the case of an atom in front of a LHM combined with a mirror [Fig. 1(b)]. One finds for a dipole moment of the atom parallel or perpendicular to the mirror

$$\Gamma^{\parallel}(D\hat{\mathbf{z}},D\hat{\mathbf{z}}) = 0, \quad \Gamma^{\perp}(D\hat{\mathbf{z}},D\hat{\mathbf{z}}) = 2\Gamma_0, \tag{18}$$

Γ_0 being the vacuum decay rate. Figure 2 shows $\Gamma^{\parallel}$ as a function of the displacement of the atom from the focal point perpendicular to the mirror. One clearly sees a complete suppression of the spontaneous emission for dipoles parallel to the mirror at the focal point. In free space, i.e., without the LHM, a similar effect only occurs if the atom is located on the surface of the mirror [10].

Remarkably the property (16), and thus the existence of subradiant and superradiant states, Eq. (17), as well as the suppression of spontaneous emission, Eq. (18), do not seem to depend on the distance between the atoms or the distance of the atom from the mirror, provided the space in between is half filled by a LHM. In particular, in contrast to the free space, the strong radiative (self) coupling exists also over large distances compared to the transition wavelength due to the vanishing optical length of all pathways between the two foci.

While the decay properties are determined only by the Greens tensor at one frequency, the dipole-dipole shift $\delta\omega$ depends on the whole spectrum of the dielectric function $\varepsilon(\omega)$ and the magnetic permeability $\mu(\omega)$ and thus a general expression cannot be given. Using various single-resonance model functions for ε and μ, which fulfill the Kramers-Kronig relations we found values of $\delta\omega$ of up to $0.5\Gamma_{11}$. If $|\delta\omega| \gg \Gamma_{11}$ could be achieved, a perfect coherent excitation transfer between two atoms at the focal points would be possible without the use of a resonator. A more detailed study of the dipole-dipole shift in LHM will be the subject of further studies.

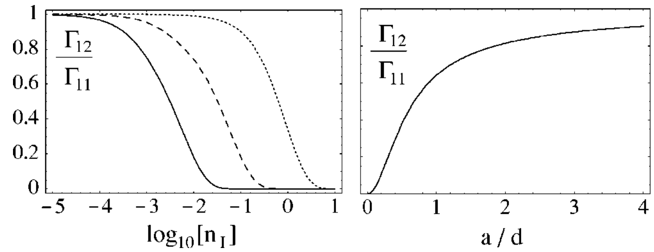

FIG. 3. Left: Γ_{12}/Γ_{11} as function of the imaginary part of the refractive index n_I for $\mathrm{Re}[n]=-1$ for different thicknesses d of the lens, $d=100\lambda/2\pi$ (solid line), $d=10\lambda/2\pi$ (dashed), and $d=1\lambda/2\pi$ (dotted). Right: Γ_{12}/Γ_{11} as a function of the transversal radius a.

Let us now discuss the limitations of the predicted effects. When the lens is not perfect, e.g., in the presence of losses, the ratio Γ_{12}/Γ_{11} decreases roughly exponentially with the increasing distance of the atoms and the subradiance and the superradiance effect disappears Fig. 3(a). The radiative coupling is also not perfect if the lens has only a limited transversal extension. It is not possible to give an analytical expression for the Greens tensor of a lens consisting of a disk of finite radius a. Also a numerical calculation of $\mathbf{G}$ for this case is quite difficult. One can, however, obtain an estimate of the effect if $d \gg \lambda$ by employing a short-wavelength or ray-optics approximation. Noting that for a lossless LHM with $n(\omega_0)=-1$, only propagating waves with $k_\perp \leqslant \omega_0/c$ contribute to $\mathrm{Im}[\mathbf{G}^{20}]$, we can model the effect of a finite transverse extension of the lens by restricting the $k_\perp$ integration in Eq. (10) to values $k_\perp \leqslant k[(a/d)/\sqrt{(1/4)+(a/d)^2}]$. The corresponding result is shown in Fig. 3(b). It is apparent that already a moderate ratio a/d is sufficient to obtain close to 100% suppression of the decay of the antisymmetric state $|a\rangle$.

There is another limitation of the predicted radiative effects even under ideal conditions, which arises solely from fundamental properties of LHM. If the slab of negative-index material has arbitrarily small losses at the frequency of interest and if it has a sufficiently large transversal extension, the previous discussion seems to suggest that subradiance and superradiance is possible for two atoms at an arbitrary distance. For causality reasons this should, of course, not be possible. Thus the question arises what is the maximum possible separation $2d$ of the atoms over which the strong radiative coupling persists? As pointed out already by Veselago [1], a lossless LHM is necessarily dispersive. The positivity of the electromagnetic energy in a lossless LHM requires that $(d/d\omega)\{\omega\,\mathrm{Re}[\epsilon(\omega)]\} \geqslant 0$, and $(d/d\omega)\{\omega\,\mathrm{Re}[\mu(\omega)]\} \geqslant 0$, which implies for $n(\omega_0)=-1$

$$\frac{d}{d\omega} n(\omega_0) \geqslant \frac{1}{\omega_0}. \tag{19}$$

As a consequence of the dispersion the frequency window $\Delta\omega$ over which $\mathbf{G}^{20}(\omega) \approx \mathbf{G}^{00}(\omega)$ narrows with increasing

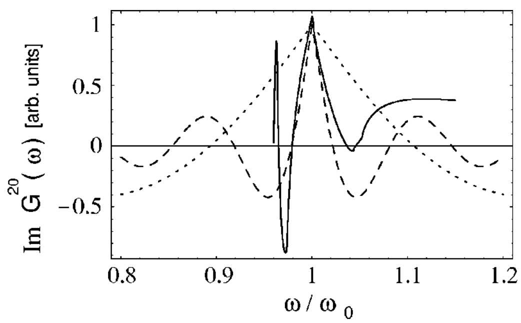

FIG. 4. Im$[\mathbf{G}^{20}(\omega)]$ following from Eq. (20) for lossless LHM with $n=-1+\alpha(\omega-\omega_0)$ for $\alpha=45/\omega_0$ for $dk_0=1$ (dashed), 0.2 (dotted). Also shown is a numerically calculated spectrum for a specific causal model for $n(\omega)$ with resonances of $\varepsilon(\omega)$ and $\mu(\omega)$ below ω_0. $n(\omega)$ was chosen such that Re$[n(\omega_0)]=-1$ and $\alpha=45/\omega_0$. The central structure is well represented by the linear-dispersion approximation (20).

thickness d of the lens. When $\Delta\omega$ becomes comparable to the natural linewidth of the atomic transition, Γ_{11}, the Markov approximation used in Eq. (2) is no longer valid. To give an estimate when this happens, we note from Eqs. (10)–(15) that for $d \gg \lambda$ the term in $\mathbf{G}^{20}$ that is most sensitive to dispersion is the exponential factor $e^{i\mathbf{K}\cdot(\mathbf{r}-\mathbf{r}')}e^{i(k_{1z}-k_z)d}$. Taking into account a linear dispersion in this factor, according to $n=-1+\alpha(\omega-\omega_0)$, with a real value of α, while keeping the resonance values in T^{TE}, T^{TM}, and R^{TE}, R^{TM}, one finds for the Greens tensor

$$\mathrm{Im}[\mathbf{G}^{20}(\omega)] = \frac{k}{8\pi}\mathrm{Re}\left[\int_0^1 d\xi(1+\xi^2)e^{i(dk_0/\xi)\alpha(\omega-\omega_0)}\right]\hat{\mathbf{1}}. \tag{20}$$

As can be seen from Fig. 4 the spectral width $\Delta\omega$ of the Greens function is in this approximation of order $\Delta\omega \approx (k_0 d\alpha)^{-1}$. Since Eq. (19) implies $\alpha \geq 1/\omega_0$, one arrives at $\Delta\omega \leq c/d$. This leads to an upper bound for the distance of the atoms, as $\Delta\omega \gg \Gamma_{11}$ implies

$$d \ll \frac{c}{\Gamma_{11}}. \tag{21}$$

This condition can easily be understood. It states that the distance between the two atoms must be small enough such that the travel time of a photon from one atom to the other is small compared to the free-space radiative lifetime.

In summary we have shown that the property of LHM to allow for a zero optical path length between macroscopically separated points in a space over a substantial range of frequencies can lead to interesting novel effects in the interaction of atoms with the quantized radiation field. Two atoms put in the focal points of an ideal, i.e., lossless Veselago-Pendry lens, exhibit perfect subradiance and superradiance as long as their distance is smaller than the propagation length of light corresponding to the free-space decay time. This effect can be used, e.g., to prepare a maximally entangled state between the two atoms in a similar way, as suggested in [11] for a cavity system. In addition an atom put at some distance from the surface of a mirror with a layer of LHM behaves as an atom directly on the mirror surface, i.e., shows complete suppression of spontaneous emission. Although negative refraction has been observed so far only in the cm-wave range [2–5] and would thus be limited to applications involving, e.g., Rydberg atoms, some ideas have been put forward to extend negative refraction to the Thz [12] and optical domain [13].

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[1] V. G. Veselago, Sov. Phys. Usp. **10**, 509 (1968).
[2] D. R. Smith, W. J. Padilla, D. C. Vier, S. C. Nemat-Nasser, and S. Schultz, Phys. Rev. Lett. **84**, 4184 (2000); R. A. Shelby, D. R. Smith, S. C. Nemat-Nasser, and S. Schultz, Appl. Phys. Lett. **78**, 489 (2001).
[3] R. A. Shelby, D. R. Smith, and S. Schultz, Science **292**, 77 (2001).
[4] A. Grbic and G. V. Eleftheriades, J. Appl. Phys. **92**, 5930 (2002).
[5] C. G. Parazzoli, R. B. Greegor, K. Li, B. E. C. Koltenbah, and M. Tanielian, Phys. Rev. Lett. **90**, 107401 (2003).
[6] J. B. Pendry, Phys. Rev. Lett. **85**, 3966 (2000).
[7] R. H. Dicke, Phys. Rev. **93**, 99 (1954).
[8] Ho Trung Dung, S. Y. Buhmann, L. Knöll, D.-G- Welsch, S. Scheel, and J. Kästel, Phys. Rev. A **68**, 043816 (2003).
[9] L. Tsang, Jin Au Kong, and R. T. Shin, *Theory of Microwave Remote Sensing* (Wiley, New York, 1985).
[10] H. Morawitz, Phys. Rev. **187**, 1792 (1969).
[11] M. B. Plenio, S. F. Huelga, A. Beige, and P. L. Knight, Phys. Rev. A **59**, 2468 (1999).
[12] T. J. Yen *et al.*, Science **303**, 1494 (2004).
[13] M. Ö. Oktel and Ö. E. Müstecaplioğlu, e-print physics/0406039.