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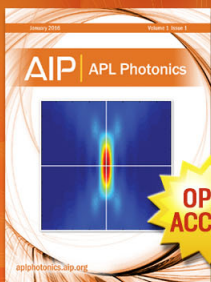
Near-field enhancement and imaging in double planar polariton-resonant structures

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It is shown that a system of two coupled planar material sheets possessing surface mode (polariton) resonances can be used for the purpose of evanescent field restoration and, thus, for sub-wavelength near-field imaging. The sheets are placed in free space so that they are parallel and separated by a certain distance. Due to interaction of the resonating surface modes (polaritons) of the sheets an exponential growth in the amplitude of an evanescent plane wave in the system can be achieved. This effect was predicted earlier for backward-wave (double-negative or Veselago) slab lenses. The alternative system considered here is proved to be realizable at microwaves by grids or arrays of resonant particles. The necessary electromagnetic properties of the resonating grids and the particles are investigated and established. Theoretical results are supported by microwave experiments that demonstrate amplification of evanescent modes. © *2004 American Institute of Physics.* [DOI: 10.1063/1.1765865]

I. INTRODUCTION

Negative refraction and subwavelength resolution in slab lenses based on backward-wave (also called *double-negative* or *Veselago*[1]) materials are the topics that have been widely discussed in the recent literature. The discussion started soon after Pendry had published his article[2] claiming a possibility to overcome the $\lambda/2$ optical resolution limit using a slab of Veselago material with the relative parameters $\epsilon_r = -1$ and $\mu_r = -1$. The result came out of the discovery that a Veselago slab[1] can "amplify" exponentially decaying evanescent modes of a source field. The slab restores the amplitudes of these modes in the image plane recovering the fine details of a source.

It was no wonder that many scientists, especially those from the optical community, found this idea of "amplification" difficult to accept.[3] Indeed, if one assumes an exponential growth of the field amplitude in a slab of a finite thickness, it seems that by increasing the slab thickness the amplitude of the outgoing field can be made arbitrarily high. However, an accurate analysis shows that if the slab relative permittivity and permeability are not exact −1 (due to inevitable losses and dispersion) that cannot happen. Indeed, for a slab of thickness d of Veselago material characterized by ϵ_r, μ_r the slab transmission coefficient can be found using the standard procedure of expressing the slab fields in terms of two oppositely propagating (decaying) waves and finding the unknown wave amplitudes by solving a system of four equations coming from the boundary conditions on the slab interfaces. For excitation by an evanescent plane wave with the tangential component of the propagation factor k_t we have

$$T = \frac{2}{2\cosh(\alpha d) + (\gamma + 1/\gamma)\sinh(\alpha d)}. \quad (1)$$

Here $\alpha = \sqrt{k_t^2 - \epsilon_r \mu_r k_0^2}$ is the decay factor in the slab, $\alpha_0 = \sqrt{k_t^2 - k_0^2}$ is the same for free space, $\gamma = \alpha_0 \mu_r / \alpha$ for TE incidence or $\gamma = \alpha / \alpha_0 \epsilon_r$ for TM case, and $k_0 = \omega\sqrt{\epsilon_0 \mu_0}$. Assuming $\epsilon_r = \mu_r = -1$, we see that $\gamma = -1$ and $T = \exp(\alpha_0 d)$ as in the ideal Pendry's case. However, if the permittivity and (or) permeability differ from that very value, then $\gamma \neq -1$ and $|\gamma + 1/\gamma| > 2$, resulting in domination of the growing exponent in the denominator of Eq. (1) when the slab thickness and (or) the incidence field spatial frequency become large enough. We see that the region where the evanescent fields are indeed amplified in a Veselago slab is limited by several factors. Some of them are even inevitable in any realistic material, e.g., losses and spatial dispersion.[4]

For the following, it is important to understand what main phenomena lead to amplification of evanescent modes in the Veselago slab. For this purpose, we will shortly review the plane wave incidence problem for an interface of free space and a half-space filled by a Veselago material. In the material we will look for a solution which, as usual, is either exponentially decaying (if the transmitted wave is evanescent) or transmitting energy *from* the interface (if the wave is propagating).

If A, B, and C denote the complex amplitudes of the incident, transmitted, and reflected wave electric field tangential components, respectively, then using the interface boundary conditions we can write:

$$A + C = B,$$

$$\frac{A - C}{\eta_0} = \frac{B}{\eta}. \quad (2)$$

Here η_0 and η are the wave impedances that connect the tangential components of electric and magnetic fields in free

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space and in the medium, respectively. The solution of Eq. (2) is, obviously,

$$C = \frac{\eta - \eta_0}{\eta + \eta_0} A, \quad B = \frac{2\eta}{\eta + \eta_0} A. \quad (3)$$

The wave impedance of propagating transmitted waves is given by $\eta = \omega\mu_0\mu_r/k_n$ for TE waves, and $\eta = k_n/(\omega\epsilon_0\epsilon_r)$ for TM waves [$k_n = \sqrt{\epsilon_r\mu_r k_0^2 - k_t^2}$ denotes the normal to the interface wave vector component; the formula applies for passive lossy materials with $\text{Im}\{\epsilon_r, \mu_r\} < 0$ (or $\text{Im}\{\epsilon_r, \mu_r\} \rightarrow -0$) if the square root branch is chosen so that $\text{Im}\{k_n\} < 0$ (or $\text{Im}\{k_n\} \rightarrow -0$); the time dependence is in the form $\exp(+j\omega t)$].

In a Veselago medium both ϵ_r and μ_r are negative. In the same time, k_n is also negative because the propagating waves are backward waves. Nothing especially interesting comes out of Eq. (3) in this case except that when $\epsilon_r = \mu_r = -1$ the interface is perfectly matched: $\eta = \eta_0$, $C = 0$, $B = A$.

But let us suppose that the incident and transmitted waves are evanescent. Then, $\eta = j\omega\mu_0\mu_r/\alpha$ for TE waves, and $\eta = \alpha/(j\omega\epsilon_0\epsilon_r)$ for TM waves. Because the transmitted wave must decay from the interface, α is positive. We see that for evanescent modes the ideal case when $\epsilon_r = \mu_r = -1$ leads to purely imaginary wave impedances such that $\eta = -\eta_0$! A resonance occurs: $C = B \rightarrow \infty$. The reason for such resonant growth of the field amplitudes when $\epsilon_r, \mu_r \rightarrow -1$ is in the excitation of a surface mode (surface polariton) associated with the interface. Indeed, if there is no incident field in Eq. (2) ($A = 0$) and $\epsilon_r = \mu_r = -1$ we can observe that for any $k_t > k_0$ (imaginary wave impedances) there is a solution $C = B \neq 0$ corresponding to a surface wave concentrated near the interface.

Based on similar considerations several authors[5,6] explained the evanescent mode amplification in the Pendry lens as the result of resonant excitation of a pair of coupled surface modes (polaritons) sitting at the slab interfaces. Under certain conditions the polariton excited at the second (output) interface is much stronger than that excited at the first interface. The exponentially decaying trail of the polariton sitting at the output interface appears as an exponential growth of the field inside the slab.

The effects taking place in the material depth (backward waves) and the properties of the slab interfaces (polariton resonances) both contribute to the Pendry's lens operation. However, it can be shown that in general the presence of a bulk material layer is not crucial. Conceptually, if one can realize a planar sheet such that traveling waves refract negatively when crossing this sheet, a system of two such sheets placed *just in free space* will focus the propagating modes of a source just like a Veselago slab. If the sheets also support surface waves for all $k_t > k_0$ then such a system will possess surface polariton resonances reconstructing the evanescent spectrum as well. We found in our recent paper[7] that a system of two phase conjugating interfaces in air behaves as a perfect lens. A possible drawback of phase conjugating design is the necessity to utilize nonlinear effects like wave mixing. In this paper we will discuss alternative possibilities to evanescent spectrum reconstruction not involving nonlinearity. The design will be based on the above-mentioned

principle: We will make use of a couple of polariton-resonant surfaces or grids placed in free space. No bulk backward-wave materials will be involved, providing more flexibility and less limitations in design.

II. ANALYSIS BASED ON TRANSMISSION MATRICES

In this section and in what follows we restrict our consideration by the evanescent spectrum only. Our purpose here will be to find such conditions on resonating sheets that lead to "amplification" of the evanescent modes in the proposed double-grid system. We will refer to the system simply as *device*. A possible name for such a device can be *near-field lens*, but we would prefer not to use the word *lens* in this context to avoid misunderstanding. Let us emphasize that our aim here is the restoration of the near-field or evanescent field picture of a source. The systems to be considered in the following do not focus propagating modes. This can be done by other well-known optical means.

We will make use of a powerful method based on so-called 2×2 wave transmission matrices, well known in the microwave circuit theory.[8] These matrices connect the complex amplitudes of waves traveling (or decaying) in the opposite directions in a waveguiding system or a system where one can determine the principal axis of propagation and measured at two reference planes:

$$\begin{pmatrix} E_2^- \\ E_2^+ \end{pmatrix} = \begin{pmatrix} t_{11} & t_{12} \\ t_{21} & t_{22} \end{pmatrix} \begin{pmatrix} E_1^- \\ E_1^+ \end{pmatrix}. \quad (4)$$

Here, $E_1^{\pm}$ and $E_2^{\pm}$ denote the tangential components of the electric field complex amplitudes of waves at the first (input) and the second (output) interfaces of a device, respectively (we restrict ourselves by planar layered structures and plane waves). The signs $\pm$ correspond to the signs in the propagator exponents $e^{\pm jk_n z}$ of these waves, and z is the axis orthogonal to the interfaces (the main axis of the system). It is known that the T-matrix of a serial connection of several devices described by their T-matrices is simply a multiplication of the matrices in the order determined by the connection.

Our purpose is to build a theoretically ideal near-field imaging device. Hence, the total transmission matrix from the source plane to the plane where the source field distribution is reconstructed must be the identity matrix

$$T_{\text{tot}} = T_{\text{sp after}} T_{\text{dev}} T_{\text{sp before}} = \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix} \quad (5)$$

for every spatial harmonic of the source field. Here, $T_{\text{sp before}}$ and $T_{\text{sp after}}$ represent the air layers occupying the space between the source plane and the device, and the space between the device and the image plane. T_{dev} is the transmission matrix of the device. From this formula it is obvious that a *complete* reconstruction of the field distribution in the source plane at a distant image plane must involve phase compensation for the propagating space harmonics and "amplification" for the evanescent ones. In other words, one needs to synthesize a device that somehow inverts the action of a free-space layer. A slab of a material with $\epsilon_r = -1$ and $\mu_r = -1$ (Veselago medium)[1,2] and a pair of parallel conjugat-

ing surfaces or sheets[7] operate as such a device. In this paper we will find other *linear* solutions working for the evanescent fields of a source.

Let us note here that condition (5) is a strict condition requiring not only the one-way transmission to be such that it reconstructs the source field picture at the image plane, but also the matching to be ideal (no reflections from the device) and the device operation to be symmetric (reversible in the optical sense). We will consider some less strict conditions later.

Let us suppose that the source and the image planes are distanced by $d/2$ from the input and the output interfaces of the device. A space layer of thickness $d/2$ has the T-matrix

$$T_{sp}(d/2) = \begin{pmatrix} \exp(-jk_n d/2) & 0 \\ 0 & \exp(+jk_n d/2) \end{pmatrix}. \quad (6)$$

To compensate the action of two such layers before and after the device and satisfy condition (5), the device T-matrix T_{dev} has to be, obviously, the inverse of the transmission matrix of these space layers:

$$T_{dev} = \begin{pmatrix} \exp(+jk_n d) & 0 \\ 0 & \exp(-jk_n d) \end{pmatrix}. \quad (7)$$

Let us study if a device modeled by this transmission matrix can be realized as a combination of two "field transformers" (e.g., thin sheets of certain electromagnetic properties) separated by a layer of free space, like it was discussed in Sec. I. This system is modeled by the transmission matrix

$$T_{dev} = T_{out} T_{sp}(d) T_{in} = \begin{pmatrix} a & b \\ c & d \end{pmatrix} \begin{pmatrix} \exp(-jk_n d) & 0 \\ 0 & \exp(+jk_n d) \end{pmatrix} \begin{pmatrix} e & f \\ g & h \end{pmatrix}. \quad (8)$$

Here, T_{in} and T_{out} are matrices with yet unknown components describing the two sheets or layers forming the device, and $T_{sp}(d)$ is the matrix of the free-space layer between the sheets. It is easy to show that if a $a=d=0$, $e=h=0$, and $bg=cf=1$, then the total device T-matrix takes form (7), i.e., it is the necessary matrix of a perfect lens. From the mathematical point of view such an amazing result is simply an effect of permutation of the matrix components under the multiplication (8). The physical question which we will need to answer later is how to realize an interface with a T-matrix of the form

$$T = \begin{pmatrix} 0 & b \\ c & 0 \end{pmatrix}. \quad (9)$$

We will show that even this ultimate case is realizable (for excitation by evanescent fields) by passive grids with specific electric and magnetic susceptibilities.

However, let us consider a couple of simpler systems also, resulting from the conditions less strict than Eq. (5). We keep the same operation principle defined by Eq. (8). If we allow a mismatch at the device interfaces still maintaining the device symmetry, the following solution can be found:

$$T_{in} = T_{out} = \begin{pmatrix} a & b \\ c & 0 \end{pmatrix}. \quad (10)$$

If $b=-c=1$, we have for the device

$$T_{dev} = \begin{pmatrix} a^2 \exp(-jk_n d) - \exp(+jk_n d) & a \exp(-jk_n d) \\ -a \exp(-jk_n d) & -\exp(-jk_n d) \end{pmatrix}, \quad (11)$$

which corresponds to the device scattering matrix (well-known S-matrix) of the form

$$S_{dev} = \begin{pmatrix} -t_{21}/t_{22} & 1/t_{22} \\ t_{11} - t_{12}t_{21}/t_{22} & t_{12}/t_{22} \end{pmatrix}$$
$$= -\begin{pmatrix} a & \exp(+jk_n d) \\ \exp(+jk_n d) & a \end{pmatrix}. \quad (12)$$

The S-matrix elements are the reflection and transmission coefficients for two "ports" of our device. One can see that the device "amplifies" evanescent modes (due to pluses in the s_{21} and s_{12} exponents) and reflects in both ports with the reflection coefficient equal to $s_{11} = s_{22} = -a$.

If the symmetry is not important but the matching is, the following solution is possible:

$$T_{out} = \begin{pmatrix} a & b \\ c & 0 \end{pmatrix}, \quad T_{in} = \begin{pmatrix} 0 & f \\ g & h \end{pmatrix}. \quad (13)$$

If $bg=cf=1$, the total T-matrix for the device becomes

$$T_{dev} = \begin{pmatrix} \exp(+jk_n d) & af \exp(-jk_n d) + bh \exp(+jk_n d) \\ 0 & \exp(-jk_n d) \end{pmatrix}, \quad (14)$$

which corresponds to the scattering matrix of the form

$$S_{dev} = \begin{pmatrix} 0 & \exp(+jk_n d) \\ \exp(+jk_n d) & bh \exp(+2jk_n d) + af \end{pmatrix}. \quad (15)$$

One can see that the device is matched for the waves coming to the first interface ($s_{11}=0$) and also "amplifies" the evanescent modes. In Sec. III we will describe the ways to realize matrices (9), (10), and (13).

III. THE USE OF IMPEDANCE SHEETS

At first we consider a simple system: a lossless isotropic grid, e.g., a conductive wire mesh (possibly loaded by certain bulk reactances inserted in every cell). If the grid supports only electric currents, and there is no effective magnetic current induced on the grid, then the grid reflection coefficient R and transmission coefficient T at the grid plane are related as

$$T = 1 + R \quad (16)$$

provided that they are defined through the electric field tangential components (for a given polarization). The corresponding T-matrix of such a grid is

$$T_g = \begin{pmatrix} \dfrac{1+2R}{1+R} & \dfrac{R}{1+R} \\ -\dfrac{R}{1+R} & \dfrac{1}{1+R} \end{pmatrix}. \quad (17)$$

It is possible to make grids supporting propagation of surface modes (also known as *slow waves* in radio engineering). For

wire meshes, for example, this phenomenon is well investigated.[9] If the tangential component of the wave vector of an incident wave coincides with the propagation factor of a surface mode, the surface mode resonance takes place. Obviously, the incident wave should be evanescent in this case to match with the propagation constant of the surface mode. At a surface mode resonance $R \to \infty$ (for evanescent modes R is not bounded by $|R| \leq 1$). Then, the grid T-matrix takes the form

$$T_g = \begin{pmatrix} 2 & 1 \\ -1 & 0 \end{pmatrix}, \tag{18}$$

which is of the necessary form (10).

For a better understanding we reformulate the above-mentioned consideration in terms of the grid impedance. If the boundary condition on the grid is given as $E_t = Z_g J$, where J is the averaged electric surface current density induced on the grid, E_t is the averaged tangential electric field in the grid plane, and Z_g is the grid impedance, the reflection coefficient can be found as[10]

$$R = -\left(1 + \frac{2Z_g}{\eta_0}\right)^{-1} \tag{19}$$

and the grid transmission matrix becomes

$$T_g = \begin{pmatrix} 1 - \dfrac{\eta_0}{2Z_g} & -\dfrac{\eta_0}{2Z_g} \\ \dfrac{\eta_0}{2Z_g} & 1 + \dfrac{\eta_0}{2Z_g} \end{pmatrix}. \tag{20}$$

The reflection coefficient (19) becomes infinite and the transmission matrix (20) reduces to Eq. (18) when ω and k_t satisfy

$$Z_g(\omega, k_t) + \frac{\eta_0(\omega, k_t)}{2} = 0, \tag{21}$$

which is the dispersion equation for surface modes on the grid surface. Because η_0 is purely imaginary for evanescent modes (it is inductive for TE waves and capacitive for TM waves) one can see that in principle there are no restrictions on realizing a capacitive or inductive grid or array possessing the necessary resonance for some value(s) of k_t.

We can see from Eq. (20) that it is enough to change the sign of the grid impedance to realize the second matrix in Eq. (13). Such a grid is not at resonance with the incident evanescent field, and it works as an additional matching layer or a load for the output grid which must experience a strong resonance in accordance with Eq. (13).

Let us now consider a more complicated grid or array that supports both electric and magnetic currents. We suppose that the electric current is excited by electric fields in the array plane and the magnetic current is due to magnetic fields at the same plane. In the presence of two currents, the tangential components of both electric and magnetic fields are not continuous across the interface:

$$E_1 - E_2 = J_m, \quad H_1 - H_2 = J_e, \tag{22}$$

where J_e and J_m stand for the averaged electric and magnetic surface current densities. The following conditions determine the current amplitudes in terms of two grid impedances Z_e and Z_m:

$$\frac{E_1 + E_2}{2} = Z_e J_e, \quad Z_m \frac{H_1 + H_2}{2} = J_m. \tag{23}$$

It is possible to show that an interface defined by the above-noted conditions has the following T-matrix:

$$T_g = \left[\frac{4Z_e}{Z_m} - 1\right]^{-1}$$
$$\times \begin{pmatrix} \left[1 - \dfrac{2Z_e}{\eta_0}\right]\left[1 - \dfrac{2\eta_0}{Z_m}\right] & 2\left[\dfrac{Z_e}{\eta_0} - \dfrac{\eta_0}{Z_m}\right] \\ -2\left[\dfrac{Z_e}{\eta_0} - \dfrac{\eta_0}{Z_m}\right]\left[1 + \dfrac{2Z_e}{\eta_0}\right] & \left[1 + \dfrac{2\eta_0}{Z_m}\right] \end{pmatrix}. \tag{24}$$

Equation (20) is a particular case of Eq. (24) when Z_m tends to zero and the magnetic current vanishes. The matrix (24) reduces to form (9) in two cases. First, this happens when the electric subsystem is at resonance: $Z_e = -\eta_0/2$ and the magnetic subsystem works as a loading: $Z_m = 2\eta_0$. In this case

$$T_g = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}. \tag{25}$$

Or, second, when the roles are interchanged: $Z_e = \eta_0/2$, $Z_m = -2\eta_0$ For this case

$$T_g = \begin{pmatrix} 0 & -1 \\ 1 & 0 \end{pmatrix}. \tag{26}$$

Let us note again that nothing forbids realization of the necessary impedances for a given k_t, because for the evanescent mode they are purely imaginary.

IV. ARRAYS OF WEAKLY INTERACTING RESONANT INCLUSIONS

In the previous sections we have shown that the surface mode resonance plays the key role in the mechanism of the evanescent field amplification. We also have proven that such a resonance is possible to realize in a passive grid or array. The resonance occurs when the (ω, k_t) pair belongs to the polariton spectrum of the grid.

Practically speaking, this means that for a given frequency ω one may realize one or at most several resonant values of the transverse propagation constant k_t. It can be enough for some purposes as, for example, for resonant extraction and "amplification" of a certain spatial harmonic of the incident field. However, for a device operating as a near field lens one should provide as wide a range of operable k_t as possible.

Mathematically (and ideally), the latter means that the dispersion equation (21) should be somehow turned at a given frequency into an identity for any $k_t > k_0$. Although that is impossible in practice, there is a good approximation for this. Let us consider a dense regular two-dimensional array of small resonant dipole inclusions. Instead of writing

the boundary condition in terms of the *total* averaged field in the array plane as before, we may solve the excitation problem directly in terms of the *external* field and the induced dipole moments. In a given external field E_{ext}, the dipole moment of each particle in the array is[10]

$$p = \chi(\omega)[E_{ext} + \beta(\omega,k_t)p]. \qquad (27)$$

Here χ is the particle polarizability, and $\beta(\omega,k_t)$ is the so-called interaction factor, which is a function of k_t. Obviously, the solution for the induced dipole moment is

$$p = \frac{E_{ext}}{\dfrac{1}{\chi(\omega)} - \beta(\omega,k_t)}. \qquad (28)$$

From here it is already seen what condition is needed for a resonance, but let us proceed a bit further. In terms of the average surface current density $J = j\omega p/S_0$ (S_0 is the unit cell area) we write

$$Z_c J = E_{ext}, \qquad (29)$$

where

$$Z_c = \frac{S_0}{j\omega}\left[\frac{1}{\chi(\omega)} - \beta(\omega,k_t) \right]. \qquad (30)$$

Here we have introduced a quantity which we call *cell impedance* Z_c. This impedance is related with the grid impedance. Indeed, because the total tangential electric field at the array plane is $E = E_{ext} - \eta_0 J/2$, we have

$$Z_g = Z_c - \frac{\eta_0}{2}. \qquad (31)$$

Comparing the last relation with the resonance condition (21) we see that the polariton resonance takes place when $Z_c = 0$. Of course, the same conclusion follows directly from looking at the denominator of Eq. (28).

The imaginary parts of $1/\chi$ and β cancel out in a lossless nonradiating array[10] (we work with evanescent fields). In order to resonate at all harmonics of the evanescent field spectrum, the resonant condition $Z_c = 0$ must be satisfied for all k_t. Because the inclusions are assumed to be small dipole particles, their polarizability χ depends on the frequency, but does not depend on k_t. On the other hand, the interaction constant β depends both on the frequency and on the transverse wave number. Thus, the only possibility to realize such grids using small inclusions is to use resonant particles (Re$\{1/\chi\} \to 0$) and minimize the field interactions between the particles in the array [provide Re $\beta(\omega,k_t) \to 0$]. If these conditions are satisfied, each particle in the grid is excited *locally* by the incident field at its position (if the field interaction in the array is negligible, the local field equals the incident field). Since at the operational frequency the particles are at resonance, *arbitrary* spatial distributions of the incident field will excite a resonance of the whole array. In other words, at this frequency the grid indeed supports polaritons with arbitrary $k_t > k_0$, as needed for evanescent field amplification and imaging.

In a real system particles always interact, and the last condition cannot be exactly satisfied. However, it is possible
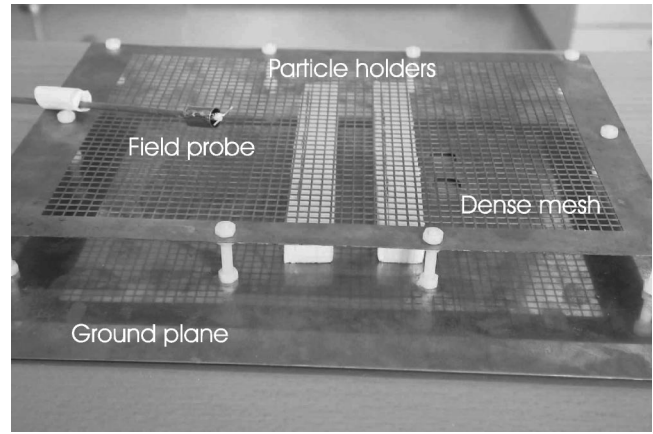

FIG. 1. Photo of the experimental setup. Resonant particles were positioned on two parallel foam holders located between two highly conducting planes. The probe used to scan the field distribution is seen on the top. A similiar antenna was used as the source. (Not shown on the photo. It was positioned between the conducting planes.)

to reduce or compensate interactions at a certain frequency. In Sec. V we will describe an experiment based on implementing an array of highly resonant but weakly interacting particles. The experiment will demonstrate evanescent field enhancement in a passive linear system. Another possibility is to use inclusions of larger sizes, and try to compensate for the spatial dispersion of the particle interaction with the spatial dispersion of the response of one inclusion. We will not explore this last possibility in this paper.

V. EXPERIMENT

In the microwave experiment (the operating frequency was close to 5 GHz), evanescent fields were generated in the space between a metal plate and a dense mesh of conducting strips forming a two-plate waveguide, see Fig. 1. Microwave absorbers were used around the system to minimize reflections from the open ends of the waveguide. The upper screen was made weakly penetrable to the fields in order to give us a possibility to measure the field distribution by a probe positioned on top of the mesh. The transmission coefficient of this mesh (for normal plane-wave incidence) at the operation frequency was about −20 dB. As a source, we used a wire dipole antenna whose length was close to $\lambda/4$. The dipole was parallel to the conducting plates, so that only evanescent modes were excited in the space between the plates (the distance between the plates was 2.5 cm, so that all the waveguide modes of this polarization were evanescent). Let us note that originally we planned to work with vertically polarized fields (excited by a dipole orthogonal to the plates). However, we soon noticed that even in the presence of absorbers around the setup the standing wave pattern created by propagating modes was quite disturbing.

As was established above, to realize a device that would "amplify" evanescent fields we need to design an array of small resonant particles that weakly interact. To validate this concept, one can minimize interactions between particles by simply increasing the distance between the particles in the array. In our first experiment, we measured fields in a system of only *two* resonant particles, which corresponds to the case
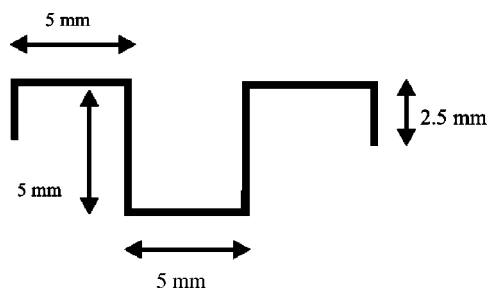
FIG. 2. Small resonant particle.

FIG. 4. The distribution of evanescent field created by a small dipole antenna in the presence of two small resonant particles. The source dipole is placed at $x=-125$ mm, $y=55$ mm and directed along the y axis. Two metal particles (Fig. 2) are placed at $x=-105$ mm, $y=55$ mm and $x=-65$ mm, $y=55$ mm. The particles are oriented along the y axis. The frequency is 5 GHz. The probe is 0.5 cm away from the top mesh of the setup. The field amplitude scale is linear.

of two parallel arrays with infinitely large periods. The particles were made of a copper wire of 0.8 mm diameter, and their shape and dimensions were as shown in Fig. 2. The wire was meandered in order to make the overall dimensions small as compared with the wavelength. The stretched wire length was close to $\lambda/2$, so the particles showed resonant response. The positions of the two particles in holders with respect to the field source are shown in Fig. 3.

The experimental results are shown in Figs. 4 and 5. It can be seen that, as expected from the theory, the first particle is very weakly excited, and a high-amplitude plasmon polariton is sitting at the second particle. The field amplitude in the "image plane" is close to that at the source position. However, since we have only one particle in each of the two "arrays," what we see in these figures is not an image in the usual meaning. Basically, these are only the field distributions in the vicinity of a couple of resonating particles. This device detects and enhances the incident evanescent field, but it cannot resolve the field distribution in the source plane. However, certain conclusions can be made even here.

First, we see that the "image" field is concentrated in the area whose size is of the same order as the size of the resonating particle itself. The size of the particle limits the device resolution at least by $k_t a < \pi$, where a is the characteristic dimension of the particle (in the direction of $\mathbf{k}_t$). Second, in the present setup the dimensions of the particles (≈ 15 mm) are "in resonance" with the source dimensions (the source $\lambda/4$ dipole is about 15 mm in length too). This provides an effective coupling because the same spatial harmonic dominates in the source field and in the field resonantly amplified by the particles.

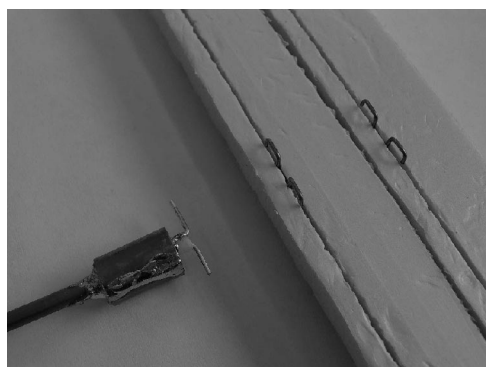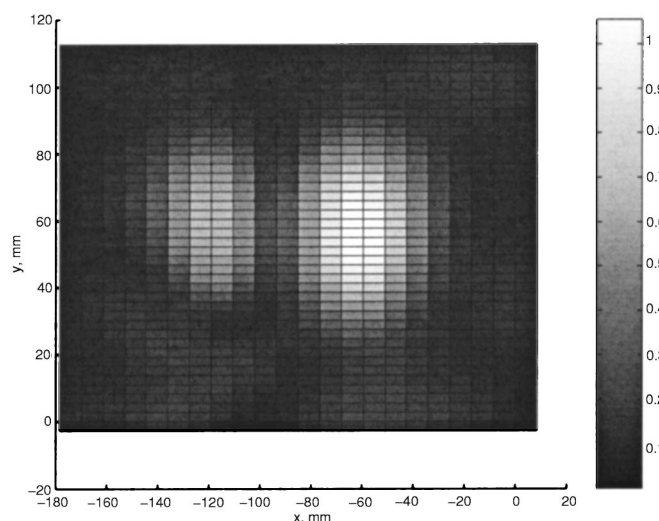Although this experiment demonstrates the validity of

the main principle of near-field enhancement and a possibility to restore the evanescent field components, grids with reasonably small periods are necessary to realize an imaging device. In such devices, the resolution is limited by the grid period: the maximum transverse wave number resolved by a pair of grids is about $k_t D \approx \pi$, where D is the grid period. This is the same limitation as in perfect lenses realized as slabs of artificial left-handed materials.

To study phenomena in such grids, we have made measurements in a system of two regular arrays of similar par-

FIG. 3. Position of two small particles in foam holders and the source dipole antenna with a quarter-wave balun.
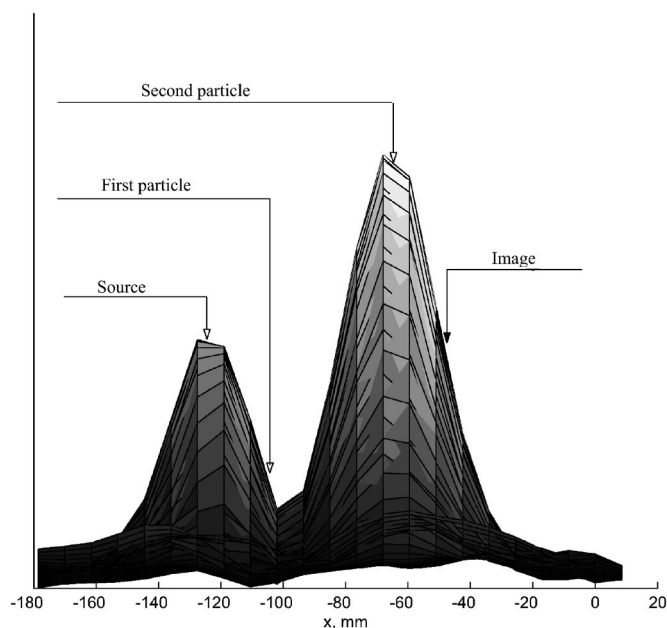

FIG. 5. Dependence of the field amplitude along the device axis for the same arrangement as in Fig. 4. This is a side view of the three-dimensional plot of the field distribution shown in Fig. 4. The key positions on the plot are indicated by arrows.
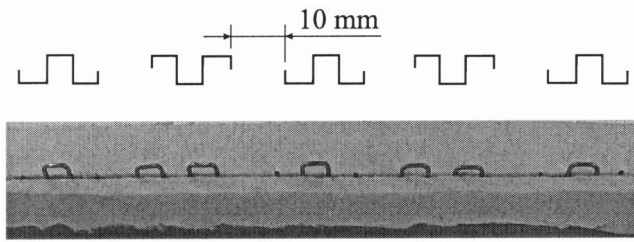
FIG. 6. A periodic array of small resonant particles. The photo shows the same array in a foam holder.

ticles. The array geometry is shown in Fig. 6. In this system, the field interaction of particles exists, meaning that maintaining polariton resonance for all transverse wave numbers is not possible. In the measurements, we first experimentally determined the eigenfrequencies of the grids. Each eigenfrequency corresponds to a certain spatial profile of the amplitudes of the dipole moments excited in the grids, i.e., to a certain transverse wave number k_t. Next, we excited the grids at each of these frequencies (two lowest frequencies were practically used) and scanned the fields. To get a complete image we superimposed the measured spatial profiles. This corresponds to reconstructing the source spatial spectrum using only a few spatial harmonics. The results for the case of excitation by two small dipole antennas are shown in Figs. 7 and 8, for two separate eigenfrequencies. The superimposed profile is given in Fig. 9. Strong excitation of the second grid is clearly visible, as well as an image of the source field behind the grids. Note that the amplitude of the field excited at the second resonant mode (Fig. 8) is much higher than that at the first one, so the total superimposed image is similar to that in Fig. 8. This last experiment should be considered as a first step only, because no effective reduction of the field interactions between the grid particles was realized, and the

FIG. 8. The same as in Fig. 7, but the field complex values are measured at 5.26 GHz (second polariton resonance).

setup had many nonidealities. However, we can conclude that the experiments successfully validate the principle of near-field enhancement in simple passive and linear resonant systems.

VI. CONCLUSIONS

In this paper we have considered a wide class of passive linear structures able to enhance evanescent fields and reconstruct the near-field image of a source. All these structures result from the idea of using a system of two parallel polariton-resonant grids or arrays separated by a certain distance and placed in free space.

The physics behind this idea is based on the known behavior of coupled resonant systems. If in a system of two resonators the first resonator is pumped by an external force and the second resonator is coupled to the first one, then under certain conditions it is possible for the amplitude of oscillations in the second resonator to be much higher than
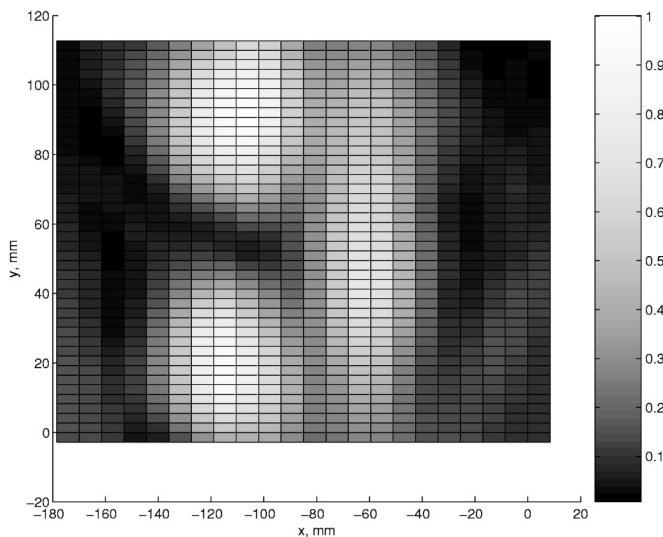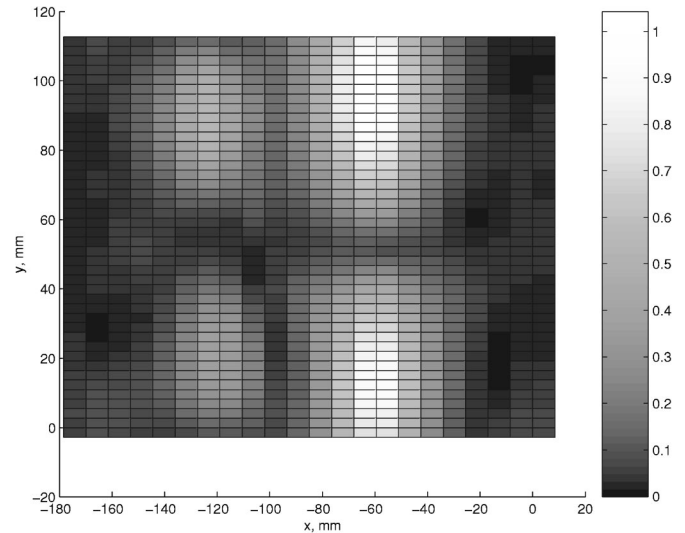

FIG. 7. The distribution of the evanescent field created by two small dipole antennas in the presence of two grids of small resonant particles. The field complex values are measured at 5.15 GHz (the first polariton resonance). Two source dipoles are placed at $x=-125$ mm, $y=20$ mm and at $x=-125$ mm, $y=90$ mm. The dipoles are oriented along the y axis. Two grids (5 particles in each) are placed at $x=-105$ mm, $y=55$ mm and $x=-65$ mm, $y=55$ mm along the y axis. The probe is 0.5 cm away from the top mesh of the setup. The field scale is linear.
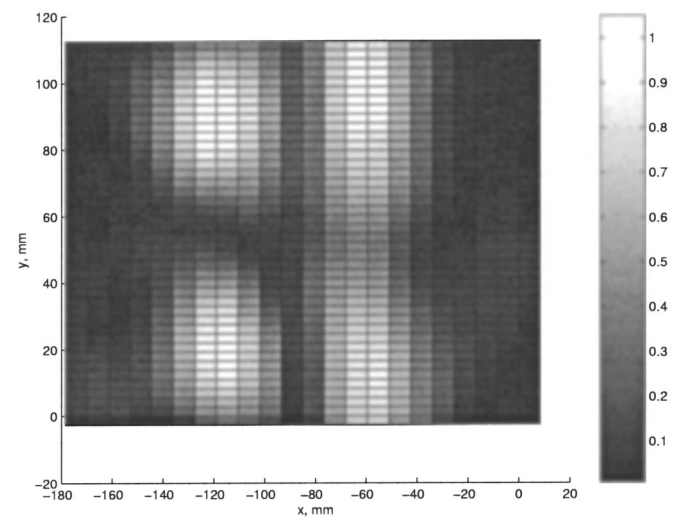

FIG. 9. The complete field profile calculated as a superposition of the data from Figs. 7 and 8.

the amplitude of the external field and the amplitude of the first resonator oscillations. A similar interpretation of the phenomena taking place in a coupled-polariton-resonant system (the Veselago slab) can be found in a recent work by Rao and Ong.[6]

The present paper continues the research we started in order to eliminate the need in backward-wave or other exotic *bulk material layers* in the design of near-field imaging devices. In our recent work[7] we showed that a system of two phase-conjugating planes or sheets placed in free space behaves as a perfect lens proposed by Pendry.[2] The obvious drawback of the phase-conjugating design is the necessity to involve nonlinear materials or devices in the structures realizing the conjugating sheets. Here we have shown that if the focusing of the propagating spectrum is not required, there exist several *linear* solutions.

We have developed a general synthesis approach based on the wave transmission matrices[8] to find the conditions under which a system of two coupled polariton-resonant grids or arrays enhances incident evanescent field. Next, the inner design of the grids has been revealed, and it has been shown that there are many possibilities arising from the use of impedance sheets. The necessary structures can be realized as arrays of weakly interacting resonant particles of different nature.

We have experimentally confirmed the possibility to use passive linear systems for evanescent field enhancement at microwaves. The experiment supports the main concepts of our theoretical findings. The resonant growth of evanescent fields in the system of two coupled polariton-resonant grids has been observed.

ACKNOWLEDGMENT

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