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 the 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 coming through 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 microwave 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.

I. INTRODUCTION

Negative refraction and sub-wavelength 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 \(\varepsilon_r = -1\) 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 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 \(\varepsilon_r, \mu_r\) 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

\[
T = \frac{2}{2 \cosh(\alpha d) + (\gamma + 1/b) \sinh(\gamma d)}
\]

for TE incidence or \(\gamma = \frac{\alpha}{\omega \varepsilon_r} \approx 1\) for TM case, and \(k_0 = \sqrt{\omega_0 \mu_r \varepsilon_r}\). Assuming \(\varepsilon_r = \mu_r = -1\), we see that \(\gamma = -1\) and \(T = \exp(\alpha d)\) 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 (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 usually, 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:

\[
\begin{align*}
A + C &= B \\
\frac{A - C}{\eta_0} &= B \eta \\
\end{align*}
\]

Here \(\eta_0\) and \(\eta\) are the wave impedances that connect the tangential components of electric and magnetic fields in free space and in the medium, respectively. The solution of (2) is, obviously,

\[
C = \frac{\eta - \eta_0}{\eta + \eta_0} A, \quad B = \frac{2\eta}{\eta + \eta_0} A
\]

The wave impedance of propagating transmitted waves is given by \(\eta = \omega \mu_r \varepsilon_r / k_n\) for TE waves, and \(\eta = k_n / (\omega \varepsilon_r \mu_r)\).
for TM waves \( k_n = \sqrt{\epsilon_r \mu_r k_0^2 - k^2} \) 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 \( \text{Im}\{\epsilon_r, \mu_r\} \to -0 \)) if the square root branch is chosen so that \( \text{Im}\{k_n\} < 0 \) (or \( \text{Im}\{k_n\} \to -0 \)); the time dependence is in the form \( \exp(+j\omega t) \).

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

But let us suppose that the incident and transmitted waves are evanescent. Then, \( \eta = j\omega\mu_0\epsilon_0/\alpha \) for TE waves, and \( \eta = \alpha/(j\omega\epsilon_0\mu_0) \) for TM waves. Because the transmitted wave must decay from the interface, \( \alpha \) 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: \( C = B \to \infty \). The reason for such resonant growth of the field amplitudes when \( \epsilon_r, \mu_r \to -1 \) is in the excitation of a surface mode (surface polariton) associated with the interface. Indeed, if there is no incident field in (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 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 \emph{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 system will posses surface polariton resonances reconstructing the evanescent spectrum as well. We found in our recent paper 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 non-linear effects like wave mixing. In this paper we will discuss alternative possibilities to evanescent spectrum reconstruction not involving non-linearity. The design will be based on the principle mentioned above: 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.

\section{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 call the system simply as \emph{device}. A possible name for such a device can be \emph{near-field lens}, but we would prefer not to use word \emph{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 \( 2 \times 2 \) wave transmission matrices, well known in the microwave circuit theory. 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 \( e^{\pm jk_z z} \) 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

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

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. \( T_{\text{dev}} \) is the transmission matrix of the device. From this formula it is obvious that a \emph{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 evanes-
cent 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 $\varepsilon_r = -1$ and $\mu_r = -1$ (Veselago medium)$^{1,2}$ and a pair of parallel conjugating surfaces or sheets$^7$ operate as such 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

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

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

$$T_{dev} = \begin{pmatrix} \exp(+jk_d) & 0 \\ 0 & \exp(-jk_d) \end{pmatrix} \tag{7}$$

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 the introduction. This system is modeled by the transmission matrix

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

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 $a = d = 0$, $e = h = 0$, 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

$$T = \begin{pmatrix} 0 & b \\ c & 0 \end{pmatrix} \tag{9}$$

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 (5). We keep the same operation principle defined by (8). If we allow a mismatch at the device interfaces still maintaining the device symmetry, the following solution can be found:

$$T_{in} = T_{out} = \begin{pmatrix} a & b \\ c & 0 \end{pmatrix} \tag{10}$$

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

$$T_{dev} = \begin{pmatrix} a^2 \exp(-jk_d) - \exp(+jk_d) & a \exp(-jk_d) \\ -a \exp(-jk_d) & -\exp(-jk_d) \end{pmatrix} \tag{11}$$

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

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

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:

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

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

$$T_{dev} = \begin{pmatrix} \exp(+jk_d) & af\exp(-jk_d) + bh \exp(+jk_d) \\ 0 & \exp(-jk_d) \end{pmatrix} \tag{14}$$

which corresponds to the scattering matrix of the form

$$S_{dev} = \begin{pmatrix} 0 & \exp(+jk_d) \\ \exp(+2jk_d) & bh \exp(+2jk_d) + af \end{pmatrix} \tag{15}$$

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 the next section 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

$$ T = 1 + R $$

(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

$$ T_g = \begin{pmatrix} 1 + 2R & R \\ R & 1 + R \end{pmatrix} $$

(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. 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

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

(18)

which is of the necessary form (10).

For a better understanding we reformulate the consideration above 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

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

(19)

and the grid transmission matrix becomes

$$ T_g = \begin{pmatrix} 1 - \frac{\eta_0}{2Z_g} & \frac{\eta_0}{2Z_g} \\ \frac{\eta_0}{2Z_g} & 1 + \frac{\eta_0}{2Z_g} \end{pmatrix} $$

(20)

The reflection coefficient (19) becomes infinite and the transmission matrix (20) reduces to (18) when $\omega$ and $k_t$ satisfy equation

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

(21)

which is the dispersion equation for surface modes on the grid surface. Because $\eta_0$ 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 (20) that it is enough to change the sign of the grid impedance to realize the second matrix in (13). Such 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 (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:

$$ E_1 - E_2 = J_m, \quad H_1 - H_2 = J_e $$

(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 $Z_m$:

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

(23)

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

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

(24)

Eq. (20) is a particular case of (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

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

(25)

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

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

(26)

Let us note again that nothing forbids realization of the necessary impedances for a given $k_t$, because for the evanescent modes 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 resonance is possible to realize in a passive grid or array. The resonance occurs when the \((\omega, k_t)\) pair belongs to the polariton spectrum of the grid.

Practically speaking, this means that for a given frequency \(\omega\) 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 range of operable \(k_t\) as possible.

Mathematically (and ideally), the last 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\(^9\)

\[
p = \chi(\omega) \left[ E_{ext} + \beta(\omega, k_t)p \right] \tag{27}
\]

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

\[
p = \frac{E_{ext}}{\chi(\omega) - \beta(\omega, k_t)} \tag{28}
\]

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

\[
Z_c J = E_{ext}, \tag{29}
\]

where

\[
Z_c = S_0 \frac{1}{j\omega \chi(\omega) - \beta(\omega, k_t)} \tag{30}
\]

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

\[
Z_g = Z_c - \frac{\eta_0}{2} \tag{31}
\]

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 (28).

The imaginary parts of \(1/\chi\) and \(\beta\) cancel out in a lossless non-radiating array\(^9\) (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 \(\chi\) depends on the frequency, but does not depend on \(k_t\). On the other hand, the interaction constant \(\beta\) 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 (\(Re\{1/\chi\} \rightarrow 0\) and minimize the field interactions between the particles in the array (provide \(Re\beta(\omega, k_t) \rightarrow 0\)). 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 to reduce or compensate interactions at a certain frequency. In the next section 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 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 Figure 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).

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 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 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 Figure 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 experimental results are shown in Figures 3 and 4. 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.

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.

To study phenomena in such grids, we have made measurements in a system of two regular arrays of similar particles. The array geometry is shown in Figure 5. 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 transverse wavenumber $k_t$. Next, we exited the grids
FIG. 5: A periodic array of small resonant particles.

FIG. 6: The distribution of 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 (1st polariton resonance) and 5.26 GHz (2nd polariton resonance) and summed up. 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. Probe is 0.5 cm away from the top mesh of the setup. The field scale is linear.

at each of these frequencies (2 frequencies were practically used) and superimposed the measured spatial profiles. This corresponds to reconstructing the source spatial spectrum using only a few spatial harmonics. The result for the case of excitation by two small dipole antennas is shown in Figure 6. Strong excitation of the second grid is clearly visible, as well as an image of the source field behind the grids. 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 set-up had many non-idealities. 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 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.

The present paper continues of 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 we showed that a system of two phase-conjugating planes or sheets placed in free space behaves as a perfect lens proposed by Pendry. 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 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 polariton-resonant systems for evanescent field enhancement at microwaves. The experiment supports the main concepts of our theoretical findings. The resonant growth of the evanescent fields coming through the system has been observed.

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