Eur. Phys. J. D **39**, 295–301 (2006)
DOI: 10.1140/epjd/e2006-00092-4

THE EUROPEAN PHYSICAL JOURNAL D

Anomalous lateral beam shift and total absorption due to excitation of surface waves in materials with negative refraction

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Received 25 January 2006 / Received in final form 17 March 2006
Published online 28 April 2006 –

Abstract. We studied electromagnetic beam reflection from layered structures that include materials with negative refraction. Excitation of *leaky surface waves* leads to the formation of anomalous lateral shifts in the reflected beams with single or double peak structures. The presence of reasonable losses within material with negative refraction, besides significant influence on manifestation of the giant lateral shifts, can lead to their total suppression and anomalously high absorption of the incident radiation. If, in addition to the resonant excitation of leaky surface waves, radiation inflow exactly compensates their irreversible damping, *total absorption* of the incoming radiation can be achieved for moderately wide beams.

PACS. 78.68.+m Optical properties of surfaces – 42.25.Bs Wave propagation, transmission and absorption – 42.79.Gn Optical waveguides and couplers

1 Introduction

Experimental fabrication [1,2] of the microstructured materials with negative refraction called left-handed metamaterials, whose peculiar optical properties have been predicted many years ago [3], have recently attracted extensive studies for the development of novel optical devices (see e.g. [4,5]). In particular, attention has been paid to the negative lateral beam shift (Goos-Häncken effect in the opposite direction) of the reflected beam that is obliquely incident upon the interface between the conventional or right-handed (RH) and left-handed (LH) media [6–8]. However, such interfaces have been shown to support surface waves of both TE and TM polarizations [9–11] for certain range of interface parameters. In order to excite surface waves it is necessary to satisfy phase-matching conditions that can only be achieved with the use of properly chosen layered structures. In that case, surface modes energy can leak out of the layered structure and thus lead to their reversible damping. This is in contrast to the usual irreversible damping that is due to energy losses within the layered structure itself. Namely, the reversibility of the leaky surface wave (LSW) damping allows for the resonant energy pumping by the incident beam.

Excitation of leaky waves is usually realized via attenuated or frustrated total reflection configurations in the two well-known geometries: (i) prism-air-dielectric called Otto configuration and (ii) prism-dielectric-air called Kretchman configuration. For the purpose of this paper we confine ourselves to the Otto configuration and replace dielectric by material with negative refraction or LH medium. In fact, similar problem has been studied by Shadrivov et al. [12,13], who showed that lateral shift can be dramatically enhanced by the resonant excitation of the leaky surface waves at the interfaces that involve LH materials. Such giant shifts are known to exist in the case of layered structures made of RH materials [14–18]. However, the surface between two RH materials may support leaky surface waves of TM polarization only, with the shift of the reflected beam in the forward direction. The use of LH materials in layered structures opens up the possibilities of giant shifts in both TM and TE polarized beams in forward, as well as backward direction. This requires more comprehensive analysis of lateral beam shifts in the case of LH materials.

The role of absorption in LH materials also calls for more extensive investigation having in mind that the plane-wave theory reveals the possibility of total absorption of the incoming radiation [19–21]. For this to occur, besides the resonant excitation of LSW, the energy of the incoming radiation must exactly compensate all irreversible dissipative processes within the layered structure materials.

From the point of view of possible applications, it is worth noting that actual LH materials are composed of unit cells of finite dimensions and thus heterogeneous. In order to consider them as continuous media and introduce effective dielectric permittivity and magnetic permeability

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that are simultaneously negative, it is necessary to observe the homogenization limit, as discussed in reference [22]. In the present paper we assume that the wavelength of the excited LSW is sufficiently greater then the size of the composite material unit cell. The same holds for the distance between the prism and LH material in the Otto configuration.

Although the geometry of the problem and some starting equations in the present paper are similar to those of references [12,13], there are essential differences in the treatment of the problem and the results. Our choice of the phase shift, that is inversely proportional to the imaginary part of the LSW wave number, is fully justified by the analytical results presented, which are in excellent agreement with numerical ones and allow for complete physical understanding of the physical processes involved. The definition of lateral beam shift used in references [12,13] as normalized first moment of the electric field of the reflected beam has entirely different physical meaning and may have essentially different values than the phase shift used in our paper. Here, we present the results for all possible cases of polarization and propagation directions and predict that the giant lateral shifts may reveal single-peak structures, i.e. double peak structures do not necessarily appear. We show that the field structure (one or two peaks for giant lateral shifts) significantly depends on the width of the gap and the incident beam, as well as on the angle of incidence. Finally, we present thorough investigation of the role of losses on the reflected beam.

Thus, the aim of the present paper is twofold: (a) to investigate comprehensively, both analytically and numerically, the manifestation of giant lateral reflected beam shifts in cases when absorption within materials with negative refraction can be neglected; (b) to study how the finite beam width will influence the effect of anomalously high absorption when reasonable losses are present. Obviously, absorption leads to high or even total suppression of the giant lateral beam shifts. Moreover, we predict the effect of resonant absorption of moderately wide beams by lossy materials with negative refraction that, under certain conditions, can be considered as total.

2 Formulation of the problem

We consider a Gaussian beam that has the beam width w and that is obliquely incident from medium 1 (usually a prism) upon a two-dimensional, two layered structure with dielectric permittivities $\varepsilon_2, \varepsilon_3$ and magnetic permeabilities μ_2, μ_3 (see Fig. 1). The angle of incidence of the beam θ_i is defined with respect to the normal to the interface so that the wave vector component along the interface is: $k_{xi} = k_1 \sin\theta_i$, where $k_1 = \omega(\varepsilon_1\mu_1)^{1/2}/c$ is the wave number in medium 1, ω is the frequency, c the speed of light in vacuum, ε_1, μ_1 are dielectric permittivity and magnetic permeability of the medium 1, respectively. If the beam is TE or s-polarized it is convenient to work with the electric field, while for TM or p-polarization one will rather use magnetic field of the beam. In both cases, the fields are oriented normally to the plane of incidence, i.e. along y-axis, and are consequently continuous across the boundaries between media.

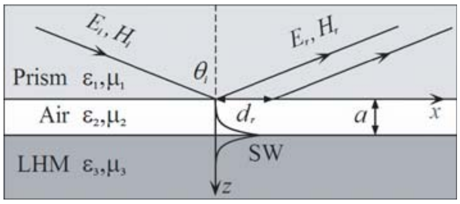

Fig. 1. Geometry of the problem.

If we choose e.g. s-polarization, the electric field of the incident beam at the interface $z = 0$ has the form $E_i(x, z = 0) = \exp(-x^2/2w_x^2 - ik_{xi}x)$, where $w_x = w/\cos(\theta_i)$, assuming the amplitude is 1. In the case of p-polarization the same expression can be used for magnetic field at $z = 0$. For the purpose of LSW excitation it is necessary that the medium 1 is optically dense ($\varepsilon_1\mu_1 > \varepsilon_2\mu_2$) and that incident angle θ_i is greater than the angle of total internal reflection $\theta_{tir} = \arcsin(\varepsilon_2\mu_2/\varepsilon_1\mu_1)^{1/2}$. Medium 2 represents a gap layer of width a, between medium 1 (usually a prism) and those two media are considered to be non-dispersive. However, medium 3 has been assumed to be a LH metamaterial with negative real parts of both ε_3 and μ_3 that are shown to be frequency dependent [2]:

$$\varepsilon_3(\omega) = 1 - \frac{\omega_p^2}{\omega^2}; \quad \mu_3(\omega) = 1 - \frac{F\omega^2}{\omega^2 - \omega_r^2}. \tag{1}$$

Here, the parameters that have been chosen to fit experimental data of reference [2] are given in reference [10] as follows: $\omega_p/\omega_r = 2.5$; $F = 0.56$. The interface between media 1 and 2 generates reflected beam and the beam transmitted in the form of evanescent fields that can couple with evanescent field of surface waves supported by the interface between media 2 and 3. The regions of existence and the properties of such surface waves have been investigated in references [9–11]. In reference [9] the reflection in the plane-wave approximation has also been calculated. The authors of reference [11] use theoretically more general expressions of ε_3 and μ_3 than those presented in equation (1), for their general analysis of the surface waves. In the present paper, however, we will use expressions (1) that have been experimentally confirmed. Let us mention that ε_3 and μ_3 are both negative in the frequency range $\omega_r < \omega < \omega_r/(1-F)^{1/2}$, and the condition $\omega_r < \omega < \omega_r/(1-F/2)^{1/2}$ is imposed for the existence of surface waves if $\varepsilon_2 = \mu_2 = 1$, i.e. if the gap is supposed to be the air. Then, for the parameter values given above $1.75 < |\varepsilon_3(\omega)| < 3.5$; $0 < |\mu_3(\omega)| < 1$. When $1.75 < |\varepsilon_3(\omega)| < 2.7$ surface waves appear to be TM polarized with the energy transfer in the forward direction, while for $2.7 < |\varepsilon_3(\omega)| < 3.5$ they are TE polarized with the energy transfer in the backward direction. Another two possible types of surface waves, TE forward and TM backward (see Ref. [10]), can exist if the gap is filled-up with the dielectric material with $\varepsilon_2 > 3.5$; $\mu_2 = 1$, or for (suitably chosen) different values of the parameters ω_p/ω_r and F.

In what follows, throughout this paper we are going to use dimensionless wave vectors (normalized to k_1) and lengths (normalized to $1/k_1$). At the same time ε_2, ε_3 are normalized to ε_1 and μ_2, μ_3 to μ_1 without loss of generality. When $k_{xi} = \sin\theta_i$ is close to the real part of the wave vector of the corresponding surface wave (polariton) LSW is resonantly excited. Energy transfer along the interface (up to the distances that are inversely proportional to imaginary part of LSW wave number) can lead to a substantial enhancement of the lateral shift of the reflected beam as compared to the well known Goos-Häncken effect. Here, we are talking about shifts that are grater than (or comparable to) the incident beam width, while the opposite is true for Goos-Häncken shifts. Moreover, the shape of the reflected beam can be far from Gaussian one and usually contains double peak structures. This, relatively simple picture, may hold also for the layered structures that contain RH materials only, for TM or p-polarized beams. The presence of LH metamaterials is distinguished by the fact that both TE and TM polarizations can be used and in both cases forward, as well as backward shifts can appear depending on the properties of the layered media, i.e. on ε_n, μ_n ($n = 1, 2, 3$). For s- or p-polarized Gaussian beam, the plane wave spectrum can be written as:

$$\{E_i(k_x); H_i(k_x)\} = \frac{w_x}{\sqrt{2\pi}} e^{-(k_x - k_{xi})^2 w_x^2/2}. \quad (2)$$

The reflected electric (magnetic) field is given by:

$$\{E_r(x); H_r(x)\} = \int_{-\infty}^{+\infty} R(k_x)\{E_i(k_x); H_i(k_x)\} e^{ik_x x} dk_x, \quad (3)$$

where $R(k_x)$ is the plane-wave reflection coefficient, which for monochromatic waves, can be written in the following form, suitable for LSW considerations:

$$R(k_x) = R_{12}\frac{G - gA_2 + i[F - g(1 + A_1)]}{G + gA_2 - i[F + g(1 + A_1)]}. \quad (4)$$

Here:

$$G = (1 + \alpha_2) - (1 - \alpha_2)A_1; \quad A_1 = \frac{(1 - \alpha_1^2)}{(1 + \alpha_1^2)} e^{-2\kappa_{z2}a};$$

$$F = (1 - \alpha_2)A_2; \quad A_2 = \frac{2\alpha_1}{(1 + \alpha_1^2)} e^{-2\kappa_{z2}a};$$

$R_{12} = (\alpha_1 - i)/(\alpha_1 + i)$, $\alpha_{1,2} = \mu_2\kappa_{z1,3}/\mu_{1,3}\kappa_{z2}$ in the case of s-polarization and $\alpha_{1,2} = \varepsilon_2\kappa_{z1,3}/\varepsilon_{1,3}\kappa_{z2}$ in the case of p-polarization, and $\kappa_{z1} = \sqrt{1 - k_x^2}$, $\kappa_{z2,3} = \sqrt{k_x^2 - \varepsilon_{2,3}\mu_{2,3}}$. In fact, R_{12} represents the reflection coefficient from a single boundary between media 1 and 2, $G = 0$ represents the dispersion equation of LSW, while F describes wave damping due to leakage into the prism region. Losses are represented by imaginary parts of dielectric permittivity and magnetic permeability ε_3'' and μ_3''. It is assumed that the prism and the gap consist of the RH lossless materials. Then, all losses in the expression (4) for the reflection coefficient appear in α_2. Thus, for the purpose of the present paper, it is convenient to introduce the complex quantity $\tilde{\alpha}_2 = \alpha_2 - ig$, where:

$$\text{(TE)}: g \simeq \alpha_2[\mu_3''/\mu_3 + \varepsilon_3\mu_3/2\kappa_3^2(\mu_3''/\mu_3 + \varepsilon_3''/\varepsilon_3)], \quad (5a)$$

$$\text{(TM)}: g \simeq \alpha_2[\varepsilon_3''/\varepsilon_3 + \varepsilon_3\mu_3/2\kappa_3^2(\mu_3''/\mu_3 + \varepsilon_3''/\varepsilon_3)]. \quad (5b)$$

Thus, expression (4) represents generalization to the case of LH materials of the results of references [19–21]. Of course, when absorption can be neglected (i.e. when $g = 0$) $|R|^2 = 1$ and the real and imaginary part of the numerator in equation (4) cannot simultaneously vanish. However, the phase shift in R can be dramatically enhanced when the resonant excitation of LSW takes place (i.e. when $G = 0$). In the presence of absorption, the approximation of plane-waves gives the possibility to achieve total absorption, provided that the two conditions are simultaneously fulfilled:

$$G = gA_2; \quad F = g(1 + A_1). \quad (6)$$

The first of those two conditions gives a small correction to the LSW dispersion relation $G = 0$, and, thus, requires resonant excitation of LSW. However, the second condition requires exact compensation of the irreversible damping by radiation inflow making perfect connection of the source and the sink.

3 Analytical considerations

It is clearly seen from equations (2, 3) that the main contribution to the integral in equation (3) comes from the region

$$(k_x - k_{xi}) \leq \frac{1}{w_x}. \quad (7)$$

Provided there are no other singularities such as branch points within this region, the essential contribution to the reflected field comes from the poles and zeros of the reflection coefficient that can be written in the following form:

$$R(k_x) = R_{12}\frac{k_x - k_n}{k_x - k_p}, \quad (8)$$

where $k_n = k_{sn} + i\beta_n$, $k_p = k_{sp} + i\beta_p$, and:

$$k_{sn} = k_{sw} - [(1 - \alpha_2)A_1 + gA_2]Q;$$
$$\beta_n = -[(1 - \alpha_2)A_2 - g(1 + A_1)]Q; \quad (9)$$

$$k_{sp} = k_{sw} - [(1 - \alpha_2)A_1 - gA_2]Q;$$
$$\beta_p = [(1 - \alpha_2)A_2 + g(1 + A_1)]Q. \quad (10)$$

Here, $Q = (d\alpha_2/dk_x)^{-1}$ taken at $k_x = k_{sw}$; the propagation constants of the surface waves k_{sw} have been evaluated in reference [10] from the dispersion relations $1 + \alpha_2 = 0$ that correspond to the infinitely wide gaps:

$$\text{(TE)}: k_{sw}^2 = \varepsilon_2\mu_2\frac{Y(Y - X)}{Y^2 - 1};$$
$$\text{(TM)}: k_{sw}^2 = \varepsilon_2\mu_2\frac{X(X - Y)}{X^2 - 1}, \quad (11)$$

where the effective parameters are given by: $X = |\varepsilon_3|/\varepsilon_2$ and $Y = |\mu_3|/\mu_2$. By using these two parameters the derivatives in equation (7) can be written as:

$$\text{(TE)}: Q = k_{sw}\frac{Y(1-XY)}{(Y^2-1)(Y-X)};$$
$$\text{(TM)}: Q = k_{sw}\frac{X(1-XY)}{(X^2-1)(X-Y)}. \quad (12)$$

As can be verified, the sign of Q is equal to the sign of the corresponding x-component of the integrated energy flux carried by LSW within the gap and LH metamaterial: $P_x = \int_0^\infty S_x dz$; where S_x is x-component of the Poynting vector ($\mathbf{S} = c/8\pi[\mathbf{E}\times\mathbf{H}]$). This can be easily understood having in mind that expressions (11) represent inverse derivatives of the dispersion function and, consequently, are inversely proportional to the group velocity of the LSW [11]. In fact:

$$P_x = P_0 Q^{-1} + O(A_1); \quad (13)$$

Here, $P_0 = c^2E_0^2\kappa_{z2}\sin\theta_i/16\pi\omega\mu_1\mu_2$ for TE mode, $P_0 = c^2H_0^2\kappa_{z2}\sin\theta_i/16\pi\omega\varepsilon_1\varepsilon_2$ for TM mode, E_0, H_0 are the field amplitudes at $z = d$, and $O(A_1)$ stands for small terms that are proportional to A_1. Since P_0 is essentially positive, the sign of P corresponds to sign of Q. Notice that $P_x \to 0$ when $Y \to 1$ for TE mode or when $X \to 1$ for TM mode, while $P_x \to \infty$ when $XY \to 1$ for both polarizations. Notice that, when $XY \to 1$ and $k_{sn} = k_{sp} \to k_{sw}$, both κ_{z2} and $\kappa_{z3} \to 0$, so that the field(s) amplitude(s) become constant throughout the region $0 < z < \infty$. Consequently, no LSW can be excited. Moreover, X and Y cannot be close to 1, because otherwise matching conditions with the prism for excitation of LSW cannot be fulfilled ($k_{sw} = \sin\theta_i < 1$). Therefore, in the present paper we confine ourselves to the range of parameters X and Y where each of them, as well as their product XY, is not too close to 1. In contrast, surprisingly, the authors of reference [13] have used the parameters: $X = 0.5$, $Y = 2$ (i.e. $XY = 1$ and consequently infinite P_x) for the excitation of the TE forward surface waves (see last paragraph on p. 487 of Ref. [13]).

For the configuration studied in the present paper excitation of four different types of the LSW are possible. To identify those we recall the requirement $\sin\theta_{tir} < \sin\theta_i < 1$ that, up to the small corrections that are proportional to B, reveals:

(a) forward TE mode for $Y(1-Z)+Z/Y < X < 1/Y$;
(b) backward TE mode for $Y(1-Z)+Z/Y > X > 1/Y$;
(c) forward TM mode for $X(1-Z)+Z/X < Y < 1/X$;
(d) backward TM mode for $X(1-Z)+Z/X > Y > 1/X$.

Here $Z = (\varepsilon_2\mu_2)^{-1}$. For example, in Figure 2 we plot Y versus X for the parameters $Z = 10$, (i.e. $\varepsilon_1 = 10$; $\mu_1 = 1$); $F = 0.56$; $\omega_p/\omega_r = 2.5$ (dashed curve) and $\omega_p/\omega_r = \sqrt{2}$ (dotted curve), assuming that the gap is the air.

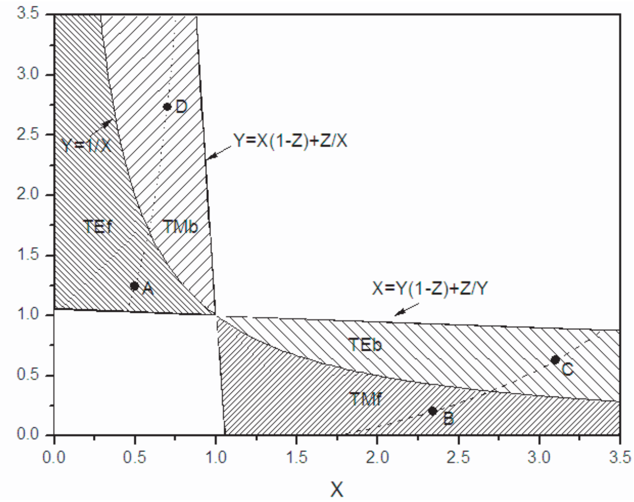

Fig. 2. Existence regions (shaded areas) for the four possible types of surface waves excitation. Dotted curve corresponds to $(\omega_r/\omega_p)^2 = 0.5$, dashed curve to $(\omega_r/\omega_p)^2 = 0.16$ in formula (1). For both curves: $F = 0.56$; $Z = 10$. Points A, B, C and D used for numerical calculations shown in the following figures.

Upon inserting (7) into (2), one can find the following expression for the reflected electric (magnetic) field:

$$\{E_r(x); H_r(x)\} = R_{12}\Big[1 + i\sigma\sqrt{2\pi}(k_p - k_n)(w_x/2) \times \exp(\gamma^2)\,\text{erfc}(\sigma\gamma)\Big]\exp(ik_{xi}x - x^2/2w_x^2) \quad (14)$$

where $\gamma = \sqrt{2}(w_x\beta_p/2 - x/2w_x + iw_x(k_{xi} - k_{sp})/2)$, $\sigma = \text{sgn}(\beta_p) = \text{sgn}(P)$, and erfc($\gamma$) is the complementary error function. This represents generalization to the case of materials with negative refraction of the results previously obtained by Shah and Tamir [16] for the case of ordinary RH materials, when only TM forward LSW can be excited. Besides the presence of the parameter σ in our formula (12), the main difference from reference [16] arises through our definitions of k_n and k_p (see Eqs. (8–11)). This leads to clear physical understanding of anomalously high lateral beam shifts in the absence of absorption (see Sect. 4) , as well as to the understanding of the effect of total absorption of moderately wide beams in the presence of reasonable losses within the material with negative refraction (see Sect. 5). Moreover, comparison with the results obtained via direct numerical integration of expression (2) using equation (3) shows not only quantitative, but also excellent quantitative agreement.

4 Lateral shifts due to excitation of leaky surface waves. Zero absorption

Within this section we neglect irreversible damping of the leaky surface waves (LSW) i.e. ε_n, μ_n are considered to be real quantities and the expression (4) becomes:

$$R(k_x) = R_{12}\frac{G + iF}{G - iF}. \quad (15)$$

In that case, $G = 0$ represents the condition for resonant excitation of LSW which leads to the following expression for a as a function of θ_i:

$$a = \frac{1}{2\kappa_2} \ln \left(\frac{1-\alpha_2}{1+\alpha_2} \frac{1-\alpha_1^2}{1+\alpha_1^2} \right)_{k_x = k_{xi}}. \qquad (16)$$

From this simple expression resonant angle of incidence θ_i can be determined for a given gap width a. Then, it can be shown that there is a simple relation between the imaginary part of the LSW wave number β_s and the lateral shift of the reflected beam d_r, defined as:

$$d_r = - \left. \frac{d\Phi_r}{dk_x} \right|_{k=k_s} = i \left. \frac{d}{dk_x} \ln \left[R\left(k_x\right) \right] \right|_{k=k_s} = \frac{2}{\beta} \qquad (17)$$

where Φ_r is the phase of the reflection coefficient, $k_s = k_{sn} = k_{sp}$ and $\beta = \beta_p = -\beta_n$ are determined from equations (8, 9). In fact, these expressions are valid, strictly speaking, in the plane-wave approximation but produce quite reasonable results for moderately wide beams. Shadrivov et al. [12,13] pointed out that in the case of narrow beams with wide spectrum, it is more appropriate to define the relative shift of the beams, by using the normalized first moment of the electric (magnetic) field of the reflected beam. Of course, exact numerical integration is needed in that case, without assumptions made above for the purpose of analytical considerations. However, our analysis shows that in the case of double peak formations, the use of equation (14) is fully justified. Really, the reflected electric (magnetic) field at $z = 0$ can be evaluated analytically from equation (14)

$$\{E_r(x); H_r(x)\} = R_{12} \times \left[1 - \sigma \frac{w_x}{d_r} 2\sqrt{2\pi} e^{\gamma^2} \operatorname{erfc}(\sigma\gamma) \right] e^{-\frac{x^2}{2w_x^2} + ik_x x} \qquad (18)$$

where: $\gamma = \sqrt{2}(w_x/d_r - x/2w_x + iw_x(k_{xi} - k_s)/2)$, $\sigma = \mathrm{sgn}(d_r) = \mathrm{sgn}(P)$, and $\mathrm{erfc}(\gamma)$ is the complementary error function of the real argument when $k_{xi} = k_s$. Here, d_r is calculated using equation (17) with help of equations (10, 16) when $g = 0$. The results obtained via equation (18) are in excellent agreement with the following numerical ones.

Direct numerical solutions of the integral (3) using expression (4) for the reflection coefficient, for four different sets of parameters that corresponds to the possible excitation of TE forward and TE backward, as well as to TM forward and TM backward LSW have been performed and the results are presented in Figures 3 and 4. In Figure 3, we present density plots of the squares of the reflected field amplitudes at $z = 0$ as a function of x/w_x and the normalized gap width a, for the angles of incidence given by equations (9–11), i.e. $k_{xi} = \sin\theta_i = k_{sp} \simeq k_{sw}$, while the width of the beam is $w_x = 20$. Notice, that for the parameters used in our calculations, small corrections to k_{sw} given by equations (9, 10) do not change the results qualitatively. Figure 3a presents TE forward, Figure 3b TM forward, Figure 3c TE backward and Figure 3d TM backward LSW. Within the certain range of the gap widths a giant lateral shifts due to resonant excitation of LSW are clearly seen in all cases of interest. Outside this range, only small shifts that correspond to the normal Goos-Häncken effect can be observed. In that case, Gaussian structure of the incident beam is preserved in the reflected beam. Inside the range of the giant lateral beam shifts, formation of single, as well as double peak structures, in the laterally shifted reflected beam can be observed. When the double peak structure exists, one peak is due to the mirror

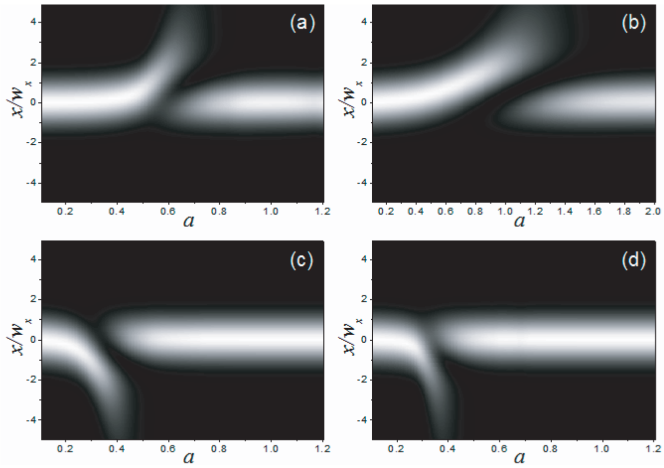

Fig. 3. Squares of the reflected field profile as a functions of gap width a for the following non-normalized parameters that correspond to the points A, B, C and D in Figure 2: (a) $\varepsilon_3 = -0.5$, $\mu_3 = -1.24$, $\theta_i = 24.402°$; (b) $\mu_3 = -0.203$, $\varepsilon_3 = -2.34$, $\theta_i = 19.528°$; (c) $\varepsilon_3 = -3.1$, $\mu_3 = -0.628$, $\theta_i = 30.417°$; (d) $\mu_3 = -2.73$, $\varepsilon_3 = -0.7$, $\theta_i = 31.86°$. Other parameters are: $\varepsilon_1 = 10$, $\varepsilon_2 = 1$, $\mu_1 = \mu_2 = 1$, $w_i = 20$. In cases (a) and (c) E_r field magnitudes for s-polarized beams are shown while cases (b) and (d) depict H_r field magnitudes for p-polarized beams.

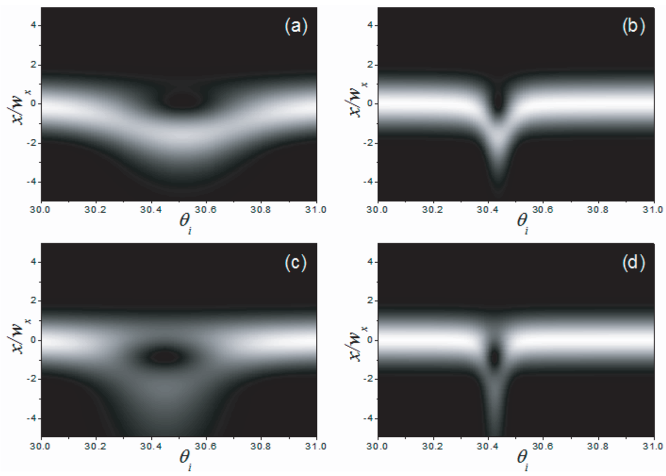

Fig. 4. Squares of field profiles for backward s-polarized beam as a function of incident angle for different gap widths and beam widths. The parameters are: (a) $w = 20$, $a = 0.33$; (b) $w = 100$, $a = 0.435$; (c) $w = 20$, $a = 0.4$; (d) $w = 100$, $a = 0.505$. Other parameters are as in Figure 3c, corresponding to point C in Figure 2.

reflection of one part of the beam, while the other peak corresponds to the resonant excitation of LSW. However, when the single peak structure appears, the whole energy of the incident beam is transferred to and shifted via LSW.

To get a complete insight in the phenomena, in Figure 4 we present density plots of the TM backward reflected squared field amplitude, as a function of x/w_x and the angle of incidence, for given values of a and different beam widths w_x. As one can see, double peak structures do not always accompany giant lateral shifts (see Figs. 4a and 4b). This follows also from Figure 3, but is much more pronounced in Figure 4. For wider beams ($w_x = 100$; Figs. 3b and 3d) we see, as expected, that the shifts are higher and more narrow than for narrow beams ($w_x = 20$; Figs. 4a and 4c).

5 Irreversible damping of leaky surface waves. Total absorption

The aim of this section is to examine how the presence of absorption will influence the lateral beam shift, as well as whether total absorption can be achieved for finite beam widths. For that purpose, we have numerically solved the integral given by equation (3) using expression (4) for the reflection coefficient. However, before presenting those numerical results, in order to understand better the physical picture of the processes involved, we will analyze what follows from equation (12).

First of all we find from equations (8, 9)

$$k_p - k_n = 2\left[g + i\left(1 - \alpha_2\right)\right] A_2 Q, \quad (19)$$

Since $\alpha_2 \simeq -1$ and $g \simeq 1$, the imaginary part in (16) is much grater than the real one. As to the argument γ that enters in equation (14), it becomes the real quantity if $k_{xi} = k_{sp}$. Of course, one can make such choice of the angle of incidence and study reflected beam as a function of the gap width a for given g and w_x. However, we are here interested in total absorption and will make slightly different choice of the angle of incidence that follows from equation (5). Namely, we choose $k_{xi} = k_{sn}$ and recall that the plane-wave approximation also requires $\beta_n = 0$, i.e.: $A_2 \simeq g/2$; for total absorption. In that case, it follows from equation (12):

$$|E_r(x); H_r(x)|^2 = \Big|\Big[1 - \sqrt{2\pi}\sigma\left(1 - ig/2\right) g w_x \times Q \exp(\gamma^2) \mathrm{erfc}\left(\sigma\gamma\right)\Big] \exp(-x^2/2w_x^2)\Big|^2, \quad (20)$$

where $\gamma = \sqrt{2}\sigma\left[g w_x Q(1 - ig/2) - x/2w_x\right]$. If,

$$g w_x Q \simeq 1, \quad (21)$$

we can neglect γ dependence on x for all reasonable values of x, i.e. when $|x| \simeq 2 g w_x^2 Q$. Then, by using asymptotic expansion: $e^{\gamma^2}\mathrm{erfc}(\gamma) \sim 1/\sqrt{\pi}\gamma$ one gets zero on the right-hand side of equation (20) and thus, total absorption. Of course, we are aware of approximations that have been made and consequently, that exactly 100% absorption cannot be achieved for finite beam widths. However, here we talk about total absorption when reflection becomes negligibly small. In fact, recalling that the amplitude of the incident field is assumed to be 1, one can calculate z-component of the integrated reflected power flux $P_{zr} = \int_{-\infty}^{+\infty} S_z dx$, to obtain:

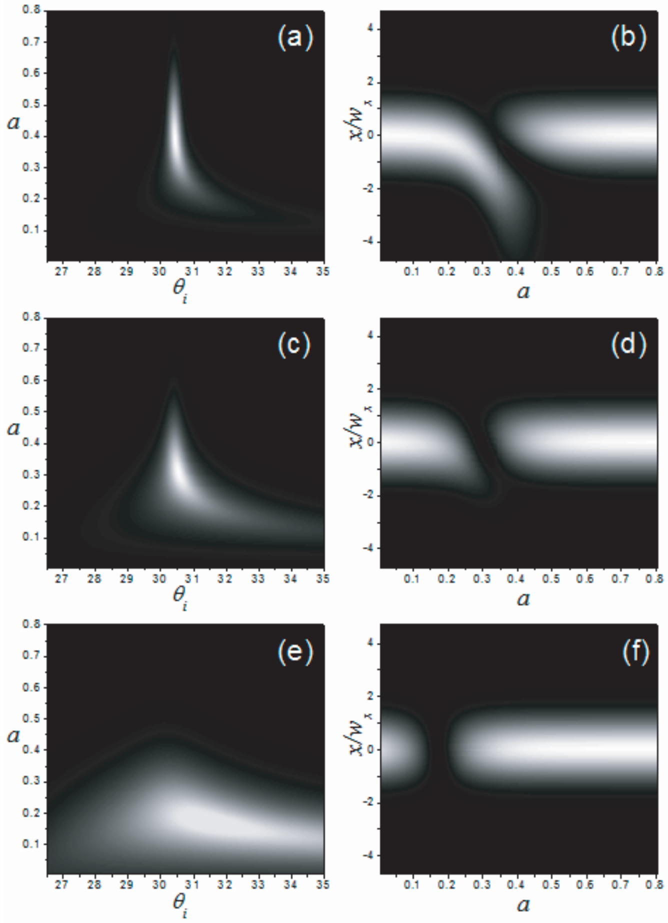

Fig. 5. Absorption as a function of incident angle and gap width a (a, c, e). Reflected field profiles as a function of gap width a (b, d, f). Parameter g takes values 0.001 (a, b), 0.01 (c, d), and 0.1 (e, f). Incident angles for (b, d, and f) refer to the maximum absorption for respective losses. Other parameters are as in Figure 3c, corresponding to point C in Figure 2.

$$P_{zr} = -\frac{c^2 w_x \cos\theta_i}{16\sqrt{\pi}\omega\mu_1}\left(\frac{g - 2A_2}{g + 2A_2}\right)^2 + O\left((g + 2A_2)^{-3} w_x^{-3} Q^{-3}\right). \quad (22)$$

This confirms that the conditions for vanishing reflection coefficient ($g \simeq 2A_2$) that follow from plane-wave approximation, is justified with high accuracy when the condition (21) is fulfilled.

Absorption is calculated as $A = 1 - |R|^2$ as a function of θ_i and a for given g. The results are presented in Figures 5a, 5c, and 5e for $g = 0.001$; 0.01 and 0.1 respectively, as pertinent samples for the parameters that allow for the excitation of s-polarized backward LSW. Our analysis shows that in all other cases of interest (s-backward and forward, p-forward) the similar results are obtained. As one can see, the resonant angle of incidence increases, while the resonant gap width a decreases for higher values of g, as expected from equation (6). Absorption peak is extremely narrow for small values of g and becomes wider for greater ones. However, total absorption can always be achieved (i.e. for any $g > 0$), when the conditions (6) are satisfied.

In order to compare the results obtained within the plane-wave approximation with those obtained for finite beam widths, we present in Figures 5b, 5d and 5f the numerical results of integration in equation (3) using expression (4) for the reflection coefficient. We present density plots of the squares of the reflected field profiles at $z = 0$ as a function of x/w_x and the normalized gap width a, for the angles of incidence given by equations (9–11), i.e. $k_{xi} = \sin\theta_i = k_{sn}$, while the width of the beam is $w_x = 20$. The values of k_{sn} have been calculated for each particular value of g, i.e. $g = 0.001$; 0.01 and 0.1 that have been used in Figures 5b, 5d and 5f respectively. As can be seen in Figure 5b, for very small losses $g = 0.001$ the picture is very similar to Figure 3d. In other words, giant lateral beam shift still appears. However, for moderate losses ($g = 0.01$) Figure 5d shows that anomalous lateral beam shift disappears. High absorption can be achieved for certain value of the gap width. Finally, for relatively high losses ($g = 0.1$), in Figure 4f one can see the absorption resonance shifted towards smaller gap widths as expected, while maximum absorption achieves ≃100%.

6 Conclusion

We have investigated the scattering of an obliquely incident Gaussian electromagnetic beams from the layered structures that contain LH metamaterials. When absorption can be neglected, resonant excitation of leaky surface waves can lead to giant lateral beam shifts in the reflected beams. We have presented comprehensive analytical, as well as numerical analysis of the phenomena in the case of Otto configuration: prism-air-LH metamaterial. Excitation of all possible TE and TM, both backward or forward, leaky surface waves leads to the formation of giant lateral shifts in the reflected beams with single or double peak structures depending on angle of incidence and width of the air gap. The presence of reasonable losses within LH metamaterial, besides significant influence on manifestation of the giant lateral shifts, can lead to their total suppression and anomalously high absorption of the incident radiation. If, in addition to the resonant excitation of LSW, radiation inflow exactly compensates their irreversible damping, nearly total absorption of the incoming radiation can be achieved. Such phenomena can occur for moderately wide incident beams and reasonable losses within LH metamaterial itself.

The authors acknowledge a financial support of the Ministry of Science and Environmental Protection of the Republic of Serbia (grant No. OI 1475).

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