Hysteretic jumps in the response of layered superconductors to electromagnetic fields

V. A. Yampol'skii,[1,2] T. M. Slipchenko,[2,3] Z. A. Mayzelis,[2,3] D. V. Kadygrob,[3] S. S. Apostolov,[2,3] S. E. Savel'ev,[1,4] and Franco Nori[1,5]

[1]*Advanced Science Institute, The Institute of Physical and Chemical Research (RIKEN), Wako-shi, Saitama, 351-0198, Japan*
[2]*A. Ya. Usikov Institute for Radiophysics and Electronics Ukrainian Academy of Sciences, 61085 Kharkov, Ukraine*
[3]*V. N. Karazin Kharkov National University, 61077 Kharkov, Ukraine*
[4]*Department of Physics, Loughborough University, Loughborough LE11 3TU, United Kingdom*
[5]*Department of Physics, Center for Theoretical Physics, Applied Physics Program, Center for the Study of Complex Systems, University of Michigan, Ann Arbor, Michigan 48109-1040, USA*

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We consider here a layered superconductor subject to an externally applied moderately strong electromagnetic field. We predict hysteretic jumps in the dependence of the surface reactance of the superconductor on the amplitude H_0 of the incident electromagnetic wave. This very unusual nonlinear phenomenon can be observed in thin superconducting slabs at not very strong ac amplitudes if the frequency of the irradiating field is close to the Josephson plasma frequency. Using the set of coupled sine-Gordon equations, we derive the expression for the phase shift χ of the reflected wave and obtain the conditions for the appearance of hysteresis in the $\chi(H_0)$ dependence.

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I. INTRODUCTION

There has been a recent surge of studies of electromagnetic waves (EMWs) propagating in artificially fabricated media (see, e.g., Refs. 1–3), including metals with modulated properties,[4] arrays of coupled waveguides,[5] left-hand materials,[6–8] and layered superconductors.[9] The excitation of these waves can produce a large variety of resonance anomalies[10] in the reflectivity, transmissivity, and absorptivity, offering different types of optical nanodevices.

The recent increase in this type of studies is related to nonlinear surface and waveguide electromagnetic modes (see, e.g., Refs. 1, 5, and 7). In this broad context, a layered superconductor is a medium favoring the propagation of nonlinear[11–13] and surface[9] waves in the (important for applications[14,15]) terahertz (THz) and sub-THz frequency ranges. Both the existence of surface waves[9] and the nonlinear effects[11–13] occur due to the gap structure (see, e.g., Ref. 16) of the spectrum of Josephson plasma waves, which was experimentally observed via Josephson plasma resonance.[17] The nonlinearity of Josephson plasma waves (JPWs) with frequency ω close to the Josephson plasma frequency ω_J becomes important even at small field amplitudes $\propto|1-\omega^2/\omega_J^2|^{1/2}$. In close analogy to nonlinear optics,[18] the nonlinear JPWs exhibit numerous remarkable features,[11–13] including the slowing down of light, self-focusing effects, and the pumping of weaker waves by stronger ones. However, the nonlinearity of EMWs in layered superconductors is quite different from optical nonlinearities. This leads one to expect very unusual phenomena in the EMW propagation in this nonlinear media.

In this paper, we predict and analyze theoretically one of such unexpected nonlinear effects in a thin slab of a layered superconductor subject to an externally applied electromagnetic wave. We show that, under specific conditions, the amplitude dependence of the phase χ of the reflected wave becomes many valued. This should result in hysteretic jumps of χ when sweeping the amplitude of the incident wave, a phenomenon which is very unusual for conductors and superconductors.

II. PROBLEM STATEMENT AND EQUATIONS FOR THE ELECTROMAGNETIC FIELD

Consider a slab of a layered superconductor of thickness d (see Fig. 1). The crystallographic **ab** plane coincides with the xy plane and the **c** axis is along the z axis. The interlayer distance D is much smaller than the thickness d of the slab.

Let the sample be irradiated by two p-polarized (transverse magnetic) plane monochromatic electromagnetic waves with the magnetic fields symmetric with respect to the middle of the sample: the plane $z=0$. Therefore, the magnetic $\vec{H}=\{0,H,0\}$ and electric $\vec{E}=\{E_x,0,E_z\}$ fields satisfy the symmetry conditions,

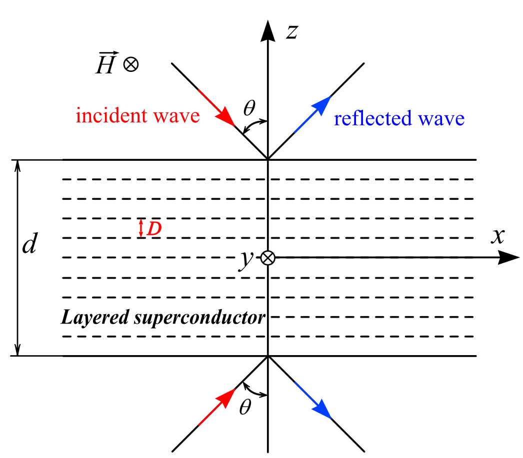

FIG. 1. (Color online) Geometry of the problem. The slab of a layered superconductor is irradiated with a p-polarized electromagnetic wave with the magnetic field symmetrical with respect to the middle of the sample.

$$H(x,z,t) = H(x,-z,t),$$

$$E_x(x,z,t) = -E_x(x,-z,t), \quad E_z(x,z,t) = E_z(x,-z,t). \tag{1}$$

Due to this symmetry, we will only consider the field distribution in the upper half space $z>0$.

The electromagnetic field in the vacuum, $z>d/2$, is the sum of the incident and specularly reflected waves. The Maxwell equations give for them

$$H^V(x,z,t) = H_0 \cos \gamma_- + H_r \cos(\gamma_+ + \chi),$$

$$E_x^V(x,z,t) = -\frac{k_z}{k}[H_0 \sin \gamma_- - H_r \sin(\gamma_+ + \chi)],$$

$$\gamma_- = k_x x - k_z z - \omega t, \quad \gamma_+ = k_x x + k_z z - k_z d - \omega t, \tag{2}$$

with $k_x = k \sin \theta$, $k_z = k \cos \theta$, and $k=\omega/c$. Here ω is the wave frequency, c is the speed of light, and θ is the angle of incidence. The value of χ is the phase shift of the reflected wave at the boundary $z=d/2$ of the slab. As is known, χ defines the surface reactance X of a sample: when neglecting the dissipation,

$$X = \frac{4\pi}{c}\tan\left(\frac{\chi}{2}\right)\cos \theta.$$

Inside a layered superconductor, the electromagnetic field is determined by the interlayer gauge-invariant phase difference φ of the order parameter. The spatial distribution of $\varphi(x,z,t)$ obeys the set of coupled sine-Gordon equations (see, e.g., Ref. 19 and 20). We consider the nonlinear JPWs with $|\varphi| \ll 1$, when the Josephson current $J_c \sin \varphi$ can be approximated by $J_c(\varphi - \varphi^3/6)$. We also assume that the gauge-invariant phase difference experiences small changes on the scale D in the z direction, and thus we can use the continuum approach. In the continuum limit, the coupled sine-Gordon equation has the form,

$$\left(1 - \lambda_{ab}^2 \frac{\partial^2}{\partial z^2}\right)\left(\frac{1}{\omega_J^2}\frac{\partial^2 \varphi}{\partial t^2} + \varphi - \frac{\varphi^3}{6}\right) - \lambda_c^2 \frac{\partial^2 \varphi}{\partial x^2} = 0. \tag{3}$$

Here λ_{ab} and $\lambda_c = c/\omega_J \varepsilon^{1/2}$ are the London penetration depths across and along layers, respectively, and $\omega_J = (8\pi e D J_c/\hbar \varepsilon)^{1/2}$ is the Josephson plasma frequency. The latter is determined by the maximum Josephson current J_c, the interlayer dielectric constant ε, and the interlayer spacing D. The spatial variations in the z direction of the fields inside the very thin superconducting layers are neglected. Here we also omit the dissipation terms related to the quasiparticle conductivity. They are controlled by the sample temperature T and can be reduced to negligibly small values. Moreover, in Eq. (3), we neglect the term with the capacitive coupling, for waves with sufficiently high $k_x = k \sin \theta \sim \omega/c \gg \beta/\lambda_c$. Here $\beta = R_D^2 \varepsilon / sD \ll 1$ is the prefactor of the capacitive coupling,[21] R_D is the Debye length for a charge in a superconductor, and s is the thickness of the superconducting layers.

The coupled sine-Gordon Eq. (3) for weakly nonlinear Josephson plasma waves was derived in Ref. 13. The reasons why dissipation is negligible in the regime $|1-\omega^2/\omega_J^2| \ll 1$ are also discussed in Ref. 13. The magnetic and electric fields in a layered superconductor are related to the gauge-invariant phase difference as

$$\frac{\partial H^S}{\partial x} = \frac{\mathcal{H}_0}{\lambda_c}\left(\frac{1}{\omega_J^2}\frac{\partial^2 \varphi}{\partial t^2} + \varphi - \frac{\varphi^3}{6}\right),$$

$$E_x^S = -\frac{\lambda_{ab}^2}{c}\frac{\partial^2 H^S}{\partial z\, \partial t}, \quad \mathcal{H}_0 = \frac{\Phi_0}{2\pi D \lambda_c}, \tag{4}$$

where $\Phi_0 = \pi c \hbar / e$ is the flux quantum and e is the elementary charge.

As was shown in Ref. 13, the nonlinearity in Eq. (3) can play a crucial role in the wave propagation for frequencies close to ω_J, i.e., for $|1-\Omega^2| \equiv |1-\omega^2/\omega_J^2| \ll 1$. Indeed, if $\varphi \sim \sqrt{|1-\Omega^2|} \ll 1$, the cubic term φ^3 in Eq. (3) is of the same order as the linear term $\omega_J^{-2}\partial^2\varphi/\partial t^2 + \varphi$.

We consider the frequency range below the Josephson plasma frequency, $\Omega < 1$, and seek a solution of Eq. (3) of the form,

$$\varphi(x,z,t) = a(z)(1-\Omega^2)^{1/2}\sin(\gamma_0 + \alpha),$$

$$\gamma_0 = k_x x - k_z d/2 - \omega t, \tag{5}$$

keeping only the first harmonics in $(k_x x - \omega t)$.

Substituting φ in Eq. (5) into Eq. (4), one obtains

$$H^S(x,\zeta,t) = -\mathcal{H}_0 \frac{(1-\Omega^2)}{\kappa} h(\zeta)\cos(\gamma_0 + \alpha),$$

$$E_x^S(x,\zeta,t) = \mathcal{H}_0 \frac{(1-\Omega^2)}{\kappa} P h'(\zeta)\sin(\gamma_0 + \alpha). \tag{6}$$

Here we introduce the dimensionless variables,

$$h(\zeta) = a(\zeta) - \frac{a^3(\zeta)}{8}, \quad \zeta = \frac{\kappa z}{\lambda_{ab}}, \tag{7}$$

and parameters

$$P = \frac{\lambda_{ab}}{\lambda_c}\frac{\kappa}{\sqrt{\varepsilon}}, \quad \kappa = \frac{\lambda_c k_x}{(1-\Omega^2)^{1/2}}, \tag{8}$$

and the prime denotes $d/d\zeta$.

Equations (3) and (5) yield the ordinary second-order differential equation for $a(z)$:

$$\left[1 - \kappa^2 \frac{d^2}{d\xi^2}\right]\left(a - \frac{a^3}{8}\right) + \kappa^2 a = 0. \tag{9}$$

Further, we also assume $\kappa \gg 1$, which is valid for not very small incident angles θ. In this case, integrating Eq. (9) with the symmetry condition

$$a'(0) = 0,$$

we obtain

$$\frac{3}{4}(a')^2 = \left(\frac{8-3a_0^2}{8-3a^2}\right)^2 - 1, \tag{10}$$

where $a_0 = a(0)$. The solution of Eq. (10) can be written in the implicit form,

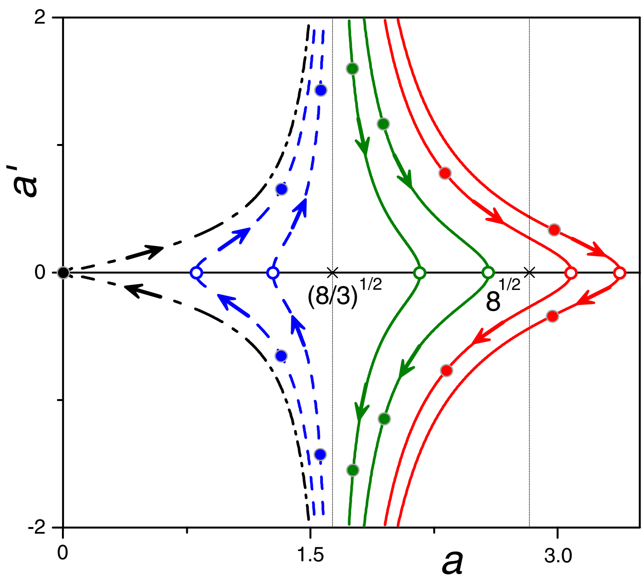

FIG. 2. (Color online) The phase diagram $a'(a)$ (for $a>0$). Recall that a is the amplitude of the gauge-invariant phase φ [see Eq. (5)] while a' is its derivative with respect to the dimensionless coordinate $\zeta=\kappa z/\lambda_{ab}$. The paths along the arrows on the phase trajectories between solid circles correspond to the change of the coordinate z inside the sample, from $z=-d/2$ to $z=d/2$. Open circles correspond to the middle of the slab ($z=0$).

$$\zeta = \sqrt{\frac{3}{4}}\int_{a_0}^{a(\zeta)} da \frac{8-3a^2}{\sqrt{(8-3a_0^2)^2-(8-3a^2)^2}}. \quad (11)$$

The phase diagram, i.e., the set of $a'(a)$ curves for different values of the constant a_0, is shown in Fig. 2. Solid circles mark the sample boundaries while open circles indicate the middle of the slab. Arrows show the direction of motion along the phase trajectories when increasing the coordinate ζ.

Thus the electromagnetic fields in the vacuum and in the superconducting slab are determined by Eqs. (2) and (6), respectively. The latter equations contain the function $a(\zeta)$ given by Eq. (11).

III. PHASE SHIFT OF THE REFLECTED WAVE

Now we find the relationship between the phase shift χ of the reflected wave and the amplitude H_0 of the incident wave. With this purpose, we join the tangential components of the electric and magnetic fields in the vacuum and in the superconducting slab at the interface $z=d/2$. Thus, separating terms with $\sin(k_x x-\omega t)$ and $\cos(k_x x-\omega t)$, we derive four equations for H_r, $H(d/2)$, χ, and α:

$$-h_0 + h_r \cos\chi = Ph'(\delta)\sin\alpha,$$

$$-h_0 - h_r \cos\chi = h(\delta)\cos\alpha,$$

$$-h_r \sin\chi = Ph'(\delta)\cos\alpha,$$

$$-h_r \sin\chi = h(\delta)\sin\alpha. \quad (12)$$

Here

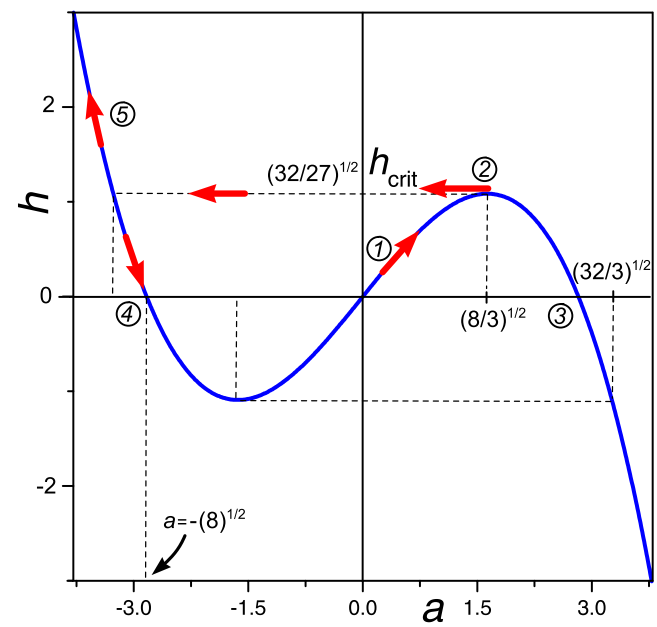

FIG. 3. (Color online) The normalized magnetic field h versus the normalized amplitude a of the gauge-invariant phase ϕ. This $h(a)$ dependence is described by Eq. (7). Arrows indicate the motion along the $h(a)$ curve when changing the normalized amplitude h_0 of the incident wave.

$$h_0 = \frac{H_0}{\mathcal{H}_0}\frac{\kappa}{(1-\Omega^2)}, \quad h_r = \frac{H_r}{\mathcal{H}_0}\frac{\kappa}{(1-\Omega^2)},$$

$\delta=\kappa d/2\lambda_{ab}$ is the value of ζ at the interface $z=d/2$. Note that, without any loss of generality, we can assume the value $h(\delta)$ of the total magnetic field at the sample surface to be positive, $h(\delta)>0$. A negative $h(\delta)$ corresponds to replacing $\chi\to\chi+\pi$ in Eq. (12).

Excluding h_r and α from Eq. (12), we find

$$h_0 = \tfrac{1}{2}\sqrt{h^2(\delta)+[Ph'(\delta)]^2}, \quad (13)$$

$$\chi = 2\arctan\left(\frac{Ph'(\delta)}{h(\delta)}\right). \quad (14)$$

Equations (13) and (14), together with Eqs. (7) and (11), give, in an implicit form, the required dependence of the phase shift χ on the amplitude of the incident wave H_0.

For further analysis of the amplitude dependence of the phase shift χ, it is very important to take into account the non-single-valued relation between the values $h(\delta)$ and $a(\delta)$. Indeed, Eq. (7) and Fig. 3 show that there exist *three values* of a that correspond to the *same value* of h if $0<h<(32/27)^{1/2}$. Taking into account Eq. (13), we conclude that three different values of the parameter $h'(\delta)/h(\delta)$ in Eq. (14) correspond to the same value of h_0. This results in the appearance of *three branches* of the dependence $\chi(h_0)$.

Below, for simplicity, we restrict ourselves to the case of small sample thicknesses, when

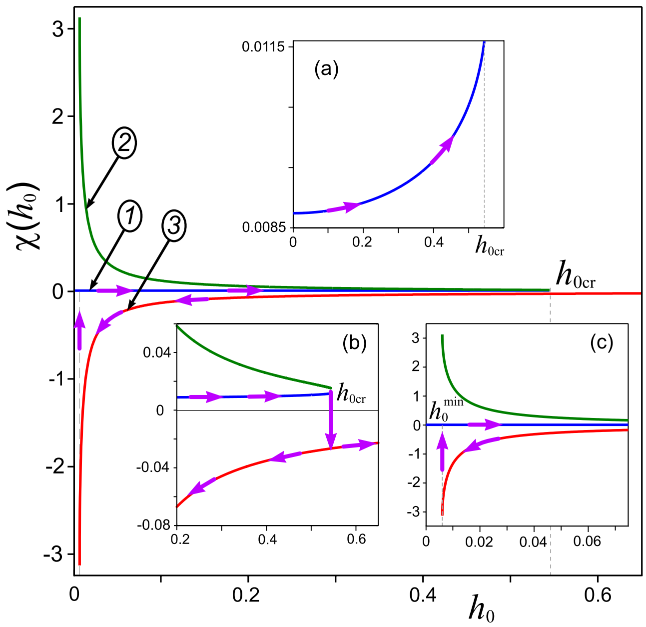

FIG. 4. (Color online) The numerically obtained dependence of the phase shift χ of the reflected wave on the dimensionless amplitude h_0 of the incident electromagnetic wave. The values of the parameters here are: $d/\lambda_{ab}=0.05$, $\lambda_c/\lambda_{ab}=200$, $\kappa=10$, and $\varepsilon=16$. The insets show magnified portions of the $\chi(h_0)$ plots. The arrows on the curves indicate how the phase shift χ changes when periodically varying the amplitude h_0. The first jump in $\chi(h_0)$ at $h_0=h_{0\text{cr}}$ is shown by the downward vertical arrow in inset (b). This produces a jump $\Delta\chi\approx-9P\delta$. The second (reverse) jump at $h_0=h_0^{\min}$ is shown in the main frame and in inset (c). Namely, the upward vertical arrow [from $\chi(h_0)|_{\text{before jump}}=-\pi$ to $\chi(h_0)|_{\text{after jump}}\approx 0$] shows the jump in $\chi(h_0)$.

$$\delta \ll 1. \tag{15}$$

In this case, $|a(\zeta)-a_0|\ll 1$, and Eq. (11) is significantly simplified,

$$a(\zeta) \simeq a_0\left(1+\frac{4\zeta^2}{8-3a_0^2}\right), \tag{16}$$

and one can easily derive the asymptotic equations for all three branches $\chi(h_0)$.

A. Low-amplitude branch $\chi(h_0)$

First, we discuss the branch of the dependence $\chi(h_0)$ that corresponds to the portion "1–2" of the $h(a)$ curve in Fig. 3. This monotonically increasing branch is shown by the blue curve (with number 1) in the main panel of Fig. 4, and also in inset (a) of Fig. 4. This blue branch is defined within the interval $(0,h_{0\text{cr}})$ of the h_0 change. The ending point $h_{0\text{cr}}$ of the branch is described by the expression,

$$h_{0\text{cr}} \simeq \sqrt{\frac{8}{27}}+\sqrt{\frac{3}{8}}P^2\delta^2. \tag{17}$$

At this point,

$$\chi(h_{0\text{cr}}) \simeq 3P\delta-\sqrt{\frac{9}{8}}P\delta^2. \tag{18}$$

The low-amplitude asymptotics of $\chi(h_0)$ dependence is

$$\chi \simeq 2P\delta\left(1+\frac{h_0^2}{2}\right), \quad h_0 \ll 1. \tag{19}$$

B. High-amplitude branch $\chi(h_0)$

The second and third branches of the $\chi(h_0)$ dependence correspond to the portions "2–3" and "4–5," respectively, of the $h(a)$ curve in Fig. 3. The second and third branches of $\chi(h_0)$ are shown by the green and red curves (indicated by numbers 2 and 3, respectively) in Fig. 4. The second branch exists in the interval

$$h_0^{\min} < h_0 < h_{0\text{cr}}, \quad h_0^{\min} \simeq \sqrt{2}P\delta. \tag{20}$$

As is seen in Fig. 4, the first and second branches of $\chi(h_0)$ almost meet at the point $h_{0\text{cr}}$. The third branch is defined for

$$h_0^{\min} < h_0 < \infty. \tag{21}$$

The lower value $h_0=h_0^{\min}$ for the second and third branches $\chi(h_0)$ corresponds to the points 3 and 4 on the $h(a)$ curve in Fig. 3. At these points, $a(\delta)=\pm\sqrt{8}$. The asymptotic expansion of the integral in Eq. (11), near $\xi=\pm\delta$, provides

$$a(\xi) = \pm\left[\sqrt{8}-\sqrt{h_0^2-2P^2\delta^2}+\frac{\delta^2-\xi^2}{\sqrt{2}}\right], \tag{22}$$

where the plus and minus signs correspond to the second and third branches $\chi(h_0)$, respectively. Using this formula [and Eqs. (13), (14), and (7)], we derive the asymptotic equations for the second and third branches near the minimum value of h_0, i.e., for $h_0 \sim h_0^{\min}$:

$$\chi(h_0) = \pm\left(\pi-\sqrt{\frac{2h_0^2}{\delta^2}-4P^2}\right). \tag{23}$$

The third branch tends to zero for $h_0\to\infty$ following the expression,

$$\chi(h_0) = -P\delta\left(\frac{4}{h_0}\right)^{2/3}, \quad h_0 \gg 1. \tag{24}$$

Now we present the parametrically defined formula for all three branches, which is valid for h_0 not very close to $h_{0\text{cr}}$:

$$h_0(a_0) = \frac{a_0|8-a_0^2|}{16}, \tag{25a}$$

$$\chi(a_0) = \frac{16P}{8-a_0^2}\delta. \tag{25b}$$

Figure 5 demonstrates the very good agreement of this formula with numerical results.

C. Hysteresis jumps in the dependence $\chi(h_0)$

From the analysis shown above, we can now describe the behavior of the phase shift χ of the reflected wave when periodically changing the amplitude h_0 of the incident wave.

When increasing h_0 from zero, the phase shift $\chi(h_0)$ increases monotonically following the first branch in Fig. 4. At

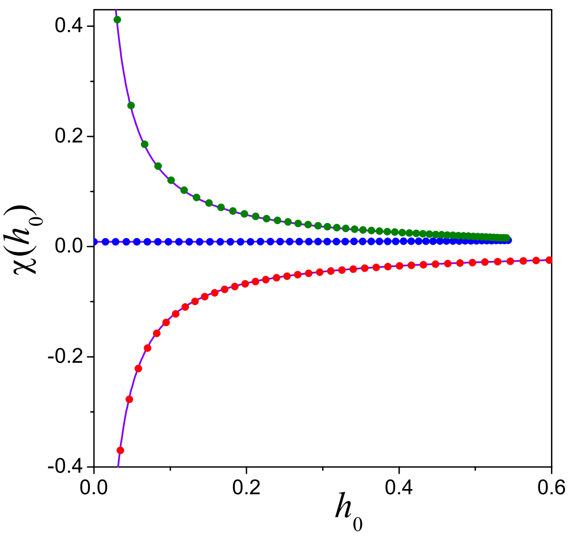

FIG. 5. (Color online) The phase shift χ of the reflected wave versus the dimensionless amplitude h_0 of the incident wave, for $d/\lambda_{ab}=0.05$, $\lambda_c/\lambda_{ab}=200$, $\kappa=10$, and $\varepsilon=16$. The solid curves are the plots of the three branches of the function given by Eq. (25) while the dots were numerically obtained from Eqs. (7), (11), (13), and (14).

the point $h_0=h_{0\mathrm{cr}}$, the first branch comes to the end, and a jump to the third branch should occur by further increasing h_0. The phase shift $\chi(h_0)$ jumps from the value

$$\chi(h_{0\mathrm{cr}}-0) \simeq 3P\delta$$

to the final value

$$\chi(h_{0\mathrm{cr}}+0) \simeq -6P\delta.$$

Thus, the jump here is

$$\Delta\chi|_{h_0=h_{0\mathrm{cr}}} \approx -9P\delta. \tag{26}$$

When decreasing h_0, the phase shift $\chi(h_0)$ decreases following the third branch, crosses the point $h_0=h_{0\mathrm{cr}}$, and only at the point $h_0=h_0^{\mathrm{min}}$ it performs a reverse jump to the first branch. Therefore, now

$$\Delta\chi|_{h_0=h_0^{\mathrm{min}}} \approx \pi. \tag{27}$$

Thus, the hysteretic jumps $\Delta\chi$ of the phase $\chi(h_0)$ of the reflected wave could be observed when periodically changing the amplitude h_0 of the incident wave.

IV. CONCLUSION

In this paper, we have predicted and theoretically analyzed unusual phenomenon for conducting media. We have shown that, due to the specific nonlinearity of layered superconductors, hysteretic jumps of the surface reactance could be observed when periodically changing the amplitude H_0 of the incident wave. A remarkable feature of the predicted phenomenon is the relatively small values of the necessary ac amplitudes H_0. The hysteretic jumps can be observed even when $\varphi \ll 1$. According to our analysis, the critical amplitude is

$$H_{0\mathrm{cr}} = h_{0\mathrm{cr}} \cdot \mathcal{H}_0 \frac{(1-\Omega^2)}{\kappa} \ll \mathcal{H}_0 \sim 20\ \mathrm{Oe}.$$

There are two small parameters here, $(1-\Omega^2)$ and $1/\kappa$. As was shown in Ref. 13, the heating effect is negligible for such ac amplitudes.

The phenomenon discussed here is another exciting example of the numerous unusual effects related to the very specific nonlinearity of layered superconductors.

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