Soliton-plasmon resonances as Maxwell nonlinear bound states

C. Milián,1,2,* D. E. Ceballos-Herrera,3 D. V. Skryabin,2 and A. Ferrando4
1Instituto de Instrumentación para Imagen Molecular (IIM), InterTech, Universidad Politécnica de Valencia, Camino de Vera S/N, Valencia 46022, Spain
2Centre for Photonics and Photonic Materials, Department of Physics, University of Bath, Bath BA2 7AY, UK
3Universidad Autónoma de Nuevo León, Facultad de Ciencias Físico Matemáticas, Avenida Universidad S/N, Cd. Universitaria, Nuevo León 66450, México
4Departament d’Òptica, Interdisciplinary Modeling Group InterTech, Universitat de València, Dr. Moliner 50, Burjassot, València 46100, Spain
*Corresponding author: cme22@bath.ac.uk

Received July 25, 2012; revised September 3, 2012; accepted September 4, 2012; posted September 5, 2012 (Doc. ID 173035); published October 5, 2012

We demonstrate that soliplasmons (soliton–plasmon bound states) appear naturally as eigenmodes of nonlinear Maxwell’s equations for a metal/Kerr interface. Conservative stability analysis is performed by means of finite element numerical modeling of the time-independent nonlinear Maxwell equations. Dynamical features are in agreement with the presented nonlinear oscillator model. © 2012 Optical Society of America

OCIS codes: 190.6135, 240.6680.

Nanoscaled plasmonic optical solitons have attracted much attention in the last few years. Recent studies report plasmonic solitons in single and double metal/dielectric interfaces [1,2], systems with gain and loss [3,4], waveguide arrays [5–7], and chains of nanoparticles [8]. Strong surface plasmon polariton (SPP) fields at the interface enhance the nonlinear effect [9].

Recently, it has been suggested that SPPs can couple to spatial solitons [10] because both have a wavenumber greater than that of the light cone (see, e.g., [11–13]) and their dispersion relations intersect [see Fig. 1(b)]. A symmetric coupled oscillator model was introduced by means of heuristic reasoning, and its dynamical properties were fully analyzed in [14]. Although surface solitons were earlier studied (see, e.g., [15,16]), they consisted of one component only, whereas the soliplasmons considered here consist of two: one peaked at the metal interface and the other one peaked far from it in the Kerr medium.

In this Letter, we prove for the first time that the soliplasmons proposed in [10] exist in the context of Maxwell equations. We compute numerically the soliplasmons as eigenstates of the nonlinear Maxwell equations for a metal/Kerr interface [Fig. 1(a)], showing that they are classified according to the relative soliton-plasmon phase $\delta = 0, \pi$. Finite element analysis modeling is used to integrate the time-independent nonlinear Maxwell equations in two dimensions to analyze soliplasmon stability. The non-self-adjoint character of the Maxwell operator (property of the vectorial nature of the system) leads to an oscillator model with asymmetric coupling between the SPP and the soliton. This model has no unknown parameters (as opposed to [10]) and therefore predicts realistic physical properties of the stationary solutions and their stability. Inclusion of ohmic losses is observed to yield a nontrivial dynamics.

We analyze the full-vector nonlinear equation for a monochromatic wave (cw) $E_\omega$ with frequency $\omega = ck$,

$$\frac{\partial^2}{\partial z^2} + k^2 \epsilon_L(x) E_\omega = \mathcal{L}_\omega(x) E_\omega - P_{NL},$$

where $\epsilon_L = \epsilon_{m(K)}$ for $x < 0$ ($x > 0$), $\mathcal{L}_\omega \equiv \nabla [\nabla \phi]$, $P_{NL} = [k^2 \chi^{(3)}/3][2|E_\omega|^2 E_\omega^* E_\omega^*]$, and $\chi^{(3)} \equiv \epsilon_0 \epsilon_\omega \epsilon_2$. Our geometry is assumed to be illuminated from the Kerr medium [see Fig. 1(a)], and diffraction along $y$ is neglected. Equation (1) can be transformed into the dynamical equations of a soliplasmon by using the variational ansatz of [10].

$$E_\omega(x, z) = [c_p(z)e_p(x) + uc_s(z)f_s(x - \alpha; c_s(z))]e^{i\beta k x z},$$

where $c_p$, $c_s$ are the complex amplitudes, $e_p(x)$ is a TM-SPP on the interface with propagation constant $\beta_p$, which is a stationary solution of Eq. (2), and hence $c_p(z) = c_p(0)e^{i\beta_p z}$. The soliton term in Eq. (2), $f_s(z) = c_s e^{i\kappa_s(x - a)}$, $\kappa_s \equiv |k n_K|^{1/2} |c_s|$, is located at a distance $a$ from the interface, such that the overlap with the SPP is small (weak coupling), and it is a solution of the stationary scalar ($\mathcal{L}_\omega = 0$) and paraxial nonlinear Schrödinger equation, $\{1/[2k n_K]\}^{1/3} + \gamma |\phi_s|^2 \phi_s = \mu_s \phi_s$ ($\gamma \equiv k \chi^{(3)}/[2 n_K]$), with $c_s(z) = c_s(0)e^{i\beta_s z}$.

![Fig. 1. (Color online) (a) Metal/Kerr structure with linear dielectric constants $\epsilon_m = -n_m^2$, $\epsilon_K = n_K^2$ and nonlinear Kerr index $n_2$. (b) Dispersion of a SPP in a lossy metal (black) and a spatial soliton (gray) owning two different amplitudes ($\beta_s = k n_K [1 + g |\phi_s|^a]$). Circles enclose the matching points, and the dashed line marks the light cone $\omega = \beta c/\sqrt{\epsilon_K}$.](image-url)
and $\beta = k n_K + \mu_s$, where $\mu_s = \gamma |c_x(0)|^2 / 2$. Note that $c_s(z)$ appears nonlinearly in Eq. (2), which prevents the soliplasmon from behaving as a linear superposition of the two modes. Considering that $a$ is constant and that the soliton is $x$-polarized, i.e., $\mathbf{u} = \hat{x}$, (as supported by Figs. 3–5), substitution of Eq. (2) in the paraxial version of Eq. (1), $-i \partial_x c_x = \hat{M} c_x$, leads to the soliplasmon equations [17] ($\Psi \equiv d\Psi / dz$)

$$-i \partial c = M c, \quad M = [\hat{M}] = \begin{pmatrix} \mu_p & q(c_x) \\ \bar{q}(c_x) & \mu_s \end{pmatrix},$$

where $|c\rangle \equiv [c_p, c_s]^T$. The origin of the coupling in Eq. (3) is the nonorthogonality relation $\int \mathbf{E}_p \cdot \mathbf{E}_s \approx 0$, and it is not symmetric in general, $\bar{q}/q \sim N_p / N_s$, where $q \equiv k / \{2 n_K N_p \} \int \mathbf{E}_p \cdot \mathbf{E}_s \sim \exp \left( -a \sqrt{k n_K} |c_s| \right)$, $N_p \equiv \int |\mathbf{E}_p|^2$ and $N_s \equiv \int |\mathbf{E}_s|^2$. This feature was not captured by previous heuristic models [10, 14]. Interestingly, $q, \bar{q}$ are proportional to the value of the soliton tail at the interface \[10\], revealing that a strong soliton drives a weak SPP ($N_s \gg N_p$) at a rate $q$.

Stationary soliplasmons of Eq. (3) are determined from the eigenvalues, $\mu_s$, and eigenvectors, $|\mu_s\rangle = c_s[|\mu_p - \mu_s|, 1]^T$ of $M (c_s \in C)$,

$$\beta = k n_K + \bar{\mu} + e^{i \delta} \sqrt{\Delta_\mu + q \bar{q}}, \quad \delta = 0, \pi.$$

$$E_s(x, z) = \left\{ \frac{q e^{i x} (x)}{\mu_p - \mu_x} + \text{sech}(\kappa_s |x - a|) \right\} e^{i \beta z},$$

where $\beta = k n_K + \mu_s$, $\bar{\mu} \equiv [\mu_p + \mu_x] / 2$, and $\Delta_\mu \equiv [\mu_p - \mu_x]/2$. Note from Eq. (4) that $\mu_p > \mu_x$, and $\mu_x < \mu_p$, so the plasmon term in Eq. (5) is $> 0$ ($< 0$) for $\delta = 0$ ($\pi$) and $\delta$ is the relative soliton–plasmon phase. The minimum value of $\mu_s - \mu_p = \sqrt{q \bar{q}} \exp(i \delta)$ imposes a maximum in the soliplasmon norm, the solution with infinite power $\mu_s = \mu_p$, being nonphysical. We verified these features by plotting the guided power $P(\mu, \omega, a) = \int_0^\ell E_s^2 H_y^2 / 2 = \omega n_0 \int_0^\ell |E_s|^2$ of the numerically computed soliplasmons in Fig. 2. At fixed $(\omega, a)$ there are two divergent branches, $\mu_0 > \mu_p$ (right) and $\mu_x < \mu_p$ (left), in agreement with Eq. (4). Far from the asymptote, both branches coalesce into the monotonically increasing soliton curve $P_s(\mu) \sim \mu^{1/2}$, but close to it they open a gap in $\mu$, which is proportional to $\sqrt{q \bar{q}} \sim \exp \left( -a \sqrt{k n_K} |c_s| \right)$. Soliplasmons are computed from Eq. (1) ($E(x, z) = E(x) e^{i \beta z}$) on a silver/glass interface ($\epsilon_{m} = 82, \epsilon_K = 2.09, n_g = 2.6 \times 10^{-20} \text{m}^2 / \text{W}$), by means of an iterative Fourier method that fixes $a$, letting the families $P(\mu; a)$ be found separately. Soliplasmons with $\delta = 0, \pi$ naturally appear (see insets of Fig. 2).

The stability of soliplasmons was checked by propagating the solutions with input noise (20% in amplitude) and without losses ($\epsilon_{p, s} \in \Omega$). Focusing a beam into the Kerr material is likely to excite soliplasmons with a strong solitonic component, so we focus below on the case $N_p / N_s \ll 1$ ($q \ll q$), in which the soliton dynamics is quasi-stationary, i.e., $|c_s| / |c_x| \ll |c_p| / |c_s|$. In this situation, the dynamics can be qualitatively predicted by rewriting Eq. (3) as two equations for the relative phase, $\phi_{sp} \equiv \phi_p - \phi_s$, and $|c_p| = |c_p| \exp (i \phi_{sp})$,

$$\dot{\phi}_{sp} = 2 \Delta_\mu + q |c_s| \cos \phi_{sp}, \quad |c_p| = q |c_s| \sin \phi_{sp}. \quad (6)$$

Perturbations introduced here induce an increase of $\phi_{sp}$ [see Figs. 4(b) and 5(b)]. The amplitude Eq. (6) predicts that in this situation $|c_p|$ will increase (decrease) if $\sin \phi_{sp} > 0 (< 0)$. These features are clear in our dynamical simulations, which integrate Eq. (1) with no approximations and permit us to evaluate $c_p, c_s$ as the peak amplitudes of the plasmon and soliton components. Figure 2 shows the propagation of a $\delta = \pi$ solution. Apart from being diffraction free, the input noise introduces fluctuations that propagate away from the soliplasmon.
The decrease of $|c_p|$ is associated to the transfer of energy from the SPP to the soliton.

Propagation of a $\delta = 0$ solution (see Fig. 4) shows a very different behavior, since the initial increase of $|c_p|$ implies that the SPP drains energy from the soliton. Remarkably, $|c_p| = 0$ at $\phi_{sp} = 0$, $\pm \pi$ and the flow of energy between the soliton and SPP is reversed [see Figs. 4(b), 4(c)], as Eq. (6) predicts. $\delta = \pi$ soliplasmons appear to be more oscillatory stable than the $\delta = 0$ ones, stability meaning a small energy transfer between the soliton and SPP components.