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Quadratic solitons as nonlocal solitons

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We show that quadratic solitons are equivalent to solitons of a nonlocal Kerr medium. This provides new physical insight into the properties of quadratic solitons, often believed to be equivalent to solitons of an effective saturable Kerr medium. The nonlocal analogy also allows for analytical solutions and the prediction of bound states of quadratic solitons.

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I. INTRODUCTION

Quadratic nonlinear (or $\chi^{(2)}$) materials have a strong and fast electronic nonlinearity, which makes them excellent materials for the study of nonlinear effects, such as solitons [1]. The main properties of quadratic solitons are well known [2] and both (1+1)-dimensional [3] and (2+1)-dimensional [4] bright spatial solitons have been observed experimentally. Unlike conventional solitons, which form due to a self-induced refractive index change, the formation of quadratic solitons does not involve any change of the refractive index. Thus the underlying physics of quadratic solitons is often obscured by the mathematical model. Only recently Assanto and Stegeman used the concepts of the cascading phase shift and parametric gain to give an intuitive interpretation of effects, such as self-focusing, defocusing, and soliton formation in $\chi^{(2)}$ materials [5].

Nevertheless certain features of quadratic solitons, such as formation of bound states, are still without a physical interpretation. The nice simple phase-shift approach of Assanto and Stegeman [5] predicts, e.g., that two dark solitons and two out-of-phase bright solitons will always repel and thus can never form a bound state, whereas it is known that such bound states of quadratic solitons do exist [6].

Here we use the analogy between parametric interaction and nonlocality and present a physically intuitive nonlocal theory, which is exact in predicting the profiles of stationary quadratic solitons and which provides a simple physical explanation for their properties including formation of bound states.

The nonlocal analogy was applied recently by Shadrivov and Zharov to find approximate bright quadratic soliton solutions [7]. The formal equivalence of bright solitons in nonlocal liquid crystals and parametric solitons was also discussed recently by Conti, Peccianti, and Assanto [8]. However, the nonlocal concept was not fully exploited in Refs. [7,8] to give a broad physical picture in the whole regime of existence and discuss, e.g., dark solitons and bound states of out-of-phase bright solitons.

We do that here and we go one step further in showing how a simple phase-sensitive nonlocal model provides a better description of the dynamics in $\chi^{(2)}$ materials than the nonlinear Schrödinger (NLS) equation obtained in the cascading limit.

We consider a fundamental wave (FW) and its second harmonic (SH) propagating along the $z$ direction in a lossless quadratic nonlinear medium under conditions for type I phase matching. The normalized dynamical equations for the slowly varying envelopes $E_{1,2}(x,z)$ are then [9]

\begin{align}
\dot{i}\partial_x E_1 + d_1 \partial^2_x E_1 + E_1^* E_2 \exp(-i \beta z) &= 0, \\
\dot{i}\partial_x E_2 + d_2 \partial^2_x E_2 + E_1^2 \exp(i \beta z) &= 0.
\end{align}

In the spatial domain $d_1 \approx 2d_2$, $d_j > 0$, and the coordinate $x$ represents a transverse spatial direction. The term $\partial^2_x E_j$ represents pulse dispersion. The parameter $\beta$ is the normalized phase mismatch and $j = 1,2$.  

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II. GENERAL NONLOCAL MODEL

Physical insight into the properties of Eqs. (1) and (2) may be obtained from the cascading limit, in which the phase mismatch is large, $\beta^{-1} \rightarrow 0$. Writing $E_2 = e_2 \exp(i\beta z)$ and assuming slow variation of $e_2(x,z)$ gives the NLS equation

$$i\partial_t E_1 + d_1 \partial_z^2 E_1 + \beta^{-1}|E_1|^2 E_1 = 0,$$  

in which the local Kerr nonlinearity is due to the coupling to the SH field $e_2 = E_1^2/\beta$. The SH is thus slaved to the FW and the widths of the SH and FW are fixed. The sign of the mismatch $\beta$ determines whether the effective Kerr nonlinearity is focusing or defocusing and thus the cascading limit predicts that bright and dark quadratic solitons exist for $\beta d_1 > 0$ and $\beta d_1 < 0$, respectively.

However, even for stationary solutions the NLS equation is inaccurate, since the term $\partial_z^2 E_2$ is neglected. Thus it predicts, e.g., modulational stability of dark quadratic solitons for all values of $d_2$, whereas this is known to depend on the value of $d_2$ [6]. It further predicts that in higher dimensions bright solitons are unstable and will either spread out or collapse [10], whereas it is known that stable quadratic solitons exist in all dimensions and that collapse cannot occur in the $\chi^{(2)}$ system (1) and (2) [11,12]. The stabilizing effect of the $\chi^{(2)}$ nonlinearity is often described as being due to saturation of the effective Kerr nonlinearity [5,11,13]. We show below that the nonlinearity is in fact nonlocal.

To obtain a more accurate model than that given by the cascading limit we assume a slow variation of the SH field $e_2(x,z)$ in the propagation direction only. Thus, neglecting only $\partial_x e_2$, we solve Eq. (2) exactly using Fourier transformation and the convolution theorem, treating $E_1^2$ as a function. The SH is still expressed in terms of the FW, but now the relation has the form of a convolution, leading to the nonlocal equation for the FW:

$$i\partial_t E_1 + d_1 \partial_z^2 E_1 + \beta^{-1} N(E_1^2) E_1 = 0,$$  

with $E_2 = \beta^{-1} N \exp(i\beta z)$. Equations (4) and (5) clearly show that the interaction between the FW and the SH is equivalent to the propagation of a FW in a medium with a nonlocal nonlinearity. In the Fourier domain (denoted with tilde) the response function $R(x)$ is a Lorentzian $\tilde{R}(k) = 1/(1 + s\sigma^2 k^2)$, where $\sigma = |d_2/\beta|^{1/2}$ represents the degree of nonlocality and $s = \text{sgn}(d_2 \beta)$. Both Eqs. (1) and (2) and (4) and (5) are trivially extended to include more transverse dimensions.

For $s = +1$, where the $\chi^{(2)}$ system (1) and (2) has a family of bright (for $d_1 > 0$) and dark (for $d_1 < 0$) soliton solutions [6], $\tilde{R}(k)$ is positive definite and localized, giving

$$R(x) = (2\sigma)^{-1} \exp(-|x|/\sigma).$$  

The cascading limit $\beta^{-1} \rightarrow 0$ is now seen to correspond to the local limit $\sigma \rightarrow 0$, in which the response function becomes a $\delta$ function, $R(x) \rightarrow \delta(x)$.

With the nonlocal analogy one can assign simple physically intuitive pictures to several important cases. When the mismatch $|\beta|$ is large Eqs. (4) and (5) reduce to the so-called weakly nonlocal equation with $\sigma \rightarrow 1$ [14]. Similarly the nearly phase-matched limit when $\beta \rightarrow 0$ corresponds to the strongly nonlocal limit with $\sigma \rightarrow \infty$, when Eqs. (4) and (5) become effectively linear [15,16].

For $s = -1$, $\tilde{R}(k)$ has poles on the real axis and the response function becomes oscillatory in nature with the Cauchy principal value

$$R(x) = (2\sigma)^{-1} \sin(|x|/\sigma).$$  

In this case the propagation of solitons has a close analogy with the evolution of a particle in a nonlinear oscillatory potential. In fact, it is possible to show that the oscillatory response function explains the fact that dark and bright quadratic solitons radiate linear waves for $s = -1$ [6].

Equations (4) and (5) show the important result that in contrast to the conventional nonlocal NLS equation, describing materials with a thermal [18] or diffusion [19] type nonlinearity, liquid crystals [20], and photorefractive crystals [21], the nonlocal response of the $\chi^{(2)}$ system depends on the square of the FW, not its intensity. Thus the phase of the FW enters into the picture and one cannot directly transfer the known nonlocal stability results for plane waves [15,17] and solitons [14,16].

However, the simple nonlocal model (4) and (5) is indeed an improved model of quadratic nonlinear materials, as compared with the even simpler NLS equation obtained in the cascading limit. Thus the nonlocal model correctly shows that the properties of quadratic solitons depends sensitively on the parameter $d_2$. For symmetric response functions, $R(x) = R(-x)$, the nonlocal model in arbitrary dimension conserves both power and Hamiltonian and it is straightforward to carry out the same analysis as for the conventional nonlocal NLS equation [16] and show that the Hamiltonian for the system (4) and (5) is also bounded from below. Thus the nonlocal model (4) and (5) also correctly predicts that collapse cannot occur in the $\chi^{(2)}$ system. Furthermore, for stationary fields $E_1(x,z) = E_1(x)$ the nonlocal model (4) and (5) represents an exact model for $\chi^{(2)}$ materials.

III. NONLOCAL QUADRATIC SOLITONS

The properties of nonlocal solitons in terms of profiles thus directly apply to quadratic solitons. Consider stationary solutions to Eqs. (1) and (2) in the form

$$E_1(x,z) = a_1 \phi_1(\tau) \exp(i\lambda z),$$  

$$E_2(x,z) = a_2 \phi_2(\tau) \exp(i2\lambda z + i\beta z),$$  

where the profile $\phi_1(\tau)$ is real, with $\tau = x|\lambda|/d_1|^{1/2}$, $a_1^2 = \lambda^2|d_2/(2d_1)|$, and $a_2 = \lambda$. This scaling reduces the number of free parameters to one and transforms Eqs. (1) and (2) into the following system [6]:
where $s_j = \text{sgn}(\lambda d_j) = \pm 1$, $\alpha = (2 + \beta/\lambda)|d_1/d_2|$, and prime denotes differentiation with respect to the argument. The properties of solitons described by Eqs. (10) and (11) are well known [6]. A family of bright (dark) solitons exist for $s_2 = s_1 = +1$ ($s_2 = -s_1 = +1$) and $\alpha > 0$. As discussed above we do not consider the combinations $s_2 = \pm s_1 = -1$, for which solitons do not exist in the whole $\alpha$ space.

Equation (11) has the formal solution $\phi_2 = \gamma N(\phi^2_1)$, with $\gamma = 1/(2\alpha)$ and the nonlocal nonlinearity $N(\phi^2_1)$ given by Eq. (5). For $\text{sgn}(s_2, \alpha) = +1$ the response function is $R(\tau) = (2\sigma)^{-1} \exp(-|\tau/\bar{\sigma}|)$, with the degree of nonlocality $\bar{\sigma} = |\alpha|^{-1/2}$. Inserting the SH into Eq. (10) then gives the exact nonlocal model for the FW in the $\chi^{(2)}$ system (10) and (11):

$$s_1 \phi''_1 - \phi_1 + \gamma \phi_1 \int_{-\infty}^{\infty} R(\tau - \xi) \phi^2_2(\xi) d\xi = 0,$$

where $\gamma$ is the strength of the nonlocal nonlinearity. Thus $\chi^{(2)}$ solitons are equivalent to nonlocal solitons.

In the weakly nonlocal case $\bar{\sigma} \ll 1$ (i.e., $|\alpha| > 1$) the response function $R(\tau)$ is much narrower than the FW intensity $\phi^2_1$. Taylor expansion of $\phi^2_1$ under the integral in Eq. (12) gives the weakly nonlocal model [14]

$$s_1 \phi''_1 - \phi_1 + \gamma (1 + \bar{\sigma}^2 \phi^2_1) \phi_1 = 0,$$

where $\phi_2 = \gamma (1 + \bar{\sigma}^2 \phi^2_1) \phi_1$. This model has exact bright soliton solutions for $s_2 = s_1 = +1$ and $\alpha > 0$ [14]:

$$\pm \tau = \text{tanh}^{-1}(\rho) + 2\bar{\sigma} \tan^{-1}(2\bar{\sigma} \rho),$$

where $\rho^2 = (a^2_1 - \bar{\sigma}^2)/(a^2_1 + 4\bar{\sigma}^2 \phi^2_1)$, $a^2_2 = 2/\gamma$ being the maximum intensity of the FW. Exact stationary black soliton solutions exist for $s_2 = s_1 = +1$ [14].

For $|\alpha| \ll 1$ the nonlocality is strong, $\bar{\sigma} \gg 1$, and we can expand the response function $R(\tau)$ in Eq. (12). For bright solitons we then obtain the linear equation for the FW:

$$s_1 \phi''_1 - \phi_1 + \gamma P_1 R(\tau) \phi_1 = 0,$$

where $\phi_2 = \gamma P_1 R(\tau)$. In this eigenvalue problem the FW power $P_1 = \int_{-\infty}^{\infty} \phi^2_1(\tau) d\tau$ plays the role of the eigenvalue, and bright solitons correspond to the fundamental mode of the waveguide structure one can associate with the exponential response function. For $s_2 = s_1 = +1$ and $\alpha > 0$, Eq. (15) has exact bright soliton solutions in the form of the Bessel function of the first kind of order $2\bar{\sigma}$ [22]

$$\phi_1(\tau) = A_1 J_{2\bar{\sigma}}(\bar{\sigma}^2 \sqrt{2P_1} R(\tau)).$$

For the single-soliton ground-state solution $P_1$ is found as the first zero of the derivative, $J'_{2\bar{\sigma}}(\sqrt{\bar{\sigma}^2 P_1}) = 0$, which assures that $\phi'_1(0) = 0$. The amplitude $A_1$ is then found from the definition of $P_1$, giving $A^2_1 = P_1/2\left[ P_1 (1/\pi^2 \bar{\sigma}^2) \right]^{1/2}$.
assures that

For the fundamental single-soliton solution

do not satisfy the self-consistency relation $Q_1 = \int_{-\infty}^{\infty} [A_1^2 - \phi_1^2(\tau)]d\tau$. Nevertheless the strongly nonlocal model is able to predict and physically explain the existence of multihump dark quadratic solitons found earlier [6].

In Fig. 3 we show the full width at half maximum of the FW intensity $\phi_1^*\frac{2}{\pi}$ of dark quadratic solitons versus the mismatch parameter $\alpha$. The dark solitons have the constant background $\phi_1^*\frac{2}{\pi}(\pm \infty) = 2\alpha$, $\phi_2^*\frac{2}{\pi}(\pm \infty) = 1$. The analytical weakly nonlocal dark soliton solution exists for $\alpha < 2$ and was taken from Ref. [14]. Unlike bright solitons, whose width is a monotonic function of $\alpha$, dark solitons are seen to have a minimum width at $\alpha = \alpha_0 \approx 3.1$.

Figure 3 confirms that for $\alpha > \alpha_0$ the weakly nonlocal model correctly predicts how the soliton width decreases when the mismatch parameter decreases. This, as well as the appearance of the minimum in the soliton width, is again elegantly explained by the nonlocal analogy: Because of the convolution in the nonlinearity in Eq. (12) representing the trapping potential or waveguide structure, the contribution from the constant background tends to contract this potential. This leads to a stronger confinement and thus a smaller width of the soliton. However, this is only true as long as the amplitude of the trapping potential is not affected by nonlocality, as in the weakly nonlocal regime. For a high degree of nonlocality (i.e., smaller value of $\alpha$) not only the width of the trapping potential, but also its amplitude is affected. In this regime the nonlocality leads to a drop in the amplitude of the potential, resulting in a weaker confinement and an increase of the soliton width. The profiles shown in Fig. 3 further illustrate the excellent agreement of the numerical results and approximate nonlocal analytical solutions in the weakly nonlocal limit $\alpha > 1$.

IV. CONCLUSION

In conclusion, we have used the analogy between parametric interaction in quadratic media and nonlocal Kerr-type nonlinearities to provide a physically intuitive theory for
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The strong nonlocal limit has enabled us to find bound states of bright quadratic solitons and explain their formation using the natural concept of the nonlocality-based attraction between out-of-phase constituent solitons.

Finally we have discussed how a simple phase-sensitive nonlocal model provides a better description of the dynamics as compared to the standard NLS equation obtained in the cascading limit.

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