One-dimensional steady-state spatial solitons in photovoltaic photorefractive materials with an external applied field

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A theory on spatial solitons in closed-circuit photorefractive-photovoltaic media is developed that gives rise to spatial solitons due to both the bulk photovoltaic effect and the spatially nonuniform screening of the external bias field. In case the photovoltaic effect can be neglected, these solitons are screening solitons. In case the external bias field is absent, these solitons are photovoltaic solitons in the closed-circuit case. We also show theoretically that the photovoltaic nonlinearity can be switched from self-defocusing to self-focusing by adding the external electric field. Under a strong bias condition, the photovoltaic nonlinearity can be switched from self-defocusing to self-focusing by changing the polarity of the external electric field or by rotating the polarization of the light.

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Photorefractive (PR) spatial optical solitons have attracted much interest in the past few years [1–15]. Thus far, quasi-steady-state solitons [1–3], screening solitons [4–8], and photovoltaic (PV) solitons [9–15] in dielectric PR crystals have been suggested in the literature. Recent work on PV solitons has shown theoretically and experimentally that adding background illumination for these solitons enables a uniform electric field across the crystal that is screened nonuniformly [12,13]. When background illumination is the same polarization as the focused beam, the PV nonlinearity in a short circuit can be controlled by the intensity of the uniform background illumination [12]. When the focused and the background beams are orthogonally polarized, the PV nonlinearity in a open circuit can be switched from self-defocusing to self-focusing by use of background illumination [13].

In this paper we show theoretically that the application of an external field to photovoltaic PR crystals enables the PV nonlinearity that gives rise to steady-state planar solitons due to both the PV effect and the spatially nonuniform screening of the external bias field. Screening solitons and PV solitons in the closed-circuit case are a special case of these solitons, respectively. Thus, screening solitons and PV solitons in the closed-circuit case studies may be integrated in these solitons studies. There are differences between using an external field and adding background illumination in PR crystals. The application of an external field enables screening solitons in nonphotovoltaic PR crystals, which result from the spatially nonuniform screening of the external bias field [7,8]. Adding background illumination for nonphotovoltaic PR crystals merely increases the background density of free carriers. If the electrical circuit consists of a PV crystal, an external source ε, and an external resistor R, an external field may screen the PV effect, which results in a change of the sign of the space-charge field. The sign of the refractive index perturbation Δn may change because $\Delta n \propto E$. On the other hand, for PV solitons, adding a background beam is equivalent to adding a term in the PV current, which results in establishing a uniform electric field across the crystal. The space-charge field is shifted upward or downward by the uniform electric field. However, only the shape and not the sign of Δn changes [13].

We start with the standard set of rate and continuity equations and Gauss's law, which describe the photorefractive effect in a medium in which electrons are the only charge carriers, plus the wave equation for the slowly varying amplitude of the optical field. In steady state and in two dimensions these equations are [9,11]

$$(s|A|^2+\beta_T)(N_d-N_d^i)-\gamma\hat{n}N_d^i=0, \tag{1}$$

$$\nabla\cdot\hat{\boldsymbol{J}}=\nabla\cdot[q\mu\hat{n}\hat{\boldsymbol{E}}+k_BT\mu\nabla\hat{n}+\kappa_{\text{eff}}(N_d-N_d^i)|A|^2\boldsymbol{s}]=0, \tag{2}$$

$$\nabla\cdot\hat{\boldsymbol{E}}+(q/\varepsilon_s)(\hat{n}+N_A-N_d^i)=0, \tag{3}$$

$$\left(\frac{\partial}{\partial z}-\frac{i}{2k}\frac{\partial^2}{\partial x^2}\right)A(x,z)=\frac{ik}{n_b}\Delta n(\hat{E})A(x,z), \tag{4}$$

$$V=-\int_{-l/2}^{l/2}dx\,\hat{E}=RS\hat{J}+\mathcal{E}, \tag{5}$$

where $\Delta n(\hat{E})=-\frac{1}{2}n_b^3r_{\text{eff}}\hat{E}$ is the perturbation in the refractive index. The independent variables are as follows: z is the propagation axis and x is the transverse coordinate. The dependent variables are as follows: $\hat{n}$ is the electron number density, N_d^i is the number density of ionized donors, $\hat{\boldsymbol{J}}$ is the current density, $\hat{\boldsymbol{E}}$ is the space-charge field inside the crystal, and A is the slowly varying amplitude of the optical field defined by $E_{\text{opt}}(x,z,t)=A(x,z)\exp(ikz-i\omega t)+\text{c.c.}$ ($k=2\pi n_b/\lambda$, where λ is the wavelength *in vacuo*, ω is the frequency, and n_b is the unperturbed refractive index). Relevant parameters of the crystal are as follows: N_d is the total donor number density, N_A is the number density of negatively charged acceptors, β_T is the dark generation rate, $\boldsymbol{s}$ is the photoionization cross section, γ is the recombination

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rate, μ is the electron mobility, κ_{eff} is the effective photovoltaic constant, ε_s is the low-frequency dielectric constant, r_{eff} is the effective electro-optic coefficient, and V is the potential measured between the crystal's electrodes separated by distance l; $-q$ is the electron charge, k_B is Boltzmann's constant, and T is the absolute temperature. S is the surface area of the electrodes, R is the external resistance, and ε is the external source potential applied to the crystal. Finally, we define the optical intensity as $I=|A|^2$, and the dark irradiance as $I_{\mathrm{dark}}=\beta_T/s$.

We seek solutions of the form

$$A(x,z)=u(x)\exp(i\Gamma z)\sqrt{I_{\mathrm{dark}}}, \tag{6}$$

where Γ is the soliton propagation constant. We limit our analysis to real $u(x)$. Since $I=|A|^2$ depends on x alone, we look for solutions in which the dependent variables $\hat{n}$, N_d^i, $\hat{\boldsymbol{J}}$, and $\hat{\boldsymbol{E}}$ depend solely on x, and the only component of $\hat{\boldsymbol{J}}$ and $\hat{\boldsymbol{E}}$ is in the x direction.

We transform the equations to dimensionless form by the substitutions $n=\hat{n}/(aN_d)$, $r=N_d/N_A$, $N=N_d^i/N_d$, $E=|\hat{\boldsymbol{E}}|qL_D/(k_BT)$, $J=|\hat{\boldsymbol{J}}|L_D/(a\mu N_d k_B T)$, and $\xi=x/L_s$, where $L_D=[k_BT\varepsilon_s/(q^2N_A)]^{1/2}$ is the Debye length, L_s is the soliton length scale defined by $L_s=1/(\pm 2kb)^{1/2}$, where $b=(k/n_b)\frac{1}{2}n_b^3 r_{\mathrm{eff}}k_BT/(qL_D)$ is the parameter that characterizes the strength and the sign of the optical nonlinearity. The sign of r_{eff} determines the sign (positive or negative) of the perturbation in the refractive index Δn [7]. We therefore introduce the dual-sign ($\pm$) notation in the definition of L_s, where the upper (lower) sign applies to the positive (negative) value of L_s (and, consequently, of Δn). The dimensionless equations are

$$n-(1+u^2)(1-N)/(rN)=0, \tag{7}$$

$$J=nE+\alpha(1-N)u^2+\hat{\varepsilon}n'=\mathrm{const}, \tag{8}$$

$$\left(N-\frac{1}{r}-an\right)-\hat{\varepsilon}E'=0, \tag{9}$$

$$u''=\pm\left(\frac{\Gamma}{b}+E\right)u, \tag{10}$$

$$\beta J+C=-\int_{-l/2L_s}^{l/2L_s}d\xi\, E. \tag{11}$$

The prime stands for the derivative with respect to the variable ξ, $a=sI_{\mathrm{dark}}/(\gamma N_A)$, $\hat{\varepsilon}=L_D/L_s$, $C=q\mathcal{E}L_D/(k_BTL_s)$, $\beta=RSq\mu aN_d/L_s$, and $\alpha=\kappa_{\mathrm{eff}}I_{\mathrm{dark}}L_D/(\mu k_B T a)=E_pL_Dq/(k_BT)$, where $E_p=\kappa_{\mathrm{eff}}\gamma N_A/(q\mu)$. Even though E can be obtained, in principle, from Eqs. (7)–(11), this task is considerably involved. However, we can conveniently derive E by taking a similar way to that of Ref. [11]. Using inequality the $N\leqslant 1/r\leqslant 1$ [11] in Eq. (7) leads to $n=(1+u^2)/rN$. Neglecting both terms an [11] and $\hat{\varepsilon}E'$ [7] in Eq. (9) leads to $N=1/r$. These approximations can be used in Eq. (8) and yield

$$J=nE+a\left(1-\frac{1}{r}\right)u^2\approx nE+\alpha u^2=(1+u^2)E+\alpha u^2=\mathrm{const}. \tag{12}$$

From Eq. (12) the space-charge field is given by

$$E=\frac{J-\alpha u^2}{1+u^2}. \tag{13}$$

The substitution of Eq. (13) into Eq. (11) leads to

$$J=-\chi\delta-\chi(1-\rho)\alpha, \tag{14}$$

where $\delta=C\eta$, $\rho=\eta l/L_s$, $\chi=1/(\eta\beta+1)$, and $\eta=1/\int_{-l/2L_s}^{l/2L_s}d\xi/(1+u^2)$. The substitution of Eq. (14) into Eq. (13) yields

$$E=-\frac{\chi\delta}{1+u^2}-\frac{\alpha[\chi(1-\rho)+u^2]}{1+u^2}. \tag{15}$$

The substitution of Eq. (15) into Eq. (10) then yields the following nonlinear wave equation that describes stationary (soliton) propagation in a photovoltaic nonlinear medium with an external applied field

$$u''=\pm\left\{\frac{\Gamma}{b}-\frac{\chi\delta}{1+u^2}-\frac{\alpha[\chi(1-\rho)+u^2]}{1+u^2}\right\}u. \tag{16}$$

We integrate Eq. (16) using quadrature, and obtain

$$p^2-p_0^2=\pm\left\{\left(\frac{\Gamma}{b}-\alpha\right)(u^2-u_0^2)+\chi[\alpha(\rho+\eta\beta)-\delta]\times\ln\left(\frac{1+u^2}{1+u_0^2}\right)\right\}, \tag{17}$$

where $p=u'$, $P_0=p(\xi=0)$, and $u_0=u(0)$.

For dark solitons one requires that the boundary conditions are $u(0)=0$, $u(+\infty)=u_\infty=u(-\infty)\neq 0$, and $u'(\infty)=u''(\infty)=0$. Substituting $\xi\to\infty$, and conditions $u_\infty\neq 0$ and $u''(\infty)=0$ into Eq. (16) leads to

$$\frac{\Gamma}{b}=\frac{\chi\delta}{1+u_\infty^2}+\frac{\alpha[\chi(1-\rho)+u_\infty^2]}{1+u_\infty^2}. \tag{18}$$

This implies that the propagation constant Γ depends on u_∞.

By substituting Eq. (18), and $p(\infty)=u'(\infty)=0$, $u(\infty)=u_\infty$, and $u(0)=0$ into Eq. (17), we obtain

$$p_0^2=\pm\chi[\alpha(\rho+\eta\beta)-\delta]\left[\frac{u_\infty^2}{1+u_\infty^2}+\ln(1+u_\infty^2)\right]. \tag{19}$$

Substituting Eqs. (18) and (19) into Eq. (17) yields

$$p^2=\pm\chi[\delta-\alpha(\rho+\eta\beta)]\left[\frac{u^2-u_\infty^2}{1+u_\infty^2}+\ln\left(\frac{1+u_\infty^2}{1+u^2}\right)\right]. \tag{20}$$

It is proved that the term inside the latter square brackets of Eq. (20) is always positive. The reality of p can be obtained for the upper sign or the lower sign according to a given

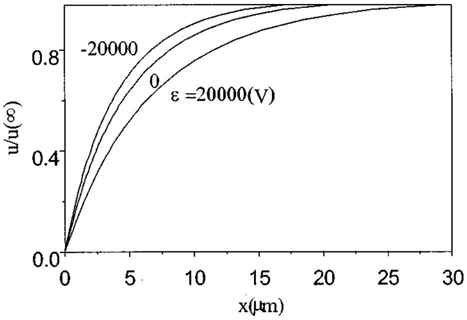

FIG. 1. Dark soliton amplitude u/u_∞ as a function of x for $u_\infty=5$ and $\varepsilon=0$, $-20\,000$, and $20\,000$ V.

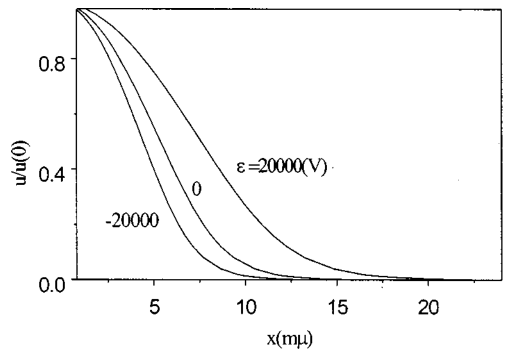

FIG. 3. Bright soliton amplitude $u/u(0)$ as a function of x for $u(0)=5$ and $\varepsilon=0$, $-20\,000$, and $20\,000$ V.

physical system $[\alpha(\rho+\eta\beta)$ and $\delta]$. When $\delta>0$ and $\delta>|\alpha(\rho+\eta\beta)|$ or $\alpha<0$ and $|\alpha(\rho+\eta\beta)|>|\delta|$, the upper sign gives dark solitons. Similar to screening solitons [7], dark solitons in biased photovoltaic PR materials require that $\Delta n(\xi)>0$ for all ξ. When $\delta<0$ and $|\delta|>|\alpha(\rho+\eta\beta)|$ or $\alpha>0$ and $|\delta|<\alpha(\rho+\eta\beta)$, the lower sign gives dark solitons. Similar to photovoitaic solitons [11], dark solitons in biased photovoltaic PR materials require that $\Delta n(\xi)<0$ for all ξ. We integrate Eq. (20) numerically for various values of u_∞, and obtain the waveforms $u(\xi)$ of dark solitons in biased photovoltaic PR materials. To illustrate our results, we consider the following examples: let $\lambda=0.5\,\mu$m, $l=1$ cm, $R=0$, and $\varepsilon=0$, $-20\,000$, and $20\,000$ V. The $LiNbO_3$ parameters are taken here to $n_b=2.2$, $r_{\rm eff}=30\times10^{-12}$ m/V, and $E_p=40$ KV cm. A particular case of $u_\infty=5$ is shown in Fig. 1 when $\varepsilon=0$, $-20\,000$, and $20\,000$ V. Figure 2 shows the half-width as a function of u_∞ when $\varepsilon=0$, $-20\,000$, and $20\,000$ V.

In the limit of low intensity ratio, i.e., $u_\infty^2\ll1$, Eq. (16) becomes

$$u''=\pm\left\{\frac{\Gamma}{b}-\chi[\delta+\alpha(1-\rho)]+\chi[\delta-\alpha(\rho+\eta\beta)u^2]\right\}u. \tag{21}$$

The solutions of Eq. (21), $u(\xi)$ with the appropriate boundary conditions, are bright and dark solitons.

For dark solitons, using boundary conditions $u(\infty)=u_\infty\neq0$ and $u''(\infty)=0$, and substituting $\xi\to\infty$ into Eq. (21) leads to $\Gamma/b=\chi[\delta+\alpha(1-\rho)]-\chi[\delta-\alpha(\rho+\eta\beta)]u_\infty^2$. Integrating Eq. (21) by quadrature and substituting conditions $p(\infty)=0$, $u(0)=0$, the expression we have just found for Γ/b, and $\xi\to\infty$ yields $p_0=\sqrt{\pm\chi[\delta-\alpha(\rho+\eta\beta)]/2}u_\infty^2$. The reality of p_0 can be obtained for the upper sign or the lower sign according to a given physical system $[\alpha(\rho+\eta\beta)$ and $\delta]$. This is identical to Eq. (20). Therefore in the limit of low intensity ratio, dark solitons in biased photovoltaic PR materials are solutions of

$$u''=\pm\chi[\delta-\alpha(\rho+\eta\beta)](u^2-u_\infty^2)u, \tag{22}$$

with $u(0)=0$ and $u'(0)=p_0=\sqrt{\pm\chi[\delta-\alpha(\rho+\eta\beta)]/2}u_\infty^2$. Integration of Eq. (22) leads to $u=u_\infty\tanh\{\sqrt{\pm\chi[\delta-\alpha(\rho+\eta\beta)]/2}u_\infty\xi\}$.

For bright solitons one requires the boundary conditions are $u(+\infty)=u(-\infty)=0$, $p(\infty)=u''(\infty)=0$, and $p(0)=0$. Substituting $\xi\to\infty$, and using conditions $u(\infty)=p(\infty)=0$ and $p(0)=0$ into Eq. (17) yields

$$\frac{\Gamma}{b}=\alpha-\frac{\chi[\alpha(\rho+\eta\beta)-\delta]}{u_0^2}\ln(1+u_0^2). \tag{23}$$

This implies that the propagation constant Γ depends on u_0.

Substituting Eq. (23) and using condition $p(0)=0$ into Eq. (17) yields

$$p^2=\pm\left\{\chi[\alpha(\rho+\eta\beta)-\delta]\left[\ln(1+u^2)-\frac{u^2}{u_0^2}\ln(1+u_0^2)\right]\right\}. \tag{24}$$

It is proved that the term inside the latter square bracket of Eq. (24) is always positive. The reality of p can be obtained for the upper sign or the lower sign according to a given physical system $[\alpha(\rho+\eta\beta)$ and $\delta]$. When $\alpha>0$ and $\alpha(\rho+\eta\beta)>|\delta|$ or $\delta<0$ and $|\delta|>|\alpha(\rho+\eta\beta)|$, the upper sign gives bright solitons. Similar to photovoltaic solitons [11], bright solitons in biased photovoltaic PR materials require that $\Delta n(\xi)>0$ for all ξ. When $\alpha<0$ and $|\alpha(\rho+\eta\beta)|>|\delta|$ or $\delta>0$ and $\delta>|\alpha(\rho+\eta\beta)|$, the lower sign gives bright solitons. Similar to screening solitons [7], bright solitons in biased photovoltaic PR materials require that $\Delta n(\xi)<0$ for all ξ. We integrate Eq. (24) numerically for various values of u_0, and obtain the waveforms $u(\xi)$ of bright solitons in biased photovoltaic PR materials. A particular case of $u_0=5$ is shown in Fig. 3 when $\lambda=0.5\,\mu$m, $l=1$ cm, $R=0$, and ε

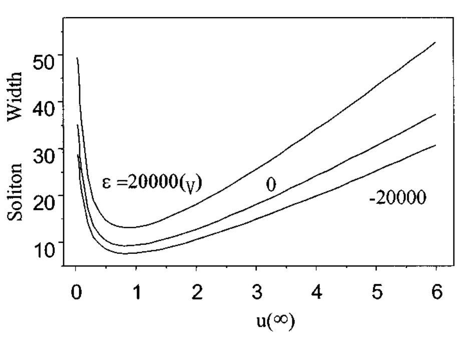

FIG. 2. Dimensionless half-width of the dark soliton as a function of u_∞ for $\varepsilon=0$, $-20\,000$, and $20\,000$ V.

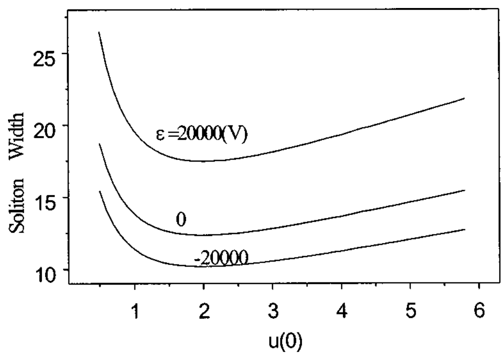

FIG. 4. Dimensionless half-width of the bright soliton as a function of u_0 for $\varepsilon=0$, $-20\,000$, and $20\,000$ V.

$=0$, $-20\,000$, and $20\,000$ V. Figure 4 shows the half-width as a function of u_0 when $R=0$, $\lambda=0.5\,\mu$m, $l=1$ cm, and $\varepsilon=0$, $-20\,000$, and $20\,000$ V. The $LiNbO_3$ parameters are taken to be the same as those considered in the previous examples.

In the limit of low intensity ratio, i.e., $u_0^2 \ll 1$, for bright solitons, integrating Eq. (21) by quadrature and substituting boundary conditions $u(\infty)=p(\infty)=0$ and $p(0)=0$ yields $\Gamma/b=\delta-\frac{1}{2}\chi[\delta-\alpha(\rho+\eta\beta)]u_0^2$. Substituting the expression we have just found for Γ/b into Eq. (21) yields

$$u''=\pm\chi[\alpha(\rho+\eta\beta)-\delta]\left(\frac{u_0^2}{2}-u^2\right)u, \tag{25}$$

with $u(0)=u_0$ and $u'(0)=0$. The integration of Eq. (25) leads to $u=u_0\,\mathrm{sech}\{\sqrt{\pm\chi[\alpha(\rho+\eta\beta)-\delta]/2}u_0\xi\}$.

Next, we discuss the relation between spatial solitons in biased photovoltaic PR crystals and both screening solitons and photovoltaic solitons in the closed-circuit case. When $\kappa_{\mathrm{eff}}=0$, i.e., $\alpha=0$, and $R=0$, i.e., $\chi=1$, our physical system becomes the physical system of previously studied screening solitons. In this case, substituting $\delta=C\eta$ into Eq. (15) leads to

$$E=-\frac{C\eta}{1+u^2}, \tag{26}$$

substituting Eq. (18) and $\delta=C\eta$ into Eq. (20) leads to

$$p^2=\pm\left[\frac{\Gamma}{b}(u^2-u_\infty^2)-C\eta\ln\left(\frac{1+u^2}{1+u_\infty^2}\right)\right], \tag{27}$$

and substituting Eq. (23) and $\delta=C\eta$ into Eq. (24) leads to

$$p^2=\pm\left[\frac{\Gamma}{b}u^2-C\eta\ln(1+u^2)\right]. \tag{28}$$

Equations (27) and (28) have the same dimensionless parameters and form as Eqs. (16) and (18) in Ref. [7], respectively, in which screening solitons are discussed for biased nonphotovoltaic-photorefractive crystals when $R=0$. On the other hand, in the absence of the external source, i.e., $\delta=0$, our physical system becomes the physical system of previously studied PV solitons in the closed-circuit case. In this case, Eqs. (14) and (15), and in the limit of the reality of p, Eqs. (20) and (24) become

$$J=-\chi(1-\rho)\alpha, \tag{29}$$

$$E=-\frac{\alpha[\chi(1-\rho)+u^2]}{1+u^2}, \tag{30}$$

$$p^2=[\chi\alpha(\rho-1)+\alpha]\left[\frac{u^2-u_\infty^2}{1+u_\infty^2}+\ln\left(\frac{1+u_\infty^2}{1+u^2}\right)\right], \tag{31}$$

$$p^2=[\chi\alpha(\rho-1)+\alpha]\left[\ln(1+u^2)-\frac{u^2}{u_0^2}\ln(1+u_0^2)\right], \tag{32}$$

respectively. By the substitution $J_p=J/\alpha$, we transform Eq. (30) to the following form:

$$E_p=-\frac{J_p-u^2}{1+u^2}, \tag{33}$$

where $E_p=E/\alpha$. By the substitution $J_p=J/\alpha$ and $\zeta=\sqrt{\alpha}\xi=x/d$, where $d=(\pm 2\,\mathrm{kg})^{-1/2}$ is the characteristic length scale $[g=(k/n_b)(1/2)n_b^3 r_{\mathrm{eff}}E_p]$, we transform Eqs. (31) and (32) to the following form:

$$\hat{p}^2=(J_p+1)\left[\frac{u^2-u_\infty^2}{1+u_\infty^2}+\ln\left(\frac{1+u_\infty^2}{1+u^2}\right)\right], \tag{34}$$

$$\hat{p}^2=(J_p+1)\left[\ln(1+u^2)-\frac{u^2}{u_0^2}\ln(1+u_0^2)\right], \tag{35}$$

where $\hat{p}=du/d\zeta$. Equations (33)–(35) have the same dimensionless parameters and form as Eqs. (13), (26), and (29) for a closed circuit in Ref. [11], respectively, in which photovoltaic solitons are discussed in photovoltaic-photorefractive crystals without an external bias field for both open and closed circuits.

Finally, we show the refractive index perturbation in biased photovoltaic PR crystals. Substitution of Eq. (6) into $I=|A|^2$ leads to $u^2=I/I_{\mathrm{dark}}$. Substituting Eq. (15), and using $u^2=I/I_{\mathrm{dark}}$ into $\Delta n(\hat{E})=-\frac{1}{2}n_b^3 r_{\mathrm{eff}}\hat{E}$ yields

$$\Delta n=\Delta n_0+\Delta n_0\chi(\rho+\eta\beta)\left[\frac{\delta}{\alpha(\rho+\eta\beta)}-1\right]\frac{I_{\mathrm{dark}}}{I+I_{\mathrm{dark}}}, \tag{36}$$

where $\Delta n_0=\frac{1}{2}n_b^3 r_{\mathrm{eff}}E_p$. The sign of Δn_0 cannot be changed in a given material [13]. On the other hand, δ may be switched from a positive value to a negative value by changing the polarity of the external electric field [8,9]. Some PV materials (for example, $BaTiO_3$) possess PV constants κ_{eff} (i.e., α) that change signs under polarization rotation [9]. When $\Delta n_0<0$, $\delta>0$, and $|\alpha(\rho+\eta\beta)|<\delta<|\alpha(\rho+2\eta\beta+1)|$, the photovoltaic nonlinearity that gives rise to spatial solitons can be switched from self-defocusing to self-focusing by adding the external electric field. Notice that only the shape and not the sign of Δn changes. When $\delta = -\alpha(\rho+\eta\beta)$, changing the polarity of the external electric field or rotating the polarization of the light enables $\Delta n = \Delta n_0$. When $\delta > \alpha(\rho+2\eta\beta+1)$, the photovoltaic nonlinearity can be switched from self-defocusing to self-focusing by changing the polarity of the external electric field or by rotating the polarization of the light.

In conclusion, a theory on spatial solitons in closed-circuit photorefractive-photovoltaic media has been developed that gives rise to spatial solitons due to both the bulk photovoltaic effect and the spatially nonuniform screening of the external bias field. Screening solitons and PV solitons in the closed-circuit case can be obtained from these solitons. Subsequently, we have shown that under the appropriate condition $[|\alpha(\rho+\eta\beta)| < \delta < |\alpha(\rho+2\eta\beta+1)|]$, the photovoltaic nonlinearity can be switched from self-defocusing to self-focusing by adding the external electric field, and under strong bias condition $[\delta > \alpha(\rho+2\eta\beta+1)]$, the photovoltaic nonlinearity can be switched from self-defocusing to self-focusing by changing the polarity of the external electric field or by rotating the polarization of the light.

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