Roles of resonance and dark irradiance for infrared photorefractive self-focusing and solitons in bipolar InP:Fe

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This paper shows experimental evidence of photorefractive steady state self-focusing in InP:Fe for a wide range of intensities, at both 1.06 and 1.55 μm. To explain those results, it is shown that despite the bipolar nature of InP:Fe where one photocarrier and one thermal carrier are to be considered, the long standing one photocarrier model for photorefractive solitons can be usefully applied. The relationship between the dark irradiance stemming out of this model and the known resonance intensity is then discussed.

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I. INTRODUCTION

Photorefractive self-focusing and spatial solitons have been up to now the object of more than a decade of growing interest, both experimentally and theoretically [1–4]. During this decade, solitary waves induced by the photorefractive effect have been shown to interact with each other showing specific behaviors upon collision [5] such as birth and fusion [6]. When a spatial solitary wave propagates in a photorefractive medium, it induces an optical waveguide that does not disappear instantly when the beam is turned off [7,8], leading to optically induced optical components such as junctions [9], couplers [10], and even optically induced soliton arrays and photonic lattices [11–13].

The attraction towards photorefractive solitons and the waveguides they induce owes a lot to their potentialities in optically inducing photonic circuitry in the bulk of a material, on the one hand, and to their reconfiguration properties, on the other hand. The former relies on the material property of retaining the waveguide once the beam is shut off, such as, for instance, LiNbO₃ [14]. The latter, on the contrary, relies on the possibility to rapidly screen out an applied field with a laser beam, to give rise to a waveguide, and to erase this waveguide as rapidly. Up to now and to our knowledge, the characteristic build-up and erase times are on the order of the second in the widely used SrₓBaₙ₋ₓNb₂O₆ [15] or sillenite Bi₁₂TiO₂₀ [16], or up to the minute in LiNbO₃ [17], all of these experiments having been conducted at visible wavelengths to which most of the cited materials are sensitive. Shorter response times around the millisecond have been obtained in semiconductor CdₓZn₁₋ₓTe [18] and around the nanosecond for visible ultrahigh intensities in Bi₁₂TiO₂₀ [19].

If one wants to target optical networking as a potential application of the dynamically reconfigurable photorefractive soliton technology, build-up and reconfiguration times below the millisecond are to be achieved ideally at infrared wavelengths compatible with optical fiber networks. Thanks to their high carrier mobility and despite their low electro-optic coefficient, photorefractive semiconductors such as InP:Fe might respond to this demand. Charge transport, in this material, has a bipolar nature with photoexcited holes coexisting with thermally excited electrons. This yields a peculiar resonance of the photorefractive two wave mixing (TWM) gain as a function of intensity [20–23].

Owing to its bipolar transport character, the physics behind spatial soliton propagation in InP:Fe is more complicated than in other materials. It has already been the subject of theoretical analysis [24]. In this work, we present a series of recent investigations on photorefractive self-focusing in InP:Fe at infrared wavelengths and we propose a model that takes into account the transport of both charge carriers. Unlike in previous work [24–26], the experimental results in our crystal samples and the corresponding theoretical prediction do not show any clear evidence of an influence of the intensity resonance of the bipolar photorefractive effect, thus indicating that photorefractive self-focusing behaves to a large extent as in the case of the well-known single-carrier model [27]. A hint is finally provided that the reason why qualitatively different behaviors are observed in different samples might be a varying oxidoreduction state.

II. EXPERIMENTAL EVIDENCE OF THE ROLE PLAYED BY DARK IRRADIANCE

The bipolar photorefractive resonance mechanism was evidenced experimentally by Mainguet and Idrissi et al. [20,21] and then defined by Picoli et al. [22,23] as the mean intensity at which the TWM gain is greatly enhanced. It corresponds to the intensity at which the intensity dependent

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hole generation rate is equal to the temperature dependent electron generation rate. The resonance intensity is thus expected to depend on the temperature and has been evidenced so. It has also been shown that this TWM enhancement corresponds to a $\pi/2$ phase shift between the index grating and the optical intensity grating, and to a corresponding enhancement of the index modulation amplitude.

As a preliminary characterization of our samples, we evidenced and measured the resonance intensity in the samples used in our experiments. The iron content of the various samples we tested varied from $8 \times 10^{16}$ to $2 \times 10^{17}$ cm$^{-3}$, as measured by secondary ion mass spectroscopy. Similar behavior was observed in all samples. The results presented in this paper were obtained with a $2 \times 10^{17}$ cm$^{-3}$ Fe doped sample.

The resonance intensity has been measured with a standard TWM setup [20] as described in Ref. [28], for various temperatures, and particularly at 20 °C, the temperature at which further self-focusing experimental results have been obtained. The measured resonance intensity at 20 °C is 300 mW/cm$^2$ at a wavelength of 1.06 \(\mu\)m and 3 W/cm$^2$ at a wavelength of 1.56 \(\mu\)m. This wavelength disparity is explained by the wavelength dependent photoexcitation cross sections.

The experiments were conducted independently at both 1.06 and 1.56 \(\mu\)m wavelengths, the latter because it corresponds to a telecom wavelength and the former for InP:Fe large sensitivity to it. The 1.06 \(\mu\)m wavelength is delivered by a solid state diode pumped Nd:YAG laser whereas the 1.56 \(\mu\)m beam is produced by a laser diode.

The beam is focused down to a 30 \(\mu\)m waist onto the entrance face of a temperature regulated InP:Fe sample. The beam propagates along the (110) direction and is linearly polarized along (110) with an electric field of 10 kV/cm applied along (001). The crystal dimensions are $5 \times 5 \times 12$ mm$^3$ along the directions (001), (110), and (110), respectively.

The steady state beam profile on the sample output face is observed through a microscope objective and a silicon CCD camera. The camera is sensitive enough for the 1.06 \(\mu\)m wavelength but requires to be equipped with a luminescent phosphorus converter to achieve 1.56 \(\mu\)m observations. The observed image is then gamma corrected to account for the conversion nonlinearity.

The output beam profile is observed at steady state for various intensities, at a temperature of 20 °C. The best result giving the strongest focusing is shown in Fig. 1(b) for both 1.06 and 1.56 \(\mu\)m wavelengths. Varying the input intensity reveals that the steady-state self-focusing phenomenon vanishes at low and high intensities and reaches an optimum focusing for an intermediate intensity, when the output beam width reaches a minimum.

This is quantified in Fig. 2 where the normalized horizontal output beam diameter is plotted as a function of the input beam peak intensity for an applied field of 10 kV/cm. Linear diffraction is observed at steady state for small and large intensities. A minimum output beam diameter is reached for an intermediate intensity of approximately 210 mW/cm$^2$ at 1.06 \(\mu\)m and 7.5 W/cm$^2$ at 1.56 \(\mu\)m. A noticeable characteristic of both curves is that the upside down bell-shape is attracted by the dotted line when it approaches it, i.e., when the output beam diameter approaches the input one. This evidences a domain of stability of the beam width (from 1 to 4 W/cm$^2$ at 1.06 \(\mu\)m and from 5 to 50 W/cm$^2$ at 1.56 \(\mu\)m). The horizontal beam width associated to these domains roughly corresponds to the input beam width. This may suggest that the beam evolution is attracted towards a solitary wave solution.

We note that the width stability intensity range is much wider at 1.56 \(\mu\)m than at 1.06 \(\mu\)m: this is due to the fact that in the former case, the intensity that fulfills soliton propagation conditions is much nearer to the minimum of the reverse bell-shaped curve. Furthermore, one has to keep in mind that the beam intensity decreases with propagation, owing to absorption—2 cm$^{-1}$ at 1.06 \(\mu\)m and 0.5 cm$^{-1}$ at

![FIG. 1. (Color online) Steady state output beam profile for a peak intensity of 190 mW/cm$^2$ at 1.06 \(\mu\)m (top) and a peak intensity of 8.7 W/cm$^2$ at 1.56 \(\mu\)m (bottom), for no applied field (a) and a 10 kV/cm horizontally applied field (b). The images total width is 200 \(\mu\)m. The non-Gaussian shape of the 1.56 \(\mu\)m beam in absence of the applied field is due to the multimode character of the laser itself. The curves (c) are the horizontal beam profiles (a.u.) at beam maximum.](image-url)
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FIG. 2. (Color online) Output beam width as a function of input beam peak intensity at 1.06 μm (top) and 1.56 μm (bottom) wavelengths. The beam diameter is normalized to the linearly diffracted output beam diameter. The horizontal dotted line represents an output beam as wide as the input one, suggesting a spatial solitonlike propagation.

1.56 μm—so that the soliton propagation conditions cannot be fulfilled precisely all the way through the crystal, thus strengthening the need for stability if a spatial soliton is to propagate.

As in many other materials, self-focusing in InP:Fe comes together with beam bending, as is evidenced by Fig. 1 and particularly Fig. 1(e). The steady-state beam shift at the output face with respect to the propagation axis, along the same horizontal direction as the applied field, is reported in Fig. 3. We observe that beam bending is negligible for sufficiently low and high intensities, that is, when self-focusing is absent. Furthermore, the maximum beam shift is obtained precisely at the intensity at which self-focusing is most efficient.

The curves reported in Fig. 2 resemble the photorefractive soliton existence curve—giving the soliton width as a function of the beam peak intensity to dark irradiance ratio—as predicted earlier [27,29] in the framework of a single-photocarrier model. The dark irradiance plays a central role in this theory since it is the ratio of the beam peak intensity to the dark irradiance which determines the steady state soliton width. The dark irradiance is the equivalent uniform optical irradiance which should be shone on the crystal to generate as many photocarriers as thermal excitation does. In other words, shining the dark irradiance on a given sample doubles its conductivity.

If we assume that this one-photocarrier model—which will be named model I in the following—describes our experiments in InP:Fe, we can use its predictions to deduce the dark irradiance in InP:Fe at both the wavelengths we have used. According to it, the minimum width soliton which corresponds to the maximum self-focusing effect is obtained for a beam peak intensity equal to three times the dark irradiance, as confirmed recently [30] in a (1+1)D model. Using this hypothesis, the dark irradiance in our sample can then be evaluated, from Fig. 2, to approximately 70 mW/cm² for the 1.06 μm wavelength and 2.5 W/cm² for 1.56 μm, at a temperature of 20 °C; the accuracy of the value measured at 1.06 μm is low because of InP:Fe strong absorption at this wavelength. Let us note here that, in both cases, the estimated dark irradiance and the resonance intensity are on the same order of magnitude.

At this stage, one may wonder why model I should describe photorefractivity in InP:Fe since both holes and electrons are known to play a significant role in this semiconductor—holes are photoexcited and electrons are thermally generated—whereas model I assumes only one carrier, both photoexcited and thermally excited. One may hastily think that electrons can simply be neglected in our experiments. This is, however, not the case since both the dark irradiance and the resonance intensity are included in our measurement range: they, respectively, correspond to the equality of the concentrations on the one hand and of the generation rates on the other hand, of holes and electrons.

Therefore, since model I assumptions do not match with our experiments, we now need to ascertain by an independent measurement that the values of the dark irradiance we have deduced from Figs. 2 and 3 are indeed correct. We have thus shone a uniform intensity on our InP:Fe sample and measured its photoconductivity. This allowed us to deduce a
50 mW/cm² dark irradiance at 1.06 µm and a 2 W/cm² one at 1.56 µm, confirming the previous evaluation at this wavelength. The slight discrepancy is accounted for by the relative precision with which Fig. 2 allows the determination of the dark irradiance on the one hand and on the other hand by the difficulty to shine a strong and uniform light on the whole sample for the independent measurement.

These recent experimental results evidence the fact that the dark irradiance plays a crucial role in photorefractive self-focusing in InP:Fe, as it does in previously developed models, despite the greater complexity of the photorefractive mechanisms in InP:Fe. However, as previously explained, model I does not allow us to explain our observations, nor do previously published ones [24–26], as will be detailed later on. Therefore, we will now tackle the theoretical problem of showing how and why the dark irradiance plays a significant role in InP:Fe.

III. SPACE CHARGE FIELD AND DARK IRRADIANCE

Let us start with the general two-carrier band transport model [20,24] reduced to one dimension $x$, which we will consider at steady state in the following. It is described by the set of equations

$$\frac{\partial E}{\partial x} = \frac{e}{\epsilon} (N_D - N_A + p - n - n_T),$$  

(1a)

$$j_n = e \mu_p n E + \frac{e}{\epsilon} \kappa_b T \frac{\partial n}{\partial x},$$  

(1b)

$$j_p = e \mu_n p E - \frac{e}{\epsilon} \kappa_b T \frac{\partial p}{\partial x},$$  

(1c)

$$\frac{\partial n}{\partial t} = e n_T - c_n p T + \frac{1}{e} \frac{\partial j_n}{\partial x},$$  

(1d)

$$\frac{\partial p}{\partial t} = e p_T - c_p n T + \frac{1}{e} \frac{\partial j_p}{\partial x},$$  

(1e)

$$\frac{\partial n_T}{\partial t} = e p_T - e n_T - c_p n T + c_n p T,$$  

(1f)

where $E$ is the electric field, $n$ and $p$ are the electron and hole densities in the respective conduction and valence bands, $n_T$=[Fe²⁺], is the density of ionized occupied traps, $p_T$=[Fe³⁺] is the density of neutral unoccupied traps, $j_n$ and $j_p$ are the electron and hole current densities, respectively. $N_T$, $N_D$, and $N_A$ are the densities of iron atoms, the shallow donors, and the shallow acceptors, respectively. The charge mobilities are given by $\mu_n$ for electrons and $\mu_p$ for holes, the electron and hole recombination rates are, respectively, $c_n$ and $c_p$, $T$ is the temperature, and $k_b$ is the Boltzmann constant. The dielectric permittivity is given by $\epsilon$ while $e$ is the elementary charge. The electron and hole generation rates $e_n$ and $e_p$ depend on both thermal and optical emission as described by

$$e_n = e_n^0 + \sigma_n I \approx e_n^0,$$  

(2a)

$$e_p = e_p^0 + \sigma_p I \approx \sigma_p I,$$  

(2b)

where the thermal contribution to the emission rate coefficient is $e_n^0$ and the optical cross section of the carriers is given by $\sigma$, $I$ is the spatially dependent intensity of light. The approximations in both above equations are valid in the case of InP:Fe as stated above, with thermally excited electrons and photoexcited holes. The need for a large applied electric field leads us to neglect carrier diffusion when compared to drift diffusion, which is accounted for by all the terms which contain $k_b T$.

Solving the differential system (1) requires defining the boundary conditions. To that aim, we will consider that the input beam is small with respect to the sample dimensions and thus that its influence is limited to its surroundings. Therefore, the boundary values for $E, n, p, n_T, p_T, j_n$, and $j_p$ are the values reached at steady state in the dark. As was previously done [20,22,23,28], they can be deduced from system (1), setting to zero all derivatives and optical intensity. Following, if the 0 subscript identifies the boundary condition $E_0 = V/d$, where $V$ is the applied voltage and $d$ the crystal width; $n_0 = e_n^0 n_{T0}/e_p T_0$ and $p_0 = e_p^0 p_{T0}/e_n T_0$, where $n_{T0} = N_D - N_A$ and $p_{T0} = N_T - n_{T0}$ are the ionized and nonionized iron densities in the dark [22,23]; $j_n = e \mu_n n_0 E_0$ and $j_p = e \mu_p p_0 E_0$.

The time evolution of the electric field $E$ is determined as follows. The time derivative of Eq. (1a) is simplified using Eqs. (1d)–(1f) to yield

$$\frac{\partial^2 E}{\partial t \partial x} = \frac{-1}{\epsilon} \frac{\partial (j_n + j_p)}{\partial x}.$$  

(3)

Making use of Eqs. (1b) and (1c) and integrating with respect to the space variable $x$, it becomes

$$\frac{\partial E}{\partial t} = \frac{-e}{\epsilon} \left[ \left( \mu_n n + \mu_p p \right) E - \left( \mu_n n_0 + \mu_p p_0 \right) E_0 \right].$$  

(4)

At steady state, the evolution equations (1d)–(1f) and (4) reduce to

$$\left( \mu_n n + \mu_p p \right) E - \left( \mu_n n_0 + \mu_p p_0 \right) E_0 = 0,$$  

(5a)

$$e_n n_T - c_n p T + \frac{e}{\epsilon} \frac{\partial E}{\partial x} - \mu_n E \frac{\partial n}{\partial x} = 0,$$  

(5b)

$$e_p (N_T - n_T) - c_p n T - \mu_p E \frac{\partial n_T}{\partial x} = 0.$$  

(5c)

To further simplify this differential system, the light intensity spatial variations will be assumed small enough so that the resulting free carriers spatial variations remain small, and the derivatives $\partial n/\partial x$ and $\partial p/\partial x$ can be neglected in Eqs. (5b) and (5c). In the same spirit, we also assume that the field derivative $\partial E/\partial x \ll e_n^0 N_T/\mu_n n_0$. 

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The left-hand side of Eq. (1a) becomes irrational within this approximation. Light intensity and temperature will again be considered not so large so that free carriers density remains small with respect to ionized iron density: $n \ll n_T$ and $p \ll n_T$, so that Eq. (1a) reduces to $n_T = n_{T_0}$, and consequently from Eqs. (5b) and (5c) we get $n = n_0$ and $p = (\sigma_f p_{T_0})/(\sigma_p p_{T_0})$. The assumptions made are valid for intensities small against $c_T n_{T_0}^2/\sigma_p p_{T_0}$, itself greater than the TW/cm$^2$, and at temperatures orders of magnitude above room temperature. Within these approximations, we obtain the electric field $E$ dependency upon the local intensity $I$:

$$E = \frac{I_d}{I_d + I_E_0}. \quad (6)$$

This expression is strictly the same as the one found before in the framework of model I [27,29] when diffusion is neglected; only the expression of the dark irradiance $I_d$ differs (the approximation below being valid in the case of InP:Fe):

$$I_d = \frac{\epsilon_0^2 n_{T_0}^2 c_\mu_p + \epsilon_0^2 n_{T_0}^2 c_\mu_n}{\sigma_\mu_p^2 n_{T_0}^2 c_\mu_p + \sigma_\mu_n^2 n_{T_0}^2 c_\mu_n} = \frac{\epsilon_0^2 n_{T_0}^2 c_\mu_p}{\sigma_\mu_p^2 n_{T_0}^2 c_\mu_p}. \quad (7)$$

Let us point out here that this has implications beyond the above equation: all the conclusions that have been driven, in model I, from an equation identical to Eq. (6), still hold, if the above value of $I_d$ is accounted for. There is, for instance, no need to derive again the soliton existence curve since the involved equations are the same. Using the predictions of model I, our observations of self-focusing in InP:Fe are partially explained. Indeed the presence of self-focusing over a broad range of intensity is predicted and no electron-hole resonance effect is observed. The beam bending is, however, not yet accounted for, which is one reason to further extend and refine our model.

**IV. RESONANCE INTENSITY**

As was said before, we felt the need to establish the theoretical analysis presented in the previous section—which we will now call model II—because we could not satisfactorily explain our observations using the theory available in the literature [24,26] for one particular reason: this theory—which we will call model III—predicts the existence of a resonance intensity around which the self-focusing phenomenon reverses itself; while we have successfully measured the resonance intensity in our samples, we were never able to measure the predicted reversal from focusing to defocusing.

The resonance intensity around which a self-focusing behavior reversal is predicted by model III is precisely the same resonance intensity at which TWM gain is maximum. As was said before, the resonance intensity is also the intensity at which the generation rates of holes and electrons are equal:

$$I_r = \frac{\epsilon_0^2 n_{T_0}^2 + \epsilon_0^2 n_{T_0}^2}{\sigma_\mu_p^2 n_{T_0}^2 + \sigma_\mu_n^2 n_{T_0}^2} = \frac{\epsilon_0^2 n_{T_0}^2}{\sigma_\mu_p^2 n_{T_0}^2}. \quad (8)$$

At resonance, the imaginary part of the space charge field reaches a maximum while the real part—the in phase component—reaches zero. Following, the real part changes sign around $I_r$ [22,26]. The link with self-focusing can then be thought of as follows: the in phase component of the space-charge field—the real part—is a local photorefractive response, responsible for self-focusing or defocusing, whereas the out of phase one—the imaginary part—is responsible for beam bending. This phenomenological interpretation has been confirmed by an experiment [26]. Later on, analytical derivations and numerical simulations [24] came to the same predictions. As stated before, however, we were never able to reproduce the experiment, for a reason we will now give hints of.

First of all, in the case of InP:Fe, there is a close and fairly simple relationship between $I_d$ and $I_r$:

$$\frac{I_d}{I_r} = \frac{\mu_p e^2 c_p}{\mu_p e_n n_{T_0}}. \quad (9)$$

We have evaluated the above ratio to 0.65 with values issued from the literature [31–33], a numerical value which is fully compatible with those we have measured at 1.56 μm, the discrepancy for the measures at 1.06 μm being accounted for by the strong absorption impairing the $I_d$ estimation from Figs. 2 and 3.

The above ratio (9) evidences the ratio $[\text{Fe}^{2+}]/[\text{Fe}^{3+}]$ in the dark ($n_{T_0}/p_{T_0}$). Therefore, the oxydoreduction state plays a significant role, as well as the doping of the samples, in the ratio of the two characteristic intensities. Therefore, a different doping but also a different $[\text{Fe}^{2+}]/[\text{Fe}^{3+}]$ ratio might be the cause of the discrepancy between our experimental results, which do not exhibit a crucial role for the resonance intensity, and those previously reported.

However, our theoretical analysis, model II, does not either bring the resonance intensity forth. A more universal model is in development in order to predict a doping dependent behavior and thus unify models II and III.

**V. CONCLUSION**

In light of recent experimental results, we have shown that despite more complex photorefractive charge transport mechanisms, steady-state photorefractive self-focusing and solitons in InP:Fe can be described by the same model as was done for the simpler one-free-carrier model, provided diffusion is neglected. We have evidenced in the process the role of the dark irradiance. Our model does not, however, evidence any crucial role for the resonance intensity, a role which is to be the subject of future studies.

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