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Abstract
We experimentally study a fiber loop laser in a figure-8 configuration and explore the dependency of the degree of polarization on controlled parameters. To account for the experimental observations, a mapping is derived to evaluate the polarization time evolution. The nonlinearity induced by the Kerr effect and gain saturation gives rise to rich dynamics. We find that the degree of polarization can be increased by tuning the system into a region where the mapping has a locally stable fixed point.

I. INTRODUCTION
Scrambling is commonly employed to lower the degree of polarization (DOP) of optical sources when effects such as polarization hole-burning and polarization mode dispersion are unwanted. On the other hand, a high DOP is desired for some applications, including the detection of magneto-optic and optomechanical effects.

In this study, we explore the DOP of a fiber loop laser in a figure-8 configuration. Mode locking (ML) in figure-8 lasers (F8Ls) has been studied in Refs. 7-11. The stability of a passively mode-locked F8L was studied in Ref. 12. Various types of vector solitary waves have been investigated in optical fibers and mode-locked fiber lasers.

Polarization evolution of solitary waves has been investigated in optical fibers and fiber lasers. Polarization evolution in optical fibers with spatially varying birefringence has been explored in Ref. 29. Polarization evolution of solitary waves in fiber lasers is complex as optical components, e.g., polarization controllers (PCs), laser gain media, optical couplers (OCs), wavelength filters (WLFs), and multiple cavity round trips, contribute to polarization evolution apart from birefringence and dispersion of optical fibers. Reference 24 has shown that the soliton polarization rotation is periodic where the period can be a multiple of the cavity round trip time in a fiber ring laser. It has been shown that soliton dynamics in a mode-locked fiber laser can produce either multiple fixed polarization states occurring periodically or no fixed polarization state at a fixed location of the laser cavity. This can lead to different DOPs depending on the soliton dynamics. F8L is promising as this fiber laser configuration can give rise to sub-picosecond laser pulses in the ML region. However, polarization instability is a challenging issue in mode-locked F8L not based on polarization-maintaining fibers.

In our experiment, we employ the rotating quarter-wave plate method to measure the DOP in both continuous wave (CW) and ML regions. We find that for both cases, the DOP can be tuned.

II. EXPERIMENTAL SETUP
A sketch of the experimental setup is shown in Fig. 1. The F8L consists of two fiber optical loops, which are connected through a 50:50 optical coupler (OC). PCs are integrated into both loops. The right loop serves as a fiber optical loop mirror (FOLM). An erbium doped fiber (EDF) of 3.6 m length is integrated into the FOLM to provide gain in the telecom band. The EDF is pumped using a 980 nm diode laser and a wavelength division multiplexer (WDM). The EDF is connected to a long nonlinear single mode fiber (NSMF) of 50 m length. Band selection is performed using a tunable wavelength filter (WLF) of 0.42 nm width and −3 dB insertion loss. A 99:1 OC in the right loop is connected to a photodetector (PD), which is probed using a radio frequency spectrum analyzer (RFSA).
For clockwise unidirectional propagation of light in the left loop, two optical isolators (OIs) are employed (labeled by arrows in Fig. 1). A 90:10 OC in the left loop is connected to an optical spectrum analyzer (OSA), a PD, an oscilloscope (OS), and a rotating quarter-wave plate polarimeter (PM) (Thorlabs PA430).

III. MEASUREMENTS

The temporal F8L output signal measured by the OS in the region of ML is shown in Fig. 2(a), Figure 2(b) shows the corresponding autocorrelation (AC) signal as a function of delay time. Figure 2(c) shows the AC signal calculated using Eq. (A4) of Appendix A. The calculated AC signal approximately imitates the coherence artifact (spike) and broad pedestal (wings) of the measured AC. The pulse width is extracted by fitting the data with Eq. (A4) (see caption of Fig. 2).

Plots of the optical spectrum measured using the OSA, the RF spectrum measured using the RFSA, the averaged optical power measured using the PD in the right loop, and the DOP measured using the PM are shown in Fig. 3 as a function of diode current $I_D$. For the measurements shown in Figs. 3(a1)–3(d1), both PCs are tuned to maximize the DOP in the ML regime, whereas the DOP in the ML regime is minimized for the measurements shown in Figs. 3(a2)–3(d2). For the first (second) case, the CW threshold occurs at a diode current of $I_D = 0.1$ A ($I_D = 0.15$ A), and the ML threshold occurs at $I_D = 0.26$ A ($I_D = 0.37$ A). As can be seen by comparing Figs. 3(d1) and 3(d2), the DOP in the ML regime can be tuned. The DOP plot in Fig. 3(d1) reveals a sharp change in the state of polarization (SOP) at the transition from the CW to the ML regime occurring at a diode current of $I_D = 0.26$ A. For this case, the DOP increases with diode current $I_D$ in both CW and ML regions. An SOP change occurring at the transition from CW to ML regions is also seen in Fig. 3(d2). However, for this case, the DOP remains very low above the transition.
IV. FOLM

To account for the experimental observations, scattering and birefringence in the FOLM are theoretically studied. The 50:50 OC is characterized by forward (backward) transmission $t$ ($t'$) and reflection $r$ ($r'$) amplitudes. Time reversal and mirror symmetries together with unitarity imply that $t' = t$, $r' = r = it|r|$, and $|t|^2 + |r|^2 = 1$.

Consider light having an SOP $|p⟩$ being injected into port a1 of the 50:50 OC. The transmitted (reflected) state $|p⟩$ ($|p⟩$) at port a2 (a1) is given by $A_T|p⟩ = g_A|p⟩$ ($g_A|p⟩$), where the total gain $g$ is real, the operator $A_T$ ($A_R$) is given by $A_T = tt'e^{-iφ /2}σ_z + rr'e^{-iφ /2}J_zσ_z$ ($A_R = tr'e^{iφ /2}σ_zJ_z + rt'e^{-iφ /2}J_zσ_z$), the nonlinear phase shift $φ$ is real, $σ_z = \text{diag}(1, -1)$ is Pauli’s $z$ matrix, and $J_z$ is the Jones matrix corresponding to circulating the FOLM loop in the clockwise (counter clockwise) direction. The assumption is that both $f_s$ and $J_z$ are unitary, i.e., $J_z^†J_z = 1$, and together with the relations

$$A_T^†A_T = \mathcal{J} \left( t'^t e^{-iφ /2} \sigma_z \right) \left( t't e^{iφ /2} \sigma_z \right) \mathcal{J}$$

$$= \mathcal{J} \begin{pmatrix} |t|^2 & -|r|^2 \\ -|t|^2 & |r|^2 \end{pmatrix} \mathcal{J} \tag{1}$$

and

$$A_R^†A_R = \mathcal{J} \left( t'r' e^{-iφ /2} \sigma_z \right) \left( t'r' e^{iφ /2} \sigma_z \right) \mathcal{J}$$

$$= |r|^2 \mathcal{J} \begin{pmatrix} 1 & 1 \\ 1 & 1 \end{pmatrix} \mathcal{J}, \tag{2}$$

where

$$\mathcal{J} = \left( e^{-iφ /2} f_s \sigma_z, e^{iφ /2} \sigma_z f_s^† \right)$$

yields

$$A_T^†A_T = (|t|^2 - |r|^2)^2 + 2|r||t| (1 - S), \tag{4}$$

$$A_R^†A_R = 2|r|^2 (1 + S), \tag{5}$$

where the operator $S$ is given by (note that $σ_z^2 = 1$)

$$S = \frac{e^{-iφ} f_s σ_z J_z + e^{iφ} σ_z J_z^† f_s^†}{2}. \tag{6}$$

These relations imply the unitarity condition

$$A_T^†A_T + A_R^†A_R = (|t|^2 - |r|^2)^2 + 4|r||t|^2 = 1. \tag{7}$$

A derivation (see Appendix B) similar to the one that yields the Heisenberg uncertainty principle is employed to derive a lower bound upon the FOLM reflectivity $P_R = \langle p | A_R^†A_R | p \rangle / \langle p | p \rangle$ (the commutation relation is denoted by $[\cdot, \cdot]$).

$$P_R \geq \frac{|r|^2 |\text{Re}(p e^{-iφ} [J_z^† σ_z, J_z σ_z]) p|}{\langle p | p \rangle}. \tag{8}$$

The relation $A_R^†A_T + A_T^†A_R = 1$ together with inequality (8) yields an upper bound upon the FOLM transmissivity $P_T = \langle p | A_T^†A_T | p \rangle / \langle p | p \rangle$. However, both $P_R$ and $P_T$ become unbounded when $[J_z^† σ_z, J_z σ_z] = 0$. In particular, no bound is imposed by inequality (8) when $J_z = σ_z J_z σ_z$. As will be shown below, the condition $J_z = σ_z J_z σ_z$ is satisfied when the SOP evolution along the FOLM depends only on the geometry of the fiber spatial curve.

V. SOP EVOLUTION

The Jones matrices $J_z$ and $J_z$ can be calculated by integrating the equation of motion for the SOP along the FOLM. Consider an optical fiber wound in some spatial curve in space. Let $s(s)$ be the arc-length parametrization of this curve, i.e., the tangent $\mathbf{s} = ds/\mathrm{ds}$ is a unit vector. The normal and binormal Serret–Frenet unit vectors are denoted by $\mathbf{ν}$ and $\mathbf{b}$, respectively. The curve torsion $t$ is defined as $\mathrm{d} \mathbf{b}/\mathrm{ds} = -\mathbf{t}$. A unit vector parallel to the electric field is denoted by $\mathbf{e}_0$. In the Serret–Frenet frame, the unit vector $\mathbf{e}_0$, which is expressed as $\mathbf{e}_0 = \mathbf{e}_0 \mathbf{ν} + \mathbf{e}_0 \mathbf{b}$, evolves according to the parallel transport equation of geometrical optics as $[^9, ^11]$

$$\frac{\mathrm{d}}{\mathrm{ds}} \begin{pmatrix} e_0(s) \\ e_ν(s) \end{pmatrix} = i \mathcal{K} \begin{pmatrix} e_0(s) \\ e_ν(s) \end{pmatrix}, \tag{9}$$

where $\mathcal{K} = \mathcal{K}_0 + \mathcal{K}_t$. The geometrical birefringence $\mathcal{K}_0$ is given by $[^3, ^2]$

$$\mathcal{K}_0 = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \tag{10}$$

whereas $\mathcal{K}_t$ is the fiber birefringence induced by elasto-optic, electro-optic, or magneto-optic effects. For a lossless fiber, the matrix $\mathcal{K}$ is Hermitian.

Consider the case where $\mathcal{K}_t$ vanishes. For this case, the solution to Eq. (9) is given by $[^2]$

$$\begin{pmatrix} e_0(s) \\ e_ν(s) \end{pmatrix} = \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} e_0(0) \\ e_ν(0) \end{pmatrix}, \tag{11}$$

where the rotation angle $\theta$ is given by the integrated torsion along the fiber curve,

$$\theta = \int_0^s \mathrm{d}s' t(s'). \tag{12}$$

For the case where the curve end points are parallel, i.e., $\mathbf{s}(s) = \mathbf{s}(0)$, the following holds good: $\theta = \Omega$, where $\Omega$ is the solid angle subtended by the closed curve $\mathbf{s}(s')$ at the origin. This geometrical rotation of polarization, which is closely related to Berry’s phase, has been experimentally measured in Refs. 42 and 44.
For the case $K_f = 0$, $J_r = \sigma f J_u \sigma$, holds good, and consequently, $S = \cos \Theta$, $A^1_{1,2} A_r = (|r|^2 - |t|^2)^2 + 4|t|^2 \sin^2(\Theta/2)$, and $A^2_{1,2} A_r = 4|t|^2 \cos^2(\Theta/2)$ (it is assumed that $f_j J_m = 1$). Hence, for this case, both reflectivity $P_R$ and transmissivity $P_T$ become independent of the input SOP. For a 3 dB OC, i.e., $|t|^2 = |r|^2 = 1/2$, and in the linear limit, i.e., for $\Theta = 0$, the FOLM for the same case where $K_f = 0$ becomes perfectly reflecting, i.e., $P_R = 0$ and $P_T = 1$.

For treating the general case (where $K_f$ cannot be disregarded and consequently FOLM’s perfect reflectivity cannot be guaranteed), it is convenient to express the matrix $J_r$ as $J_r = f_m J_u f_s J_u$, where $J_m$ is unitary, i.e., $f_m^2 J_m = 1$. Using this notation, one finds that

$$P_R = 4|t|^2 (1 - \eta),$$

$$P_T = (|t|^2 - |r|^2)^2 + 4|t|^2 \eta,$$

where

$$\eta = \sin^2 \frac{\Theta}{2} - \frac{\text{Re}(e^{-i\phi}(p|J_m f_r|p))}{2(p|p)}$$

and $V_m = \sigma_c (J_m - 1) \sigma$. Note that $J_m = 1$ and $V_m = 0$ for the case $K_f = 0$.

VI. MAPPING

In this section, a cycle to cycle mapping is derived to analyze the SOP evolution. Stability analysis of the mapping allows DOP evaluation in the steady state. The input SOP $|p_i|$ after $n$ cycles is denoted by $|p_n|$. The mapping between $|p_n|$ and $|p_{n+1}|$ is given by

$$|p_{n+1}| = g f J_r A_r |p_n|,$$

where $f$ is an SOP, for which the amplification factor $F(|p|)$ is maximized (for a given $\Theta$). If, in addition, $|p_f|$ is an eigenvector of the mapping operator $g f J_r A_r$, then the DOP is expected to be relatively high in the steady state, provided that $|p_f|$ is a locally stable fixed point of the mapping Eq. (16). For the case shown in Fig. 4(a), for which the operator $A_r$ is given by $A_r = i \sin(\Theta/2) \sigma_2$, the SOP $|p_f|$ is an eigenvector of $g f J_r A_r$ for all $\Theta$ [note that $|\lambda_{1,2}| = \lambda_{1,2}^{1/2} = \sin(\Theta/2)$ for this case]. The equality $|\lambda_{1,2}| = \lambda_{1,2}^{1/2}$ [see Fig. 4(a)] is violated by the transformation $J_m$ [see Fig. 4(b)], and the equality $|\lambda_{1}| = |\lambda_{2}|$ [see Figs. 4(a) and 4(b)] is violated by the transformation $J_1$ [see Fig. 4(c)]. The effect of violating the condition $|t|^2 = |r|^2$ is demonstrated in Fig. 4(d).

The drop in PDO [see Figs. 3(c1) and 3(c2)] and the DOP [see Figs. 3(d1) and 3(d2)] occurring in the transition from CW to ML is attributed to the change in peak power, which, in turn, affects SOP evolution due to nonlinearity of the cycle to cycle mapping equation (16) (induced by both Kerr effect and gain saturation). The drop in PDO represents the change in the averaged optical loss per cycle, which depends on the FOLM reflectivity $P_R$.

A bifurcation diagram example of the mapping Eq. (16) is shown in Fig. 5. The parameters that are used for the calculation are listed in the figure caption. To account for saturation, the gain $g$ is assumed to be given by

$$g = \frac{g_0}{1 + \frac{I}{I_s}}.$$  

where $g_0$ is the small signal gain and $I_s$ is the saturation intensity.

The Kerr effect induced nonlinear phase $\Theta$ is assumed to be given by $\Theta = \zeta I$, where $\zeta$ is a positive constant and where $I = |p|/\sigma$. The equality $|\lambda_{1,2}| = \lambda_{1,2}^{1/2}$ [see Fig. 4(a)] is violated by the transformation $J_m$ [see Fig. 4(b)], and the equality $|\lambda_{1}| = |\lambda_{2}|$ [see Figs. 4(a) and 4(b)] is violated by the transformation $J_1$ [see Fig. 4(c)]. The effect of violating the condition $|t|^2 = |r|^2$ is demonstrated in Fig. 4(d).
the region $4.28 \leq g_0 \leq 4.51$, where the period time of the mapping equation (16) becomes much longer than the loop fundamental period.

Note that the mapping equation (16) strongly depends on the unitary matrices $J_u$, $J_{in}$, and $J_{out}$. In the experiment, this dependency can be explored by tuning both PCs. The dependency shown in Fig. 3, of both PDO and DOP on diode current $I_{D0}$, can be significantly modified by retuning both PCs. Similarly, the bifurcation diagram shown in Fig. 5 can be significantly modified by changing the assumed unitary matrices. In a fiber-based setup, it is very difficult to independently measure each of these unitary matrices for any given paddle configuration of both PCs. However, when these unitary matrices are treated as fitting parameters, good agreement between theory and experiment can be obtained.

VII. SUMMARY

In summary, DOP tunability is experimentally demonstrated in an F8L, in both CW and ML regions. The mapping equation (16) is employed to derive the SOP time evolution. A high DOP can be obtained in the regions where the mapping has a locally stable fixed point. The nonlinearity of the cycle to cycle SOP mapping equation (16) gives rise to complex dynamics. The current experimental setup, which is based on the rotating quarter-wave plate method, does not allow monitoring of the relatively fast SOP cycle to cycle evolution since the fiber ring period time is much shorter than the plate rotation period time. Future experiments, which will employ a faster polarimeter, will be devoted to the cycle to cycle complex dynamics of this system.

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AUTHOR DECLARATIONS

Conflict of Interest

The authors have no conflicts to disclose.

Author Contributions

Banoj Kumar Nayak: Investigation (lead); Methodology (equal); Data curation (lead); Writing – original draft (lead). Ciji Mathai: Methodology (equal). Dmitry Panna: Methodology (equal). Eyal Buks: Supervision (lead); Formal analysis (lead); Writing – original draft (lead).

DATA AVAILABILITY

The data that support the findings of this study are available from the corresponding author upon reasonable request.

APPENDIX A: AUTOCORRELATION

In a Michelson interferometer with a second harmonic signal, the detector signal $I_C(t)$ corresponding to second-order interferometric correlation can be written as a function of delay time $\tau$ as [see Eq. (9.6) of Ref. 45] as

$$I_C(t, E(t)) = \int_{-\infty}^{\infty} dt \left| (E(t) + E(t - \tau))^2 \right|. \quad (A1)$$

where $E(t) = e^{i\omega_0 f_{\alpha}^{1/2}}(t)$ is the electric field of a Gaussian partial pulse at time $t$ [see Eq. (2.39) of Ref. 46], $\omega_0$ is the carrier angular frequency, and the normal distribution function having a standard deviation $\sigma$ and a vanishing expectation value is given by

$$f_\sigma(x) = \frac{1}{\sigma \sqrt{2\pi}} \exp \left( -\frac{x^2}{2\sigma^2} \right). \quad (A2)$$

The pulse full width at half maxima is $2\sqrt{2\ln(1/2)}\sigma_c$.

A fitting based on Eq. (A1) could not accurately replicate the broad pedestal (wings) in our experimental AC data shown in Fig. 2(b). Better agreement can be obtained from a double-pulse model, which assumes pulse splitting with a varying separation time $t_s$ (see Fig. 4.8 of Ref. 46). The splitting gives rise to the so-called coherence artifact. In this model, the electric field $E_c(t, t_s)$ is assumed to be given by

$$E_c(t, t_s) = e^{-i\omega_0 f_{\alpha}^{1/2}(t - \frac{\pi}{2} + t_s)} + e^{i\omega_0 f_{\alpha}^{1/2}(t + \frac{\pi}{2} - t_s)}. \quad (A3)$$

The double-pulse assumption is supported by Ref. 47, which demonstrates transition from the single-pulsing to the double-pulsing regime with increasing pump current beyond the single-pulsing instability threshold. In our case, the time of separation $t_s$ is...
considered to have a normal distribution with standard deviation $\sigma$. The averaged AC signal, which is denoted by $\langle I_C(t) \rangle$, is given by

$$\langle I_C(t) \rangle = \int_{-\infty}^{\infty} dt \, f_\sigma(t) I_C(t, E_x(t, t)).$$  \hspace{1cm} (A4)

The values of both standard deviations $\sigma$ and $\sigma_\tau$ are extracted by data fitting (see the caption of Fig. 2).

**APPENDIX B: INEQUALITY (8)**

The positive definite operator $A_{\not{\tau}}^+ A_{\not{\tau}}$ can be expressed as

$$A_{\not{\tau}}^+ A_{\not{\tau}} = \left[ A_{\not{\tau}}^+ A_{\not{\tau}} - A_{\not{\tau}} A_{\not{\tau}}^+ \right]/2,$$  \hspace{1cm} (B1)

where $\left[ A_{\not{\tau}}^+ A_{\not{\tau}} - A_{\not{\tau}} A_{\not{\tau}}^+ \right]$ is anti-Hermitian and $\left[ A_{\not{\tau}}^+ A_{\not{\tau}} + A_{\not{\tau}} A_{\not{\tau}}^+ \right]$ is Hermitian; thus,

$$\left| p A_{\not{\tau}} A_{\not{\tau}}^+ |p\rangle \right|^2 \geq \frac{\left| p \left( A_{\not{\tau}}^+ A_{\not{\tau}} + A_{\not{\tau}} A_{\not{\tau}}^+ \right) |p\rangle \right|^2}{2\langle p |p \rangle}.$$  \hspace{1cm} (B2)

The relation $A_{\not{\tau}} = tr e^{\rho_{\not{\tau}}/2} \sigma J \cdot \sigma + rt' e^{-\rho_{\not{\tau}}/2} f_{\not{\tau}} \Delta \sigma z$ yields (recall that it is assumed that $f_{\not{\tau}} f_{\tau} = f_{\not{\tau}} f_{\not{\tau}} = 1$)

$$\left| A_{\not{\tau}}^+ A_{\not{\tau}} \right|_{|rt|} = e^{\rho_{\not{\tau}}} \left| p \left( f_{\not{\tau}} f_{\not{\tau}} \right) \sigma_{\not{\tau}} \sigma_{\not{\tau}} \right| |p\rangle \right| + e^{\rho_{\not{\tau}}} \left| p \left( f_{\not{\tau}} f_{\not{\tau}} \right) \sigma_{\not{\tau}} \sigma_{\not{\tau}} \right| |p\rangle \right|,$$

and then, the inequality equation (8) holds good.

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