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Dissipative Solitons in Passively Mode-Locked Fiber Lasers

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School of Electrical & Electronic Engineering

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Statement of Originality

I hereby certify that the work embodied in this thesis is the result of original research and has not been submitted for a higher degree to any other University or Institution.

Date

Wu Xuan

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Summary

Passively mode-locked fiber lasers due to their capability of generating ultrashort pulses, low cost and compact size have attracted considerable attentions. Conventionally, to generate ultrashort pulses the mode-locked fiber lasers were operated in the anomalous dispersion regime, where the natural balance between the fiber dispersion and optical Kerr effect shapes the mode-locked pulses into optical solitons. Recent studies have further shown that optical solitons could even be formed in the normal dispersion regime, where the soliton formation is a result of the mutual interaction among the cavity dispersion, fiber nonlinearity, laser gain saturation and gain bandwidth filtering. The solitons formed in the mode-locked fiber lasers are a type of dissipative solitons.

The dynamics of dissipative solitons formed in a fiber laser is governed by the Ginzburg-Landau equation. Different from the conventional nonlinear Schrödinger equation (NLSE) type of solitons, dissipative solitons possess a number of new properties, such as scalable pulse energy, strong frequency chirp. Although there have been considerable theoretical and numerical studies on the properties and features of the dissipative solitons, no systematic experimental studies on them have been made, since such dissipative solitons were only experimentally observed recently. The current PhD thesis presents results of extensive experimental investigations on the features and dynamics of the dissipative solitons formed in passively mode-locked fiber lasers.

Firstly, an all-normal-dispersion (ANDi) erbium-doped fiber (EDF) laser passively mode-locked by the nonlinear polarization rotation (NPR) technique was set up. The formation and properties of the scalar dissipative solitons were experimentally investigated. The NPR initiates the mode locking of the laser. The mutual action of the normal cavity dispersion, cavity nonlinearity and the gain bandwidth filtering shapes the mode-locked pulses into the dissipative solitons. Furthermore, due to the existence of an intracavity polarizer, which fixes the polarization of light in the cavity, only the scalar solitons could be formed in the laser. The advantage of operating a fiber laser in the dissipative soliton regime is that large energy mode-locked pulses could be generated. It was shown that through optimizing the cavity design, dissipative solitons with maximum single pulse energy as large as 7.2 nJ could be obtained in an EDF laser.

Dissipative solitons have not only been observed in the ANDi cavity fiber lasers, but also in the dispersion-managed (DM) cavity fiber lasers. However, depending on the net cavity dispersion dissipative solitons could exhibit different features. In the large anomalous cavity dispersion regime dissipative solitons are characterized by the Kelly sidebands formation, while in the large normal cavity dispersion regime dissipative solitons are characterized by spectra with sharp steep edges. Firstly in my experiments it was demonstrated that either dissipative solitons with clear Kelly sidebands or sharp steep spectral edges could appear in a DM fiber laser with net near zero cavity dispersion. Moreover, stretched pulses, which are normally observed at around zero cavity dispersion, were also observed in a large net normal dispersion cavity fiber laser. It is understood that in a DM cavity fiber laser a soliton pulse is compressed in

the anomalous dispersion fiber and stretched in the normal dispersion fiber. The effective cavity dispersion actually varies with the concrete cavity design. The observed experimental results therefore are a direct reflection of this interesting DM laser cavity feature.

Dissipative soliton resonance is a special feature of the dissipative solitons formed in a fiber laser occurred under certain special laser parameters settings. Unlike the dissipative solitons beyond the resonance range, which have soliton energy quantization, the solitons in resonance have unlimited soliton pulse energy and clamped peak power. Dissipative soliton resonance provides an effective way to generate large energy mode-locked pulses in fiber lasers. Using a specially designed long-cavity ANDi EDF laser, we have first experimentally confirmed the dissipative soliton resonance effect. It was shown that the energy of a near resonance dissipative soliton could be increased continuously. A record-high single pulse energy of 281.2 nJ was generated, which was only limited by the available pump power. The dissipative solitons observed at near resonance have exactly the properties as those predicted theoretically, namely a rectangular pulse with fixed peak power and variable pulse width or a flat-top soliton.

Different from the soliton propagation in a single mode fiber, a soliton circulating in the laser cavity is also subject to the actions of the discrete cavity components, which shape the dissipative solitons. The intracavity dynamics of the dissipative solitons in a DM fiber laser with large normal cavity dispersion has been firstly investigated. The results showed that due to the de-chirping of the anomalous fiber in the cavity, the dissipative soliton no longer experienced pulse width narrowing in the normal

dispersion gain fiber as in the non-DM fiber lasers, but had a similar evolution as that of parabolic pulses only with nonlinear chirp in virtue of the finite gain bandwidth.

Apart from the period-doubling bifurcation and period-doubling route to chaos, experimentally a new type of the deterministic soliton dynamics known as intermittency was further observed in a mode-locked fiber laser. In the intermittency state it was observed that the regular soliton pulse intensity of the dissipative solitons was interrupted by the bursts of large intensity solitons, and the average durations of the pulse bursts decreased with the pump power. Numerical simulations have well reproduced our experimental observations.

Formation of bound states of solitons is now a well-known phenomenon for solitons formed in the anomalous dispersion fiber lasers. Different from the conventional solitons, dissipative solitons have large frequency chirp. Whether dissipative solitons could also form bound states is unclear. In a DM fiber laser with large normal dispersion we have firstly observed bound states of dissipative solitons. It was found that in spite of the frequency chirp and strong pulse stretch and compression in the DM laser, the bound dissipative solitons still have fixed pulse separation and phase difference. Bound states of dissipative solitons in an EDF laser mode-locked by carbon nanotubes (CNTs), which has different saturable absorption properties from NPR, have also been studied. Various bound states of dissipative solitons obtained in the same fiber laser makes a systematic experimental study on them and a comparison of their properties possible. Our results further show that the tightly bound soliton states are characterized by the fixed phase difference, discrete pulse separations, and ultra stability, while the loosely bound soliton states no longer possess these features.

Replacing the NPR mode locking technique with the polarization insensitive saturable absorption method, such as CNTs or semiconductor saturable absorber mirror (SESAM) method, vector dissipative solitons (VDSs), which are solitons with two orthogonal polarization components, were also experimentally obtained and investigated. Depending on the cavity birefringence the polarization locked or the polarization rotating VDSs were observed in spite of the fact that dissipative solitons are strongly chirped. Experimentally it was observed that the period of polarization rotation of the polarization rotating VDSs was either lock to the cavity roundtrip time or multiple of it. Multiple VDSs could also be formed. Apart from two sets of VDSs with orthogonal polarization orientations, two sets of VDSs with parallel orientations were also achieved in our laser since the gain competition was effectively suppressed when the pulse interval was comparable or even larger than the relaxation time of the EDF. In addition, bound state of VDSs were also experimentally observed and numerically verified.

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List of Acronyms

Abbreviations	Full Expressions
ANDi	All-Normal-Dispersion
CNT	Carbon Nanotube
CW	Continuous Wave
DCF	Dispersion-Compensating Fiber
DFWM	Dissipative Four-Wave-Mixing
DM	Dispersion-Managed
DSR	Dissipative Soliton Resonance
EDF	Erbium-Doped Fiber
FBG	Fiber Bragg Grating
FWHM	Full Width at Half Maximum
GLE	Ginzburg-Landau Equation
GVD	Group Velocity Dispersion
GVLVS	Group-Velocity-Locked Vector Soliton
MBE	Molecular Beam Epitaxy
MI	Modulation Instability
MOVPE	Metal-Organic Vapor-Phase Epitaxy
NLSE	Nonlinear Schrödinger Equation
NPR	Nonlinear Polarization Rotation

PC	Polarization Controller
PDL	Polarization Dependent Loss
SA	Saturable Absorber
SESAM	Semiconductor Saturable Absorber Mirror
SMF	Single Mode Fiber
SPM	Self-Phase Modulation
SWCNT	Single-Wall Carbon Nanotube
TOD	Third-Order Dispersion
VDS	Vector Dissipative Soliton
WDM	Wavelength Division Multiplexer
XPM	Cross Phase Modulation
YDF	Ytterbium-Doped Fiber

Chapter 1

Introduction

1.1 Background and Motivation

Ultrafast science has developed vigorously in the past two decades and become one of the most active fields of optics today. Ultrashort pulses with high energy have attracted considerable attentions in virtue of their widespread applications in various scientific fields and industries. Notable examples of their applications include high-density optical data storage [1], material processing [2], fabrication of nanostructures in transparent dielectrics such as silica glass [3], precise laser surgery without destructive effects of the surrounding tissues [4], dissection of biological samples [5], and noninvasive manipulation of the structural machinery of cells [6]. They could also be applied to generate exceptionally wideband, or supercontinuum light [7], which has found numerous applications in biomedical imaging [8]. The promising applications seem to cover all areas of human activities. Hence, progress in the generation of ultrashort optical pulses with high energy possesses great practical significance, as well as promoting the advances in theoretical research.

Bulk solid-state lasers and fiber lasers are the two common and notable devices employed to produce high intensity ultrashort pulses today [9-12]. A great deal of work related to the laser systems has been done to improve the properties of the output pulses, such as higher peak power and shorter pulse width. The classical solid-state lasers have dominated the market and are still the preferred choice as they generate the shortest pulses with highest peak power. However, these bulk solid-state lasers have their own shortcomings and limitations. They are in general very bulky, excessively complicated, and expensive. They are sensitive to the misalignment of bulk optics and unreliable, which require continuous maintenance. All above cumbersome features limit them to the usage of research laboratories, and hamper their further practical application development.

Fiber lasers overcome above mentioned drawbacks of the solid-state lasers and have become strong candidates that make ultrafast lasers applicable beyond research laboratories. Fiber lasers are well known for their prominent advantages of compact design, low sensitivity to alignment, great efficiency, and low production cost. However, they still have their own drawbacks that they normally generate less powerful pulses than the solid-state lasers. Recent years, notable efforts have been made on improving the output pulse energy of fiber lasers so that their performance becomes comparable to their solid-state counterparts.

When an intense ultrashort pulse propagates in an anomalous dispersion optical fiber, the balanced interaction between the anomalous group velocity dispersion (GVD) and nonlinearity always lead to the formation of an optical soliton [13]. Soliton here refers to a narrow, high peak power optical pulse, whose pulse profile and energy remain

invariant during the propagation. The existence of the soliton in optical fibers was first theoretically predicted by Hasegawa and Tappert in 1973 [13] and experimentally observed in single mode fibers (SMFs) by Mollenauer et al. in 1980 [14]. Optical solitons have been extensively investigated as the information carriers for the ultra-high speed telecommunications since their first realization in optical fibers. Theoretically, an optical soliton is a stable solution of the nonlinear Schrödinger equation (NLSE) which describes the optical pulse propagation in a Hamiltonian system with Kerr nonlinearity and anomalous dispersion.

The mode-locked erbium-doped fiber (EDF) lasers have been generally operated in the anomalous cavity dispersion regime, where a mode-locked pulse could be automatically shaped into an optical soliton. A laser operating in the soliton formation parameter regime is referred to as a soliton laser [15]. Soliton operation of fiber lasers has been extensively investigated [16-18]. Although a soliton fiber laser can emit stable ultrashort optical pulses, the output pulse energy is limited. For a given dispersion and soliton pulse width, the soliton pulse energy is ascertained by the soliton area theorem. As a result, although the pulse duration achieved from a soliton fiber laser could be as short as sub-100 fs, the maximum single pulse energy available only reaches ~ 0.1 nJ.

To achieve larger pulse energy, mode-locked fiber lasers with dispersion-managed (DM) cavity, where the cavity is made of alternatively positive and negative dispersion fibers, have been proposed [19, 20]. Pulses propagating in a DM cavity have their pulse width stretched and compressed periodically. Benefit from the fact that a pulse circulating in a DM cavity would have less net nonlinear phase shift, a

larger energy mode-locked pulse could be generated in the fiber lasers. It has been shown that in spite of the periodic variation of the fiber dispersion within the cavity, on average the features of the stable mode-locked pulses formed still follow the soliton theory. For this reason, the nonlinear pulses formed in this type of lasers are known as the DM solitons. The shape of the DM solitons is closer to a Gaussian profile rather than the “sech” shape associated with the standard solitons [21]. A DM soliton fiber laser was first demonstrated by Tamura et al. in 1993 [19]. Under strong pumping and near zero net cavity dispersion, a DM fiber laser could also operate in an interesting state, where the optical spectrum of the mode-locked pulse has a complicated profile and a bandwidth broader than the gain bandwidth. Such a state of fiber laser operation was called the stretched-pulse operation [19]. Ultrashort mode-locked pulses with pulse energy roughly an order of magnitude higher than that of the conventional solitons could be obtained in the stretched-pulse fiber lasers. Until now, ultrashort pulses with energy as high as 2.7 nJ and pulse width of 100-fs have been reported for passively mode-locked fiber lasers operating in the stretched-pulse state [22], and the shortest pulse width achieved in a stretched-pulse EDF laser is 37.4 fs [23].

To pursue even higher pulse energy, the research interests have been further shifted to normal dispersion region. Recent researches reveal that solitary waves could even be formed in fiber lasers with positive dispersion [24-28]. In a normal dispersion fiber laser, no balance of the nonlinearity and anomalous fiber dispersion occurs, and hence no formation of the conventional NLSE type soliton is possible. Solitons are formed in the lasers as a result of the mutual nonlinear interaction among the positive cavity dispersion, self-phase modulation (SPM), effective gain bandwidth filtering, and gain

saturation. Considering that the gain-loss balance plays an essential role in the pulse shaping, the formed solitons are also known as the dissipative solitons. The dissipative solitons formed in fiber lasers are described by the extended Ginzburg-Landau equations (GLEs).

Bélanger et al. first theoretically predicted the existence of such a type of soliton in the amplified nonlinear dispersive media [24]. They pointed out that the solitary pulses had a hyperbolic secant amplitude profile. Akhmediev et al. have conducted a series of theoretical and numerical investigations on the properties and kinetics of the dissipative solitons formed in fiber lasers [29-33]. It was shown that the dissipative solitons had abundant and distinctive features, such as the scalable pulse energy and formation of stable asymmetric soliton pairs. The experimental observation of dissipative solitons in fiber lasers with normal dispersion was reported recently. Zhao et al. firstly reported the observation of solitons in an all-normal-dispersion (ANDi) passively mode-locked EDF laser [25]. To highlight the effect of the gain bandwidth filtering on the soliton formation, they also named the observed soliton as “gain-guided soliton”. Chong et al. reported the observation of the dissipative solitons in an ANDi ytterbium-doped fiber (YDF) laser with a spectral bandpass filter in cavity [34], where the intracavity spectral filter performed a similar function as the gain bandwidth filtering since the YDF had a much broader gain bandwidth than that of the EDF. Under the dissipative soliton operation much larger pulse energy could be achieved in a fiber laser. The highest pulse energy achieved in both of the ultrafast EDF and YDF lasers has been over 20 nJ [35, 36].

Energy dissipation is a general feature of all practical physical systems. Therefore,

study on the formation and dynamics of dissipative solitons is of great practical significance. Passively mode-locked fiber lasers provide an ideal test-bed for studying the general properties of the dissipative solitons. Although solitons formed in fiber lasers have been explored extensively, until now the main emphasis has been given to the solitons formed in anomalous dispersion lasers. Despite that solitons formed in an anomalous dispersion cavity fiber laser also experience laser gain and loss, due to the influence of other effects such as the natural balance between the dispersion and SPM, and the pulse peak clamping, the formed solitons have normally a much narrower spectral bandwidth than that of the laser gain, so the gain spectral filtering effect is generally not obvious. Consequently, the solitons dominantly exhibit the NLSE soliton features rather than the dissipative soliton features. In contrast, for the solitons formed in normal dispersion fiber lasers the gain bandwidth filtering always plays an essential role in their pulse shaping in virtue of the absence of the natural soliton formation mechanism.

1.2 Objectives

Dissipative solitons, as a new type of solitons, attract great attention and are under intensive research nowadays. As addressed above, the dynamics of dissipative solitons formed in a fiber laser is governed by the GLE. Numerous theoretical and numerical investigations on the GLE have been conducted previously, which revealed that the dissipative solitons possessed plenty of unique and interesting features, such as the strong frequency chirp and scalable pulse energy. However, since such a new type soliton was only experimentally observed recently, most of these features lack the experimental confirmations, especially for the existence of the flat-top dissipative

solitons and dissipative soliton resonance, the formation of bound states of dissipative solitons with steep spectral edges and even vector dissipative solitons subjected to large frequency chirp. Considering that majority of the predicted properties of the dissipative solitons still remain experimentally untouched, a systematic experimental study on the features and properties of the dissipative solitons formed in fiber lasers appears to be necessary and meaningful.

On the other hand, the pulse propagation in a fiber laser with normal dispersion gains a strong positive chirp which broadens its pulse width heavily and significantly lowers the peak power. Such a pulse could tolerate larger nonlinear phase shift. As a result the energy of the dissipative solitons formed in a normal dispersion fiber laser can be greatly enhanced as compared to that of the conventional NLSE type solitons. Hence, through optimizing the cavity design and operation condition, large energy mode-locked pulses can be obtained in fiber lasers.

Note that the nonlinearity still limits the highest pulse energy achievable for dissipative solitons. It is important to understand the diversiform nonlinear effects and figure out their impacts on the performance of the fiber laser systems.

The main objective of the thesis is through conducting extensive experimental investigations on the optical dissipative solitons formed in passively mode-locked fiber lasers to gain a comprehensive understanding on the features of the dissipative solitons, and to pursue the dissipative solitons with large pulse energy by optimizing the design of the fiber laser cavity.

1.3 Major Contributions of the Thesis

The original contributions of this thesis are listed as follows:

1. Investigate experimentally the fundamental properties of the scalar dissipative solitons formed in a nonlinear polarization rotation (NPR) mode-locked fiber laser. Through optimizing the laser design in terms of the net cavity dispersion and output coupling ratio, we show that dissipative solitons with 7.2 nJ pulse energy could be achieved.
2. Explore the features of the dissipative solitons formed in DM fiber lasers under different cavity dispersion. We show that solitons with either Kelly sidebands or steep spectral edges could appear in the same DM fiber laser with around zero cavity dispersion, and the stretched pulses could exist in a DM fiber laser with large normal cavity dispersion. The underlying mechanism is also studied, and it turns out to be the effective dispersion shift induced by the nonlinearity.
3. Demonstrate the first observation of the dissipative soliton resonance (DSR) and the flat-top solitons in fiber lasers by employing a specially designed long-cavity ANDi EDF laser. The dissipative solitons at near resonance are observed to exhibit exactly the properties predicted theoretically. Single pulse energy with a record high 281.2 nJ has been achieved. The result is limited by the available pump power.
4. Present the first investigation on the intracavity dynamics of the dissipative

solitons in a DM fiber laser with large normal cavity dispersion. The results show that on account of the strong pulse stretch and compression in a DM cavity, their evolution is different from the dissipative soliton evolution in an ANDi fiber laser, but similar to that of the parabolic pulses with the nonlinear chirp in virtue of the finite gain bandwidth.

5. Demonstrate the first experimental observation of soliton intermittency in a mode-locked fiber laser. We show that under such a soliton dynamic state the regular soliton pulse intensity is interrupted by the bursts of large intensity solitons, and the average duration of the soliton pulse bursts decreases with the pump power. Numerical simulations further confirm the experimental results.

6. Demonstrate the first observation of the bound state of dissipative solitons in a DM fiber laser with large normal dispersion. We show that in spite of the large frequency chirp and strong pulse stretch and compression within the cavity, the bound dissipative solitons could still have fixed pulse separation and phase difference. Numerical simulations well reproduce our experimental observations. Bound states of dissipative solitons in an EDF laser mode-locked by carbon nanotubes (CNTs), a new type mode locker with different saturable absorption properties from NPR, have also been firstly studied. Various bound states of dissipative solitons are obtained in the same fiber laser, making a systematic experimental study on the bound dissipative solitons possible.

7. Investigate the formation and properties of the vector dissipative solitons (VDSs). We show that both of the polarization-locked and polarization-rotating VDSs

could be formed in spite of the fact that the dissipative solitons are strongly chirped. The polarization rotation dynamics of the latter is also studied and it turns out that the period of polarization rotation of the group-velocity-locked VDSs can either lock to the cavity roundtrip time or multiple of it. Bound states of the VDSs are also experimentally observed and numerically verified.

1.4 Organization of the Thesis

The rest of the thesis is organized as the following:

Chapter 1 provides a brief introduction to the research work of this thesis.

Chapter 2 introduces the fundamental principles of the mode-locking of the fiber laser, and presents briefly the background knowledge on the pulse propagation in optical fibers and corresponding mathematical equations. Our numerical method used to simulate the dissipative soliton operation of the fiber laser, the pulse tracing method, is also illustrated.

Chapter 3 studies the formation and characteristics of the scalar dissipative solitons in an ANDi EDF laser mode-locked by NPR. A dissipative soliton with ~ 7.2 nJ pulse energy is achieved, taking advantage of the large normal cavity dispersion. Dissipative solitons in an extended NPR mode-locked fiber laser is also studied and the underlying mechanism of the extended NPR mode-locking is analyzed.

Chapter 4 explores the impact of the effective cavity dispersion shift induced by the nonlinearity on the dissipative soliton operation in a DM fiber laser, and the resulted features of the dissipative solitons in DM fiber lasers with different dispersion.

Chapter 5 investigates the novel phenomena of the DSR and flat-top solitons. The main advantage concerning energy scaling is present, and a dissipative soliton with ~ 281.2 nJ pulse energy and ~ 18.5 ns pulse duration is achieved as a result.

Chapter 6 explores the intracavity dynamics of the dissipative solitons in a DM fiber laser with large normal cavity dispersion, and compares it with the result attained in an ANDi fiber laser. The occurrence and characteristics of the intermittency phenomenon in a fiber laser are also investigated experimentally. Numerical simulations confirm our experimental results.

Chapter 7 discusses the formation and properties of the bound state of the dissipative solitons formed in a DM fiber laser with large normal dispersion. The bound dissipative solitons still exhibit fixed pulse separation and phase difference, although they experience strong pulse stretch and compression in the cavity. A systematic investigation on the bound state of the dissipative solitons in a fiber laser mode-locked by CNT is also presented.

Chapter 8 investigates the properties of the VDSs. Depending on different cavity birefringence, various VDSs, such as the polarization-locked and polarization-rotating VDSs, are formed in a fiber laser. The polarization dynamics of the latter is analyzed further, and it turns out that the period of polarization rotation could either lock to the

cavity roundtrip time or multiple of it. Bound states of VDSs are also observed.

Finally, Chapter 9 gives the conclusions of the research and the recommendation for future work.

Chapter 2

Theory of Dissipative Soliton in a Mode-Locked Fiber Laser

Short-duration optical pulses have extensive applications in the fields of digital communications, diagnostics of ultrafast processes, and ablation of materials. Mode-locking is a well known technique for the intense short pulse generation. It produces a train of equally spaced optical pulses in temporal domain. Since a mode-locked pulse comprises many frequency components and generally possesses high intensity, its propagation in optical fibers is simultaneously affected by the fiber dispersion and nonlinearity, which in mathematics is well described by the nonlinear Schrödinger equation (NLSE). Provided the dispersion is balanced by the nonlinearity, the mode-locked pulses could be further shaped into optical solitons. When analyzing the pulse propagation in practical optical fibers, the fiber birefringence caused by an imperfect-circular fiber core or stress induced anisotropy needs to be considered as well, so the governed equations are replaced by the coupled NLSEs. Theoretically describing the pulse evolution in a mode-locked fiber laser is much more complicated, since besides optical fibers, the gain and gain dispersion introduced by the EDF

amplifier and the saturable absorption effect also need to be taken into account. Then the dominant equations are extended to the Ginzburg-Landau equations (GLEs) from the NLSEs. Instead the dissipative solitons would be formed provided the mutual interaction among the dispersion, nonlinearity, gain and loss reaches a balance.

The objective of this chapter is to present the basic theoretical background and main equations that govern the dissipative soliton pulse propagation in a mode-locked fiber laser. Numerical modeling to simulate its evolution in the laser cavity is also introduced. Section 2.1 gives a brief introduction of the basic theory about the mode-locked fiber laser. Section 2.2 provides a mathematic description of nonlinear pulse propagation in optical fibers. NLSE and the soliton solutions are introduced. Section 2.3 provides a mathematic description of nonlinear pulse propagation in a gain fiber. GLE and the dissipative soliton solutions are introduced. The last section of this chapter presents the numerical modeling that we use to simulate the dissipative soliton operation of a fiber laser, the so-called pulse tracing technique.

2.1 Fundamentals of Mode-Locked Fiber Lasers

Mode-locking is a widely used technique to generate ultrashort pulses [37-39]. The pioneering work in this area was conducted by W. E. Lamb in 1964, who investigated the frequency locking of three modes [40]. After the theoretical prediction of mode-locking [41], Hargrove et al. successfully demonstrated the first mode-locked laser in the same year [42]. With employing an acousto-optical modulator as the mode locker, pulses with 2.5 ns pulse width and 56 MHz repetition rate were achieved. The first stable passively mode-locked laser was reported by Shank et al., and pulse width

of 1.5 ps was achievable in a dye laser [43]. Mode-locked fiber lasers have been under intensive investigations since its first realization in 1983 [44].

2.1.1 Principle of Mode-Locking

The principle of the mode-locking technique is to lock the phase of longitudinal modes in a laser. It is well known that if the pump is strong enough, a laser cavity may support many different modes oscillating simultaneously within the gain bandwidth. The angular frequency spacing among the modes is given by

$$\Delta\omega = \frac{\pi c}{L} \quad (2.1)$$

where L is the optical length during one round trip inside the cavity. The output of such a laser is a function of time depending on the relative phases, frequencies and amplitudes of the modes, and the total electrical field can be written as

$$E(t) = \sum_{n=0}^{N-1} E_n \exp[i(\varphi_n + \omega_n t)] \quad (2.2)$$

where E_n , φ_n and ω_n are respectively the amplitude, phase, and angular frequency of the n_{th} mode. Usually these parameters are all time varying, so that the modes are incoherent. Assuming that N modes have equal amplitude E_0 for simplicity, then the total intensity is given by:

$$I(t) = |E(t)|^2 = E_0^2 \left| \sum_{n=0}^{N-1} \exp[i(\varphi_n + \omega_n t)] \right|^2 \quad (2.3)$$

This value may exhibit small fluctuations if a few modes happen to be in phase at some time, but for large N it does not vary significantly from the average value. In this case, the laser produces a randomly distributed, averaged output of the intensity

in time domain, which is called the continuous wave (CW) state.

Figure 2.1: Different optical output from (a) CW state and (b) Mode-locking state

However, when the N modes have the same relative phase to one another, i.e. the laser is mode locked, defining ω_n as $\omega_n = \omega_{N-1} - n\Delta\omega$, the total intensity would be given by:

$$I(t) = |E(t)|^2 = E_0^2 \frac{\sin^2(N\Delta\omega t/2)}{\sin^2(\Delta\omega t/2)} \quad (2.4)$$

This expression of the intensity is a periodic function with the maximum value occurring at $\frac{\Delta\omega t}{2} = 0, \pi, 2\pi, \dots, n\pi$. It displays clearly that if the laser is mode-locked,

the laser output is a sequence of short pulses with the periodicity of the cavity round trip. The pulse width is $1/N$ times of the pulse intervals. The peak power of each pulse equals to N times the average power (or N times the power of the same laser with the modes uncoupled) NE_0^2 . Thus to obtain high-power short-duration pulses, there should be a large number of modes; this depends on a broad laser transition and a long laser cavity. Figure 2.1 shows the different output of the CW and mode-locking state.

2.1.2 Active and Passive Mode-Locking

Generally speaking, the mode-locking techniques could be classified into two sorts: the active and passive mode-locking techniques.

- **Active Mode-Locking**

Active mode-locking of fiber laser is a preferred choice for producing ultralow-jitter pulses at high repetition rates. The first actively mode-locked fiber laser was realized by employing a bulk acousto-optic modulator in the mid 1980s [45]. In active mode-locking, a modulator, either an amplitude or phase one, modulated by an external signal source must be applied [46, 47]. When the modulation frequency is the integral times of the laser cavity frequency, the laser can be mode-locked and output pulse train is synchronized to the external modulation signal. Through harmonic active mode-locking, the repetition rate over 1000 times of the fundamental laser frequency and as high as tens of GHz has been realized.

Pulses from actively mode-locked fiber lasers have the advantages of high repetition rate and narrow line-width, and drawbacks of broad pulse width, low peak power, and requiring an expensive modulator to implement the mode-locking. The soliton operation, which possesses the born preponderance: good stability, low time jittering, short pulse width and high peak power, is also difficult to be realized in actively mode-locked fiber lasers. The reason is that harmonic mode-locking is normally implemented, so multiple pulses are generated, which share the laser cavity energy. Hence each mode-locked pulse has low energy. In combination with the initial broad pulse width determined by the active mode-locking technique, it results in the fact that only weak nonlinear self-phase modulation (SPM) is generated by the pulse, which makes them difficult to shape into solitons. Fortunately, passive mode-locking might overcome those shortcomings.

- **Passive Mode-Locking**

Passively mode-locked fiber lasers do not require any externally modulated media or devices [43], but only employ a saturable absorber (SA) in the cavity to realize the mode-locking. The SA refers to an optical device or medium that exhibits an intensity-dependent transmission [48], which selectively absorbs low intensity light and transmits light with sufficiently high intensity, as shown in Figure 2.2. In a passively mode-locked fiber laser, a small intense spike arising from the random intensity fluctuation will be transmitted preferentially by the SA. As this process repeats, the intense spike is amplified and the low-intensity light is absorbed continuously in the cavity, which eventually builds up a train of pulses and starts the mode-locking of the laser.

Figure 2.2: The function of an ideal SA

The repetition rate of passively mode-locked fiber laser is generally low (typically 10~100 MHz) due to the long laser cavity, so it is not a good choice for the source of high bit-rate data transmission systems. Harmonic mode-locking can be used to increase the repetition rate, but it reduces the laser stability [49]. However, compared to actively mode-locked fiber laser, passively mode-locked fiber lasers can produce much shorter and more intense optical pulses simultaneously with a low cost. Considering that the soliton operation is easily obtained with passive mode-locking technique, we mainly focus on the passively mode-locked fiber lasers in this dissertation.

2.1.3 Saturable Absorbers for Passive Mode-Locking

As mentioned above, a SA, known as any device transmitting higher intensities of light at expense of lower intensities, is the key component to realize the passive mode-locking. Nonlinear polarization rotation (NPR), semiconductor saturable absorber mirror (SESAM) and carbon nanotubes (CNTs) are three mainly recognized SAs. NPR is an artificial one, while another two are real ones. Their basic

configuration, fundamental properties and respective merits and drawbacks are listed in the following.

- **Nonlinear Polarization Rotation**

The NPR technique is a method exploiting an artificial SA effect to achieve self-started mode-locking, which is realized by the mutual interaction of a polarizer and nonlinear fiber birefringence. It was first used to mode lock the fiber lasers as a self-sustaining mechanism by Hoffer et al. [50], where mode-locking could not self-start due to the high birefringence and short fiber length used. One year later Matsas et al. are successful to employ it to realize the first self-started mode locking in fiber laser [51]. Since then this technique has been widely used in passively mode-locked fiber lasers to generate ultrashort pulses.

Figure 2.3 shows the basic setup for the NPR technique, which consists of a birefringent optical fiber and two polarizers at both its input and output. The later polarizer is also referred to as an analyzer. The light is transformed into a linearly polarized light after the polarizer, and then splits into two components along two principal polarization axes of the birefringent fiber after entering it. At output of the fiber before the analyzer, the light becomes elliptically polarized due to the fiber birefringence. When the light intensity is low, the final polarization ellipse is entirely determined by the linear fiber birefringence and the polarizer orientation, which is independent on the intensity. When considering the high intensity light, nonlinearity becomes so strong that it introduces a polarization-dependent, intensity-dependent refractive index. The shape and orientation of the final polarization ellipse now

become the light intensity dependent, in combination with the analyzer, which ultimately determines the light transmission. If the analyzer orientation is appropriately selected, a situation that the stronger the light, the larger transmission it experiences after analyzer would be realized, which is similar as the SA effect.

Figure 2.3: Basic setup for the NPR technique [52]

Ultrashort mode-locked pulses could be achieved by using NPR techniques, since now the SA effect comes from the Kerr effect, which has a fast recovery time of fs level. The NPR technique also possesses the advantages of variable modulation depth and saturation energy, and high damage threshold. However, it still has its own disadvantages, such as high mode-locking threshold. Since NPR mode-locking critically depends on the polarization state, it is unstable to the environmental perturbations (especially the birefringence), which is harmful to its commercial applications. Moreover, the peak power clamping effect is found always existing in NPR mode-locked fiber lasers, which significantly limits the pulse energy.

- **Semiconductor Saturable Absorber Mirror**

SESAM, as an intensity and wavelength-dependent mirror constructed using a semiconductor material, is the most widely used SA in commercial passively mode-locked fiber lasers [53-56]. It concretely consists of a Bragg-mirror on a semiconductor substrate, and a single quantum well absorber layer near the surface. Figure 2.4 shows its typical structure. It shows that GaAs/AlAs on the GaAs substrate serves as the Bragg mirror, which causes the wavelength dependent response. The InGaAs quantum well on top of the mirror serves as the SA material, responsible for the intensity-dependent reflection.

Figure 2.4: Schematic structure of the SESAM [57]

Along with the light absorption, the electron-hole pairs are generated in the medium continuously. Limited by that only a finite number of electron-hole pairs could be created, the absorption is saturated at last. Then the electron-hole pairs recombined non-radiatively, and it is possible for the medium to absorb photons again only after a period of time, which is known as the recovery time of SAs. Ideally, shorter recovery time is always preferred provided mode locking operation is still ensured. A typical relaxation time caused by the spontaneous photon emission in a semiconductor is around 1 ns. Extra complicated techniques such as a suboptimal low-temperature

growth environment or high-energy heavy-ion implantation are required to further reduce it to the ps region through introducing defect sites. Bi-temporal recovery time is a distinct property of SESAM, which also makes it a perfect mode locker. The slower one enables the self-started mode-locking, and the faster one, generally around ps or sub-ps, allows the ultrashort pulse generation. Noticeably, the pulse width achieved in fiber lasers mode-locked by SESAMs could be significantly narrower than the recovery time of the SESAM. The reason is that once soliton is formed, its pulse width is only determined by the soliton theory rather than the absorber recovery time. The latter only plays the role in initiating mode-locking of the laser and suppressing the background noise then.

The mode-locking achieved by SESAM is stable and robust against small environmental perturbations since the work principle does not involve the concrete polarization state, which makes them especially favored in commercial applications. However, SESAM still has some shortcomings. For example, its damage threshold is generally low due to the semiconductor SA layer. Expensive and complex clean-room-based fabrication systems such as Molecular Beam Epitaxy (MBE), Metal-Organic Vapor-Phase Epitaxy (MOVPE) are always required for the growth or the recovery time reduction of SESAM. The operation bandwidth of SESAM is also very narrow. Consequently, there is always strong demand for seeking new materials to replace SESAM, achieving the goal of perfect SA with wider operation range, faster response time and lower cost.

- **Carbon Nanotubes**

CNTs, especially single-walled carbon nanotubes (SWCNTs), have been of great interest nowadays due to their distinct physical properties [58]. After they were found to exhibit the Saturable absorption property, their potential application as optical switches has been proposed [59, 60]. Later, a SWCNT-based SA was successfully employed to suppress the optical noise of ultrafast optical pulses in the picoseconds regime [61]. The first Q-switched and mode-locked fiber laser using SWCNTs were reported by S. Y. Set et al. successively [62, 63].

An ideal CNT is a seamless cylinder (Figure 2.5) rolled by a hexagonal network of carbon atoms, whose ends are typically capped with half of a fullerene molecule [58, 64]. Its diameter is normally several nanometers and length is a few microns. A SWCNT can be classified into semiconducting, metallic or semimetallic one in terms of its chirality [65]. Normally, SWCNTs are the hybrid of both semiconductor and metallic tubes, as current synthesis methods can not accurately control the desired chirality or diameter of the tube. Statistically, $2/3$ are semiconducting, while $1/3$ are metallic ones. Excitonic absorption in semiconducting nanotubes has an ultrafast recovery time and plays a key role in the saturable absorption property of SWCNT. Metallic nanotubes play the role of the recombination centers, similar as the defects in SESAM, and are responsible for the ultrafast recovery time of SWCNTs. Semiconducting SWCNTs are direct band-gap materials and their band gaps depend on the nanotube structures, so SWCNTs with a certain diameter and chiral angle only absorb the light with a particular wavelength [66]. By mixing SWCNTs with large number of different diameters and chiral angles, a broadband saturable absorption mode locker could be realized.

Figure 2.5: Schematic model of a chiral nanotube [58]

The superior properties of the SWCNTs, such as intrinsically ultrafast recovery time (subpicosecond), wide operating spectral range, the easiness of fabrication and low cost, all make it a better choice for mode locker than conventional ones. A lot of achievements have been indeed gained in this area. Mode-locking at different wavelengths ranging from C band to L band has been realized by using different SWCNTs-based SA [67, 68]. A wideband-tunable EDF laser with tunable span of 40 nm mode locked by SWCNTs was also reported [69]. With a SWCNTs mode-locked fiber laser, the shortest pulse width achieved was less than 200fs [70], and the highest repetition rate of the pulses was 17.2GHZ up to now [71]. However, the SWCNTs mode locker still suffers intrinsic drawbacks: SWCNTs tend to form bundles that finish up as scattering sites, so the coexistence of SWCNTs with different diameters introduces extra linear losses, making the mode locking of a laser difficult.

2.2 Pulse Propagation in a Single Mode Fiber

After mode-locking of the fiber laser, optical pulse train is formed in the laser cavity. An intense pulse propagating in a single mode fiber is always subject to all the actions of the fiber dispersion, birefringence and nonlinearity. In the following, I will introduce

the mathematical expressions of these items respectively and further present the equation which describes the nonlinear pulse propagation in optical fibers.

2.2.1 Fiber Dispersion and Birefringence

Light of different wavelengths propagates at different speed in fibers. This effect is known as dispersion, which is an intrinsic property of optical fibers. In fact, there are generally two sources of dispersion in single mode optical fibers: material dispersion and waveguide dispersion. Material dispersion comes from a frequency-dependent response of the materials made of a fiber to waves, while waveguide dispersion arises from the wavelength dependence of the propagation constant of a mode to the fiber core diameter and the difference in refractive index between the fiber core and cladding. For a single mode fiber, the material dispersion is generally difficult to be changed, but the total dispersion could still be managed through altering the waveguide dispersion, which corresponds to the special design of the fiber refraction index structure. Therefore, single mode optical fibers of different total dispersion could be realized and are now commercially available. For example, the zero dispersion wavelength of a standard single mode fiber is around 1.3 μm , while that of a dispersion-shifted single mode fiber is about 1.55 μm .

Dispersion is a linear optical property, and it always plays a critical role in the propagation of short optical pulses in optical fibers. Mathematically, the effect of fiber dispersion on the propagation of an optical pulse is described by expanding the mode-propagation constant β in a Taylor series about the pulse's carrier frequency ω_0 :

$$\beta(\omega) = n(\omega)\frac{\omega}{c} = \beta_0 + \beta_1(\omega - \omega_0) + \frac{1}{2}\beta_2(\omega - \omega_0)^2 + \frac{1}{6}\beta_3(\omega - \omega_0)^3 + \dots \quad (2.5)$$

where

$$\beta_m = \left(\frac{d^m \beta}{d\omega^m} \right)_{\omega=\omega_0} \quad (m = 0, 1, 2, \dots) \quad (2.6)$$

The parameters β_0 and β_1 are the propagation constant at the carrier frequency and the inverse of the group velocity of the pulse respectively. The parameter β_2 represents dispersion of the group velocity and is responsible for the optical pulse broadening when it propagates in a single mode fiber, which is also known as the group velocity dispersion (GVD) parameter.

Owing to the technical limitation in fabricating a perfect-circular fiber core and/or the random mechanical stresses, a single mode fiber is generally not perfectly symmetric everywhere along the fiber. As a result, practical single mode fibers always support two polarization eigenmodes, or in other words exhibiting a certain birefringence. The strength of birefringence of an optical fiber is described by the polarization mode beat length L_b :

$$L_b = \frac{\lambda}{n_x - n_y} \quad (2.7)$$

where n_x and n_y are respectively the refractive index of the two polarization eigenmodes of the fiber. A weak CW light propagating in a linearly birefringent single mode fiber is described by:

$$\begin{aligned} \frac{dA_x}{dz} &= i\beta_x A_x \\ \frac{dA_y}{dz} &= i\beta_y A_y \end{aligned} \quad (2.8)$$

where β_x and β_y are the propagation constants along two polarization axes of the fiber, respectively, A_x and A_y are the light field amplitude along each polarization axes.

Obviously, when the light propagates for a distance of L , two polarization

components will accumulate a phase difference of $\Delta\Phi = \frac{2\pi(n_x - n_y)L}{\lambda} = \frac{2\pi L}{L_b}$,

suggesting that the polarization of the light varies along the fiber. For a propagation distance of multiple of L_b , the phase difference between two polarization components becomes multiple of 2π so that the polarization of the light returns to its initial state. Therefore, physically, L_b is the length over which the same polarization state appears during the light propagation in optical fibers.

The speed of light propagating in fibers depends on the refractive index, which is different along two polarization axes, so light polarized along different polarization axes have different phase as well as group velocities. Assuming $n_x > n_y$ for a birefringent fiber, the light polarized along the X axis then moves slower than that along the Y axis, so the X-axis is called the slow axis and the Y-axis is the fast axis. Different single mode fibers have different beat length. For standard single mode optical fibers, a typical value of beat length is around 1 meter, but for strongly birefringent fibers, the beat length could be as small as several millimeters.

2.2.2 Fiber Nonlinearity

Nonlinear effects arise in optical fibers as optical power level increases. Therefore, when describing an intense light propagation in a single mode fiber, the effect of fiber nonlinearity should be taken into account as well. Fundamentally, the fiber

nonlinearity originates from the nonlinear response of polarization to strong light field. Considering that optical fibers normally do not exhibit second-order nonlinear effects, since they are made of SiO₂, which appears to be centro-symmetric at mesoscopic scale, the lowest-order nonlinear effects in optical fibers is the third order effect, which includes the third-harmonic generation, four-wave mixing, stimulated Raman and Brillouin scattering, and optical Kerr effect. The occurrence of third-harmonic generation and four-wave mixing needs to fulfil the phase matching condition, so these effects are always not efficient in single mode optical fibers. The most important nonlinear effect in optical fibers is the optical Kerr effect, which originates from the intensity dependence of the refractive index:

$$n(I) = n_0 + n_2 I \quad (2.9)$$

where n_0 is the refractive index of the fiber without light, and n_2 is the nonlinear-index coefficient related to the third-order nonlinear property of the fiber. I is the optical intensity inside the fiber.

A typical value of the nonlinear refractive index coefficient n_2 of optical fibers is $2.2\sim 3.4 \times 10^{-20} \text{ m}^2/\text{W}$ [72], which is quite small. However, due to the small core of single mode fibers, normally with a diameter $<10 \text{ }\mu\text{m}$, the light is confined to propagate in the fiber core with almost ignorable loss, which makes a very long distance of light-matter interaction possible. Therefore, optical fibers still exhibit strong nonlinearity in spite of the small value of n_2 as a result of the high efficiency of the nonlinear optical process.

2.2.3 Nonlinear Schrödinger Equation and the Soliton Solutions

When an intense optical pulse propagates in a single mode optical fiber, both of the fiber dispersion and nonlinear optical Kerr effect must be considered. As we know, the most basic mathematical method used to accurately model the propagation of optical fields in the media is the Maxwell's equations. Taking account of the concrete expressions describing the dispersion and Kerr nonlinearity of the optical fiber, as shown in 2.2.1 and 2.2.2, the well-known NLSE would be derived from the Maxwell's equations after a complicated process and a few approximations. The most important approximations among them are the nonresonance approximation, the slowly varying envelop approximation and the weak wave-guiding approximation. In fact, the derivation process of the NLSE from Maxwell's equations is well covered in numbers of textbooks such as Ref. [73, 74], and the final equation format is given in the following.

$$\frac{\partial A}{\partial z} + \frac{\alpha}{2} A + i \frac{\beta_2}{2} \frac{\partial^2 A}{\partial T^2} - \frac{\beta_3}{6} \frac{\partial^3 A}{\partial T^3} = i\gamma \left[|A|^2 A + \frac{i}{\omega_0} \frac{\partial}{\partial T} (|A|^2 A) - T_R A \frac{\partial |A|^2}{\partial T} \right] \quad (2.10)$$

Here T is the local time, and z is the propagation distance. $A = A(T, z)$ is a slowly varying field amplitude of the optical pulse. α denotes the gain/loss coefficient, β_2 and β_3 are the GVD and third-order dispersion (TOD) coefficient in Eq. (2.5) respectively. High order dispersion coefficients are omitted here due to their negligible effects in most cases. The nonlinear coefficient γ is defined as

$$\gamma = \frac{n_2 \omega_0}{c A_{\text{eff}}} \quad (2.11)$$

n_2 is the nonlinear index coefficient, c is the velocity of light in vacuum, and A_{eff} is

the effective fiber core area. $\frac{i}{\omega_0} \frac{\partial}{\partial T} (|A|^2 A)$ and $T_R A \frac{\partial |A|^2}{\partial T}$ terms are higher-order

nonlinear terms, which are responsible for the effects of self-steepening, and the Raman self-frequency shift respectively. Considering that the TOD and higher-order nonlinear terms are generally negligible for a pulse with a relatively broad pulse width larger than 100 fs and a low peak power (valid for most of fiber lasers), then Eq. (2.10) is reduced to

$$\frac{\partial A}{\partial z} + \frac{\alpha}{2} A + i \frac{\beta_2}{2} \frac{\partial^2 A}{\partial T^2} = i\gamma |A|^2 A \quad (2.12)$$

Nowadays, optical fiber could be fabricated with extremely low loss so that the gain/loss α could even be neglected without affecting the calculation of the pulse propagation within the fibers. Then Eq. (2.12) is further reduced to the famous NLSE

$$\frac{\partial A}{\partial z} + i \frac{\beta_2}{2} \frac{\partial^2 A}{\partial T^2} = i\gamma |A|^2 A \quad (2.13)$$

Eq. (2.13) includes only two phase modulation terms induced by the GVD and nonlinearity effects. Although the GVD effect alone always broadens the pulse width during its propagation, and consequently the pulse becomes linearly frequency chirped, under the coaction of the fiber Kerr effect, the situation becomes more complicated and totally different.

Here, we would like to point out that a special phenomenon, known as soliton, would occur provided the fiber dispersion is anomalous, where the pulse propagates in the medium without changing its shape. To be more specific, γ is always positive for a silica fiber, so the nonlinear optical Kerr effect always generates a positive frequency chirp in the center part of a pulse, while the GVD could induce either a positive or negative frequency chirp depending on the sign of the GVD parameter.

Provided that the pulse propagates in the anomalous dispersion regime, the GVD of the fiber will generate a negative linear frequency chirp, which is possible to counteract the positive frequency chirp induced by nonlinearity, and eventually lead to the pulse propagating in fibers without any broadening and distortion. Moreover, the chirp induced by the nonlinear Kerr effect is the pulse intensity dependent, which can adjust its intensity to automatically balance the chirp induced by the dispersion. Namely, the generation of optical solitons is actually a generic property of fibers in the anomalous dispersion regime. Until now, it shows obviously that two conditions are necessary for the soliton formation. One is that the light intensity must be strong enough so that the nonlinear phase modulation is generated, and the other is that the fiber has to provide the anomalous dispersion.

The soliton solution could also be achieved from the Eq. (2.13) analytically. As a matter of fact, Eq. (2.13) is integrable. In 1971, Zakharov and Shabat first reported that by employing the inverse scattering method, the NLSE Eq. (2.13) could be solved in analytical form and had a stable localized solution, referred to as the fundamental soliton solution, in the case of $\beta_2 < 0$. Such solitons are derived to possess a characteristic hyper-secant functional form, which is given in the following

$$A(T, z) = \sqrt{\frac{|\beta_2|}{\gamma T_0}} \text{Sech}\left(\frac{T}{T_0}\right) \exp\left(\frac{iz}{2}\right) \quad (2.14)$$

where T_0 is the soliton pulse width, and $T_{FWHM} \approx 1.763T_0$. The peak power required to form above fundamental soliton in optical fibers is

$$P_0 = \frac{|\beta_2|}{\gamma T_0^2} \approx \frac{3.11|\beta_2|}{\gamma T_{FWHM}^2} \quad (2.15)$$

First experimental observation of solitons in a single mode optical fiber with anomalous dispersion was reported by Mollenauer et al. in 1980 [14].

2.3 Pulse Propagation in a Gain Fiber

In a fiber laser, one of the essential components is the gain fiber. When an optical pulse propagates in a gain fiber, in addition to the dispersion and nonlinearity, the effects of the gain is necessary to be taken into account.

2.3.1 Gain Profile of Erbium-Doped Fiber

The gain medium is the heart of a laser. In fiber lasers, the rare-earth doped fibers are always exploited as the gain fibers which are manufactured by doping the rare-earth ions in the single mode fiber core. Different rare-earth ions are chosen to provide gain at different wavelength region, such as ytterbium for wavelength around 1.06 μm , and thulium for wavelength around 2.0 μm . The erbium-doped fibers (EDFs) among them have attracted special attentions and been widely used since they provide a broadband gain around 1.55 μm in the optical fiber communication window.

Generally, a fiber amplifier is characterized by the parameters of its small signal gain, gain saturation and gain bandwidth. For a homogeneously broadened gain medium such as EDF, the gain coefficient can be written as [74]:

$$g(\omega) = \frac{G}{1 + (\omega - \omega_a)^2 T_2^2 + P/P_s} \quad (2.16)$$

where G is the small signal gain, ω is the frequency of the incident signal, ω_a is the

resonant frequency of the dopant energy levels. T_2 is the dipole relaxation time, and its typical value is ~ 100 fs for fiber amplifiers. P is the power of the incident light, P_s is the saturation power of the rare-earth-doped fiber, which depends on dopant parameters. In the case of three-level pumping scheme, the saturation power even depends on the pumping strength.

Eq. (2.16) shows that the gain of a pulse depends on the ratio of the pulse energy over the saturation energy. The gain could even saturate with the increase of this ratio. In fact, the typical value of the saturation energy for EDF is around $10 \mu\text{J}$, which is much larger than the single pulse energy so the gain saturation effect is always negligible over the duration of a single pulse. However, in the case of a pulsed fiber laser, the pulse circulates in the laser cavity and it is the average power of the light that saturates the gain and determines the saturated gain value. In more detail, the lifetime of the excited state of erbium is ~ 10 ms, while the typical pulse interval is around $10\sim 100$ ns. The fact that thousands of pulses propagating through the gain fiber within the excited state lifetime will easily generate the strong saturation effect, which also leads to the result that the actual saturation energy for each single pulse is much lower and easily to be satisfied. Therefore, the gain saturation in fiber lasers is almost always present.

However, in the case of only considering the gain with small incident signal, the ratio P/P_s could be neglected, so Eq. (2.16) becomes:

$$g(\omega) = \frac{G}{1 + (\omega - \omega_a)^2 T_2^2} \approx G \left[1 - \frac{(\omega - \omega_a)^2}{\Omega_g^2} \right] \quad (2.17)$$

Ω_g is the gain bandwidth, which is defined as the full width at half maximum (FWHM)

of the gain spectrum $g(\omega)$. Eq. (2.17) shows that the gain exhibits a Lorentzian profile, which is the characteristic of homogeneously broadened systems. It also shows the gain is frequency dependent, and reaches its maximum value when the incident frequency coincides with the atomic transition frequency ω_a . For a Lorentzian profile, the bandwidth Ω_g is inversely related to the dipole relaxation time of the dopant T_2 . Typical absorption and emission spectra of EDFs are shown in Figure 2.6. The gain bandwidth of an EDF is generally 40 nm, which can approximately support the pulses as short as 60 fs. However, due to the effects of the nonuniform gain profile and fiber birefringence, not the whole part of the gain could be efficiently utilized in practice.

Figure 2.6: The absorption and emission spectra of the EDFs [74].

Pumping the fiber from one end will produce a non-uniform pump light, which further leads to a non-uniform gain for the pulse propagating in the fiber. Such a nonuniform population inversion can be calculated by rate equations. However, provided the EDF is short, the nonuniformity of the gain along z direction is small and can be neglected to simply the numerical model, so the gain is still considered to be unchanged along

the EDF.

2.3.2 Ginzburg-Landau Equation and the Dissipative Soliton

Solutions

As demonstrated in above section, when an optical pulse propagates in an EDF, the effect of the light amplification must be considered. Generally, the interaction of the doped ions with the light is governed by the Maxwell-Bloch equations [75]. However, in the case of that the pulse width is large compared to the dipole relaxation time T_2 of the Er^{3+} ions ($T_2 < 0.1$ ps), the rate equation approximation can be made and consequently the gain term can be simply added into the NLSE.

Considering the gain is frequency dependent, its expression can be approximately written as:

$$g(\omega) = g(\omega_0) + \frac{\partial g}{\partial \omega} \Big|_{\omega=\omega_0} (\omega - \omega_0) + \frac{1}{2} \frac{\partial^2 g}{\partial \omega^2} \Big|_{\omega=\omega_0} (\omega - \omega_0)^2 + \dots \quad (2.18)$$

where ω_0 is the central gain frequency. Without loss of the generality, assuming that the pulse carrier frequency coincides with the central gain frequency, the equation describing optical pulse propagation in the doped fibers becomes:

$$\frac{\partial A}{\partial z} + \frac{i}{2} (\beta_2 + igT_2^2) \frac{\partial^2 A}{\partial T^2} = i \left(\gamma + \frac{i}{2} \alpha_2 \right) |A|^2 A + \frac{g}{2} A \quad (2.19)$$

where g is the peak gain coefficient, T_2 is the life time of the excited states, α_2 is the two-photon absorption coefficient. On account of the gain saturation, the peak gain coefficient g is time dependent and its variation with time is governed by [74]

$$\frac{\partial g}{\partial t} = \frac{G - g}{T_1} - \frac{g|A|^2}{E_s} \quad (2.20)$$

where G is the small signal gain, and E_s is the saturation energy. T_1 is the population decay time, which is around 10 ms for erbium doped ions. Considering that the pulse width is always much narrower than it, the term T_1 could be ignored during the pulse amplification, and $g(t)$ then becomes:

$$g(t) = G \exp \left[-\frac{1}{E_s} \int_{-\infty}^t |A|^2 dt \right] \quad (2.21)$$

For EDFs, the typical value of the saturation energy E_s is about 10 μJ , which is much larger than the normal pulse energy, so the gain saturation is negligible over the duration of a single pulse. However, in the case of a pulsed fiber laser, the typical pulse interval is around 10~100 ns, which is much shorter than the life time of the excited state of the erbium ion 10 ms. Therefore, strong saturation effect could still be generated as a result of thousands of pulses passing through the gain fiber within the life time of the excited states. That also means that the saturation energy for each single pulse becomes much lower now.

Through transforming T_2 into the frequency domain, Eq. (2.19) then can be written as

$$\frac{\partial A}{\partial z} + \frac{i}{2} \left(\beta_2 + i \frac{g}{\Omega_g^2} \right) \frac{\partial^2 A}{\partial T^2} = i \left(\gamma + \frac{i}{2} \alpha_2 \right) |A|^2 A + \frac{g}{2} A \quad (2.22)$$

where Ω_g is the gain bandwidth. In reality, the fiber loss (~ 0.2 dB/km) should also be taken into account for long distance light propagation in fibers. Then Eq. (2.22) further becomes

$$\frac{\partial A}{\partial z} + \frac{i}{2} \left(\beta_2 + i \frac{g}{\Omega_g^2} \right) \frac{\partial^2 A}{\partial T^2} = i \left(\gamma + \frac{i}{2} \alpha_2 \right) |A|^2 A + \frac{1}{2} (g - \alpha) A \quad (2.23)$$

where α represents the fiber loss. When $\alpha = 0$ and $g = 0$, the equation reduces back to NLSE. Eq. (2.23) is known as the cubic Ginzburg-Landau equation (GLE). By introducing three dimensionless variables

$$\xi = \frac{z|\beta_2|}{T_0^2}, \quad \tau = \frac{T}{T_0}, \quad u = \left(\frac{\gamma T_0^2}{|\beta_2|} \right)^{1/2} A \quad (2.24)$$

Eq. (2.23) can also be written in a dimensionless form

$$i \frac{\partial u}{\partial \xi} = \frac{i}{2} \mu u + \frac{1}{2} (s + id) \frac{\partial^2 u}{\partial \tau^2} - (1 + i\mu_2) |u|^2 u \quad (2.25)$$

where $s = \text{sgn}(\beta_2) = \pm 1$, and other parameters are defined as

$$d = \frac{g T_0^2}{|\beta_2|}, \quad \mu = (g - \alpha) \frac{T_0^2}{|\beta_2|}, \quad \mu_2 = \frac{\alpha_2}{2\gamma} \quad (2.26)$$

Different from the NLSE, which is integrable and has exact solitary wave solutions, the GLE is non-integrable. For $s = -1$, i.e. when the fiber has anomalous dispersion, the general solution of Eq. (2.25) is a chirped hyper-secant profile that could be written as following [76]

$$\mu(\xi, \tau) = N_s \text{sech}(p\tau) \exp[iK_s \xi - iq \ln(\cosh p\tau)] \quad (2.27)$$

where N_s , p , q , K_s are constants depending on the parameter d . This solution of GLE can be categorized as a dissipative soliton, which is defined as a localized structure existing for an extended period of time in spite of parts of it experiencing gain and loss of energy. Obviously, different from the solitons in undoped fibers, these solitons are frequency chirped as a result of the influence of the fiber gain.

2.4 Modeling of the Soliton Fiber Laser

To gain a thorough understanding of the properties of the dissipative solitons in fiber lasers and the underlying mechanism, it is necessary for us to numerically simulate the nonlinear pulse propagating in the mode-locked fiber lasers. In this dissertation, we exploit a method called pulse tracing technique to implement this task. The key

idea and process are presented briefly in the following.

As introduced above, traditionally the nonlinear pulse propagating in a fiber laser was modeled by the GLE [77] or the master equation [78], which considered the gain, loss, and SA effect of a laser, but no laser cavity effects. Although these equations could give a good qualitative description on the soliton features of a fiber laser, they couldn't reproduce some detailed soliton features experimentally observed, such as the soliton sideband generation on the laser soliton spectrum [52]. The reason is that apart from propagating in fibers, a soliton circulating in a laser cavity is also subjected to the actions of the cavity components and feedback. To describe the soliton operation of a fiber laser more accurately, the laser cavity effects must also be included.

The pulse tracing method is an improvement of the traditional modeling techniques. The main idea of the method is that one starts the calculation with an arbitrary initial weak pulse and let the pulse circulate in a specific laser cavity. The pulse propagation in the cavity fibers is described by the coupled GLEs in the following:

$$\begin{aligned}\frac{\partial u}{\partial z} &= i\beta u - \delta \frac{\partial u}{\partial t} - \frac{ik''}{2} \frac{\partial^2 u}{\partial t^2} + \frac{ik'''}{6} \frac{\partial^3 u}{\partial t^3} + i\gamma \left(|u|^2 + \frac{2}{3} |v|^2 \right) u + \frac{i\gamma}{3} v^2 u^* + \frac{g}{2} u + \frac{g}{2\Omega_g^2} \frac{\partial^2 u}{\partial t^2} \\ \frac{\partial v}{\partial z} &= -i\beta v + \delta \frac{\partial v}{\partial t} - \frac{ik''}{2} \frac{\partial^2 v}{\partial t^2} + \frac{ik'''}{6} \frac{\partial^3 v}{\partial t^3} + i\gamma \left(|v|^2 + \frac{2}{3} |u|^2 \right) v + \frac{i\gamma}{3} u^2 v^* + \frac{g}{2} v + \frac{g}{2\Omega_g^2} \frac{\partial^2 v}{\partial t^2}\end{aligned}\quad (2.28)$$

where u and v are the normalized envelopes of the optical pulses along two orthogonal polarized modes of the optical fiber. $2\beta = 2\pi\Delta n/\lambda$ is the wave-number

difference between the two modes. $2\delta = 2\beta\lambda/2\pi c$ is the inverse group velocity difference. k'' is the second-order dispersion coefficient, k''' is the third-order dispersion coefficient and γ represents the nonlinearity of the fiber. g is the saturable gain coefficient of the fiber and Ω_g is the bandwidth of the laser gain. For undoped fibers $g = 0$, and for the erbium-doped fiber, we have also considered its gain saturation as

$$g = G \exp \left[-\frac{\int (|u|^2 + |v|^2) dt}{P_{sat}} \right] \quad (2.29)$$

where G is the small signal gain coefficient and P_{sat} is the normalized saturation energy. The standard symmetric split-step Fourier method is applied to numerically solve the coupled GLEs to get the u and v . The stability of the numerical results by selecting different step sizes has also been checked later.

Then whenever the pulse encounters a discrete cavity component, such as a SA or polarization controller (PC), the effect of the cavity component is then accounted by multiplying its Jones matrix to the optical field of the pulse. SA is the key cavity component for self-start mode-locking of the fiber laser, which provides the pulse-shortening effect due to the intensity dependent loss. Practical SAs could be classified into two sorts: the real and the artificial type, and their mathematical descriptions are different. In the following, taking SESAM and NPR as examples, the concrete modeling for these two types of SA is described.

- **Modeling of SESAM**

SESAM is a real SA. Its transmission property could be directly modeled by the following equation provided it has a finite recovery time [83]:

$$\frac{\partial l_s}{\partial t} = \frac{l_s - l_0}{T_{rec}} - \frac{P(t)}{E_{sat}^{SA}} l_s \quad (2.30)$$

where T_{rec} is the absorption recovery time, l_s is the saturation absorption of the absorber, l_0 is the initial absorption of the absorber, $P(t)$ is the temporal power of the pulse and E_{sat}^{SA} is saturation energy of the SA.

For SAs with recovery time much faster than the pulse duration, the loss can be modeled as [79]:

$$l_s = \frac{l_0}{1 + P(t)/P_{sat}} \quad (2.31)$$

where P_{sat} is the saturation power of the absorber. In fact, this model is also suitable for the new saturable absorber CNT.

- **Modeling of NPR**

Figure 2.7: A simplified model for NPR: PC, polarization controller; SMF, standard single mode fiber; EDF, erbium-doped fiber; CGLE, complex Ginzburg-Landau equation.

NPR is an artificial SA, so the model appears more complicated than that of the SESAM. Without losing its fundamental features, a NPR mode-locked fiber laser could be theoretically simplified to the model shown in Figure 2.7. It consists of three parts, the optical fiber, an intracavity polarizer, and a PC. Starting from the polarizer, any light passing through it becomes a linearly polarized one. Assume that the polarization of the polarizer is aligned with the x-axis of a laboratory reference coordinate system, which has an angle of θ with respect to the fast principal polarization axis of the fiber, the Jones matrix of the polarizer is written as

$$J_p = \begin{bmatrix} 1 & 0 \\ 0 & 0 \end{bmatrix} \quad (2.32)$$

After the polarizer, the pulse enters the cavity fiber and propagates inside it. Assume that after traversing the cavity fiber the phase shift of the light along two principal polarization directions are Φ_u and Φ_v , respectively, the Jones matrix of the cavity fiber is

$$J_F = \begin{bmatrix} e^{i\Phi_u} & 0 \\ 0 & e^{i\Phi_v} \end{bmatrix} \quad (2.33)$$

The values of Φ_u and Φ_v could be found from u and v , determined by Eq. (2.28). The desired function of the PC in the laser is to fine-tune the linear birefringence of the cavity. However, in the reality it also introduces a rotation of the fiber principal axes with respect to the laboratory reference coordinate system. In addition, due to bending and stress etc. generally the principal axes of the cavity fibers could also experience a rotation [80]. These values are experimentally difficult to measure. But one could combine them together and simply use a parameter of the net rotation of the fiber

principal axes to describe them. Therefore, the Jones matrix of the PC could be simply written as

$$J_{PC} = \begin{bmatrix} e^{i\delta} & 0 \\ 0 & 1 \end{bmatrix} \quad (2.34)$$

where δ is the linear phase shift between the two fiber polarization modes induced by the PC, known as the linear cavity phase delay bias. Combining the effects of the intracavity polarizer, the cavity fiber and the PC, the Jones matrix of the cavity is

$$J_{cavity} = J_P \cdot R(\varphi) \cdot J_{PC} \cdot J_F \cdot R(\theta) \cdot J_P \quad (2.35)$$

where φ is the angle of the fast principal polarization axis of the fiber with respect to the x axis of the laboratory reference system after the rotation of the fiber principal axes, and

$$R(\theta) = \begin{bmatrix} \cos(\theta) & \sin(\theta) \\ -\sin(\theta) & \cos(\theta) \end{bmatrix} \quad (2.36)$$

and

$$R(\varphi) = \begin{bmatrix} \cos(\varphi) & \sin(\varphi) \\ -\sin(\varphi) & \cos(\varphi) \end{bmatrix} \quad (2.37)$$

In the experiments the values of φ , θ , and δ are controlled by the PC, while the values of Φ_u and Φ_v are determined by the cavity design and the intensity of light in the cavity. Substituting all the expressions into the right hand side of the Eq. (2.35), the cavity transmission is then given by

$$T = \sin^2(\theta)\sin^2(\varphi) + \cos^2(\theta)\cos^2(\varphi) + \frac{1}{2}\sin(2\theta)\sin(2\varphi)\cos(\Phi_u + \delta - \Phi_v) \quad (2.38)$$

Based on the cavity transmission function of a fiber laser mode-locked by the SESAM or NPR technique obtained above, the result after one roundtrip of calculation will then be used as the initial pulse for the next round of calculation. The calculation is continued until a steady state is reached. The steady state will be considered as the solitary emission of the laser under the particular laser cavity parameters and operation condition. With this simulation technique not only the nonlinear pulse propagation in the cavity fibers, the laser gain and loss effects are considered, but also the influences of the cavity feedback and the features of a concrete cavity design on the pulse are included. It is found that this numerical method could faithfully reproduce all features of the soliton operations of fiber lasers experimentally observed, such as the soliton sidebands and multiple soliton generation, which confirms the reliability of our simulation technique.

2.5 Summary

In this chapter, the fundamental theory in mode-locked fiber lasers was briefly introduced. The principle of laser mode locking, and the mainly used mode-locking techniques such as the active and passive mode-locking were introduced. Three different SAs for passive mode-locking and the comparison among them were described. After mode-locking, pulses are formed in a laser cavity. Therefore, the mathematical descriptions of nonlinear pulse propagation in optical fibers and gain fibers were presented. As a result, the NLSE and its soliton solutions, the GLE and its dissipative soliton solutions were given. At the end of this chapter, the numerical modeling technique we used to simulate the dissipative soliton operation of fiber

lasers, the so-called pulse tracing technique, was also illustrated.

Chapter 3

Scalar Dissipative Solitons

Conventional solitons are formed due to the balance between anomalous fiber dispersion and self-phase modulation (SPM), whose properties are determined by the nonlinear Schrödinger equations (NLSEs), while dissipative solitons are formed as a result of the mutual nonlinear interactions among the normal cavity dispersion, SPM, effective gain bandwidth filtering, and gain saturation. Properties of the dissipative solitons are determined by the Ginzburg-Landau equations (GLEs). As a result of different soliton formation mechanism the dissipative solitons exhibit remarkably distinct properties from those of the conventional solitons.

Solitons whose dynamics are described by one-dimensional GLE are known as the scalar dissipative solitons. They have the simplest form. Experimentally solitons formed in the nonlinear polarization rotation (NPR) mode-locked fiber lasers are regarded as scalar solitons. The reason is that to implement the NPR technique, a polarizer is inserted in the laser. The polarizer fixes the polarization state of the formed solitons at the cavity position. Consequently no polarization dynamics of the solitons could be observed. NPR technique was first successfully employed to mode

lock fiber lasers in 1992. Since then solitons formed in fiber lasers mode-locked by NPR technique have been intensively investigated [81-85]. A great number of related interesting properties have been revealed and the corresponding physical mechanisms have been explained, such as Kelly sidebands [16, 86], multiple soliton formation [18, 52], soliton energy quantization [17, 52], and soliton interaction [30, 87], etc.

In this chapter, the properties and features of the scalar dissipative solitons formed in an all-normal-dispersion (ANDi) erbium-doped fiber (EDF) laser passively mode-locked by NPR technique are studied. Section 3.1 studies the fundamental and principal properties of the scalar dissipative solitons. Through analyzing the mechanism and occurrence conditions of the noise-like state in normal dispersion and multiple dissipative solitons, the guideline of optimal cavity design for achieving large pulse energy dissipative solitons is given. A 4.8 nJ dissipative soliton is obtained as a result of the cavity optimization. Section 3.2 investigates the dissipative solitons in an extended NPR mode-locked fiber laser, which are realized by replacing the intracavity polarizer with a large polarization dependent loss (PDL) cavity component. Taking advantage of the effectively reduced intracavity loss, a dissipative soliton with the highest pulse energy of 7.2 nJ is attained. Section 3.3 shows the conclusion.

3.1 Scalar Dissipative Solitons in a NPR Mode-Locked Fiber Laser

3.1.1 Properties of the Scalar Dissipative Solitons

Figure 3.1: Schematic of the ANDi fiber laser. WDM, wavelength division multiplexer; EDF, erbium-doped fiber; DCF, dispersion-compensating fiber; ISO, isolator; PC, polarization controller.

To experimentally investigate properties of the scalar dissipative solitons, an ANDi EDF laser passively mode-locked by NPR technique was constructed. Figure 3.1 shows the schematic of the cavity configuration. The ring cavity consists of 3.5-m EDF with doping concentration of 2880 ppm and group velocity dispersion (GVD) parameter of $41 \text{ ps}^2/\text{km}$, and 8.5-m dispersion-compensating fiber (DCF) with GVD parameter of $5 \text{ ps}^2/\text{km}$. The net cavity dispersion is estimated to be $\sim 0.19 \text{ ps}^2$. A polarization-independent isolator was inserted into the cavity to ensure unidirectional operation of the ring. An inline polarizer together with a polarization controller (PC) was inserted to the cavity to realize the NPR mode-locking technique. The laser was pumped by a 1480 nm high power fiber Raman laser source (BWC-FL-1480-1) through a wavelength division multiplexer (WDM), and the laser emission was output via a 10% fiber coupler. It is to emphasize that the isolator, WDM and fiber output coupler were specially made with the DCF. An optical spectrum analyzer (Ando AQ-6315B) and a 350 MHz oscilloscope (Agilent 54641A) together with a 2 GHz

photodetector were used to simultaneously monitor the output signal in the frequency domain and time domain, respectively. A commercial optical autocorrelator (FR-103MN) was used to measure the pulse width of the mode-locked pulses.

Figure 3.2: (a) Typical optical spectrum and (b) corresponding autocorrelation trace of the dissipative soliton.

Like operated in the case of anomalous cavity dispersion, with the appropriate settings of orientations of the PC, mode-locked pulses were easily obtained by simply increasing the pump power beyond the mode-locking threshold, which is ~ 150 mW in our laser. Figure 3.2 shows a typical mode-locking state observed. The optical spectrum exhibits steep spectral edges with an effective gain-bandwidth-limited

bandwidth rather than the Kelly sidebands [16, 86], which is the typical feature of the dissipative solitons formed in normal dispersion [24, 25, 34]. The autocorrelation trace shows the pulse width is about 19 ps (FWHM), assuming the soliton has a sech-form profile. Cooperating with the 3 dB spectral bandwidth of 21 nm, it gives a time-bandwidth product of 50, suggesting that the achieved dissipative soliton is a strongly chirped pulse. The oscilloscope trace further confirms that there is only one soliton pulse circulating in the laser cavity now.

Based on above dissipative soliton state, keeping the orientations of the PC unchanged, which corresponds to fixing the linear phase delay bias in theory [52], but only carefully increasing the pump strength, the intensity and edge-to-edge spectral bandwidth of the dissipative soliton were observed to be enhanced accordingly. Figure 3.3 shows the variation of the spectrum with respect to the pump power. A maximum edge-to-edge width change of ~ 4 nm was observed in our experiment. Further increasing the pump power, a noise-like state appeared.

Figure 3.3: Variation of the optical spectrum with pump power. Legend shows the corresponding pump strength.

Figure 3.4: The noise-like state: (a) Optical spectrum; (b) Oscilloscope trace; (c) Autocorrelation trace.

Figure 3.4 shows a typical noise-like state observed in our experiment. Figure 3.4(a) displays clearly that the optical spectrum no longer exhibits steep edges but still possesses a broad and smooth profile, indicating that the noise-like state is also a mode-locked state. Its oscilloscope trace shows that the mode-locked pulse still repeats at the cavity repetition rate, and only one pulse circulates in the cavity. From the autocorrelation trace in Figure 3.4(c), it is easy to identify that the mode-locked pulse profile contains complicated fine-structures, and details of the fine-structures vary randomly as the mode-locked pulse moves in the cavity, which is similar as those noise-like pulse emissions observed in the anomalous dispersion [88-91].

However, we would like to point out that the formation mechanism of a noise-like state in normal dispersion is different from that in anomalous dispersion. In anomalous dispersion regime, a generic property of the soliton propagation in the gain fiber is the explosive increase in amplitude and decrease in pulse width, which eventually results in the soliton collapse [92]. It is this soliton collapse that leads to the generation of the noise-like pulses. There is another effect, known as the cavity peak power clamping [52], existing in the cavity. It is an intrinsic feature of the soliton fiber lasers mode-locked with the NPR technique independent on the sign of the cavity dispersion, and its peak power clamping value is determined by the linear cavity phase delay bias, which corresponds to the settings of the orientations of the PC in the experiment [52]. Noise-like state in anomalous dispersion could only be formed when the soliton collapses before its peak power getting clamped, which means the linear cavity phase delay should be so biased that the peak power of the pulse only clamps at high value. Otherwise, multiple solitons would be formed instead, which will be discussed later.

Nevertheless, it was numerically found that the noise-like pulse emission in normal dispersion was caused by the peak power clamping effect [93]. When the peak power of the dissipative soliton is clamped, further increasing the laser gain will not increase the soliton energy but amplify the dispersive wave. As a result of the normal dispersion reducing the group velocity difference between the linear and nonlinear waves, the interference between the dissipative soliton and dispersive wave eventually leads to the noise-like emission. Obviously, noise-like state in normal dispersion could only be formed when the peak power of the dissipative soliton gets clamped. To avoid the noise-like pulse emission, which is harmful for the high pulse energy attainment, optimizing the operation condition to lift the peak power clamping threshold becomes a vital guideline when we design an experiment to pursue large pulse energy in normal dispersion.

It was experimentally found that with different settings of the PC, increasing the pump power could also lead to the state of multiple dissipative solitons in addition to the noise-like state. Depending on different operation conditions, such as the pump power and orientations of the PC, multifarious modes of multiple dissipative solitons were observed. The stable irregular distributed dissipative solitons, where solitons are scattered randomly and stay far apart from each other with stable relative positions, the dissipative soliton bunching, where several solitons bunch tightly and move as a whole with the fundamental repetition rate, and even the complicated dissipative soliton motion were all obtained in our experiment. Starting from the stable irregular distributed dissipative solitons, carefully tuning the orientations of the PC the dissipative solitons could even automatically arrange themselves equally spaced in the

cavity, which is known as the harmonic mode locking state [49, 83]. As an example, Figure 3.5 shows a case where 5 dissipative solitons were equally spaced in the cavity with the repetition rate of the laser emission 5 times of its fundamental value. Other harmonic mode locking states of the dissipative solitons were also obtained by simply changing the pump strength.

Figure 3.5: Oscilloscope trace of a harmonic mode-locking state of dissipative solitons. 5 solitons coexist in cavity.

Similar as the noise-like pulses, the formation mechanism of multiple dissipative solitons is also distinct from that of the conventional solitons. The generation of multiple conventional solitons in a NPR mode-locked fiber laser is found mainly ascribed to the peak power clamping effect [52]. As introduced above, the dispersive wave is amplified with the increase of the gain after the peak power of the soliton being clamped. However, different from in the normal dispersion, the linear and nonlinear waves now have very different group velocities, which results in the amplified dispersive waves quickly separating from the soliton and being reshaped

into a new dissipative soliton.

In contrast, numerical studies show that the formation of multiple dissipative solitons in normal dispersion is still a result of the peak power clamping effect, and the narrow effective gain bandwidth also plays a key role [94]. The reason is that after the peak power being clamped, a soliton with a narrower spectral bandwidth is easier to separate spatially from the dispersive waves than that with a broader bandwidth, while the spectral bandwidth of the dissipative soliton formed in normal dispersion is determined by the effective gain bandwidth. Hence, the multiple dissipative solitons could only be formed provided the effective gain bandwidth is narrow enough. Otherwise, the noise-like pulses are prone to be generated. Obviously, minimizing the effective gain bandwidth limitation could be effectively avoid the generation of multiple dissipative solitons, which is harmful for the high pulse energy attainment, and becomes another crucial guideline of the optimal cavity design to achieve large pulse energy in normal dispersion.

It was found that no matter what mode of multiple dissipative solitons was formed, all the dissipative solitons within the cavity always had exactly the same pulse height in the oscilloscope trace. This shows exactly the same feature as that of the soliton energy quantization effect observed for the conventional solitons in fiber lasers with anomalous dispersion [52]. In addition, it was also experimentally found that starting from the dissipative soliton state, the laser pump power could be reduced to a much lower value than the mode-locking threshold while the laser still maintained the soliton operation, which was known as the pump power hysteresis of the soliton operation [77]. Furthermore, provided the laser operated in the soliton operation and

the orientations of the PC were appropriately set, new dissipative solitons were even observed to be generated one by one with careful pump power enhancement. The threshold for a new soliton pulse generation was much lower than the mode-locking threshold, which indicated that the generation and annihilation of each individual dissipative soliton also exhibited pump power hysteresis. Similar phenomena have also been observed for conventional solitons in NPR fiber lasers with negative GVD dispersion previously [52].

Taking advantages of our ring cavity consisting of only elements with normal GVD, the laser could be mode locked with the net (positive) cavity GVD an order of magnitude higher than that of fiber lasers with dispersion maps, while the pulse energy is theoretically supposed to increase rapidly with increasing GVD [95]. In addition, the ANDi fiber laser is more tolerant of large nonlinear phase shifts than conventional ones. Our laser cavity is potential to achieve large pulse energy. Under appropriate settings of PC, a single pulse could be obtained in the cavity and the pulse energy increased with the pump power until certain pumping threshold. Once above threshold, the single pulse became unstable, and multiple solitons or noise-like pulses were formed in the laser. The highest average soliton output power obtained in our system was 6.2 mW with the pump strength of 500 mW, corresponding to the maximum single pulse energy of 0.4 nJ, which is significantly larger than the single-pulse energy obtained in the same type of fiber laser operating in the anomalous dispersion regime.

3.1.2 Optimization of the Dissipative Soliton Operation

It is known that in a normal dispersion amplifier the pulses could propagate in a wave-breaking free manner but accumulate a large nonlinear phase shift [96]. When they enter the following DCF, such a large nonlinear phase shift leads to the occurrence of optical wave-breaking and hence the instability of the soliton pulses, which should be avoided. The output coupler could reduce the peak power of the pulses, so it was placed immediately after the gain fiber, which significantly lowers the peak power of the pulses entering the DCF after the gain fiber, and effectively suppresses the wave-breaking. Limited by the practical splicing requirement, the length of the DCF between the EDF and the coupler was finally kept to be ~ 20 cm in our experiment. Such a position of the output coupler is also in favor of high output energy attainment. The output-coupling ratio is also a crucial parameter. In principle, a large output-coupling ratio is favorable for a higher pulse energy collected and smaller intracavity nonlinearity. However, it is against the realization of the mode-locking of a laser. After attempting different output coupling ratio, the 90% output turns out to be the optimal value and is eventually adopted in our experiment.

Based on the previous discussion, minimizing the effective gain bandwidth limitation could effectively suppress the generation of multiple dissipative solitons and hence is another crucial guideline of the cavity design optimization. Considering that pulse energy is theoretically expected to increase rapidly with increasing GVD, long cavity with large normal dispersion is favorable for achieving large pulse energy. However, it should also be noticed that with the increase of the cavity length, the cavity birefringence is enhanced accordingly. The bandwidth of the cavity birefringent filter

becomes smaller, which is equivalent to narrowing the effective gain bandwidth and should be avoided in the cavity design. Therefore, the adoption of the optimal cavity length is a trade-off between above two processes. After various attempts, it was found that the cavity length of 15 m performed well with other parameters in current settings, and was finally adopted for our experiment.

After changing the output coupling ratio into 90%, cutting the DCF pigtail between EDF and coupler to 20 cm, and adding a length of 3 m DCF between the PC and WDM, a maximum single pulse energy of 4.8 nJ was achieved, which was improved almost an order of magnitude than that before optimization (~ 0.4 nJ).

3.2 Scalar Dissipative Solitons in an Extended NPR Mode-Locked Fiber Laser and Large Energy Pulses

Conventionally, in order to implement the NPR technique in a ring fiber laser, a polarizer is inserted in the laser cavity, which in combination with the nonlinear fiber cavity forms a nonlinear birefringent filter in the laser. Under appropriate cavity linear phase delay bias, the transmission of the nonlinear birefringent filter increases with the light intensity, forming an artificial saturable absorption effect in the cavity that further results in the mode locking of a fiber laser. So far almost every NPR mode-locked fiber laser has a polarizer in the cavity. It gives the impression that the intracavity polarizer is a characteristic of the NPR mode locking. In this section, we show that under strong pumping, even without the existence of an intracavity polarizer, the NPR mode locking could still occur in a fiber laser due to the relatively large PDL of a cavity component. It is pointed out that without a polarizer in the

cavity could effectively reduce the intracavity loss so that dissipative solitons with much higher pulse energy are attainable.

3.2.1 Experimental Setup and Results

We still use the setup as shown in Figure 3.1. The only difference is that the intracavity polarizer was replaced by a piece of DCF. Hence, the cavity is still purely made of fibers with normal dispersion and the cavity length is 12 m. It is to emphasize that different from the conventional fiber lasers mode locked with the NPR technique, no polarizer was inserted in the cavity.

In spite of the fact that no polarizer was inserted in the cavity and no other mode locking technique was adopted, it was observed that under strong pumping mode locking could still self-start in the laser. Figure 3.6 shows a typical mode-locked state observed with the pump power of 272 mW. Figure 3.6(a) displays its optical spectrum. The spectrum also exhibits sharp edges with a gain-bandwidth limited bandwidth, which is the characteristic of the dissipative soliton in normal dispersion. Figure 3.6(b) is the corresponding autocorrelation trace, which shows the mode-locked pulse width is 21.9 ps (FWHM), if a sech^2 -form pulse profile is assumed. Figure 3.7(a) shows the mode-locked pulse train. It reveals that the pulses repeat with a fundamental cavity repetition rate of 16.7 MHz.

Figure 3.6: The dissipative soliton state experimentally observed: (a) optical spectrum; (b) the corresponding autocorrelation trace.

Obviously, due to the nonlinear propagation of the mode-locked pulses in the cavity, they have been shaped to the dissipative solitons. To find out the underlying mode-locking mechanism, the features of these dissipative solitons were further investigated. The polarization features of the solitons obtained were studied first. Figure 3.7 shows the comparison of the laser output before and after passing through an external cavity polarizer. It reveals that the polarizations of the solitons are locked, since after the polarizer all the pulses in the trace measured have equal intensity, as shown in Figure 3.7(b). It was experimentally found that no matter how the laser

operating conditions, such as the pump strength and the orientations of the PC, were changed, no polarization-rotating vector solitons could be observed. The observed polarization features of the solitons are exactly the same as those of the NPR mode-locked fiber lasers, where due to the existence of a polarizer in the cavity, the polarization of the output solitons cannot freely evolve but is fixed by the cavity boundary condition. Detailed experimental studies further reveal that other features of the dissipative soliton operation of the laser also have no obvious difference from those of a NPR mode-locked fiber laser, e.g. states like the noise-like pulse mode locking, multiple dissipative solitons and even the period-doubling of the soliton [97] could also be obtained. On account of the experimental results, we then measured the PDL of all the cavity components. It is found that the PDL value of the WDM used has a large value of about 3.8 dB at the central wavelength 1584 nm of the mode-locking state. The PDL of the used WDM is much larger than that of a common WDM, which is around 1.1 dB at 1584 nm wavelength. To confirm that the observed mode locking of the laser was caused by the large PDL of the WDM, replacing it with a low PDL one, it turned out that no self-started mode locking was obtained under the maximum available pump strength. The experimental result shows that a large PDL cavity component is sufficient to induce NPR mode locking. Here, the PDL value employed to achieve NPR mode-locking in our cavity seems to be slightly large, which is probably due to the fact that in the normal dispersion regime, it is more difficult to achieve the mode-locking state than in the anomalous dispersion regime.

Figure 3.7: The oscilloscope traces (a) before the external cavity polarizer; (b) after the external cavity polarizer.

It is important to note that the above mode-locking was only obtained in our laser under a high pumping strength. Under low pumping power, no matter how the PC is tuned, no mode-locking could be observed. The mode-locking threshold of the current laser was about 251 mW. For the purpose of comparison, a polarizer, which has a PDL of 30 dB, was inserted into the fiber laser with the same cavity parameters. It was found that the mode-locking threshold was reduced to 126 mW, which is much lower than that of the PDL induced mode locking. The loss of many practical passive optical components varies with the polarization state of light. Generally, a polarization discrimination component, such as a polarizer, has a very large PDL value, at least above 20 dB, while a polarization-independent component, such as a

WDM, or a fiber coupler, has a minimum PDL value, normally <0.15 dB. However, occasionally due to the technical limitations and/or other factors, a conventionally polarization independent component could possess a relatively large PDL value. Based on our above experimental result, special care needs to be taken when using these components in the laser cavity to avoid complications caused.

Finally, we point out that the PDL induced laser mode locking is independent on the cavity dispersion. A fiber laser with the WDM unchanged but the rest of the cavity fibers replaced by the anomalous dispersion fibers was constructed further. Self-started mode locking and soliton operation of the laser could be easily obtained. An example is shown in Figure 3.8. Figure 3.8(a) shows the soliton spectrum obtained in the laser, which has clear spectral sidebands, a characteristic of the soliton operation in anomalous dispersion cavity fiber lasers. Figure 3.8(b) is the corresponding autocorrelation trace, which gives the pulse width of the soliton about 972 fs (FWHM). Figure 3.8(c) is the pulse train of the soliton. Again the soliton operation features of the laser have no difference from those of a NPR mode-locked laser with a polarizer in cavity. It comes to our attention that, the PDL value of the WDM at the central wavelength 1563 nm of the current mode-locking state is around 1.4 dB, which is quite a low value but is sufficient to cause the NPR mode-locking.

Figure 3.8: The conventional soliton state experimentally observed: (a) optical spectrum; (b) the corresponding autocorrelation trace; (c) the oscilloscope trace.

3.2.2 Numerical Simulations

To confirm the PDL induced mode-locking in the fiber lasers, numerical simulations have also been conducted on the operation of the lasers by employing the pulse tracing method as demonstrated in Chapter 2.

We follow the process as described in section 2.4 to conduct the simulation. Note that the current cavity employs no polarizer but a component with large PDL (WDM in our experiment), whose function is like an imperfect polarizer. Therefore, in the numerical simulation the only difference is that a partial polarizer is adopted to replace the perfect polarizer. The Jones matrix of a polarizer in conventional NPR model Eq. (2.32) was replaced by the Jones matrix of a partial polarizer in the following

$$J_{PP} = \begin{bmatrix} p_1 & 0 \\ 0 & p_2 \end{bmatrix} \quad (3.1)$$

where p_1 and p_2 denote respectively the transmissions in two orthogonal principal polarization directions of the partial polarizer, whose values are determined by the PDL value of the intracavity component.

To make the simulation comparable to the experiment, the following cavity parameters were adopted for the calculation: For the all normal dispersion cavity, $\gamma = 3 \text{ W}^{-1}\text{km}^{-1}$, $k''_{\text{EDF}} = 41 \text{ ps}^2/\text{km}$, $k''_{\text{DCF}} = 5 \text{ ps}^2/\text{km}$, $k''' = -0.13 \text{ ps}^3/\text{km}$, gain saturation energy $P_{\text{sat}} = 100 \text{ pJ}$, gain bandwidth $\Omega_g = 12 \text{ nm}$, the principal polarization of the

intracavity partial polarizer to the fast axis of the birefringent fiber $\theta = 0.152\pi$, and $\varphi - \theta = \pi/2$, $L = 8.5_{\text{DCF}} + 3.5_{\text{EDF}} = 12$ m, and birefringence length of the fibers $L_b = L$. When p_1 is set to be one, and p_2 to be zero, a polarizer is contained in our cavity, and the situation is exactly the same as that of the conventional NPR mode-locking. Soliton operation could be easily achieved as demonstrated before by changing the gain and cavity linear phase delay bias, which corresponds to varying the pump strength and orientations of the PC in the experiment. p_1 was then reduced from 1 to $\sqrt{0.5}$ but kept $p_1^2 + p_2^2 = 1$ satisfied, which denotes the decrease of the PDL value. It was found that even with a small PDL value as low as 3.7 dB, corresponding to $p_1 = \sqrt{0.7}$, $p_2 = \sqrt{0.3}$, soliton operation could still be obtained in the laser provided the laser gain is strong enough. Figure 3.9 shows a typical soliton state numerically calculated. To obtain the state, the cavity linear phase delay bias was set as $\delta = 1.5\pi$, and a large pump strength of $G = 12000 \text{ km}^{-1}$ was used. Figure 3.9(a) is the optical spectrum of the pulse. It has steep spectrum edges with a gain bandwidth limited spectral width. Figure 3.9(b), (c) show the pulse profile and evolution with the cavity round-trips. The fact that the mode-locked pulse keeps its shape unchanged with different cavity round-trips demonstrates that the soliton state is stable. Once the dissipative soliton state is obtained, further increasing the gain, either multiple dissipative solitons or noise-like state could be achieved. However, the smaller p_1 is reduced to, the more difficult is it to achieve mode-locking, and the larger is the gain required to realize the mode-locking, which is qualitatively in agreement with the experimental observation.

Figure 3.9: The dissipative soliton state numerically calculated. Transmissions in two orthogonal principal polarization directions of the partial polarizer $p_1 = \sqrt{0.7}$, $p_2 = \sqrt{0.3}$, pump strength $G = 12000 \text{ km}^{-1}$. (a) the soliton spectrum; (b) pulse profile; (c) pulse evolution with the cavity round-trips.

Figure 3.10: The conventional soliton state numerically calculated. Transmissions in two orthogonal principal polarization directions of the partial polarizer $p_1 = \sqrt{0.6}$, $p_2 = \sqrt{0.4}$, pump strength $G = 2000 \text{ km}^{-1}$: (a) the soliton spectrum; (b) pulse profile; (c) pulse evolution with the cavity round-trips.

For the fiber laser with the anomalous dispersion cavity, the following parameters were adopted for the simulation: $\gamma = 3 \text{ W}^{-1}\text{km}^{-1}$, $k''_{\text{EDF}} = -13 \text{ ps}^2/\text{km}$, $k''_{\text{SMF}} = -23 \text{ ps}^2/\text{km}$, $k''_{\text{DCF}} = 5 \text{ ps}^2/\text{km}$, $k''' = -0.13 \text{ ps}^3/\text{km}$, gain saturation energy $P_{\text{sat}} = 100 \text{ pJ}$, gain bandwidth $\Omega_g = 20 \text{ nm}$, the principal polarization of the intracavity partial polarizer to the fast axis of the birefringent fiber $\theta = 0.125\pi$, and $\varphi - \theta = \pi/2$, $L = 8_{\text{SMF}} + 1_{\text{EDF}} + 3_{\text{DCF}} = 12 \text{ m}$, and birefringence length of the fibers $L_b = L$. Mode-locked state could still be achieved when $p_1 = \sqrt{0.6}$, $p_2 = \sqrt{0.4}$, which corresponds to the PDL value of 1.8 dB. Figure 3.10 shows the spectrum, corresponding pulse profile and its time evolution with the cavity roundtrips, calculated at $G = 2000 \text{ km}^{-1}$, and $\delta = 0.4\pi$. The spectrum exhibits clear soliton sidebands and the pulse possesses narrower pulse width, displaying the characteristics of the soliton formed in an anomalous dispersion cavity.

Considering that replacing a polarizer by a large PDL component to realize the NPR mode-locking in the fiber laser could reduce the cavity loss effectively, it should be helpful to enhance the pulse energy attainable. Therefore, further optimize our extended NPR mode-locked fiber laser following the similar procedure and guidelines given in section 3.1.2. Under current situation, the optimal output coupling ratio turns out to be 50% and the best cavity length is $\sim 20 \text{ m}$. After changing the coupler and adding a length of 7 m DCF between PC and WDM, the highest single pulse energy of 7.2 nJ was finally achieved with the pump power of 573 mW and appropriate settings of the orientations of the PC.

3.3 Conclusion

In this chapter, formation and properties of the scalar dissipative solitons formed in ANDi EDF lasers mode-locked by NPR technique were studied systematically. It was experimentally found that in spite of different pulse shaping mechanism from conventional solitons, multiple solitons, soliton energy quantization, noise-like state and the pump hysteresis could still be observed for the dissipative solitons, but the underlying mechanisms are distinct. Through analyzing the principle and occurrence conditions of the noise-like state and multiple dissipative solitons in normal dispersion, lifting the peak power clamping threshold and minimizing the effective gain bandwidth limitation were proposed to be the guideline of optimizing cavity design to achieve high pulse energy dissipative solitons. A procedure of the cavity design optimization was also provided, and after that a large pulse energy dissipative soliton with 4.8 nJ was finally achieved, benefiting from the dissipative soliton features.

Dissipative solitons were even observed in a fiber laser without a polarizer and any other real SA. It was revealed that the underlying mechanism was that under strong pumping, even without the existence of an intracavity polarizer, NPR mode locking could still occur in a fiber laser due to the relatively large PDL of a cavity component, which further lead to the formation of dissipative solitons. Numerical simulations verified the experimental results. Such an extended NPR mode-locking could effectively reduce the intracavity loss. Benefit from this, a dissipative soliton with the highest pulse energy of 7.2 nJ was eventually available in our experiment.

Chapter 4

Dissipative Solitons in Dispersion-Managed Fiber Laser with Different Cavity Dispersion

In addition to be formed in all-normal-dispersion (ANDi) cavity fiber lasers, dissipative solitons could also be formed in the dispersion-managed (DM) cavity fiber lasers [26, 98-101]. A striking feature of the DM fiber lasers is that in terms of the dispersion map, its net cavity dispersion could be set to different values. It is to note that soliton formation in fiber lasers with large anomalous cavity dispersion has been extensively investigated previously [102-105]. Solitons formed in the fiber lasers are characterized by the appearance of symmetric discrete spectral sidebands, also known as the Kelly sidebands whose positions are determined by the cavity dispersion and pulse duration [86], on their soliton spectra. The solitons also have narrow pulse width and little frequency chirp. Soliton formation in fiber lasers with large normal cavity dispersion has also been experimentally observed recently [26, 106]. Zhao et al. and Chong et al. showed that these solitons had different features from those formed in fiber lasers with large anomalous dispersion. The optical spectra of the solitons are

characterized by steep spectral edges without the Kelly sidebands. In addition, the soliton pulses have much broader pulse width and large frequency chirp.

For the DM fiber lasers, cavity with near zero dispersion could be easily realized. Tamura et al. first experimentally investigated operation of such a fiber laser and demonstrated a state of the so-called stretched-pulse emission [19]. As shown by Tamura et al., the stretched pulses could have large pulse energy and very broad optical spectra with complicated spectral profiles. A stretched pulse could even be compressed significantly. However, the autocorrelation traces of stretched pulses could have clear pedestals, which are different from those of the solitons formed in the large anomalous or large normal cavity dispersion regimes. Later, it was shown that the origin of these differences could be traced back to the competition between two soliton shaping mechanisms in the fiber lasers [107].

An interesting feature of the DM fiber lasers is that the net dispersion of the cavity varies with the peak intensity of the formed solitons. J. H. B. Nijhof et al. first numerically studied the phenomenon [108]. They have numerically shown that when pulses propagate in a DM cavity, interaction between the group velocity dispersion (GVD) and nonlinearity makes them possessing different spectral bandwidths in the normal and anomalous dispersion fiber segments, narrower in the former one and broader in the latter one, which leads to the result that the effective dispersion of the system shifts to the anomalous side.

In this chapter, we study the dissipative solitons formed in DM fiber lasers with different cavity dispersions. Section 4.1 studies the dissipative solitons in a DM fiber

laser with large normal cavity dispersion. It shows that the stretched pulses, which are normally observed at around zero cavity dispersion, could also be observed in the cavity. Section 4.2 investigates the dissipative solitons in a DM fiber laser with around zero cavity dispersion. It demonstrates that either the dissipative solitons with Kelly sidebands or steep spectral edges could be observed in the DM laser. Further numerical simulation reveals that the experimental observations are a direct result of the effective dispersion shift effect induced by the nonlinearity in a DM fiber laser. Section 4.3 shows the conclusion.

4.1 Dissipative Solitons formed in a DM Fiber Laser with Large Normal Cavity Dispersion

Previous studies on the dissipative soliton operation of the DM fiber lasers have revealed that there is a cavity dispersion regime where both the stretched pulses and the solitons with steep spectral edges could be observed [107]. To our knowledge, the physical mechanism behind the effect was unexplained. In this section, a DM fiber laser with large normal cavity dispersion was constructed. The effective dispersion shift and features of the dissipative solitons formed in such a DM cavity fiber laser have been experimentally investigated. We show that under strong pumping the stretched-pulse operation could still be obtained in a DM fiber laser with fairly large positive net cavity dispersion. However, under the same cavity dispersion a non-DM fiber laser can only generate solitons with steep spectral edges. Numerical simulations based on a model of coupled nonlinear Ginzburg-Landau equations (GLEs) have confirmed the result. It further showed that in addition to the strong pumping, the cavity linear phase delay must be so biased that practically the peak power of the

formed pulse has not been clamped. The results suggest that the effective cavity dispersion of a DM fiber laser could significantly shift with the cavity nonlinearity. The stronger the cavity nonlinearity, the larger the effective cavity dispersion shift is.

4.1.1 Experimental Setup and Observations

The fiber laser used in the experiment is schematically shown in Figure 4.1. It has a ring cavity configuration and the total cavity length is around 8.43 m. A piece of 5-m erbium-doped fiber (EDF) with a doping concentration of 2880 ppm was used as the gain fiber. The EDF has a GVD parameter of about $42 \text{ ps}^2/\text{km}$. The rest of the ring was made of the standard single mode fiber (SMF), whose GVD parameter is about $-24 \text{ ps}^2/\text{km}$. The NPR mode locking technique was used to achieve the self-started mode-locking of the laser. To this end, a polarizer together with two polarization controllers (PCs), one consisting of two quarter-wave plates and the other of two quarter-wave plates and one half-wave plate, was employed for adjusting the polarization of light in the cavity. Note that the cavity is a ring, so the polarizer here plays not only the role of a polarizer, which transforms a light with arbitrary polarization state into a linearly polarized one, but also the role of an analyzer, which introduces an intensity-dependent cavity transmission, and furthermore under suitable linear cavity birefringence bias, an artificial saturable absorption effect into the laser. Under influence of the artificial saturable absorption effect the laser could become passively mode locked. Besides, a polarization-independent isolator was also inserted into the cavity to enforce the unidirectional operation of the ring. The wave-plates, polarizer and isolator were mounted on a fiber bench, which couples light from one fiber to another across a 7-cm open-air gap, and has an insertion loss (excluding the

bulk-optic components) of about 1.2 dB. In our experiment, the laser was pumped by a 1480 nm high power Raman fiber laser (KPS-BT2-RFL-1480-60-FA) through a 1480/1550 wavelength division multiplexer (WDM), and the laser light was coupled out via a 90% fiber coupler.

Figure 4.1: Schematic of the DM soliton fiber laser. PI, polarization independent isolator; P, polarizer; $\lambda/4$, quarter-wave plate; $\lambda/2$, half-wave plate; WDM: wavelength division multiplexer; EDF, erbium-doped fiber.

The laser cavity has a typical DM configuration. The net cavity dispersion was deliberately set to be a fairly large positive value of 0.128 ps^2 at 1550 nm by controlling the length of the SMF in the cavity. The length of the SMF of the output coupler outside the cavity was cut very short to eliminate its influence on the results. Self-started mode locking could be established in the laser with a launched pump power of 501 mW. As reported previously on the large positive dispersion DM fiber lasers, dissipative solitons were always obtained when the laser was operated not far away above the mode-locking threshold. In addition, the dissipative soliton operation also exhibited the pump hysteresis. Namely, once the mode locking was achieved, the pump power could be reduced to a much lower value while the laser still remained

mode locking. In the current experiment, as a powerful pump source was used, the laser could be operated at a pump power level that was several times above the laser threshold.

Figure 4.2: The two typical mode-locking operation states experimentally observed. (a) The stretched-pulse mode-locking operation. (b) The DM dissipative soliton operation.

It came to our attention that under strong pumping and depending on the orientations of the waveplates, the stretched-pulse mode-locking state, which was normally

obtained in the DM fiber lasers with near zero cavity dispersion, could also be obtained in the laser. Figure 4.2(a) shows a typical stretched-pulse state experimentally measured. The optical spectrum of the pulse has a broad bandwidth with a complicated structure, which is a typical characteristic of the stretched pulses [19, 21, 109]. The FWHM of the pulse is about 146 fs, if a sech^2 -form pulse profile is assumed. The autocorrelation trace of the pulse has a clear pedestal, indicating the existence of some weak uncompressible components in the pulse. The measured single stretched-pulse energy was about 2.3 nJ, which gives a high peak power of 15.7 kW.

Fixing all the other laser operation conditions but adjusting the orientations of the waveplates, which corresponds to changing the linear phase delay bias of the cavity, the dissipative solitons could also be obtained in the laser, as shown in Figure 4.2(b). Comparing with the stretched pulse spectrum shown in Figure 4.2(a), the spectrum exhibits sharp leading and trailing edges with a 23 nm 3dB-bandwidth (measured from the center of the spectrum). Assuming that the solitary wave has a sech^2 -form profile, the autocorrelation trace gives a pulse width of about 5 ps (FWHM), which is about 33 times broader than that of the stretched pulse obtained. The time-bandwidth product of the dissipative soliton is 15, indicating a large frequency chirp. Under strong pumping, multiple dissipative solitons were always obtained in the laser. However, through carefully decreasing the pump strength, single dissipative soliton operation could also be achieved. The single pulse energy of the dissipative soliton measured was about 2.16 nJ, corresponding to a peak power of 432 W.

The net cavity dispersion of current laser is larger than that of the dissipative soliton fiber lasers reported by Zhao et al., and much larger than that of the stretched-pulse fiber lasers of Tamura et al., but both of the dissipative soliton and stretched pulse were obtained. To find out the underlying physical mechanism, a fiber laser with all normal-dispersion fibers was further set up and operated under a variety of positive cavity dispersions and pump strengths. With the net cavity dispersion even smaller than 0.128 ps^2 , and pump strength even larger than 2.5 W , no stretched-pulse operation could be obtained, suggesting that the net cavity dispersion value is not a critical factor on the stretched pulse formation in current laser. It should be a special feature related to the DM of the laser cavity.

4.1.2 Numerical Simulations

To understand the experimental results and determine the underlying physical mechanism, numerical simulation of the operation of this laser has been further conducted by using the pulse tracing method. In accordance with the experimental laser configuration, the following fiber and cavity parameters were used in the simulation: $\gamma = 3 \text{ W}^{-1}\text{km}^{-1}$, $k''_{\text{SMF}} = -24 \text{ ps}^2/\text{km}$, $k''_{\text{EDF}} = 42 \text{ ps}^2/\text{km}$, $k''' = -0.13 \text{ ps}^3/\text{km}$, gain saturation energy $P_{\text{sat}} = 1 \text{ nJ}$, gain bandwidth $\Omega_g = 15 \text{ nm}$, the polarization of the intracavity polarizer to the fast axis of the birefringent fiber $\theta = 0.152\pi$, and $\varphi - \theta = \pi/2$, $L = 1.63 L_{\text{SMF}} + 5 L_{\text{EDF}} + 1.8 L_{\text{SMF}} = 8.43 \text{ m}$, and birefringence length of the fibers $L_b = L/4$. In current simulation, the experimentally observed soliton operation features of the DM fiber lasers have been well reproduced again. Namely, it was numerically observed that if the net cavity dispersion of a DM fiber laser was large negative, solitons with spectrum characterized by Kelly sidebands would be obtained; when the net cavity

dispersion was large positive, solitons with spectrum characterized by steep edges were generally obtained; while near the zero cavity dispersion, the DM solitons or the stretched pulses under strong pumping were obtained [107].

Figure 4.3: The two typical mode-locking operation states of the laser numerically calculated with large positive net cavity GVD. (a) The stretched-pulse mode-locking operation. Cavity linear phase delay bias $\delta=1.9\pi$, $G=3000 \text{ km}^{-1}$. (b) The DM dissipative soliton operation. Cavity linear phase delay bias $\delta=1.5\pi$, $G=1100 \text{ km}^{-1}$.

To compare with the experimental results, we have focused on the operation of the DM fiber lasers with large net cavity dispersion. Numerically, it was found that even under the net positive cavity dispersion as large as that of our laser, provided that the laser gain is strong enough and the formed pulse peak is unclamped by the laser cavity, the stretched-pulse operation of the laser could still be obtained. Figure 4.3(a) shows a typical stretched pulse numerically calculated. To obtain the state, the cavity linear phase delay bias was set as $\delta = 1.9\pi$, and a large pump strength of $G = 3000 \text{ km}^{-1}$ was used. At current laser parameters setting, the cavity linear phase delay bias corresponds to the fact that the laser cavity has a very large peak power clamping value [52]. Therefore, the peak of the pulse is practically unlimited. In this case a narrow pulse with high peak power is formed. The nonlinear propagation of the pulse in the negative dispersion SMF causes strong spectral broadening.

Under the same pump strength but the cavity linear phase delay bias is decreased to $\delta = 1.5\pi$, either the noise-like pulses or multiple dissipative solitons were obtained. However, carefully decreasing the pump strength single dissipative soliton could always be obtained. Figure 4.3(b) shows the optical spectrum and the corresponding pulse profile of a calculated dissipative soliton at $G = 1100 \text{ km}^{-1}$. It has steep spectral edges and a gain bandwidth limited edge-to-edge width. Under the dissipative soliton operation the pulse is significantly broader, which is in agreement with the experimental observation.

Both of the experimental and numerical results show that the formation of the stretched pulse in the laser is closely related to the high peak pulse formation and its nonlinear propagation in the cavity, which reminds us the relationship of the

phenomenon to the effective dispersion shift caused by the nonlinear pulse propagation in the DM system. Hence, it was speculated that under strong cavity nonlinearity, the effective cavity dispersion of the fiber laser could shift to the negative side, and eventually the natural soliton shaping mechanism still occurs, leading to the formation of the stretched pulses.

Different from an ANDi fiber laser, the effective cavity dispersion of a DM fiber laser depends on the concrete nonlinear pulse propagation in the cavity. For the fiber lasers passively mode-locked with the NPR technique, there is an intrinsic cavity pulse peak clamping effect [52]. As the cavity pulse peak clamping power is further determined by the cavity linear phase delay bias setting, and the cavity clamping is a discrete effect in the real laser, depending on the laser cavity design and operation conditions, the formed pulse could have different effective peak intensity. Generally, when the cavity linear phase delay bias δ is set close to the peak clamping point or the laser gain G is chosen small, the peak power of the stable pulse is low. When the pulse propagation in the cavity has only low nonlinearity, the effective cavity dispersion is almost the same as the set cavity dispersion. On the other hand, increasing the pump strength G and lifting the cavity pulse peak clamping level, a pulse with high peak power would be generated. Accordingly, as it propagates in the cavity, it generates strong nonlinearity. A large nonlinearity induces a large shift in the cavity dispersion towards the anomalous dispersion. When the nonlinearity becomes strong enough, the effective cavity dispersion could even become net negative, and eventually the stretched pulse is formed in the laser. Obviously, in a DM fiber laser there is no clear-cut cavity dispersion boundary between the dissipative soliton and the stretched pulse soliton. Instead depending on the concrete cavity nonlinearity, over a large

cavity dispersion range both types of pulses could be obtained in the same laser.

4.2 Dissipative Solitons in a DM Fiber Laser with around Zero Cavity Dispersion

We further shift the net cavity dispersion of a DM fiber laser to around zero and investigate experimentally the effects of effective cavity dispersion shift on the features of the dissipative solitons. Surprisingly, either the dissipative solitons with steep spectral edges, which are normally obtained in large normal cavity dispersion regime, or dissipative solitons with clear Kelly sidebands, which are normally observed in large anomalous cavity dispersion regime, are experimentally obtained in the same cavity dispersion value. The operation of the laser can be switched between the two types of dissipative solitons by simply tuning the orientations of the intracavity PC. The physical mechanism of the phenomenon can again be traced back to the effective cavity dispersion shift of a DM fiber laser.

4.2.1 Experimental Observations

The fiber laser used in our experiment is schematically shown in Figure 4.4. The total cavity length is around 11 m. A piece of 1.7 m EDF with a doping concentration of 2880 ppm was used as the gain fiber. The EDF has a GVD parameter of about 42 ps²/km. The rest of the cavity consists of the SMF with a GVD coefficient of ~ -24 ps²/km, and the DCF with a GVD coefficient of ~ 8 ps²/km. The self-started mode-locking of the laser is achieved by employing a single-wall carbon nanotube (SWCNT) saturable absorber (SA). Such a SWCNT SA is simply constructed by

sandwiching a SWCNT thin film between two fiber-pigtailed FC/PC connectors, with index-matching oil filled in. A polarization-independent isolator was inserted into the cavity to enforce the unidirectional operation of the ring. Besides, an inline PC was also exploited in our cavity to fine tune the cavity birefringence and optimize the mode-locking condition. The laser was pumped by a 1480 nm high power Raman fiber laser (P48LD-20) with a maximum output of 220 mW through a 1480/1550 WDM, and the backward pump scheme was exploited to eliminate the influence of the residual pump power. The laser light was coupled out via a 10% fiber coupler, and the SMF pigtail of the coupler outside the cavity was cut very short to eliminate its influence on the laser emission.

Figure 4.4: Schematic of the fiber laser mode-locked by SWCNTs. WDM, wavelength division multiplexer; EDF, erbium-doped fiber; ISO, isolator; PC, polarization controller; SWCNTs, single-wall carbon nanotubes.

Figure 4.5: The typical dissipative soliton state in normal dispersion experimentally observed: (a) Optical spectrum; (b) autocorrelation trace; (c) pulse train.

The laser cavity has a typical DM configuration, and the net cavity dispersion was deliberately set to be positive but near zero at 1550 nm by controlling the length of the SMF in the cavity. Self-started mode locking was easily achieved by increasing the pump power beyond the mode-locking threshold, which is as low as 3.5 mW for our laser. The mode-locked pulses were further shaped into optical solitons. The low mode-locking threshold and single pulse energy of the formed soliton is a direct consequence of the fact that our net cavity dispersion is around zero. Figure 4.5 shows such a soliton state observed experimentally. Figure 4.5(a) is the optical spectrum. It exhibits sharp leading and trailing edges with a gain-bandwidth-limited bandwidth, which are the characteristics of the spectra of the dissipative solitons formed in normal dispersion fiber lasers. The corresponding autocorrelation trace is shown in Figure 4.5(b). The FWHM of the soliton pulse is measured to be about 3.6 ps, if a sech^2 -form pulse profile is assumed. Cooperating with the 3dB bandwidth of 12 nm, the time-bandwidth product is 5.4, suggesting a large positive frequency chirp. Figure 4.5(c) further shows the soliton pulse train. The pulses repeat with a repetition rate of 18.2 MHz as determined by the cavity length, denoting that only a single pulse circulates in the laser cavity now.

The solitons achieved here exhibited all the properties of the dissipative solitons observed in normal dispersion fiber lasers previously [25, 26, 34]. Carefully increasing the pump strength, the edge-to-edge bandwidth of the dissipative soliton broadened accordingly until reaching a certain critical value. Then further increasing the pump will not broaden the bandwidth but lead to the pulse breaking, and multiple solitons were formed in the cavity as a result. A maximum edge-to-edge bandwidth change of ~ 5 nm was achieved in the experiment. Besides, the energy quantization

[17] and pump hysteresis [18] phenomena were also observed. Moreover, the central wavelength of the optical spectrum of the dissipative soliton was observed to be slightly tunable depending on the settings of the orientations of the PC, shifting from 1558.7 nm to 1560.2 nm.

Figure 4.6: The typical soliton state in anomalous dispersion experimentally observed: (a) Optical spectrum; (b) autocorrelation trace; (c) pulse train.

However, in addition to above solitons, another distinct solitary state was also experimentally observed by simply tuning the orientations of the PC. Figure 4.6 shows such a soliton state. Rather than the steep spectral edges shown in Figure 4.5(a), the soliton spectrum is characterized by the appearance of clear Kelly sidebands, which are the typical feature of the solitons formed in anomalous dispersion fiber lasers. The 3-dB spectral bandwidth is around 6 nm. The corresponding autocorrelation trace in Figure 4.6(b) shows obviously that the soliton has a much narrower pulse width than the former one, possessing a FWHM of 710 fs. The time-bandwidth product is around 0.53, indicating that the pulse has some frequency chirp. Figure 4.6(c) shows the stable pulse train of such a soliton state.

The solitons exhibit all the properties of the solitons observed in the anomalous cavity dispersion fiber lasers previously, such as the soliton energy quantization, pump hysteresis, and multiple soliton operation under strong pump power [52]. The central wavelength of the soliton spectrum was tunable from 1557.4 nm to 1562.8 nm with the variation of the orientations of the PC. However, it is noticed that when the soliton spectrum has a relative shorter central wavelength, the spectrum is symmetric and displays clear Kelly sidebands on both sides. When the soliton spectrum was tuned to a longer central wavelength, it appears to be asymmetric, which captured our attention. To be more specific, the spectrum only exhibits the Kelly sidebands on the short wavelength side. In contrast, the longer wavelength side obviously experiences an extra loss, and shows a steeper slope without sidebands as a result, which is probably caused by the limitation of the effective gain bandwidth. From this point of view, such

solitons should be dissipative solitons. Its relative large frequency chirp of ~ 0.7 also supports this standpoint.

Figure 4.7: The smooth transition between two type dissipative solitons

Starting from the above soliton spectrum with only Kelly sideband on the short wavelength side, keeping all other operation parameters the same but only varying the orientations of the paddles of the PC, the spectrum could be continuously tuned to the one with steep spectral edges and a gain-bandwidth-limited bandwidth. Such a continuous and smooth transition between these two spectra reflects indirectly that these two types of solitons have the same nature and belong to the same category, further confirming that the solitons characterized by the spectrum with only Kelly sideband on the short wavelength side are also dissipative solitons. Therefore, the smooth transition between the typical dissipative solitons in normal and anomalous dispersion could be easily realized in our fiber laser by simply changing the

orientations of PC. Figure 4.7 shows an experimental example of the transition between the two types of dissipative solitons.

4.2.2 Numerical Simulations

The previous studies have shown that the solitons characterized by the spectrum with steep spectral edges could only be observed in the large positive cavity dispersion, while those characterized by the spectrum with Kelly sidebands could only be observed in the large negative cavity dispersion regime. However, why both of the two types of solitons could simultaneously be observed in our fiber laser? To investigate this problem and try to find out the underlying mechanism, we numerically simulated the operation of the laser using the technique described in chapter 2.

In accordance with the experimental laser conditions, the following fiber parameters were adopted for the simulation: $\gamma = 3 \text{ W}^{-1}\text{km}^{-1}$, $k''_{\text{SMF}} = -24 \text{ ps}^2/\text{km}$, $k''_{\text{DCF}} = 8 \text{ ps}^2/\text{km}$, $k''_{\text{EDF}} = 42 \text{ ps}^2/\text{km}$, $k''' = -0.13 \text{ ps}^3/\text{km}$, $P_{\text{sat}} = 1 \text{ nJ}$, $l_0 = 0.15$ and $T_{\text{rec}} = 6 \text{ ps}$. In the simulation, keeping the cavity length fixed at $L = 11 \text{ m}$, only the ratio of the length of SMF and DCF was varied to manage the net cavity dispersion. It turns out that different from previous results, both of above two types of solitons could be obtained in fiber lasers with near zero dispersion. The only difference is that the solitons characterized by the steep spectral edges could only be achieved in the slightly positive dispersion regime, while those characterized by Kelly sidebands could only be achieved in the slightly negative dispersion. Figure 4.8 shows a typical soliton state numerically calculated with the net cavity dispersion of $+0.0048 \text{ ps}^2$. To obtain the

state, the absorber saturation energy was set as $E_{\text{sat}} = 50$ pJ. A gain bandwidth of $\Omega_g = 2$ nm and a pump strength of $G = 195 \text{ km}^{-1}$ were used. Obviously, the achieved soliton indeed displays sharp spectral edges and relative broad pulse width. In contrast, Figure 4.9 shows a typical soliton state numerically calculated with the net cavity dispersion of -0.01 ps^2 . The absorber saturation energy was set as $E_{\text{sat}} = 1$ pJ. A gain bandwidth of $\Omega_g = 24$ nm and a pump strength of $G = 115 \text{ km}^{-1}$ were used. The obtained soliton display Kelly sidebands and a narrower pulse width. Gain bandwidth here in fact refers to the effective gain bandwidth, which corresponds to the narrowest spectral bandwidth limitation in the cavity. When the operation condition (pump power or orientations of the pc) is different, the birefringent filter is tuned, as well as the dynamical gain and loss, which eventually leads to the narrowest spectral bandwidth limitation different. Hence, the gain bandwidth value could also be different. On the other hand, it is difficult to know the accurate value of the absorber saturation energy in experiment. Hence, different saturation energies are tried in the simulation just to investigate the formation of the stable solitons in fiber lasers.

Figure 4.8: The soliton state characterized by steep spectral edges numerically calculated with the cavity dispersion of $+0.0048 \text{ ps}^2$: (a) Optical spectrum; (b) Pulse profile.

Figure 4.9: The soliton state characterized by Kelly sidebands numerically calculated with the cavity dispersion of -0.01 ps^2 : (a) Optical spectrum; (b) Pulse profile.

Based on above simulation results, through carefully analyzing the laser cavity, it turns out that in practice it is difficult to accurately determine the net cavity GVD of

our fiber laser since many other effects need to be considered in the case of a DM cavity with near zero net dispersion, and moreover its value might differ under different mode-locking states. For example, it is well known that the cavity dispersion is wavelength dependent. As shown above, the central wavelength of the mode-locked pulse was experimentally observed to shift with the settings of the orientations of the PC, and hence the dispersion varied as well. Further considering that the dispersion is around zero, it is probable that the wavelength dependent dispersion variation lead the value varying from the positive one to negative one, or vice versa. Figure 4.10 shows a corresponding example observed experimentally. It shows clearly that the dissipative solitons with steep spectral edges and with Kelly sidebands have different central wavelengths.

Figure 4.10: Two type solitons observed with different central wavelength of the spectrum.

Figure 4.11: Two type solitons observed with almost the same central wavelength of the spectrum

Besides the wavelength dependent effect, the nonlinearity could also induce the dispersion shift in a laser [108]. Specifically, different from an ANDi fiber laser, the effective cavity dispersion of a DM fiber laser depends on the concrete nonlinear pulse propagation in the cavity. It has been demonstrated previously that under strong cavity nonlinearity, the effective cavity dispersion would shift to the anomalous side, and the stronger the cavity nonlinearity, the larger the effective cavity dispersion shift is [108]. In this experiment, changing the directions of the PC corresponds to varying the dynamical loss of the mode-locked pulses, which leads to the peak power of the solitons different. Hence, under different operation conditions, the dispersion shift induced by the nonlinearity is not the same. When the dispersion is near zero, and the nonlinearity is strong enough, it is possible that this nonlinearity induced dispersion shift makes the effective cavity dispersion switching to the negative value from positive one. In addition, the gain dispersion is also the nonlinearity dependent, which

further influences the effective net cavity dispersion. Figure 4.11 shows an experimental example, where the solitons with steep spectral edges and with Kelly sidebands have almost the same central wavelengths, which is probably a result of the effective dispersion shift induced by the nonlinearity.

4.3 Conclusion

In summary, an important and distinct feature of DM fiber lasers, the nonlinearity induced effective dispersion shift effect, and its impact on the dissipative soliton operations have been investigated experimentally. With the influence of this effect, the variations of the features of the dissipative solitons in a DM fiber laser with different cavity dispersion have been discussed in detail. The dissipative solitons in a DM fiber laser with large normal cavity dispersion have been studied. It was shown that under strong pumping the stretched-pulse operation, which was normally formed in around zero dispersion, could even be obtained in such a cavity. Numerical simulations confirmed the result and further showed that in addition to the strong pumping, the cavity linear phase delay must be so biased that practically the peak power of the formed pulse has not been clamped. The results suggest that the physical mechanism of the phenomenon is the effective cavity dispersion shift caused by the non-equal nonlinear pulse propagation in the normal and anomalous fiber segments of the cavity, which is the same mechanism responsible for the DM soliton formation in the net slightly positive DM systems.

The dissipative solitons in a DM fiber laser with around zero cavity dispersion have been investigated. It was first demonstrated experimentally that both of the dissipative

solitons with the steep spectral edges, which were normally achieved in large normal dispersion regime, and with clear Kelly sidebands, which were normally observed in large anomalous dispersion regime, could appear in the same DM fiber laser with net near zero cavity dispersion. These two types of dissipative solitons could be continuously changed from one to the other by only tuning the orientations of PC. With the help of numerical simulations, the physical mechanism of the phenomenon was studied. Our result shows that the nonlinearity induced effective dispersion shift is a generic feature of the DM lasers. It is important for the better understanding on the dynamics of the DM fiber lasers and thus designing the DM cavity for practical application. It also makes the realization of the changing among different soliton operations more easily, by simply varying the operation condition without modifying the dispersion map.

Chapter 5

Dissipative Soliton Resonance and Flat-Top Solitons

The complex Ginzburg-Landau equation (CGLE) has been proved to be an accurate model to describe the operation of passively mode-locked fiber lasers. Previous studies have shown that the CGLE supports a wide variety of dissipative solitons with highly complicated dynamics, which do not exist for the nonlinear Schrödinger equation (NLSE). These include the pulsating solitons [110-112], bound states of solitons [113-115], and soliton explosions [112, 116, 117].

Recently, Chang et al. have further numerically demonstrated another feature of the CGLE. They have shown that in some special parameters regime the energy of the dissipative solitons governed by the CGLE could increase indefinitely while the peak power of the soliton pulse remain a constant value. Such a feature of the dissipative solitons was named as “dissipative soliton resonance (DSR)” since it resembles the resonance phenomenon of the classic oscillators [32]. They also pointed out that the resonance mostly occurred in the regime of normal dispersion, but could be shifted to the anomalous dispersion region with appropriate parameter combinations.

The prediction of DSR effect attracts considerable interest since it provides a roadmap for achieving large energy mode-locked pulses in the fiber laser. Although the pulse energy achievable in a mode-locked fiber laser has been significantly improved in the past few years using the cascaded pulse amplification technique, the technique itself is complicated and expensive. The DSR would offer the possibility of achieving large energy pulses directly from a single laser oscillator, which is simple and easy to implement.

Considerable theoretical researches have been done on the DSR. Kalashnikov has shown the energy scaling property of the chirped dissipative solitons [118]. Chang et al. generalized the parameter space of the CGLE for DSR [119]. They further identified the parameter regime of DSR for the chromatic-dispersion-free case and anomalous dispersion case, which indicates that large energy pulses could also be generated in mode-locked fiber lasers with anomalous cavity dispersion [120]. Moreover, DSR was also found to exist in laser models with parameters management [121], which suggests that such a phenomenon could also be observed in mode-locked fiber lasers with dispersion-managed (DM) cavity. Until now, the research on DSR has mainly been focused on the numerical studies. No experimental observation of DSR has been reported.

It is worthwhile to mention that in the DSR region, the dissipative solitons will only broaden their pulse widths, but keep their amplitudes fixed, exhibiting as flat-top solitons. Apart from the well-known sech^2 -profile solitons, flat-top solitons were also predicted by the CGLE [29, 122, 123]. However, to our knowledge, no experimental

observation of such a flat-top dissipative soliton in mode-locked fiber lasers has been reported.

In this chapter, through special design of an all-normal-dispersion (ANDi) fiber laser we have first experimentally demonstrated the DSR and flat-top solitons. Taking advantage of the DSR, we have generated mode-locked pulses in the erbium-doped fiber (EDF) laser with a record high energy of 281.2 nJ. The energy of the mode-locked pulse is experimentally limited by the available maximum pump power injected into the cavity.

5.1 Experimental Setup and Results

Motivated by theoretical predictions on the DSR, we have designed an EDF laser to realize the DSR effect and flat-top solitons. An ANDi fiber laser was constructed for the experiment. Different from the conventional mode-locked fiber lasers, a long cavity was used in the laser to reduce the cavity fundamental repetition rate and simultaneously increase the normal dispersion. The nonlinear polarization rotation (NPR) mode-locking technique was adopted to mode lock the fiber laser. Based on our experience under peak power clamping the mode-locking technique provides the required strong self-amplitude modulation saturation term, which benefits the appearance of the DSR.

The EDF laser has a similar cavity configuration as that described in chapter 3 (Figure 3.1). The fiber laser cavity has a length of around 163.2 m. A segment of 3.7-m EDF with group velocity dispersion (GVD) parameter of 41 ps²/km was employed as the

gain fiber, which has a peak absorption of 80 dB/m at near 1530 nm, mode field diameter of 4.3 μm at 1550 nm. The rest of the laser cavity consists of dispersion-compensating fibers (DCFs) with GVD parameter of 5 ps^2/km , whose mode field diameter and attenuation at 1550 nm are respectively 8.1 μm and ≤ 0.21 dB/km. A 50% fiber coupler was used to output the laser emission. The backward pump scheme was exploited to eliminate the influence of the residual pump power. The laser is an ANDi all-fiber laser.

Mode locking of the laser is always self-started. Depending on the settings of the linear cavity phase delay bias and pump strength, various modes of laser operation were observed, including the conventional dissipative solitons [25] and dark solitons [124]. DSR was observed under strong pumping, as shown in Figure 5.1. Figure 5.1(a) shows the typical oscilloscope traces of the observed square pulse emission. It is noticed that the rising and falling edges of the measured square pulses are limited by the response time of our measurement systems. The square pulse duration increases with the pump strength, while the peak of the pulse almost remains constant as the pump power varies. Figure 5.1(b) shows the optical spectra of these square pulses. Their spectra has no characteristic sharp spectral edges as those of the dissipative solitons in normal dispersion, but resembles that of the amplified spontaneous emission of the EDF with the spectral center shifted to the longer wavelength side. Figure 5.2 shows the experimentally measured square pulse width and output power variation with the pump power. The narrowest stable square pulse experimentally obtained was about 632 ps, while the longest pulse width could become as long as tens of ns, which was limited by the available pump power injected into the cavity. The highest laser output power attained is 344.5 mW with the pump power of 1186

mW in this experiment. No output power saturation was observed. As a WDM at a slightly higher pump power has been burned, the maximum pump strength has been kept at the current level. At the maximum pump power the square pulse also reached the maximum duration of 18.5 ns. Associated with the pulse width increase the spectral bandwidth of the pulses also slightly increased, as shown in Figure 5.1(b). The square pulse laser emission was stable, as shown in the inset of Figure 5.1(a). The pulse train has a repetition rate of 1.225 MHz (corresponding to a pulse interval of 816.3 ns) as determined by the cavity length.

Figure 5.1: Square pulse emission of the laser. (a) Zoom-in high speed oscilloscope traces under different pump power; inset: oscilloscope trace of a square pulse train. (b) Optical spectra of the square pulses under different pump power.

Figure 5.2: The experimentally measured average output power and pulse width versus the pump power injected into the cavity.

A remarkable feature of the square pulses is that their pulse energy could increase to a very large value without pulse breaking and pulse shape distortion. In our experiment the maximum pulse energy was as large as 281.2 nJ, which was limited by the pump power injected into the cavity. At the maximum pulse energy the pulse still has a square shape, with almost the same peak power it has before but with pulse duration increased to 18.5 ns. The square pulses have also been experimentally investigated with the high-speed oscilloscope combined with a commercial autocorrelator. No internal fine structures were observed within the pulses, as shown in Figure 5.3.

Figure 5.3: The measured autocorrelation trace of the square pulse achieved in our laser.

Dissipative solitons in normal dispersion, characterized by their sharp spectral edges, were also observed in the laser. Comparatively, these dissipative solitons have not only much narrower pulse width (maximum duration about 36 ps) but also much smaller maximum pulse energy (~ 8.2 nJ). The maximum dissipative soliton pulse energy was limited by the effective gain bandwidth of the laser rather than the available pump power. In the experiment as the energy of a dissipative soliton increased beyond above maximum value, it broke into two solitons. Multiple dissipative solitons were always obtained under strong pumping. In fact, the square pulse operation of an ANDi EDF laser has been reported in our previous work [125]. However, in that laser since a large output coupling (90%) coupler and a fiber bench with free space waveplates were used, the laser cavity had very large linear losses. Therefore, with the available pump power (700 mW) only very low energy square pulses could be obtained.

Based on the pulse parameters measured, the peak power of the square pulse inside the cavity was estimated to be about 30 W. Obviously the pulse propagation in the fiber cavity is strongly nonlinear. Therefore, the stable square pulse emission of the laser suggests that the observed square pulse could be another solitary wave formed in the laser. Indeed, flat-top solitons were theoretically predicted to exist in the extended CGLE [29, 122], formed as the kink-antikink bound state. Such a soliton was previously observed in a binary fluid convection experiment [126].

5.2 Discussion

The square pulse obtained in our laser displayed exactly the DSR feature predicted, indeed increasing its pulse width and energy with the pump power while keeping its amplitude constant. Especially, the linear dependence of the pulse energy on the pump power shows a well agreement with the recent result reported in Ref. [127] qualitatively. In [127], Ding et al. theoretically studied the features of the DSR with a gain saturation model instead of the previous constant gain model [120], which described the practical laser system more precisely, and they pointed out that when the gain saturation effect was taken into account, the pulse energy increased with the pumping power linearly rather than exponentially. The only small discrepancy between our experimental observation and theoretical prediction comes from the optical spectrum. No steep spectral edges with sharp side peaks were observed on the spectral profiles. We believe this discrepancy could be due to that with the current laser parameter selection, at the maximum experimentally available pump power just a weak resonance was reached. Note that as the pump power was increased, the

experimentally observed soliton spectral profile did exhibit the tendency of edge steeping. Therefore, it is expected that if higher pump power is allowed, stronger resonance would be further achieved, and a closer agreement with those of the theoretical predictions would be obtained. Finally, it is noticed that square pulse emission had also been observed in the anomalous dispersion fiber lasers [51], indicating that its formation is independent on the sign of the cavity dispersion. Nevertheless, in the anomalous dispersion fiber lasers due to the intrinsic soliton formation, fine internal structures could also appear in the square pulses and complicate the phenomenon.

5.3 Conclusion

In summary, the laser emission of an ANDi fiber laser mode-locked by NPR technique with a relative long cavity length has been investigated. It was found that a sequence of large energy square pulses in addition to the normal dissipative solitons had also been observed in such a fiber laser, and this stable nonlinear square-pulse emission was further demonstrated to manifest the flat-top dissipative soliton operation predicted. Moreover, the width of the square pulse, as well as its pulse energy was experimentally observed to increase to quite a large value with the pump strength while keeping their intensity almost constant. As a result, pulses with record high energy of 281.2 nJ and duration of 18.5 ns were finally achieved in the experiment without pulse breaking, which is a clear qualitative experimental evidence of the existence of the DSR. To the best of our knowledge, this is the first experimental observation of the typical DSR effect and flat-top solitons in mode-locked fiber lasers. The large energy square shape soliton of the laser could be

used for large pulse energy ns level mode-locked pulse generation, which could find extensive application in micromachining, material processing, remote sensing, and light detection systems [128-130].

Chapter 6

Dynamics of Dissipative Solitons

Different from propagating in a single mode fiber, a soliton circulating in a laser cavity is also subject to the actions of discrete cavity components, such as the output loss and spectral filter, which contribute essentially to the pulse shaping of the dissipative solitons other than the conventional solitons. Hence, when a dissipative soliton circulates in a cavity, what are the intrinsic functions of these elements and how they practically affect the soliton evolution deserve to be studied. The investigation on how the parameters of dissipative solitons vary along the cavity within one roundtrip also appear necessary for better understanding the dissipative soliton operation in fiber lasers.

Chong et al. have systematically studied the propagation and intracavity dynamic behavior of the dissipative solitons in ytterbium-doped all-normal-dispersion (ANDi) fiber lasers [27, 28, 36]. They verified both numerically and experimentally that the spectral filtering of a chirped pulse was the intrinsic pulse shaping mechanism. They also revealed that the dissipative solitons presented distinct characteristics and properties at different positions along the cavity, and identified that the spectral filter

bandwidth, nonlinear phase shift and GVD were the key cavity parameters determining them. Zhao et al. have also investigated the intracavity dynamics of dissipative solitons in an erbium-doped fiber (EDF) laser made of purely normal dispersion fibers [131]. Different from a ytterbium-doped fiber (YDF), which always has a large bandwidth so an external physical spectral filter with a localized function and fixed filtering strength is necessary for the chirped pulse generation in YDF lasers, an EDF possesses a narrow gain bandwidth and itself could function as a spectral filter, whose action lasts the whole length of the gain medium and filtering strength depends on the intrinsic feature of the gain medium and pump power. On account of markedly different properties of the spectral filters, which shape the solitons, the intracavity evolution of the dissipative solitons in EDF lasers presents different features from that in YDF ones. Through investigating the variation of the dissipative soliton parameters along the cavity, Zhao et al. showed that the effects of the gain and saturable absorption process on the soliton shaping of the pulse became evident. They also pointed out that the dissipative soliton formation was a mutual interaction and balance among cavity dispersion, fiber nonlinearity, laser gain saturation and spectral filtering within one cavity round trip, so the dissipative soliton was intrinsic the average soliton like conventional solitons.

Apart from ANDi fiber lasers, dissipative solitons have also been observed in the dispersion-managed (DM) fiber lasers with large normal cavity dispersion [26, 98-101]. However, different from propagating in a non-DM laser cavity, a pulse circulating in a DM fiber laser experiences strong stretch and compression, correspondingly frequency chirping and dechirping. Consequently, their evolution will of course exhibits distinct characteristics from that in a non-DM fiber laser. What

exactly these differences are, and how could a dissipative soliton be formed in a DM laser? These questions have not been addressed to our knowledge so far. It is worth noting that a parabolic pulse can also be formed in the DM fiber laser, and its spectrum also has steep edges [95]. What are the relationship and differences between the pulses? To this end a DM fiber laser with large normal dispersion is constructed to investigate these problems.

Chaos and transitions to chaos are the universal behaviors of a nonlinear system, which exist in broad areas including physics, chemistry, and biology [132-134]. Three distinct transition routes leading to this deterministic chaos from regular motion in a nonlinear dissipative dynamical system are the period-doubling route [135], quasiperiodicity route [136], and intermittency route [137, 138]. The intermittency among them is characterized by a signal, which behaves regularly (or laminarly) in time becoming interrupted by statistically distributed periods of irregular motion (intermittent bursts), and was firstly studied by Pomeau and Meanneville [138]. The average number of these bursts was found to increase with the variation of an external control parameter until the motion becomes completely chaotic. After its first observation in the hydrodynamical system [137], the intermittency phenomenon has been extensively observed in gas and semiconductor laser systems [139-141].

Strong nonlinearity of the fiber, in combination with the cavity gain and loss, makes a mode-locked fiber laser a typical nonlinear dissipative system. Furthermore, it also exhibits an essentially infinite propagation distance with identical conditions for each period, which also provides an excellent test bed for the experimental study on the nonlinear dynamics. In fact, all of the chaos and another two transition routes to chaos

have been indeed experimentally observed in mode-locked fiber lasers, except for the intermittency. Chaotic behaviors of the mode-locked fiber lasers were reported previously [142, 143]. Tamura et al. also observed the period doubling and tripling phenomena in fiber lasers with an intracavity broad filter [85]. Subsequently, Zhao et al. even revealed a complete route of period doubling to chaos in a passively mode-locked fiber soliton laser [144]. Recently, the quasi-periodicity behavior of the mode-locked pulses has also been revealed in an additive-pulse mode-locked F-center [145]. However, to the best of our knowledge, no experimental observation of such an intermittency phenomenon in fiber laser systems has been reported so far.

Chapter 6 focuses on the dynamics issue of the dissipative solitons in passively mode-locked fiber lasers. It is organized as follows. Section 6.1 investigates the intracavity dynamics of the dissipative solitons in a DM fiber laser, which is found presenting different features from that in a non-DM cavity. Section 6.2 reports the first experimental observation of the intermittency phenomenon in fiber lasers. The simulation confirms the experimental results. Section 6.3 shows the conclusion.

6.1 Dynamics of Dissipative Solitons in a DM Fiber Laser with Large Normal Cavity Dispersion

In this section, a DM fiber laser with large normal cavity dispersion was constructed to investigate the evolution of dissipative solitons within it.

6.1.1 Experimental Results

The cavity has a similar configuration as that employed in chapter 4.1, which is shown in Figure 6.1. The cavity consists of 2.7 m EDF with group velocity dispersion (GVD) of about ~ -32 (ps/nm)/km and standard single mode fibers (SMFs). Through carefully controlling the SMF, the net cavity dispersion was purposely set to be a large positive value ~ 0.0553 ps² with a cavity length of $L = 1.4_{\text{SMF}} + 2.7_{\text{EDF}} + 1.0_{\text{SMF}} = 5.1$ m so that the dissipative soliton operation could be achieved. The nonlinear polarization rotation (NPR) technique was used to mode lock the fiber laser. Different from other cavities we constructed before, two 10% fiber output couplers were specially introduced in the laser to extract the pulses from the cavity at different positions for comparison. One was located at the position 0.9 m before the EDF (port 2) along the clockwise, and the other was situated 0.3 m after the EDF (port 1).

Figure 6.1: Schematic of the DM fiber laser. I, polarization-independent isolator; P, polarizer; $\lambda/4$, quarter-wave plate; $\lambda/2$, half-wave plate; SMF, standard single-mode fiber; WDM, wavelength division multiplexer; EDF, erbium-doped fiber.

Self-started mode locking was obtained through appropriately selecting the wave-plate orientations. Figure 6.2 shows a typical mode-locked state obtained with a pump power of 309 mW. The corresponding optical spectra and autocorrelation traces measured from two ports are presented. The dissipative soliton operation of the laser is supported by the features of the measured spectra. Through comparing the results from two ports, it is found that the spectral profile changes only slightly as well as the pulse width, which varies from 3.56 ps (Output 1) to 3.02 ps (Output 2). In contrast, the corresponding average output power changes dramatically, which varies from 6.820 dBm (Output 1) to 0.062 dBm (Output 2). The static insertion loss of the fiber bench in our experiment is measured to be ~ 3 dB, in combination with the loss introduced by the polarizer, indicating that the NPR has chopped away most of the pulse energy, which verifies that the saturable absorption is a strong dissipative process. In addition, the top of the measured spectral profiles is quite humped compared with those of the parabolic pulses [95, 146, 147], and the height of the hump is experimentally found to vary with the wave-plate orientations and the pump power through investigating the dissipative solitons under different operation conditions, which suggests that the cavity peak clamping effect [52] could have played an important role in the formation of the soliton. In fact, the observed experimental results have no obvious difference from those obtained in the non-DM fiber lasers [131]. To gain an insight into the detailed intra-cavity pulse evolution, the numerical simulations have further been conducted based on the experimental system.

Figure 6.2: Experimentally measured autocorrelation traces and optical spectra of the dissipative soliton at the two cavity output ports of the DM fiber laser. Solid lines: results obtained at the port-1 position. Dashed lines: results obtained at the port-2 position.

6.1.2 Simulation Results and Discussions

Again, the same model as described in chapter 2 was used. To make the simulation comparable with the experiment, the following simulation parameters were adopted: nonlinear fiber coefficient $\gamma = 3 \text{ W}^{-1}\text{km}^{-1}$, EDF gain bandwidth $\Omega_g = 20 \text{ nm}$, third-order fiber dispersion $D''' = 0.1 \text{ (ps}^2/\text{nm)}/\text{km}$, beat length $L_b = L/2$, orientation of the intra-cavity polarizer to the fiber fast birefringent axis $\Phi = 0.152\pi$, and the gain saturation energy $P_{\text{sat}} = 1 \text{ nJ}$. The 3 dB insertion loss of the fiber bench was considered, and two output ports have also been set at the same cavity position as those in experiment for comparison.

We confirmed numerically that the dissipative soliton could indeed be obtained in the fiber laser. Figure 6.3 shows such a typical state obtained with the linear cavity phase delay bias of 1.4π and the small-signal gain of $G = 2000 \text{ km}^{-1}$. Both of the pulse

profiles and optical spectra agree well with the experimental results shown in Figure 6.2. From port 1 to port 2, the calculated soliton pulse presents a peak spectral intensity jump of ~ 5.5 dB, which is also close to the experimental measurement value of ~ 6.1 dB. A small discrepancy between the simulation and the experiment is that the top of the numerically calculated pulse spectra is less humped than those experimentally observed, which is probably caused by the reason that the practical complex gain profile was simplified by a parabolic gain profile in our simulation. Varying the small-signal gain or linear cavity phase delay bias in the model, which corresponds to changing the pump strength or orientations of the waveplates in experiment, different dissipative solitons were achieved. Theoretically different linear cavity phase delay biases further affect the effective saturable absorption strength of the mode-locking and the cavity pulse peak clamping level. Our numerical simulations faithfully reproduced all the features of the dissipative solitons experimentally observed.

Figure 6.3: Numerically calculated pulse profiles and optical spectra of the dissipative soliton at the two output port positions as in the experiment. Linear cavity phase delay bias $\Phi = 1.4\pi$, pump strength $G = 2000 \text{ km}^{-1}$. Solid lines: results obtained at the port-1 position. Dashed lines: results obtained at the port-2 position.

After verifying the model, we further monitored the relevant parameters of the formed soliton along the resonator to get a deeper insight into the intracavity pulse shaping dynamics. Figure 6.4 shows the variations of the pulse width and the time-bandwidth product along the cavity for the soliton operation shown in Figure 6.3, which starts from the end terminal of the fiber bench. The time-bandwidth product in fact gives information about the frequency chirp of the soliton. Obviously, the pulse is positively chirped at the beginning, and it is dechirped in the first segment of SMF due to the anomalous dispersion. Before dechirped into a transform-limited one, the pulse enters the EDF with normal dispersion and accumulates the positive frequency chirp then, and becomes dechirped again when it enters the second segment of SMF. Although the cavity is DM, obviously different from the stretched-pulse regime, where the pulse is transform-limited at two positions in the resonator, the chirp of the dissipative soliton is always large and positive throughout the cavity as a result of the large net positive cavity GVD. The magnitude of the soliton frequency chirp was further numerically found to be the cavity linear phase delay bias and pump strength dependent. The higher the peak power of the formed soliton, the larger the absolute value of the frequency chirp was.

Figure 6.4: Variations of the dissipative soliton pulse width and the soliton time-bandwidth product along the cavity. The slight vertical offsets of the starting and ending points of the curves are due to the pulse shaping caused by the cavity mode-locking effect.

Figure 6.4 also shows that the pulse first narrows down in the SMF due to the dechirping process, and subsequently broadens monotonically in the gain fiber as a result of the positive frequency chirp accumulation, and narrows down again in the second SMF on account of dechirping. A slightly pulse width drop of ~ 0.4 ps also occurs when passing through the intracavity polarizer, which completes the cavity round trip, indicating that the saturable absorption effect of the NPR temporally shorten the pulse duration. For a highly chirped pulse, such a temporal compression process also provides a spectral narrowing mechanism and shapes the dissipative soliton. Note that although the cavity is DM, different from the stretched pulse stretching and compressing twice per cavity round-trip, now the pulse only stretches and compresses once in one round-trip, which resembles that of the intra-cavity pulse dynamics of the parabolic pulse [95].

Obviously, the intra-cavity pulse evolution is also different from that of the dissipative soliton in ANDi fiber lasers. Figure 6.5 shows a typical intracavity dynamics of the dissipative soliton in an ANDi EDF laser reported in Ref. [131]. It is noticed that although the pulse is positively chirped and the gain fiber has a normal GVD, there is a slight pulse compression in the gain fiber, which is the strong evidence of the soliton shaping of the gain bandwidth filtering on highly chirped pulses. However, no such a pulse width narrowing was observed in our experiment. We would like to point out that despite that no pulse narrowing occurs within the gain fiber in a DM cavity due to the de-chirping of the anomalous fiber, the gain bandwidth filtering still plays a role in the formed soliton pulse.

Figure 6.5: Pulse width variation of a dissipative soliton along the laser cavity numerically calculated [131].

Numerically it was found that with all other laser parameters fixed but decreasing the gain bandwidth, i.e. from 20 nm to 12 nm, the soliton could even be de-chirped to a transform-limited pulse in the SMF in the case of strong solitons. The same effect was

also observed under the 20 nm gain bandwidth but with small net positive cavity dispersion, e.g. less than 0.025 ps^2 , which was understood as a result of the complete balance between the positive and negative chirps generated in the cavity. Like the case of the conventional DM soliton [108], when a pulse has broader bandwidth in the anomalous fiber segment than in the normal fiber segment, it will experience more effective anomalous linear dispersion than normal linear dispersion, and consequently form DM solitons even in the net positive dispersion map. Based on this understanding the result we observed then suggests that depending on the gain bandwidth, the effective positive chirp of the formed dissipative soliton is also different. Gain with narrower bandwidth will generate dissipative solitons with smaller chirp.

It is to see that although there is no pulse narrowing in the gain fiber in the DM fiber laser, the gain filtering has still played a role in the formed soliton pulse. Therefore, in the case of large net positive cavity dispersion DM lasers, the dissipative solitons can be formed. The difference and the relation of the dissipative solitons to the parabolic pulses also become clear now. In a DM laser both pulses have the similar intra-cavity pulse evolution, however, in the case of a parabolic pulse the gain bandwidth is so broad that it could be considered as infinite. Hence it has no bandwidth limiting effect on the formed pulses, and consequently the formed pulses have a linear frequency chirp [95].

6.2 Intermittency

In contrast to the intracavity dynamics of the soliton pulses discussed above, we shift

our focus to the dynamical behavior of the pulse train of the dissipative solitons. In this section, we employ an EDF mode-locked laser, and concentrate on the behavior of the pulse train of output dissipative solitons in a time scale much larger than the round-trip time to search for the intermittency phenomenon.

6.2.1 Experimental Observations

A fiber laser with the similar configuration as that in section 4.1 (Figure 4.1) was constructed for our experiment. Now the cavity is made of purely normal dispersion fibers, consisting of 5-m EDF with a doping concentration of 2880 ppm, and 8.5-m dispersion-compensating fibers (DCFs), and the laser light was coupled out via a 50% fiber coupler.

Once the orientations of the wave plates were set appropriately, mode locking of the laser self-started as the pump power was increased above the mode-locking threshold of 112 mW. Dissipative solitons were further formed in the cavity after the nonlinear propagation of the mode-locked pulses. Their spectra are characterized by steep spectral edges with a gain-bandwidth limited bandwidth, indicating they are dissipative solitons in normal cavity dispersion. The measured autocorrelation trace shows the pulse width is 26.8 ps (FWHM), if a sech^2 -form pulse profile is assumed. These formed dissipative solitons are highly chirped, suggested by a large time-bandwidth product of ~ 20 . Figure 6.6 shows a typical pulse train of the stable dissipative soliton state achieved, and the inset is the zoom-in picture. The dissipative solitons are identified to emit at the fundamental repetition rate of 14.8 MHz, as determined by the cavity length, and only one pulse circulates in the cavity. The

output now is a uniform pulse train. All the dissipative solitons obtained have exactly the same light intensity for each roundtrip, as shown in Figure 6.6, which could be well explained by the average soliton theory of the laser.

Figure 6.6: The oscilloscope trace of the dissipative soliton state. Inset: the pulse train of the dissipative soliton state for smaller time scale.

Starting from above stable dissipative soliton state, carefully increasing the pump power or tuning the orientations of the waveplates within a certain range, the soliton peak power was observed to increase slightly, while the pulse train still remains uniform. In detail, all the properties of the laser output pulse, such as the light intensity and pulse width, still keep identical for each roundtrip. During this process, with the increase of pump power, the nonlinearity that the soliton pulse experienced becomes stronger consequently, but still not strong enough to destroy the periodicity and regularity of the dissipative soliton evolution. Hence, the laser is still operated in

the fundamental soliton state only with different soliton parameters. Until now, all the phenomena and features are exactly the same as those of all the dissipative soliton states observed previously.

However, with the pump power further increased, typically above 157 mW, it is observed that the output soliton pulse train of the laser becomes nonuniform and irregular, suggesting the laser is not on the stable soliton emission state any more, but transit to another state. Figure 6.7 shows a typical output pulse train of such a state observed in our fiber laser. Rather than the periodic output pulse train with uniform pulse intensity in the stable dissipative soliton emission state, the regularly pulsing of the intensity (“laminar phase”) now is interrupted by the bursts of larger peaks, which exhibits the exact features and properties of the intermittency observed in gas and semiconductor laser systems previously [139-141]. The growth of the output light intensity was observed clearly. When the light intensity reaches a high value, the increase of the intensity becomes rapid. Eventually the signal loses its regularity and a turbulent burst appears. Immediately after this burst, there is a reappearance of the regular behavior, corresponding to a return to the preceding periodic behavior, with a given light intensity. The whole process repeats continuously. During this burst, the pulse energy undergoes dramatic changes, and all these bursts have similar features, but not identical.

Figure 6.7: The oscilloscope trace of the intermittency behavior

Figure 6.7 also displays the randomness of the interval between two adjacent bursts. In fact, it is found experimentally that the average interval of two bursts depends on the strength of the pump power, and the settings of the orientations of the waveplates. With the orientations of the waveplates fixed, the average interval of two bursts decreases with the enhancement of the pump power, and finally transits into the chaos. The similar results were obtained by only varying the orientations of the waveplates under certain pump power. The reversibility of the intermittent behavior of the fiber laser was also observed. When the pump power was reduced, or the orientation of the waveplates was set back to the previous position from the intermittency behavior, the laser emission will return to the soliton state once again. This indicates that strong nonlinearity is necessary for the occurrence of the intermittency. In addition, other

nonlinear dynamic behaviors, such as period-doubling bifurcations and chaos were also observed in the same fiber laser only by changing the pump power or orientations of the waveplates.

6.2.2 Simulation Results and Discussions

To verify the experimental observation of the intermittency in a mode-locked fiber laser, we numerically simulated the laser operation using the pulse tracing method.

Figure 6.8: Numerically simulated result of the soliton state.

Figure 6.9: Numerically simulated result of the intermittency behavior. (a) gain $G = 392 \text{ km}^{-1}$; (b) gain $G = 393 \text{ km}^{-1}$; (c) gain $G = 394.5 \text{ km}^{-1}$.

It is numerically found that, once the cavity beat length L_b is set appropriately, which corresponds to suitably choosing the orientations of the waveplates in experiment, the stable soliton state could be easily achieved in the laser cavity by simply increasing the gain G . Figure 6.8 shows, for example, a stable soliton state numerically obtained when $L_b = 6$, and $G = 340 \text{ km}^{-1}$. The evolution of the mode-locked pulses with the cavity round-trips displays clearly that the achieved solitons have exactly the same light intensity for each roundtrip with a well-shaped pulse profile. Based on this soliton state, keeping all other parameters unchanged, but only increasing the gain G , it is observed that a bifurcation behavior of the output mode-locked pulses begins to emerge when G is above 360 km^{-1} . More specifically, at the bifurcation state, the output mode-locked pulses still have periodic behavior in a time scale larger than the round-trip time, whereas this period is not commensurate with the round-trip time, but several times of it. Then with the further increase of the gain G , the intermittency phenomenon is observed. Figure 6.9 shows the intermittency state numerically calculated in the mode-locked fiber laser when the gain is selected as a) $G = 392 \text{ km}^{-1}$, b) $G = 393 \text{ km}^{-1}$, and c) $G = 394.5 \text{ km}^{-1}$, respectively. Figure 6.9 displays clearly that the regularly pulsing of the intensity of the mode-locked pulses is interrupted by the bursts of larger peaks in the intermittency, and the interval between two adjacent bursts is random, which leads to a mean turbulent state. With the increase of the gain G , the bursts occur more and more frequently, which results in the shorter and shorter mean interval of two bursts, and finally the system reaches a complete chaotic state. All these features are in well agreement with our experimental observations.

A mode-locked fiber laser is a nonlinear dissipative system. Its operation is well known to be described by the complex Ginzburg-Landau equation (CGLE), which takes into account of all the laser gain, loss, dispersion, and nonlinearity. Theoretical studies on the CGLE reveal that it presents a variety of solutions under different laser parameters, which correspond to the complicated and rich laser outputs in experiment. Soliton, as one of the well-known solutions, and refers to a stable self-localized structure. Solitons in fiber lasers mode-locked by various techniques, such as NPR, semiconductor saturable absorber mirror (SESAM), and carbon nanotubes (CNTs) have been investigated extensively. In recent years, Soto-Crespo et al. demonstrated that, apart from the solitons with perfectly periodic behaviors, another type of pulsating solution, known as soliton explosion, was also the solution of the complex GLE under another range of the equation parameters [111], and it has been observed in a Ti:sapphire mode-locked laser [148], but still lacks experimental evidence in fiber lasers. The intermittency is characterized by the regularly pulsing of the intensity randomly interrupted by the turbulent bursts, in contrast to the well-periodic dynamics of the soliton. The average duration of the laminar phases is the control parameter dependent. The pulse train observed in our passively mode-locked fiber laser displays exactly all those features of the intermittency. Therefore, we claim that this is the first experimental observation of the intermittency in a mode-locked fiber laser. Our numerical simulation further confirms the observation of the intermittency. Based on some similar properties and characteristics, our results might also be the experimental evidence of the soliton explosion in fiber lasers. However, whether they are exactly the same effect needs more explorations.

6.3 Conclusion

We investigated the intra-cavity dynamics of the dissipative solitons in a DM fiber laser, and showed that due to the de-chirping of the anomalous fiber in the cavity, the dissipative soliton no longer experienced pulse width narrowing in the normal dispersion gain fiber as in the non-DM fiber lasers, but had an intra-cavity pulse evolution similar to that of the parabolic pulses. The main difference of the dissipative solitons to the parabolic pulses in the lasers is that the dissipative solitons have nonlinear chirp due to the effect of the finite gain bandwidth, while the parabolic pulses have linear chirp as in lasers where the parabolic pulses are formed, the gain bandwidth of the lasers is much broader than the spectral bandwidth of the formed pulses so that its limiting effect is ignorable.

We have observed intermittency phenomenon in a passively mode-locked fiber laser both experimentally and numerically. To the best of our knowledge, this is the first observation of such a nonlinear dynamical behavior in a fiber laser system. In the process of intermittency, the regularly pulsing of the intensity was observed being interrupted by the bursts of larger peaks rather than the conventional periodic output pulse train with uniform pulse intensity. The average duration of the laminar lengths of the intermittency was also found to be the control parameter dependent, either the pump power or the orientations of the waveplates. As intermittency and route to chaos are typical features of the nonlinear dynamic systems, this experimental result shows that the soliton propagation in the laser cavity itself is a nonlinear dynamic process. Numerical results well reproduced our experimental observations and further confirmed the existence of the intermittency.

Chapter 7

Bound State of Dissipative Solitons

Theoretically, dynamics of solitons formed in mode-locked fiber lasers is governed by the Ginzburg-Landau equation (GLE) [77, 114, 149]. Malomed firstly theoretically predicted that interaction of two GLE solitons could lead to the formation of bound states of solitons [113]. He pointed out that as the GLE soliton had an oscillating tail, when two GLE solitons slightly overlapped, spatial local minima could be formed on the interaction potential, which lead to the formation of bound states of solitons. In particular, the bound solitons have fixed discrete soliton separations, and a fixed phase difference of 0 or π [150]. A series of the experimental observations of such bound solitons with fixed discrete pulse separations and phase difference of π in nonlinear polarization rotation (NPR) mode-locked fiber lasers have been reported subsequently [87, 151, 152]. Akhmediev et al. and Afanasjev et al. have also theoretically predicted the existence of bound states of solitons with a $\pi/2$ phase difference [31, 114, 115]. Later, bound solitons with a phase difference of $\pi/2$ have also been experimentally observed in a fiber laser by Grellu et al. [153, 154]. Very recently, Zavyalov et al. further theoretically predicted the existence of bound state of solitons with independently varying phases [155]. Moreover, apart from two-soliton

bound states, even the three- and multiple-soliton bound states were also experimentally achieved, and their characteristics have been explored [153, 156-159].

Although a great deal of theoretical and experimental work has been conducted on the bound state of solitons, they mainly focus on the anomalous dispersion regime. No studies related to bound state of dissipative solitons in large normal dispersion characterized by the spectrum with sharp spectral edges were addressed. Considering that such dissipative solitons have completely different soliton pulse shaping mechanism from those formed in anomalous dispersion laser and have large frequency chirp, whether these dissipative solitons could still form bound state is worthy to be explored. It would also be interesting to find out how the frequency chirp influences the interactions between solitons. If the bound states of such dissipative solitons exist, do they still display the fixed soliton separations and phase difference, or do they have any other distinct features not existing for those bound solitons observed before? We will focus on these issues in this chapter.

Apart from NPR, other mode-locking techniques such as semiconductor saturable absorber mirror (SESAM), carbon nanotubes (CNTs) or graphene could also be employed to passively mode lock the fiber lasers. Single-wall carbon nanotubes (SWCNTs) are a new saturable absorber and have received considerable attention nowadays due to their superior properties such as easy fabrication, controllable modulation depth, and potentially broadband saturable absorption [59, 160-163]. They have quite different saturable absorption properties from those of NPR: e.g. the modulation depth of the SWCNTs is a certain fixed value, while that of the NPR is adjustable within a certain range, and the saturation recovery time of the SWCNTs is

around 1ps, much slower than that of the NPR, which is generally considered as < 10 fs. Considering that the SWCNTs present distinct saturable absorption properties from those of NPR, it is interesting to know whether it will affect the formation of the bound state of dissipative solitons and further lead to distinct features not existing for bound solitons obtained with NPR. To this end a fiber laser mode-locked with SWCNTs saturable absorber was constructed to study this problem.

Chapter 7 focuses on the subject of bound states of dissipative solitons. In section 7.1, the formation and properties of the bound state of dissipative solitons in a DM fiber laser with large normal dispersion mode-locked by NPR has been investigated. In Section 7.2, a systematic study on the bound state of solitons in a fiber laser mode-locked by SWCNTs saturable absorber has been conducted for the first time. Section 7.3 presents the conclusion.

7.1 Bound State of Dissipative Solitons in a DM Fiber Laser with Large Normal Cavity Dispersion Mode-Locked by NPR

Formation of a bound state of pulses in a passively mode-locked fiber laser with normal net cavity dispersion was reported by Grelu et al. [164]. However, as the net cavity dispersion of their laser is close to zero cavity dispersion, where dispersion-managed (DM) solitons [108] and stretched pulses [19] can be formed, their result cannot be representative for lasers with large normal net cavity dispersion, where no DM solitons but dissipative solitons with sharp spectral edges [26] can be generated. When the net cavity dispersion of a laser is largely normal, depending on

the gain bandwidth either dissipative solitons with sharp spectral edges or parabolic wave-breaking-free pulses [95] can be formed. The parabolic pulses have linear frequency chirp and exhibit self-similar evolution in the cavity. The formation of the parabolic pulse requires that no frequency filtering effect exists in the cavity. On the contrary, the dissipative solitons with sharp spectral edges has nonlinear frequency chirp, and its formation is a result of the narrow gain bandwidth filtering effect [24, 25]. Bound parabolic pulses have already been experimentally observed in an ytterbium-doped fiber (YDF) laser [157] and numerically confirmed to exist in a laser operating in a wave-breaking-free regime subsequently [165]. To our knowledge no bound state of dissipative solitons with sharp spectral edges has been reported so far.

In this section, a DM erbium-doped fiber (EDF) laser with large normal cavity dispersion mode-locked by NPR is investigated. It is shown that, like the solitons in anomalous dispersion, under direct pulse interaction, dissipative solitons with steep spectral edges can also bind together to form bound states in spite of their large frequency chirp. It is also numerically found that despite the fact that in a DM laser the circulation of a dissipative soliton in a cavity experiences strong pulse stretching and compression, the pulse separation between the bound dissipative solitons remains the same along the cavity, and so does their relative phase difference, even though due to the pulse compression the pulse separation in some segments of the cavity becomes temporally far larger than the pulse width.

7.1.1 Experimental Observations

The fiber laser constructed for our experiment has a typical DM cavity, with the

configuration similar as that shown in Figure 4.1. A 1.72 m EDF with group velocity dispersion (GVD) of about -50 (ps/nm)/km was used as the gain fiber, which exhibits the gain bandwidth filtering effect much easier than YDF. The rest of the cavity consists of 0.3 m dispersion-compensating fiber (DCF) with GVD of -0.196 (ps/nm)/km and standard single mode fiber (SMF) with GVD of about 18 (ps/nm)/km. Dispersion management is achieved by changing the length of the SMF. The NPR technique is used to achieve self-started mode locking, and the laser was output via a 10% fiber output coupler.

The dispersion caused by the bulk components is negligible. With 2.8 m total length of SMF, the net cavity dispersion was set to be large positive value $\sim 0.046\text{ps}^2$ to ensure that no DM solitons and no stretched pulses but only dissipative solitons with steep spectral edges were observed. This was experimentally confirmed by the fact that no Gaussian pulse spectrum was obtained under any of the laser operation conditions. Self-started dissipative soliton operation was obtained in the laser through appropriately setting the orientations of the waveplates. Under strong pumping, multiple dissipative solitons were also generated in the laser. Figure 7.1(a-c) shows as an example a case of two dissipative solitons far apart in the cavity obtained with the pump power of 296 mW. The dissipative soliton feature of the pulses is evident from their characteristic steep spectral edge and the pump-dependent edge-to-edge spectral width. Unlike dissipative solitons formed in fiber lasers made of purely normal dispersion fibers [25, 34, 166], the short-wavelength side of the optical spectrum exhibits a Gaussian-like curve, which shows the influence of the cavity dispersion management on the dissipative soliton features. Figure 7.1(b) shows the corresponding autocorrelation trace. It shows that the FWHM of the pulse width is

about 2.06 ps if a Gaussian profile is assumed. The two dissipative solitons are far apart in the cavity as shown in Figure 7.1(c), and their peak power is measured as about 6.1W. Experimentally we have also observed a harmonic-three-mode locking state of the dissipative solitons.

Figure 7.1: (a), (d) Optical spectrum; (b), (e) autocorrelation trace; (c), (f) oscilloscope trace of a typical state of two dissipative solitons with large pulse spacing and a bound state of two dissipative solitons.

The detailed pulse parameters, such as the peak power and pulse width, vary with the pump power and the orientation settings of the waveplates. Through rotating the waveplates a state of bound dissipative solitons was also experimentally obtained. Figure 7.1(d-f) shows the bound state of the dissipative solitons observed. The optical spectrum of the state is shown in Figure 7.1(d), which has clear spectral modulations.

Unlike in Figure 7.1(a), obvious spectral spikes also appeared on the spectral edges. A high-resolution zoom-in on the spectrum is also displayed in Figure 7.1(d). Figure 7.1(e) shows the corresponding autocorrelation trace of the bound dissipative solitons. It shows three peaks with a height ratio of 1:2:1, indicating that the two dissipative solitons are identical and stay close to each other. Figure 7.1(f) is the corresponding oscilloscope trace, revealing only one group of two bound dissipative solitons existing in the cavity. The autocorrelation trace shows that the pulse width is about 183 fs if a Gaussian profile is assumed, which is far narrower than the pulse width measured when the dissipative solitons are far apart in the cavity as shown in Figure 7.1(b). It also gives out the pulse spacing of 6.1 ps. The peak power now is measured to be around 89.1 W. This bound state of dissipative solitons is stable. Once the bound state is formed, slightly changing the laser operation conditions, such as the pump strength or the orientations of the waveplates, does not change the pulse separation but only change the pulse peak power, suggesting that these dissipative solitons are tightly bound.

7.1.2 Simulation Results and Discussions

At first sight this result is out of our expectation, as the ratio between the measured pulse separation and the pulse width is 34.4, which excludes the possibility of effective direct pulse–pulse interaction between the dissipative solitons. To understand the formation of the bound dissipative solitons in the fiber laser, especially their strong coupling under large pulse separation, we conducted numerical simulations. The same model described in chapter 2 was used to do the simulation and the following parameters were adopted in accordance with the experimental

parameters: the GVD coefficients of the EDF and SMF are -50 and 18 (ps/nm)/km, respectively; the GVD of the DCF is -0.196 (ps/nm)/km; the cavity length $L = 0.5_{\text{SMF}} + 0.3_{\text{DCF}} + 0.1_{\text{SMF}} + 1.7_{\text{EDF}} + 2.2_{\text{SMF}} = 4.8$ m; output is 10%; and the gain bandwidth $\Omega_g = 16$ nm.

Stable self-started mode locking can always be obtained in the laser. Especially with a narrow gain bandwidth selection, multiple dissipative solitons were generated in our simulations as the pump strength was increased [94]. Like the experimental observations, the pulse parameters of the stable dissipative soliton numerically calculated depend strongly on the selection of the linear cavity phase delay bias, which corresponds to experimentally changing the orientations of the waveplates and the level of the cavity-induced pulse peak clamping. When the pulse peak is clamped at a low value, the formed dissipative soliton has relatively broader pulse width, while if it is clamped at a high value, the pulse width could become significantly narrow; in addition, spectral spikes also appear on the edges of the pulse spectrum, indicating that the spikes are formed due to strong pulse nonlinearity. Bound states of the dissipative solitons were numerically obtained under strong pulse nonlinearity. Figure 7.2 shows such a case obtained with the linear cavity phase delay bias set at 1.9π and the small signal gain coefficient $G = 3100 \text{ km}^{-1}$. We plotted the states at the cavity positions where the dissipative soliton has the maximum and the minimum pulse width. Due to dispersion management, the pulse width varies from about 1.8 ps to about 125 fs, while the pulse separation is fixed at about 4.9 ps.

Figure 7.2: Numerically simulated bound state of two dissipative solitons at the cavity position with (solid curve) the minimum pulse width and (dashed curve) the maximum pulse width

Based on the numerical simulation it becomes clear why tightly bound dissipative solitons could still be formed in the laser even when the experimentally measured pulse separation is far larger than 5 times the pulse width. Obviously, in other positions of the cavity, the pulse width is substantially broader than that measured at the laser output, which is near the position of minimum pulse width. The pulse profiles are directly overlapped at those cavity positions, which consequently results in direct pulse interaction. The dissipative soliton is essentially a chirped pulse, though along the cavity the pulse chirp also varies. Nevertheless, our numerical simulation shows that in spite of the pulse chirp variation and the strong pulse stretching and compression, the pulse separation between the bound dissipative solitons remains constant. In the case of our simulation the relative phase between the two bound dissipative solitons is fixed at about $\pi/4$ along the cavity. Our numerical simulation suggests that like the soliton formation in a laser, which is a result of the mutually balanced interaction among different effects over the cavity round-trip, the formation of the bound dissipative solitons is a result of direct pulse interaction over the cavity.

Finally, we note that the experimentally measured autocorrelation trace [Figure 7.1(e)] shows a clear pedestal, which is an indication of the partial pulse breaking frequently observed in DM fiber lasers under strong pulse nonlinearity. The feature is also reproduced in our simulation. With a large linear cavity phase delay bias setting, a high-peak-power dissipative soliton is formed in the laser. It has strong pulse nonlinearity and large nonlinear chirp. When the soliton is compressed in the negative dispersion fiber, pulse breaking occurs. Nevertheless, the pulse breaking is just a local effect; due to the cavity boundary condition, the pulse profile is then restored, and on average the pulse displays dissipative soliton features.

7.2 Bound State of Dissipative Solitons in a DM Fiber Laser with Large Anomalous Dispersion Mode-Locked by SWCNTs

Depending on the strength of the soliton interaction, bound states of solitons could be further categorized into two sorts: 1) tightly bound solitons formed under strong soliton-soliton interaction. These states of bound solitons have the characteristics that their soliton separations exhibit fixed discrete values, and the bound solitons have a fixed π or $\pi/2$ phase difference [113, 150, 151, 164]; 2) loosely bound solitons formed under weak soliton-soliton interaction. The loosely bound solitons, in contrast, have large soliton pulse separations and unfixed soliton phase differences. The states of tightly bound solitons were also called soliton molecules as they have the ability to reconstitute equilibrium separation after perturbation [167, 168]. Moreover, it was found that the tightly bound solitons can also act as a unit to form other bound states of

solitons.

In last section, we discussed the bound state of dissipative solitons in fiber lasers with large normal cavity dispersion mode-locked by NPR, while in this section we shift our attention to the formation and properties of the bound state of dissipative solitons in fiber lasers mode-locked with SWCNTs. Mode-locking fiber lasers with SWCNTs or their polymer derivatives is a new laser mode locking technique recently developed by Set et al. [105]. It has been shown that SWCNTs, as a new saturable absorber, have the advantages of easy fabrication, controllable modulation depth, and potentially broadband saturable absorption [59, 160-163]. Soliton operation of various SWCNTs mode-locked fiber lasers has been extensively studied [70, 162, 169, 170]. However, to our knowledge, no bound state of solitons has been observed in this type of mode-locked fiber lasers. In the following, a fiber laser mode-locked with SWCNTs saturable absorber was constructed to experimentally seek for the bound dissipative solitons and explore their related properties.

7.2.1 Experimental Observations

A similar fiber laser described in section 4.2 (Figure 4.4) was used for our experiment. The only difference is that the SMF has a length of 2.1 m and the DCF in the cavity is 7.2 m. Hence the net cavity dispersion was estimated to be around -0.01 ps^2 .

The fiber laser has a mode-locking threshold as low as 5 mW. Self-started mode-locking was easily achieved by increasing the pump power beyond the mode-locking threshold. Immediately after mode locking initially a single soliton was

always obtained. Figure 7.3 shows the optical spectrum and the autocorrelation trace of a typical soliton observed in the laser. The spectrum is characterized by the clear Kelly spectral sidebands, confirming its soliton nature. The 3 dB spectral bandwidth is measured to be 5.2 nm. The autocorrelation trace shows that the pulse has a sech-form profile with a FWHM of 782 fs. The time-bandwidth product is calculated to be 0.51, indicating a large frequency chirp, which denotes that the formed soliton is the dissipative soliton in anomalous dispersion regime. Checked with a high-speed photodetector (40 GHz) and sampling oscilloscope (50 GHz), we confirmed that there was only one soliton existing in the cavity.

Figure 7.3: Optical spectrum of a single-pulse soliton state in the fiber laser. The inset is the corresponding autocorrelation trace.

From the single soliton operation state, as the pump power was carefully further increased, a bound state of solitons then appeared. Figure 7.4 shows one of the bound solitons. The optical spectrum of the bound solitons is strongly modulated due to that two solitons are closely spaced in the time domain. The period of the spectral

modulation is 6.2 nm, according to the Fourier transformation it corresponds to a soliton separation of about 1.3 ps in the time domain. A strong spectral dip always appeared at the center of the spectrum, which indicates that the phase difference between the bound solitons is π . The insert of Figure 7.4 shows the autocorrelation trace of the state. Different from the oscilloscope trace, the autocorrelation trace could only provide the temporal information indirectly and roughly. Three peaks with a height ratio of 1:2:1 in the autocorrelation trace actually corresponds to the real case in time domain that two pulses with identical intensity and duration stay close to each other [151, 154], so Figure 7.4 denotes a bound state of two solitons. Partially overlapping of the three peaks is just because the two solitons are too close to each other. The autocorrelation trace also shows that the pulse separation is about 1.3 ps, which is in agreement with the spectral modulation period of 6.2 nm. Considering that the soliton has a pulse width of 782 fs, the soliton separation is about 1.7 times of the pulse width.

Figure 7.4: Optical spectrum of a state of tightly bound solitons observed in the fiber laser (The dotted curve is the single pulse spectrum for comparison). The inset is the corresponding autocorrelation trace. The pulse separation is 1.3 ps.

Depending on the laser operation conditions, e.g. the pump strength and orientations of the polarization controllers (PCs), bound states of solitons with another two soliton separations were also observed in the laser, as shown in Figure 7.5 and Figure 7.6. The soliton separations in the two states are 1.8 ps and 2.5 ps respectively, which are 2.3 and 3.2 times of the soliton pulse width. Here, we would like to emphasize that different from other tightly bound states of solitons achieved in our laser, whose two central maxima of the spectrum are identical, indicating that the phase difference is π , the two central maxima of the optical spectrum in Figure 7.6 are clearly unequal. According to the results shown in Ref. [154], this would indicate that the phase difference of the tightly bound solitons in Figure 7.6 is closer to $\pi/2$ rather than π .

Figure 7.5: Optical spectrum of another state of tightly bound solitons observed in the fiber laser (The dotted curve is the single pulse spectrum for comparison). The inset is the corresponding autocorrelation trace. The pulse separation is 1.8 ps.

Figure 7.6: Optical spectrum of another tightly bound soliton state in the fiber laser (The dotted curve is the single pulse spectrum for comparison). The inset is the corresponding autocorrelation trace. The pulse separation is 2.5 ps.

The states of bound solitons shown in Figure 7.4-7.6 have the common features that the phase difference of the bound solitons has a fixed value of π or $\pi/2$, and their peak-to-peak soliton separation is smaller than 5 times of the soliton pulse width. They belong to the category of so called tightly bound states of solitons observed in soliton fiber lasers mode-locked with other techniques. In our experiments no matter how the experimental conditions, such as the pump power, and the orientations of the PC, were changed, once the bound states of solitons were obtained, their soliton separations always fall into one of the three soliton separation values. No other values of soliton separation have been obtained, which indicates that the bound states of solitons have fixed discrete soliton separations. All of the above bound states of solitons were extremely stable. A tiny change in the pump power or orientations of the PC did not vary the pulse separation and their phase difference of π or $\pi/2$. If the laser operation condition was continuously changed so that the variation could not be treated as a perturbation, the current tightly bound soliton states would be destroyed

and jump to another tightly bound soliton state (with another pulse separation) or to a multiple soliton state where solitons are far apart separated in the cavity. It is this jump behavior between the tightly bound states of the solitons, more precisely between the discrete pulse separations, that reinforces the strong stability of the tightly bound state of solitons from the side.

Figure 7.7: Optical spectrum of a bound state of two bound solitons in the fiber laser (The dotted curve is the single pulse spectrum for comparison). The inset is the corresponding autocorrelation trace.

Figure 7.8: Optical spectrum of a three-pulse tightly bound solitons state in the fiber laser (The dotted curve is the single pulse spectrum for comparison). The inset is the corresponding autocorrelation trace. The pulse separation between two adjacent solitons is 2.5 ps.

Moreover, like the tightly bound solitons observed in other lasers, these tightly bound solitons can also behave as a unit to form other bound states of them. Figure 7.7 shows for example such a case. Two identical tightly bound solitons with a soliton separation of 2.5 ps, which is one of the above three values, have formed a bound state of them with a separation of 9 ps. Apart from the two-pulse tightly bound solitons, even states of three-pulse tightly bound solitons were occasionally observed in the laser, as shown in Figure 7.8. The autocorrelation trace shows that three identical solitons with equal pulse separations of 2.5 ps between the neighboring solitons are bound together tightly. In contrast to the central dip in the spectrum of two-pulse bound solitons with phase difference of π , now the spectrum exhibits a small hump in the center, which indicates that the phase difference between the adjacent pulses is still π [171]. To our surprising, although multiple or multi-pulse tightly bound solitons could be formed in the laser cavity, it was found that all these tightly bound solitons had the same pulse separation between the adjacent solitons. So far we have not observed the situation where two tightly bound solitons with different pulse separations coexist in our laser. The experimental result suggests that under a concrete laser operation condition only one state of the tightly bound solitons is stable.

Figure 7.9 shows for example a loosely bound state of two solitons observed in the laser. Different from the states of tightly bound solitons, solitons in the states always have relatively large peak-to-peak separations, e.g. the soliton separation in Figure 7.9 is about 10 ps, which is in good agreement with the period of spectral modulation of 0.8 nm, and is almost 13 times of the soliton pulse width. With large soliton

separation the direct soliton-soliton interaction becomes weak. The states of bound solitons are known as the loosely bound solitons. For the loosely bound solitons, the phase difference between the solitons is no longer fixed at the π or $\pi/2$, but varied from case to case. As we have shown in Figure 7.9, as the laser operation conditions are slightly shifted, other phase relationship between the solitons could also be obtained. Moreover, experimentally these bound states are found to be weakly stable. They could be easily destroyed by any mild environmental perturbations. When the pump power was carefully changed, we found that the soliton separation also changed almost continuously with the variation of the pump strength, which indicates that the loosely bound solitons do not have fixed discrete soliton separations.

Figure 7.9: Optical spectra of two states of loosely bound solitons with different phase relationship. The inset on left is the zoom-in of the optical spectra. The spectral center is 1558.6 nm. The inset on right is the corresponding autocorrelation trace. The pulse separation is about 10 ps.

7.2.2 Discussions

Although two types of bound solitons with distinct properties and features were observed, we believe that they could have essentially the same physical origin. They are formed due to the direct soliton interaction in the fiber laser. It is well known that the strength of soliton interaction decreases exponentially with the soliton separation. For the tightly bound solitons observed, the soliton interval is always smaller than 5 times of the soliton pulse width. It means that the solitons have strong direct interaction force, which predominately determines the feature of the bound solitons. For the loosely bound solitons, the soliton interval is more than 10 times of the soliton pulse width. The bound solitons have relatively much weaker direct interaction force. We note that in a real fiber laser apart from the direct soliton interaction, other effects, such as the dispersive waves mediated soliton interaction, environmental perturbations etc. always exist and influence the direct soliton interaction. These effects are of random nature and normally weak. Under strong direct soliton interaction, the influence of these effects could be ignored. However, when the strength of the direct soliton interaction becomes not sufficiently strong, their effects could no longer be ignored but clearly affect the features of the formed bound solitons.

7.3 Conclusion

In this chapter, bound states of dissipative solitons were firstly experimentally observed in a DM EDF laser with large normal cavity dispersion. We have also shown that although they are chirped pulses and the pulses stretch and being

compressed strongly along the cavity, the bound dissipative solitons observed still exhibit a fixed pulse separation, which is invariant with the pump power. Numerical simulation well reproduced our experimental observations, and revealed that the relative phase difference between dissipative solitons in the bound state remains unchanged in the cavity. More important, our study indentified that it was sufficient to form tightly bound pulses, provided that over part of the cavity round-trip the pulse profiles are overlapped.

On the other hand, we have experimentally observed bound states of dissipative solitons in a fiber laser mode locked with SWCNTs for the first time. Both states of the tightly and loosely bound solitons are observed. Like the bound soliton states observed in other mode-locked fiber lasers, the states of tightly bound solitons are characterized by that the bound solitons have a fixed phase difference of π or $\pi/2$, and fixed discrete pulse separations, and the states are ultra stable, while the states of loosely bound solitons no longer possess these features. Their soliton separation and soliton phase difference vary with the experimental conditions, and the states are sensitive to the environmental perturbations. Our experimental results also show that by employing appropriate SWCNTs as the mode locker, various bound soliton states could be obtained in the same fiber laser. This makes a systematic experimental study on them and a comparison of their properties possible, which is generally difficult to realize in experiment. It is well known that based on nonlinear Schrödinger equation (NLSE), two solitons could not bind together to form a stable bound state of solitons. However, Malomed pointed out that the case was different for GLE solitons. Since the GLE soliton has an oscillating tail, when two GLE solitons slightly overlap with each other, their interaction potential will have spatial local minima, which eventually

makes the formation of stable bound states of solitons possible. Considering that the formation and dynamics of solitons in mode-locked fiber lasers are well described by the GLE, we come to the conclusion that the bound soliton emission is an intrinsic feature of the mode-locked soliton fiber lasers, and our result is important for the better understanding on the bound states of solitons.

Chapter 8

Vector Dissipative Solitons

Properties and kinetics of the dissipative solitons have been elaborately investigated in chapter 3 to chapter 7. However, emphasis has been mainly given to the scalar dissipative solitons. It is well known that owing to the technical limitation in fabricating a perfect-circular fiber core and/or the random mechanical stresses, practical single mode fibers (SMFs) always support two polarization eigenmodes. Except when polarization-maintaining fibers are used, a fiber laser cavity is always weakly birefringent, and soliton generation and propagation are subject to the influence of cavity birefringence. For fiber lasers mode-locked with nonlinear polarization rotation (NPR) technique, since a polarizer is inserted in the cavity for achieving the artificial saturable absorption effect, the formed solitons have fixed polarization state once they pass the polarizer. Consequently, the polarization states of solitons are fixed by the polarizer and no polarization dynamics could be displayed. Provided that no polarization sensitive component is employed in the cavity, such as fiber lasers mode-locked by semiconductor saturable absorber (SESAM), carbon nanotubes (CNTs) or graphene, the polarization states of the formed soliton could freely evolve and instead vector solitons are achieved directly as a result of the cavity

birefringence. A vector soliton differs from a scalar soliton by that it consists of two orthogonal polarization components, which are nonlinearly coupled.

Vector solitons capture great attentions nowadays on account of their promising applications in fiber-optical telecommunications as information carrier for beyond a binary system, and in optical signal processing for the collision-based information switches [172]. The additional polarization freedom further makes them presenting richer and more fascinating features than scalar solitons. Depending on the strength of cavity birefringence, two orthogonal polarization components could couple each other in distinct ways so that different types of vector solitons are formed, such as the polarization-locked [173], polarization-rotating vector solitons [174] and even the soliton trapping [175, 176]. All of them are indeed observed in mode-locked fiber lasers [177-179]. Dissipative solitons formed in the normal dispersion are different from nonlinear schrödinger equation (NLSE) solitons in anomalous dispersion not only by their completely different formation mechanism, but also the distinct soliton properties. They are strongly frequency chirped. It would be interesting to find out whether the group-velocity-locked and even the polarization-locked vector dissipative solitons (VDSs) could still be formed in fiber lasers in spite of the strong frequency chirp. Moreover, how the dissipative nature further affects the polarization states of these VDSs, and even the interaction forces between them also deserves to be explored.

Chapter 8 focuses on the formation and features of VDSs in fiber lasers with normal dispersion, and it is organized as following. In section 8.1, we show that in spite of the large frequency chirp, both the polarizations locked and rotating VDSs could still be

obtained in fiber lasers experimentally and numerically. Section 8.2 further explored the polarization dynamics of these polarization-rotating VDSs, and also emphasized their difference from those of NLSE solitons in anomalous dispersion. Multiple VDSs are also achieved with stronger pumper power, and under appropriate cavity birefringence, they could even bound together to form bound state of VDSs. The properties of these bound states of VDSs are present in section 8.3, which is followed by a conclusion in section 8.4.

8.1 Polarization-Locked VDSs and Polarization-Rotating VDSs

Taking account of the fiber birefringence, two orthogonal polarization modes of a pulse generally have unequal group velocities as it propagates, which leads to temporal pulse splitting and tends to prevent the formation of a soliton with energy in both polarization modes. However, Menyuk pointed out that the nonlinear coupling between these two modes could balance above polarization-mode induced pulse-broadening, and finally a stable pulse composed of both polarization components could propagate as a single vector soliton unit in an optical fiber overcoming their group velocity difference [175, 176]. Such a soliton was also known as a group-velocity-locked vector soliton, since two components have nonlinearly induced common group velocities. Later, Akhmediev et al. also demonstrated that for certain values of fiber birefringence and pulse energy, polarization-locked vector solitons could even be formed as a result of the fact that two polarization modes have the same phase velocities [180, 181]. Polarization-locked vector solitons refer to the vector solitons whose polarization states remain unchanged during the propagation,

and such solitons are formed because the nonlinear birefringence created by the unequal amplitudes of two components counteracts the linear fiber birefringence, which further results in the dynamic equalization of two phase velocities. In addition, polarization-rotating vector solitons were also theoretically predicted [174].

Although the polarization-locked vector solitons [181], polarization-rotating vector solitons [182], and polarization-locked black–white vector solitons [183], were theoretically predicted in weakly birefringent SMF, except for the soliton trapping, most of the vector solitons in SMF were not experimentally confirmed. A challenge for experimentally observing the polarization-locked vector solitons in SMFs is that the birefringence of the fibers must be kept small over a long distance, which in practice is difficult to realize. In contrast, both the polarization-locked vector solitons and the polarization-rotating vector solitons were experimentally observed in mode-locked fiber lasers [177, 178, 184]. A main difference of soliton propagation in a fiber laser cavity from that in a SMF is that the cavity boundary condition should also be satisfied. A consequence of the difference is that the features of the laser solitons are not determined by the localized cavity parameters but by the averaged cavity parameters. Experimentally one can easily control the averaged birefringence of a fiber laser cavity to be near zero. In this part, in order to gain an insight into the VDS operations, we constructed a fiber laser with large normal cavity dispersion mode-locked by SESAM to do the experiments.

8.1.1 Experimental Observations

The schematic of the fiber laser is shown in Figure 8.1. It has a dispersion-managed (DM) cavity with a net positive cavity dispersion of about 0.04 ps^2 . The ring cavity has a total length of around 12 m, consisting of 2.7 m EDF with a group velocity dispersion (GVD) parameter of about -32 (ps/nm)/km , 4.3 m standard SMF with a GVD parameter of about 18 (ps/nm)/km , and 5 m dispersion-compensating fiber (DCF) with a GVD parameter of about -3.4 (ps/nm)/km . Now a SESAM was employed to mode lock the laser instead of the NPR technique to allow the formation of vector solitons. It has a saturable absorption modulation depth of 8%, and a recovery time of 2 ps. The central absorption wavelength of it is around 1550 nm. A 3-port polarization independent circulator was used to incorporate the SESAM into the cavity and simultaneously force the unidirectional operation of the ring. Note that within one cavity round-trip the pulse propagates twice in the fiber between the circulator and SESAM. A polarization controller (PC) is inserted in the cavity to fine-tune the cavity birefringence. The laser emission was output through a 10% fiber coupler.

Figure 8.1: Schematic of the vector soliton fiber laser. WDM: wavelength division multiplexer; EDF: erbium-doped fiber; ISO: isolator; PC: polarization controller; SESAM: semiconductor saturable absorber mirror.

In fact, the different net cavity GVD of the laser could be realized by controlling the length of the SMF used. We first keep the length of SMF relative long so that our laser has a large net negative cavity GVD. After mode-locking of the fiber laser, solitons were immediately formed in our cavity, whose spectra were characterized by clear Kelly sidebands confirming their soliton nature and the anomalous cavity dispersion. With the help of an external polarizer outside the cavity, the solitons were found to exhibit all the typical features of the vector solitons reported in optical fibers [185]. Then we cut back the length of the SMF to manage the cavity dispersion. With the decrease of SMF length, the net cavity GVD of our laser was finally shifted to a large positive value. Consequently, dissipative solitons are formed instead of the conventional solitons. Figure 8.2 shows such an example obtained with the pump power of 68 mW, which displays the common features of the dissipative solitons in normal dispersion as obtained in all previous chapters. The soliton spectrum exhibits steep spectral edges with a gain-band-limited bandwidth. Due to the large normal dispersion, the soliton broadens so much that its pulse width is not suitable to be measured by the autocorrelator any more. A 50 GHz high-speed oscilloscope together with a 45 GHz photodetector was used to measure the pulse profile of the vector soliton, and it gives out the pulse width of 50 ps. The 3dB-bandwidth of the soliton is 13.6 nm, and the time-bandwidth product of the soliton is 85, indicating a strong frequency chirp. However, different from the dissipative solitons formed in the fiber lasers mode-locked with the NPR technique, the soliton consists of two orthogonal polarization components.

Figure 8.2: A typical state of a polarization-rotating VDS emission state of the laser.
 (a) Spectrum; (b) Pulse profile; (c) Pulse train after passing through a polarizer.

To highlight the vector nature of the soliton, we let the soliton pulse train pass through a rotatable external cavity polarizer and compared the features of the pulse train before and after passing through the polarizer, either with a high speed oscilloscope or an RF-spectrum analyzer. It turns out that the soliton shown in Figure 8.2 is

polarization-rotating vector soliton. The polarization rotation of the soliton could be easily identified e.g. by the oscilloscope trace measurement. Without passing through the external cavity polarizer, the soliton pulse had identical pulse intensity on the oscilloscope trace for each cavity roundtrip, while after passing through the polarizer the intensity became nonuniform with the cavity roundtrips as shown in Figure 8.2(c). It indicates that the polarization of the soliton rotated along the cavity.

Controlling the linear cavity birefringence through an intra-cavity PC, polarization-locked vector solitons were also obtained. Figure 8.3 shows such an example. A polarization-locked vector soliton has the characteristic that it has a fixed polarization during circulation in the laser cavity. Consequently, after passing through an external cavity polarizer the pulse heights of such vector solitons on the oscilloscope trace still have identical values as shown in Figure 8.3(c). For a polarization-locked vector soliton, its optical spectra and pulse profiles along the long and short polarization ellipse axes could be measured by the following method. We monitored the signal after the external cavity polarizer. Guided by the pulse intensity in the oscilloscope, we first orientated the external polarizer to a direction where the pulse intensities have maximum values. We denoted the direction as zero degree polarizer orientation and recorded the spectrum and pulse profile as one component. With all other parameters fixed we then turned the external polarizer by 90° and recorded the spectrum and pulse profile as the orthogonal component. Figure 8.3(a) and (b) showed the results. Different from the polarization resolved vector solitons formed in the net negative cavity GVD fiber lasers, no four-wave-mixing spectral sidebands could be identified, which was believed to attribute to the large frequency chirp of the soliton. The two resolved spectra have comparable spectral intensities and

almost the same central wavelengths, but clearly different spectral distributions. The fact that no polarization evolution frequency was experimentally observed further confirms that the two components were polarization locked [185].

Figure 8.3: A typical state of a polarization-locked VDS emission state of the laser. (a) Total and polarization resolved optical spectra; (b) Total and polarization resolved pulse profile; (c) Pulse train along two orthogonal polarization directions.

Depending on different settings of the orientations of PC, the central wavelength of the VDSs was observed to be tunable from 1568 nm to 1575 nm. The longer the central soliton wavelength is, the less stable the formed VDS turns into. This could attribute to the fact that the central absorption wavelength of the SESAM is around 1550 nm, which favors the formation of the VDS centered near 1550 nm. Multiple VDSs with identical soliton parameters were also obtained under strong pumping.

8.1.2 Numerical Simulations

To gain an insight into the VDS formation, we further conducted the numerical simulation with the pulse tracing method. To make the simulation comparable with the experimental results, we used the following experimental cavity parameters: $\gamma = 3 \text{ W}^{-1}\text{km}^{-1}$, $\Omega_g = 16 \text{ nm}$, $P_{\text{sat}} = 50 \text{ pJ}$, $k''_{\text{SMF}} = -23 \text{ ps}^2/\text{km}$, $k''_{\text{EDF}} = 41 \text{ ps}^2/\text{km}$, $k''_{\text{DCF}} = 2.6 \text{ ps}^2/\text{km}$, $k''' = -0.13 \text{ ps}^3/\text{km}$, $E_{\text{sat}} = 35 \text{ nJ}$, $l_0 = 0.2$, $T_{\text{rec}} = 2 \text{ ps}$, cavity length $L = 23.2 \text{ m}$, and beat length $L_b = 100$.

With current selection of cavity parameters, it was numerically found that the formation of VDS in the laser is severely the gain bandwidth dependent. VDS is unable to be formed for large gain bandwidth, such as 24 nm, and the spectrum of the mode-locked pulses has a Gaussian profile in that case. Stable VDS could only be formed with the gain bandwidth narrower than 16 nm. The smaller the gain bandwidth used for calculation, the narrower the spectrum bandwidth of the VDS achieved. Figure 8.4(a) shows a typical case of the calculated VDS evolution with the cavity roundtrips. It was observed that the total pulse intensity remained unchanged for each

round trip, but intensity of the horizontal and the vertical component exhibited out-of-phase intensity variation, indicating that coherent energy exchange between them still existed [186]. Such an out-of-phase intensity variation between the vector soliton components disappeared provided that the four-wave-mixing terms were moved out from Eq. (2.26). The calculated pulse width is about 20 ps and its spectrum bandwidth is about 7 nm, suggesting that the formed VDS is strongly frequency chirped.

(a)

(b)

Figure 8.4: (a) Combined pulse intensity evolution; (b) Corresponding optical spectra numerically calculated.

The result shown in Figure 8.4 was obtained with the cavity birefringence $L/L_b = 0.232$. The VDSs formed have two orthogonal polarization components with coincident central wavelength and comparable intensity. The intensity difference between two polarization components was numerically observed to become larger with the increase of the cavity birefringence. Wu et al. have used the pattern traced out by the electric-field vector tip as a function of time in a fixed transverse plane to describe the polarization state of the soliton pulse in Ref. [187] and demonstrated its reliability. Here we employ the same method to numerically investigate the polarization evolutions of the formed VDSs. Both of the polarization-locked and polarization-rotating VDSs were numerically obtained in terms of different cavity birefringence. The calculation based on the peak point of the VDS showed that a polarization-locked VDS fixed its polarization ellipse orientation as it circulated in the cavity, while the polarization ellipse of a rotating VDS changed its orientation. For a polarization-locked VDS, the phase between its two polarization components was always fixed at $\pi/2$ at the pulse peak, but deviated from $\pi/2$ at other points of the pulse, exhibiting the influence of the frequency chirp. To further determine how the frequency chirp of the solitons affects their polarization, the polarization states at different points of a pulse have also been investigated, e.g. at the pulse's tail and FWHM points. It turned out that the orientations of the polarization ellipses calculated at different points were slightly different, but the difference remained the same along the propagation, which denoted that although the polarization-locked VDSs were strongly chirped, their two orthogonally polarized components were still phase locked.

8.2 Polarization Dynamics of the Polarization-Rotating VDSs

Although in section 8.1, we have shown clearly that both of polarization-locked and polarization-rotating VDSs could still be formed in fiber lasers mode-locked by SESAM in spite of the large frequency chirp, their elaborate properties such as the polarization rotation dynamics of the observed polarization-rotating VDSs have still undressed. Note that different from propagating in single mode fibers, light propagation in a laser cavity must also satisfy the cavity boundary condition. Previous studies on solitons formed in lasers have shown that cavity boundary condition could synchronize the soliton pulsation dynamics [188]. In fact, recently Zhao et al. have shown that the vector soliton formed in anomalous dispersion cavity can rotate its polarization with a period locked at multiple of cavity roundtrip time. They further demonstrated that if multiple vector solitons were formed, their instantaneous polarization ellipse orientations are orthogonal to each other as a result of the gain competition in EDF [189]. Considering that a dissipative soliton is strongly frequency chirped in contrast to a conventional soliton, which may weaken the laser gain competition, it would be interesting to find out whether the period of the polarization rotation of the VDS could still be locked to multiple of cavity roundtrip time and the polarization ellipse orientation of the newly formed VDS is orthogonal or parallel to that of the previous one.

8.2.1 Experimental Observations

We employed a similar fiber laser as that in section 8.1 (Figure 8.1) to do the

experiment. Now the cavity has a length of 37.4 m, consisting of 1.5 m EDF with GVD of -32 ps/km/nm, SMF with GVD of 18 ps/km/nm and DCF with GVD of -2 ps/km/nm. Mode-locking of the laser is still achieved with a SESAM, which has a saturable absorption modulation depth of 5%, and a recovery time of 10 ps. The laser was output through a 10% fiber coupler. The fiber laser has a DM cavity, and DCF with different lengths was inserted between the circulator and the SESAM to control the net cavity dispersion. To achieve dissipative soliton operation, a length of 32.4 m DCF was finally adopted so that the net cavity GVD was a large positive value $\sim 0.0652 \text{ ps}^2$. The cavity birefringence has also been elaborately managed to obtain the polarization-rotating vector soliton.

Solitons were easily achieved by increasing the pump power above the mode-locking threshold. To investigate the vector nature of the soliton, we still employed a rotatable external cavity polarizer and monitor the output pulses before and after passing through it by a high speed oscilloscope or RF spectrum analyzer, as we did in section 8.1. After fine tuning the PC, which corresponds to elaborately varying the cavity birefringence, it was observed after passing through the external polarizer the original uniform pulse train became nonuniform, and simultaneously the strong extra sidebands excluding harmonics of the cavity repetition frequency began to appear in the RF spectrum, indicating that the polarization-rotating vector solitons were formed in our cavity. Multiple solitons were always obtained once the pump power increased above mode-locking threshold. Carefully decreasing the pumping strength, the solitons can be reduced one by one. Figure 8.5 shows for example the case of only one pulse in the cavity. Figure 8.5(a) shows the autocorrelation trace and optical spectra. The dissipative soliton feature is confirmed by the steep spectral edges, and the large

frequency chirp, resulting from a time-bandwidth product of 9.45. Figure 8.5(b) and (c) show the oscilloscope traces of the soliton pulses measured before and after the external polarizer respectively. Obviously before the polarizer, all the pulses in the trace have equal intensity, while after passing through it the pulse intensity then alternated between two values, indicating that the polarization of the soliton is rotating in the cavity and the period of the soliton polarization rotation is 2 (P2) times of the cavity roundtrip time.

Figure 8.5: Optical spectra measured without passing and passing through a polarizer and autocorrelation trace (a) and corresponding oscilloscope traces (b) without polarizer; (c) with polarizer.

Starting from the single P2 VDS, further increasing the pump power leads to the formation of multiple VDSs, and their polarization rotation was also observed to be synchronized with 2 times of the cavity roundtrip time. In other words, once the polarization rotation of the vector soliton synchronizes with the cavity length, increasing the pumping strength would not destroy the synchronization but only change the number of the solitons. Figure 8.6(a) and (b) show the cases of two VDSs coexisting in the cavity after passing through the external polarizer. The state in Figure 8.6(a) was obtained by carefully increasing pump strength from the state in Figure 8.5. It shows clearly that these two solitons have the same polarization rotation with a common periodicity of 2 times of the cavity roundtrip time, reflected by their pulse intensity returning to the previous value after every 2 cavity-roundtrips. However, two different pulse intensities within one cavity roundtrip time shown in Figure 8.6(a) indicate that such two vector solitons actually have different polarization ellipse orientations. To further determine the relationship between two polarization ellipse orientations, we studied the relative soliton intensity variation by rotating the external polarizer. It turned out that when the pulse intensity of one soliton increased, the pulse intensity of the other decreased. When the pulse intensity of one soliton reached the maximum, the pulse intensity of the other was almost at the minimum, indicating that the polarization ellipses orientations of these two solitons were almost orthogonal to each other. Although two vector solitons with orthogonal polarization ellipse orientations coexist, the relative positions of the VDSs are stable in the cavity, indicating that all solitons move with exactly the same group velocity. Keeping all other experimental parameters unchanged, but only fine tuning the PC, the state of harmonics-2 VDSs was obtained, as shown in Figure 8.6(b). It shows clearly that the

periodicities of the polarization rotation of two solitons are still the same and locked to 2 times of the cavity roundtrip time. However, different from Figure 8.6(a), equal pulse intensity of two solitons within one roundtrip suggests that they have almost the same polarization ellipses orientation. The fact that the pulse intensity of these two solitons increased or decreased synchronously with rotating the external polarizer, and reached the extremum almost at the same time further confirmed our result.

Figure 8.6: Oscilloscope traces of two VDSs passing through a polarizer with their instantaneous ellipse orientations almost (a) orthogonal, and (b) parallel.

The similar case in Figure 8.6(a) has been observed previously in fiber lasers with anomalous dispersion by Zhao et al. and the explanation of the formation mechanism was also given out [189]. They pointed out that as the gain was polarization independent in the fiber laser, the polarization rotation of a vector soliton along the cavity, which was actually a kind of swing of the polarization ellipse with respect to a certain direction in the fiber cross section [187], would cause a spatial gain hole burning in the cross section. When two vector solitons were formed in cavity, as a result of the laser gain competition new formed soliton tended to have polarization ellipse orientations orthogonal to the initial ones. Note that in our cavity the angle of the polarization ellipse orientations between the newly formed pulse and the original one is not exactly 90° but has a tiny deviation compared to that in anomalous dispersion fiber lasers. We believe this is because VDS is strongly chirped, so the frequency component of the pulses varies with the time, which results in the gain competition effect weaker than that of transform-limited vector solitons and further leads to such an angle deviation.

Apart from the state in Figure 8.6(a), two polarization rotation locked VDSs with almost the parallel polarization ellipse orientations have also been achieved in our cavity, e.g. Figure 8.6(b). No such a case has been observed previously. We believe that the formation of such a state in our cavity may be related to the weak gain competition between the dissipative solitons. In addition, the pulse interval (~ 93 ns) between the solitons shown in Figure 8.6(b) is much larger than that (~ 33 ns) of the solitons in Figure 8.6(a), which could also lead to the weaker gain competition than that of the state in Figure 8.6(a) where two VDSs have orthogonal orientations of their polarization ellipses.

From the states in Figure 8.6(a) and (b), reducing the pump power, they could return to P2 state with only one soliton in the cavity as shown in Figure 8.5. However, further reducing the pump strength carefully, it could jump from the current P2 state to P1 state, represented by that all the solitons had identical pulse intensities in the oscilloscope trace either with or without passing through the polarizer. It is well known that the polarization rotation of a VDS along a weak birefringence fiber is caused by both the linear and nonlinear birefringence of the fiber. Moreover, the nonlinear polarization rotation is proportional to the soliton intensity. Therefore, the polarization rotation is pump power strongly dependent. Reducing the pump power leads to less nonlinearity polarization rotation accumulated and thus the meta-stable P2 state can possibly turn into the stable P1 state.

Experimentally, soliton polarization rotation locking to the 3 (P3) and even 4 (P4) times of the cavity roundtrip time have also been observed. The states were obtained by carefully tuning the PC, corresponding to changing the net linear cavity birefringence in physics. Like the case of P2 states, increasing pump strength did not destroy the polarization rotation locking state but increase the number of solitons in cavity. P3 and P4 state with multiple solitons have also been observed in our laser.

8.2.2 Numerical Simulations

To gain an insight into polarization dynamics of the polarization-rotating VDSs, we also numerically simulated the operation of the laser. To make the simulation comparable with the experimental results, we used the following experimental cavity

parameters: $\gamma = 3 \text{ W}^{-1}\text{km}^{-1}$, $\Omega_g = 16 \text{ nm}$, $P_{\text{sat}} = 50 \text{ pJ}$, $k''_{\text{SMF}} = -23 \text{ ps}^2/\text{km}$, $k''_{\text{EDF}} = 41 \text{ ps}^2/\text{km}$, $k''_{\text{DCF}} = 2.6 \text{ ps}^2/\text{km}$, $k''' = -0.13 \text{ ps}^3/\text{km}$, $E_{\text{sat}} = 35 \text{ nJ}$, $l_0 = 0.2$, $T_{\text{rec}} = 2 \text{ ps}$, and cavity length $L = 5.2 \text{ m}$. Under the current simulation parameters, the numerically calculated spectrum exhibits the characteristic steep spectral edges, one typical feature of the dissipative solitons formed in normal dispersion. The FWHM width of the pulse obtained is about 20 ps and spectrum bandwidth is about 6 nm, which matches well with our experimental observation and suggests that the formed dissipative soliton is strongly chirped.

Figure 8.7: Polarization ellipses of VDSs (including the peak, FWHM and tail points' polarization ellipses) at the same location inside the ring cavity under different cavity birefringence: (a) $L/L_b=0.05$, and (b) $L/L_b=0.5$.

It was numerically found that the polarization rotation depended strongly on the cavity birefringence. When the cavity birefringence was set to be a small value, the polarization-locked VDS was numerically obtained, which was also known as P1 state. Figure 8.7(a) shows such a state calculated with the birefringence $L/L_b = 0.05$. The electrical-field tips of the peak only trace out one polarization ellipse, denoting that the polarization state of the VDS is the same every round. Considering that VDS is strongly chirped, to gain an insight into the influence of the chirp on the polarization rotation, we also investigated the polarization states of the VDS at other points besides the peak such as its FWHM and tail. Figure 8.7(a) shows that the electrical-field tips of the FWHM point trace out a concentric ellipse with the same orientation as that of the peak while the ellipse orientation of the tail points differs slightly. It can be inferred that not only the peak point but every point of VDS is P1 state. The chirp exerts little impact upon the period of polarization rotation of P1 VDS.

With other cavity parameters unchanged, but only increasing the cavity birefringence, the VDS with the polarization rotation period locked to two times of the cavity roundtrip was achieved. Figure 8.7(b) shows such a P2 state calculated with the birefringence $L/L_b = 0.5$. It displays clearly that the electrical-field tips trace out two different polarization ellipses in odd and even rounds respectively. Except that the angles of the orientations of these two polarization ellipses are supplementary to each other, these two polarization ellipses are exactly the same, suggesting that polarization ellipse only changes its orientation every one round trip but return to the previous state every two rounds. This case agrees well with the experimental results shown in

Figure 8.5. Further increasing the cavity birefringence, the VDSs, whose polarization rotation period was locked to multiple cavity roundtrip time such as two, five, six and even fifteen rounds, have also been achieved. Starting from the case in Figure 8.7(b), when the cavity birefringence is increased to $L/L_b = 5$, electrical-field tips obtained depict chaotic shape, which is impossible to determine the number of roundtrip time needed for the polarization state recurrence to its previous orientation. Moreover, the decrease of the laser gain can also transfer P2 state to P1 state, revealing that reduction of the laser gain produces the same qualitative trend as lowering the cavity birefringence on the polarization rotation. This is in good agreement with our experimental observation.

8.3 Bound State of VDSs

Studies in last two sections of this chapter show clearly that in spite of the large frequency chirp, both of the polarization-locked and polarization-rotating VDSs could still be formed in fiber lasers, similar as that in anomalous dispersion regime. Considering that bound state of vector solitons in anomalous dispersion and bound state of the dissipative solitons in normal dispersion have both experimentally observed in fiber lasers previously, whether the VDSs could also bind one another to form stable bound state of VDSs evokes our interests. As a matter of fact, the multiple-pulse operation of a Ti:sapphire laser mode-locked by a SESAM has been investigated by Lederer et al. [190]. They reported the observation of the multiple pulses with either large separation or close separation, as well as having their relative phase rotating or near 0 or π . However, the emphasis was given to the explanation of these observations as well as of the mechanisms involved in the transitions from

single to multiple-pulse states using the Ginzburg-Landau master equation, while the vector properties of these multiple-pulses remain untouched. In this section, we focus our attention on searching for the existence of this distinct bound state of VDSs.

We still employed the setup in section 8.1 to do the experiment. Through carefully varying the orientations of PC to control the cavity birefringence, polarization-locked VDSs were easily obtained after mode locking of the fiber laser. Multiple polarization-locked VDSs were always generated in cavity with the increase of the pump power. These multiple polarization-locked VDSs were observed to form various stable patterns in terms of their different separations. They could either distribute themselves in the cavity randomly with stable relative positions, or with equal temporal intervals, which is also known as the harmonic state.

Figure 8.8 shows an example of a case of two VDSs far apart in the cavity. The steep spectral edges with a gain-bandwidth-limited bandwidth confirm the dissipative soliton nature in normal dispersion. The FWHM of the pulse width is measured to be around 52 ps, if a sech^2 -form profile is assumed. The time-bandwidth product is 85, indicating a strong frequency chirp, which is another characteristic of the dissipative solitons formed in normal dispersion. Their vector nature was verified by consisting of two orthogonal polarization components. The fact that after passing through an external cavity polarizer the soliton pulses in the oscilloscope trace still possess equal intensities indicates that the VDSs are polarization-locked ones. Figure 8.8(a) and (b) display the polarization resolved spectra and the corresponding pulse profiles along its long and short polarization ellipse axes respectively, which are measured by the same method as described in section 8.1. The spectra show clearly that the two polarization

components have almost the same central wavelengths, but significantly different spectral intensities and profiles. The pulse train of two orthogonal polarization components confirms that there are indeed two solitons coexisting in the cavity and staying far away from each other.

Figure 8.8: (a) Optical spectrum (total and polarization resolved); (b) Pulse profile (total and polarization resolved); (c) Pulse train of a typical state of two VDSs with large pulse spacing

It was experimentally observed that by carefully tuning the orientations of the PC, the separation of the VDSs could be changed. Under appropriate settings of the pump power and PC, two VDSs could stay so close to each other that they finally form a stable bound state of VDSs. Figure 8.9 shows the bound state of the VDSs observed. Both of the total and polarization resolved optical spectra exhibit a similar spectral profile with those in Figure 8.8(a) except for a spectral modulation, resulting from the fact that two solitons stay close to each other in time domain. The insert in Figure 8.9(a) further shows the high resolution zoom in of the spectral modulation. The modulation period is 0.09 nm, corresponding to a pulse separation of 93 ps in time domain. Figure 8.9(b) shows the high-speed oscilloscope traces. Along each of the two orthogonal polarization directions two identical pulses with about 93 ps separation are clearly visible. The pulse separation well matches the period of the high frequency spectral modulation shown in Figure 8.9(a). Due to the broad pulse width reaching the measurement limitation of the autocorrelator (FR-103MN), the pulse width of each soliton (FWHM) was still measured by the high-speed oscilloscope, which was given as ~ 36 ps. The peak-to-peak soliton separation is smaller than 5 times of the soliton pulse width. The experimental results suggest that the state shown in Figure 8.9 is a case of the bound state of VDSs, where two polarization-locked VDSs as shown in Figure 8.8 are bound together. Figure 8.9(c) confirms that only one group of two bound VDSs exists in the cavity. This bound state of VDSs is stable. Once the bound state is formed, slightly changing the laser operation conditions, such as the pump strength or the orientations of the PC, does not change the pulse separation but only change the pulse intensity, suggesting that these VDSs are tightly bound.

Figure 8.9: (a) Optical spectrum (total and polarization resolved); (b) Pulse profile (total and polarization resolved); (c) Pulse train of a typical bound state of two VDSs

8.4 Conclusion

In this chapter, we showed that by replacing the NPR mode-locking technique with SESAM, VDSs were experimentally observed in a fiber laser with large positive net cavity dispersion. It was found that in spite of the large frequency chirp of the dissipative solitons formed in the laser, the polarization-locked and polarization-rotating VDSs could still be formed in terms of different cavity birefringence. Formation of multiple VDSs with identical soliton parameters was also experimentally obtained. The polarization dynamics of the polarization-rotating VDS has also been investigated for the first time. It turned out that their periodicities of polarization rotation could be locked to the cavity roundtrip time or multiple of it due to the cavity boundary condition, and consequently the solitons only exhibited a fixed polarization state or alternate discrete polarization states at a fixed cavity position. The numerical simulation confirmed our results. It was also found that two sets of VDSs with orthogonal polarization orientations could be formed on account of the laser gain competition. When the pulse interval was comparable or even larger than the relaxation time of the EDF, which nullified the gain competition to some extent, two sets of VDSs with parallel orientations could even be formed as a result. In addition, the bound state of VDSs was also experimentally observed.

Chapter 9

Conclusions and Future Work

Dissipative solitons attract great attentions nowadays due to their new soliton pulse shaping mechanism, potential to achieve large pulse energy and recent experimental realization in fiber lasers. In this dissertation, a comprehensive study on the properties and dynamics of the dissipative solitons in passively mode-locked erbium-doped-fiber (EDF) lasers has been presented both experimentally and numerically. This chapter summarizes the achievements in the research and provides some suggestions for the future work.

9.1 Conclusions

An all-normal-dispersion (ANDi) EDF laser mode-locked by nonlinear polarization rotation (NPR) technique was set up to investigate the formation and properties of the scalar dissipative soliton. It was experimentally found that in spite of different pulse shaping mechanism from conventional solitons, multiple solitons, soliton energy quantization, noise-like state and the pump hysteresis could still be observed for dissipative solitons, but the underlying mechanisms are distinct. Through analyzing

the principle and occurrence conditions of the noise-like state and multiple dissipative solitons in normal dispersion, lifting the peak power clamping threshold and minimizing the effective gain bandwidth limitation were proposed as the guideline of optimizing cavity design to achieve high pulse energy dissipative solitons. Benefit from the dissipative soliton features, 4.8 nJ dissipative soliton was achieved after the cavity design optimization. An extended NPR mode-locking scheme was also proposed, where a cavity component with a relative large polarization dependent loss (PDL) replacing the polarizer was used to induce the NPR mode-locking under high pump power. This method effectively reduces the intracavity loss and could be applied to increase pulse energy. With it, 7.2 nJ high energy dissipative soliton was eventually achieved in our experiment. In fact, the ultimate limitation of the output pulse energy is the pump conversion efficiency of the gain fiber (for most EDF lasers the pump conversion efficiency is around 30-40% at 1.55 μm and some can reach 50%), in combination with the maximum available pump power in principle. However, in practical experiments, due to some other effects such as strong nonlinearity, narrow effective gain bandwidth and peak power clamping effect etc., pulse breaking always occurs in fiber lasers, which dominantly limits the largest single pulse energy achievable. Therefore, reducing the nonlinearity, broaden the effective gain bandwidth and lifting the peak power clamping threshold become the main and effective means to enlarge the single pulse energy achieved.

The dissipative solitons formed in dispersion-managed (DM) fiber lasers with different net cavity dispersion have been elaborately studied. It was first observed that dissipative solitons either with clear Kelly sidebands or with sharp steep spectral edges could appear in a DM fiber laser with net near zero cavity dispersion. Moreover,

stretched pulses, which normally appear at around zero cavity dispersion, were also observed in a large net normal dispersion cavity fiber laser. Numerical simulations confirmed the results and revealed that the underlying mechanism was that the non-equal nonlinear pulse propagation in the normal and anomalous fiber segments of a DM cavity induced the effective cavity dispersion shifting to anomalous side. The stronger the cavity nonlinearity, the larger the shift is. Our result shows that such an effect is a generic feature of the DM lasers, and is important for the better understanding on the dynamics of the DM fiber lasers.

Dissipative soliton resonance (DSR) and the flat-top soliton, two novel and unique features of the dissipative solitons, have been first experimentally observed in a specially designed long-cavity ANDi EDF laser. Remarkably, the pulse energy of a near resonance dissipative soliton was observed to increase continuously. A record-high single pulse energy of 281.2 nJ was experimentally achieved, which was only limited by the available pump power. The dissipative solitons observed at near resonance exhibit exactly the properties as those predicted theoretically, namely a rectangular pulse with fixed peak power and variable pulse width or a flat-top soliton. The largest pulse width achieved was 18.5 ns. Such large energy ns level mode-locked pulses have potential applications in laser targeting and lidar systems.

The intracavity dynamics of the dissipative solitons in a DM fiber laser with large normal cavity dispersion has been firstly investigated. The results showed that due to the de-chirping of the anomalous fiber in the cavity, the dissipative soliton no longer experienced pulse width narrowing in the normal dispersion gain fiber as in the non-DM fiber lasers, but had a similar evolution as that of parabolic pulses only with

nonlinear chirp in virtue of the finite gain bandwidth. Apart from the intracavity evolution, the pulse train of the dissipative solitons also exhibits complicated dynamical behavior in a time scale much larger than the round-trip time. Besides the period-doubling bifurcation, a new type of the deterministic soliton dynamics known as intermittency was first experimentally observed in a mode-locked fiber laser. In the intermittency, the regular soliton pulse intensity of the dissipative solitons was observed to be interrupted by the bursts of large intensity solitons. It showed the average duration of the pulse bursts decreased with the pump power enhancement or orientations of waveplates. Numerical simulations confirmed the experimental results.

Bound state of dissipative solitons in a DM fiber laser with large normal dispersion was first experimentally observed. It was found that in spite of the frequency chirp, and strong pulse stretch and compression within the cavity, the bound dissipative solitons still exhibited fixed pulse separation and phase difference. Numerical simulations well reproduced the experimental observations. Bound states of dissipative solitons in an EDF laser mode-locked by carbon nanotubes (CNTs), which has different saturable absorption properties from NPR, have also been studied. Various bound states of dissipative solitons obtained in the same fiber laser made a systematic experimental study on them and a comparison of their properties possible. Our results further showed that the tightly bound soliton states were characterized by the fixed phase difference, discrete pulse separations, and ultra stability, while the loosely bound soliton states no longer possessed these features.

Formation and properties of vector dissipative solitons (VDSs) have been studied in a fiber laser mode-locked by semiconductor saturable absorber mirror (SESAM). It was

found that in spite of the large frequency chirp of the dissipative solitons, both of the polarization-locked and polarization-rotating VDSs could still be formed. The polarization dynamics of the latter one was also studied and it turned out that its periodicities of polarization rotation could either be locked to the cavity roundtrip time or multiple of it. Multiple VDSs could also be formed. Apart from two sets of VDSs with orthogonal polarization orientations, two sets of VDSs with parallel orientations were also achieved in our laser since the gain competition was effectively suppressed when the pulse interval was comparable or even larger than the relaxation time of the EDF. Bound state of the VDSs was also experimentally observed and numerically verified.

In fact, dissipative solitons in fiber lasers mode-locked by different saturable absorbers (SAs) have also been investigated. It was noticed that different kinds of SAs, such as NPR, SESAM, CNTs and graphene, indeed had distinct saturable absorption properties. Generally speaking, the saturation recovery time of NPR (generally considered as <10 fs) is much faster than those of three others. The bandwidth of graphene is intrinsic broader than those of NPR, SESAM and CNTs, and the bandwidth of NPR is adjustable in a specific cavity. However, for conventional NLSE type solitons, the SA only plays the role of self-starting the mode-locking of the laser, and it does not participate in the soliton pulse shaping process. The parameters of the soliton are only determined by the cavity characteristics (dispersion and nonlinearity) and the input power. Different from that, for dissipative solitons, the saturable absorption, which is a dissipative process, contributes to the soliton formation. In spite of this, the parameters of the dissipative soliton are still determined by the concrete cavity condition (the balance of dispersion, nonlinearity, gain and loss)

rather than the property of the SA itself.

9.2 Recommendations for Future Work

As summarized above, a number of accomplishments have been achieved in our research, and they also inspire some interesting topics that deserve further investigations. They are briefly summarized as the following.

Chirped pulses with ultrahigh-repetition-rate are attractive since they could be effectively amplified and further compressed to form large energy ultrashort pulses with ultrahigh-repetition-rate, which have potential applications in ultrahigh speed optical communication systems. Through employing a Lyot-type filter in a laser to induce the dissipative four-wave-mixing (DFWM) mode locking, with the effect of Modulation instability (MI) as well, Zhao et al. realized a high-repetition-rate soliton pulse train with 111 GHz. Considering that the occurrence of DFWM mode locking is independent on the cavity dispersion, through appropriately selecting a spectral filter and the laser parameters, ultrahigh-repetition-rate dissipative solitons, i.e. chirped pulses with ultrahigh-repetition-rate, are also possible to be generated by DFWM mode locking in normal dispersion. Fabry-Perot Etalons and chirped fiber Bragg grating (FBG) are good candidates for the spectral filter.

Theoretical studies have shown that, taking the fiber birefringence into account, the cross-phase modulation (XPM) could also lead to the soliton formation, which is known as the induced soliton. Recently, such induced solitons were first observed in a SESAM mode-locked fiber laser. They were found to present either the same or

different central frequency to the inducing soliton, but always the same group velocity with it, so that they formed vector solitons. Although both group-velocity-locked and polarization-locked VDSs have been achieved, no induced VDSs were observed. Considering that the formation of the dissipative soliton always requires larger nonlinearity than the conventional soliton, it would be interesting to find out whether an induced VDS exists and whether it has any other different features from that of conventional soliton.

So far all the dissipative solitons obtained from single-wall carbon nanotubes (SWCNTs) mode-locked fiber laser are VDSs since the SWCNTs saturable absorber (SA) used has random tube alignment, which presents an isotropic optical absorption property. It is known that through specific technique, the tubes could be well aligned (such as vertical aligned), and then they exhibit a polarization-dependent optical absorption [191, 192]. It would be interesting to know whether a scalar dissipative soliton would be achieved in this case. Furthermore, if all the tubes in the SA are aligned in a pattern rather than in the same vertical direction what would happen? Whether VDSs could still be obtained but with less freedom dimension between two polarization components? It seems to be possible to realize a systematic and comprehensive study on the dissipative soliton with various freedom dimension and relationship between two polarization components (from scalar soliton to vector soliton) by only managing the order of the alignment of SWCNTs mode locker.

The phase difference of tightly bound soliton states were theoretically predicted to be 0, π or $\pi/2$. However, for all the tightly bound soliton states observed, their phase difference was either π or $\pi/2$, no 0 phase difference was experimentally achieved.

The reason is believed to be the fact when the phase difference is 0, the interaction between solitons in the bound state is the repulsion. If specific technique is employed to reduce this repulsive force in the laser cavity, is it possible to observe the bound state of solitons with phase difference of 0? Our numerical simulations show a positive answer to it. Besides this, our simulation also predicts the existence of the bound state of VDSs with two anti-phased polarization components. It would be interesting to experimentally confirm them. The studies are especially essential in enriching the family of the bound state of solitons.

Graphene is a new mode locker with the superior SA properties, including the intrinsic broadband absorption and high damage threshold. On the other hand, the ytterbium-doped fiber (YDF) provides more efficient gain and much broader gain bandwidth than the EDF. Taking advantages of both sides, the YDF laser mode-locked by graphene becomes the most competitive candidate to the large energy ultrashort pulses generator. Until now, graphene mode-locked YDF laser was first and only reported by Zhao et al. [193]. However, the pulse energy achieved is still limited to be quite small. Hence, managing the SA properties of the graphene, in combination with optimizing the YDF laser cavity design to achieve the single pulse energy as large as possible is another significant topic.

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