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# Intra-Channel Collision of Dual-Power Law Optical Solitons

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The intra-channel collision of optical solitons, with dual-power law nonlinearity, is studied in this paper by the aid of quasi-particle theory. The perturbation terms that are considered in this paper are the nonlinear gain and saturable amplifiers along with filters. The suppression of soliton-soliton interaction, in presence of these perturbation terms, is acheived. The numerical simulations support the quasi-particle theory.

**KEY WORDS:** optical solitons; quasi-particle theory; soliton perturbation; dual-power law; nonlinear Schrodinger's equation.

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## 1. INTRODUCTION

The theoretical possibility of existence of optical solitons in a dielectric dispersive fiber was first predicted by Hasegawa and Tappert (1973). A couple of years later Mollenauer *et al.* (1980) succesfully performed the famous experiment to verify this prediction. Important characterestic properties of these solitons are that they posses a localized waveform which remains intact upon interaction with another soliton. Because of their remarkable robustness, they attracted enormous interest in optical and telecommunication community. At present optical solitons are regarded as the natural data bits for transmission and processing of information in future, and an important alternative for the next generation of ultra high speed optical communication systems.

The fundamental mechanism of soliton formation namely the balanced interplay of linear group velocity dispersion (GVD) and nonlinearity induced

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self-phase modulation (SPM) is well understood. In the pico second regime, the nonlinear evolution equation that takes into account this interplay of GVD and SPM and which describes the dynamics of soliton is the well known nonlinear Schrödinger's equation (NLSE). The NLSE, which is the ideal equation in an ideal Kerr media, is in its original form found to be completely integrable by the method of Inverse Scattering Transform (IST) and tremendous success has been acheived in the development of soliton theory in the framework of the NLSE model.

However, communication grade optical fibers or as a matter of fact any optical transmitting medium does posses finite attenuation coefficient, thus optical loss in inevitable and the pulse is often detoriated by this loss. Therefore, optical amplifiers have to be employed to compensate for this loss. When the gain bandwidth of the amplifier is comparable to the spectral width of the ultrashort optical pulse, the frequency and intensity dependent gain must be considered. Another hindrance to the stable propagation in a practical system is the noise induced Gordon-Hauss timing jitter. An important aspect that has not been addressed with proper perspective is the fact that due to its nonsaturable nature, Kerr nonlinearity is inadequate to describe the soliton dynamics in the ultrahigh bit rate transmission. For example, when transmission bit rate is very high, for soliton formation the peak power of the incident field accordingly become very large. On the other hand higher order nonlinearities may become significant even at moderate intensities in certain materials such as semiconductor doped glass fibers. Under circumstances, as mentioned above, non-Kerr law nonlinearities come into play changing essentially the physical features of optical soliton propagation. Therefore when very high bit rate transmission or transmission through materials with higher nonlinear coefficients are considered, it is necessary to take into account higher order nonlinearities. This problem can be addressed by incorporating the dual-power law nonlinearity in the NLSE.

It has been realized that the Gordon-Haus timing jitter can be reduced by introducing bandpass filtering. Stabilazition of soliton propgation with the aid of nonlinear gain or under combined operation of gain and saturable absorption was recommended by (Kodama *et al.*, 1992; Kodama and Wabnitz, 1991, 1993a,b, 1994). Thus, in order to model these features in the soliton dynamics, in a practical situation, the NLSE should be modified by incorporating additional terms. Thus the concept of control of soliton propagation described by the NLSE with dual-power law nonlinearity is new and important developments in the application of solitons for optical communication systems. Because the NLSE with dual-power law is not integrable, perturbation methods or numerical techniques have to be applied. Therefore, the control of soliton and interction of two neighboring solitons incorporating perturbation terms like nonlinear gain, saturable amplification and filtering are going to be addressed in this paper.

## 2. MATHEMATICAL MODEL

The propagation of pulses through an optical fiber in an optical communication system, with dual-power law nonlinearity, is governed by the Nonlinear Schrodinger's Equation (NLSE) (Biswas, 2003). The dimensionless form of the NLSE is given by

$$i\frac{\partial q}{\partial Z}+\frac{1}{2}\frac{\partial^2 q}{\partial T^2}+\left(|q|^{2p}+\nu|q|^{4p}\right)q=0 \tag{1}$$

Here $q$ is the normalized efective amplitude of the wave electric field while $Z\&T$ are the independent variables. Here $Z$ represents the distance along the fiber while $T$ is the time and $\nu$ is a constant that accounts for the strength of the dual-power law.

The particularly relevant solutions to (1) are called solitons. In most cases, the interest is confined to a single pulse described by the 1-soliton solution of the NLSE. However, in this paper, the effects of the perturbation terms in NLSE on two soliton interaction will be studied. It is necessary to launch the solitons close to each other for enhancing the information carrying capacity of the fiber. If two solitons are placed close to each other then it can lead to its mutual interaction thus providing a very serious hinderance to the performance of the soliton transmission system. However, as elucidated in the previous section, the presence of appropiate perturbation terms the NLSE can lead to soliton control and suppression of this interaction. Therefore, in view of acheiving soliton control and suppression of the interaction, the perturbed NLSE that is going to be studied in this paper is given by

$$i\frac{\partial q}{\partial Z}+\frac{1}{2}\frac{\partial^2 q}{\partial T^2}+(|q|^{2p}+\nu|q|^{4p})q=i\epsilon R[q,q^*] \tag{2}$$

where

$$R=\delta|q|^{2m}q+\sigma q\int_{-\infty}^{T}|q|^2d\tau+\beta\frac{\partial^2 q}{\partial T^2} \tag{3}$$

Here, $\delta$ is called the nonlinear gain, while $\sigma$ is the coefficient of saturable amplifiers. Also, in (3), $m$ could be 0, 1 or 2. If $m=1$, it is called the linear gain while for $m=2$, it is called the gain saturation. The coefficient of $\beta$ is called the bandpass filtering term.

The quasi-particle theory (QPT) of soliton-soliton interaction (SSI) has been investigated (Biswas, 1999; Hasegawa and Kodama, 1995) and it is proved by virtue of it that the interaction can be supressed due to these perturbation terms. Also it has been proved that the sliding frequency guiding filters (Hasegawa and

Kodama, 1995) leads to the suppression of the SSI. Here, it will be proved, by virtue of the QPT, that the SSI can be suppressed due to the NLSE given by (1) in presence of the perturbation terms as in (2).

Equation (2) is not integrable by IST as it is not of Painleve type (Ablowitz and Segur, 1981; Ablowitz and Clarkson, 1993). However, (2) supports solitary waves of the form (Biswas, 2003)

$$q(Z,T) = \frac{\eta}{[1 + a\cosh\{\zeta(T - vZ - T_0)\}]^{\frac{1}{2p}}} e^{i\{-\kappa T + \omega Z + \sigma_0\}} \tag{4}$$

where

$$\kappa = -v \tag{5}$$

$$\omega = \frac{\eta^{2p}}{2p+2} - \frac{\kappa^2}{2} \tag{6}$$

$$\zeta = \eta^p \left(\frac{2p^2}{1+p}\right)^{\frac{1}{2}} \tag{7}$$

$$a = \sqrt{1 + \frac{v\zeta^2}{2p^2} \frac{(1+p)^2}{1+2p}} \tag{8}$$

Here $\eta$ is the amplitude of the soliton, $\zeta$ is the width of the soliton, $v$ is its velocity, $\kappa$ is the soliton frequency and $\omega$ is the wave number while $T_0$ and $\sigma_0$ are the center of the soliton and the center of the soliton phase respectively. For dual-power law nonlinearity solitons exist for

$$-\frac{2p^2}{\zeta^2} \frac{1+2p}{\zeta^2 (1+p)^2} < v < 0 \tag{9}$$

Also, the 2-soliton solution of the NLSE (1) takes the asymptotic form

$$q(Z,T) = \sum_{l=1}^{2} \frac{\eta_l}{[1 + a_l \cosh\{\zeta_l(T - vZ - T_l)\}]^{\frac{1}{2p}}} e^{i(-\kappa_l T + \omega_l Z + \sigma_{0_l})} \tag{10}$$

with

$$\kappa_l = -v_l \tag{11}$$

$$\omega_l = \frac{\eta_l^{2p}}{2p+2} - \frac{\kappa_l^2}{2} \tag{12}$$

$$\zeta_l = \eta_l^p \left(\frac{2p^2}{1+p}\right)^{\frac{1}{2}} \tag{13}$$

$$a_l = \sqrt{1 + \frac{\nu\zeta_l^2}{2p^2}\frac{(1+p)^2}{1+2p}} \tag{14}$$

In the study of SSI, the initial pulse waveform is taken to be of the form

$$q(0,T) = \frac{\eta_1}{[1 + a_1\cosh\{\zeta_1(T - \frac{T_0}{2})\}]^{\frac{1}{2p}}}e^{i\phi_1} + \frac{\eta_2}{[1 + a_2\cosh\{\zeta_2(T + \frac{T_0}{2})\}]^{\frac{1}{2p}}}e^{i\phi_2} \tag{15}$$

which represents the injection of 2-soliton like pulses into a fiber. Here, $T_0$ represents the initial separation of the solitons. It needs to be noted that for $T_0 \to \infty$(4) represents exact soliton solutions, but for $T_0 \sim O(1)$ it does not represent an exact 2-soliton solution. In this paper, the case of in-phase injection of solitons with equal amplitudes will be studied. Thus, without any loss of generality the choice $\eta_1 = \eta_2 = 1$ and $\phi_1 = \phi_2 = 0$ is considered so that (15) modifies to

$$q(0,T) = \frac{1}{\left[1 + a_1\cosh\left\{\zeta\left(T - \frac{T_0}{2}\right)\right\}\right]^{\frac{1}{2p}}} + \frac{1}{\left[1 + a_2\cosh\left\{\zeta\left(T + \frac{T_0}{2}\right)\right\}\right]^{\frac{1}{2p}}} \tag{16}$$

where

$$\zeta = \sqrt{\frac{2p^2}{1+p}} \tag{17}$$

## 3. QUASI-PARTICLE THEORY

The QPT dates back to 1981 since the appearance of the paper by Karpman and Solov'ev (1981). The mathematical approach to the SSI using the QPT will be studied in this paper. Here, the solitons are treated as particles. If two pulses are separated and each of them is close to a soliton they can be written as the linear superposition of two soliton like pulses (Biswas, 1999; Hasegawa and Kodama, 1995)

$$q(Z,T) = q_1(Z,T) + q_2(Z,T) \tag{18}$$

with

$$q_l(Z,T) = \frac{A_l}{[1 + a_l\cosh[B_l(T - T_l)]]^{\frac{1}{2p}}}e^{-iB_l(T-T_l)+i\delta_l} \tag{19}$$

where $l = 1, 2$ and $A_l$, $B_l$, $T_l$ and $\delta_l$ are functions of $Z$. Also, it is important to note that $A_l$ and $B_l$ do not represent the amplitude and the frequency of the full wave form. However, they approach the amplitude and frequency respectively for large separation namely as $\Delta T = T_1 - T_2 \to \infty$, $A_l \to \eta_l$ and $B_l \to \kappa_l$. Since the waveform is assumed to remain in the form of two pulses, the method is called the quasi-particle approach. The equations for $A_l$, $B_l$, $T_l$ and $\delta_l$ will be derived first using the soliton perturbation theory. Substituting (18) into (2), yields

$$i\frac{\partial q_l}{\partial Z} + \frac{1}{2}\frac{\partial^2 q_l}{\partial T^2} = i\epsilon R[q_l, q_l^*] - \left[\sum_{r=0}^{p}\binom{p}{r} q_1^{p-r} q_2^r\right]\left[\sum_{r=0}^{p}\binom{p}{r}(q_1^*)^{p-r}(p_2^*)^r\right](q_1+q_2) - \nu\left[\sum_{r=0}^{2p}\binom{2p}{r} q_1^{2p-r} q_2^r\right]\left[\sum_{r=0}^{2p}\binom{2p}{r}(q_1^*)^{2p-r}(q_2^*)^r\right](q_1+q_2) \tag{20}$$

where $l = 1, 2$ and $\bar{l} = 3 - l$ with the definition

$$\binom{n}{r} = \frac{n(n-1)\cdots(n-r+1)}{r(r-1)\cdots 3.2.1} \tag{21}$$

Here, the separation

$$|q|^{2p}q + \nu|q|^{4p}q = \left[\sum_{r=0}^{p}\binom{p}{r} q_1^{p-r} q_2^r\right]\left[\sum_{r=0}^{p}\binom{p}{r}(q_1^*)^{p-r}(q_2^*)^r\right](q_1+q_2) + \nu\left[\sum_{r=0}^{2p}\binom{2p}{r} q_1^{2p-r} q_2^r\right]\left[\sum_{r=0}^{2p}\binom{2p}{r}(q_1^*)^{2p-r}(q_2^*)^r\right](q_1+q_2) \tag{22}$$

was used based on the degree of overlapping. By the soliton perturbation theory (Hasegawa and Kodama, 1995; Hasegawa and Matsumoto, 2003) the evolution equations are

$$\frac{dA_l}{dZ} = F_1^{(l)}(A, \Delta T, \Delta\phi; \nu, p) + \epsilon M_l \tag{23}$$

$$\frac{dB_l}{dZ} = F_2^{(l)}(A, \Delta T, \Delta\phi; \nu, p) + \epsilon N_l \tag{24}$$

$$\frac{dT_l}{dZ} = -B_l - F_3(A, \Delta T, \Delta\phi; \nu, p) + \epsilon Q_l \tag{25}$$

$$\frac{d\delta_l}{dZ} = \frac{A_l^{2p}}{2p+2} + \frac{B_l^2}{2} + F_4(A, \Delta T, \Delta\phi; \nu, p) + \epsilon P_l \tag{26}$$

where the functions $F_1^{(l)}$, $F_2^{(l)}$, $F_3$ and $F_4$ formulate on using the SPT in (20), with the right side being treated as perturbation terms. Also,

$$M_l = h_1(A_l)\int_{-\infty}^{\infty} \Re\left\{\hat{R}[q_l, q_l^*]e^{-i\phi_l}\right\} \frac{d\tau_l}{(1 + a_l\cosh\tau_l)^{\frac{1}{2p}}} d\tau_l \tag{27}$$

$$N_l = h_2(A_l)\int_{-\infty}^{\infty} \Im\left\{\hat{R}[q_l, q_l^*]e^{-i\phi_l}\right\} \frac{\sinh\tau_l}{(1+a_l\cosh\tau_l)^{\frac{1}{2p}}}d\tau_l \tag{28}$$

$$Q_l = h_3(A_l)\frac{1}{A_l^2}\int_{-\infty}^{\infty} \Re\left\{\hat{R}[q_l, q_l^*]e^{-i\phi_l}\right\} \frac{\tau_l}{(1+a_l\cosh\tau_l)^{\frac{1}{2p}}}d\tau_l \tag{29}$$

$$P_l = h_4(A_l)\frac{1}{A_l}\int_{-\infty}^{\infty} \Im\left\{\hat{R}[q_l, q_l^*]e^{-i\phi_l}\right\} \frac{(1-a_l\tau_l\sinh\tau_l)}{(1+a_l\cosh\tau_l)^{\frac{1}{2p}}}d\tau_l \tag{30}$$

Here, the functions $h_j(A_l)$ for $1 \leq j \leq 4$ are by virtue of (23)-(26) and $\Re$ and $\Im$ stands for the real and imaginary parts respectively. Also, the following notations are used

$$\begin{aligned}\hat{R}[q_l, q_l^*] &= R[q_l, q_l^*] + i|q_l|^{2p}q_l + i\nu|q_l|^{4p}q_l \\ &-i\left[\sum_{r=0}^{p}\binom{p}{r}q_1^{p-r}q_2^r\right]\left[\sum_{r=0}^{p}\binom{p}{r}(q_1^*)^{p-r}(q_2^*)^r\right](q_1+q_2) \\ &-i\left[\sum_{r=0}^{2p}\binom{2p}{r}q_1^{2p-r}q_2^r\right]\left[\sum_{r=0}^{2p}\binom{2p}{r}(q_1^*)^{2p-r}(q_2^*)^r\right](q_1+q_2)\end{aligned} \tag{31}$$

$$\tau_l = A_l(T - T_l) \tag{32}$$

$$\phi_l = B_l(T - T_l) - \delta_l \tag{33}$$

$$\Delta\phi = B\Delta T + \Delta\delta \tag{34}$$

$$\Delta T = T_1 - T_2 \tag{35}$$

$$\Delta\delta = \delta_1 - \delta_2 \tag{36}$$

$$A = \frac{1}{2}(A_1 + A_2) \tag{37}$$

$$B = \frac{1}{2}(B_1 + B_2) \tag{38}$$

$$\Delta A = A_1 - A_2 \tag{39}$$

$$\Delta B = B_1 - B_2 \tag{40}$$

Moreover, it was assumed that

$$|\Delta A| \ll A \tag{41}$$

$$|\Delta B| \ll 1 \tag{42}$$

$$A\Delta T \gg 1 \tag{43}$$

$$|\Delta A|\Delta T \ll 1 \tag{44}$$

From (37) to (40) one can now obtain

$$\frac{dA}{dZ} = \epsilon M \tag{45}$$

$$\frac{dB}{dZ} = \epsilon N \tag{46}$$

$$\frac{d(\Delta A)}{dZ} = F_1^{(1)}(A, \Delta T, \Delta\phi; \nu, p) - F_1^{(2)}(A, \Delta T, \Delta\phi; \nu, p) + \epsilon \Delta M \tag{47}$$

$$\frac{d(\Delta B)}{dZ} = F_2^{(1)}(A, \Delta T, \Delta\phi; \nu, p) - F_2^{(2)}(A, \Delta T, \Delta\phi; \nu, p) + \epsilon \Delta N \tag{48}$$

$$\frac{d(\Delta T)}{dZ} = -\Delta B + \epsilon \Delta Q \tag{49}$$

$$\frac{d(\Delta\phi)}{dZ} = \frac{A_1^{2p} - A_2^{2p}}{2p+1} + \frac{\Delta T}{2}\left(F_2^{(1)} + F_2^{(2)}\right) + \epsilon B \Delta Q + \epsilon \Delta P \tag{50}$$

where

$$M = \frac{1}{2}(M_1 + M_2) \tag{51}$$

$$N = \frac{1}{2}(N_1 + N_2) \tag{52}$$

and $\Delta M$, $\Delta N$, $\Delta Q$ and $\Delta P$ are the variations of $M$, $N$, $Q$ and $P$ which are written as for example

$$\Delta M = \frac{\partial M}{\partial A}\Delta A + \frac{\partial M}{\partial B}\Delta B \tag{53}$$

assuming that they are functions of $A$ and $B$ only, which is, in fact, true for most of the cases of interest, otherwise, the equations for

$$T = \frac{1}{2}(T_1 + T_2) \tag{54}$$

and

$$\phi = \frac{1}{2}(\phi_1 + \phi_2) \tag{55}$$

would have been necessary. In presence of the perturbation terms, as given by, (2), the dynamical system of the soliton parameters, by virtue of soliton perturbation theory, are

$$
\begin{aligned}
\frac{dA}{dZ} &= \frac{4\epsilon\delta A^{2m+2}}{2^{\frac{m+1}{p}} a^{\frac{m+1}{p}} D} F\left(\frac{m+1}{p}, \frac{m+1}{p}, \frac{m+1}{p} + \frac{1}{2}; \frac{a-1}{2a}\right) B\left(\frac{m+1}{p}, \frac{1}{2}\right) \\
&\quad + \frac{2\epsilon\sigma A^4}{D^2} \int_{-\infty}^{\infty} \frac{1}{(1+a\cosh\tau)^{\frac{1}{p}}} \left(\int_{-\infty}^{\tau} \frac{ds}{(1+a\cosh s)^{\frac{1}{p}}}\right) d\tau \\
&\quad - \frac{4\epsilon\beta A^2}{D a^{\frac{1}{p}} 2^{\frac{1}{p}}} \left[ D^2 F\left(2+\frac{1}{p}, \frac{1}{p}, \frac{3}{2}+\frac{1}{p}; \frac{a-1}{2a}\right) B\left(\frac{1}{p}, \frac{3}{2}\right) \right. \\
&\quad \left. + B^2 F\left(\frac{1}{p}, \frac{1}{p}, \frac{1}{2}+\frac{1}{p}; \frac{a-1}{2a}\right) B\left(\frac{1}{p}, \frac{1}{2}\right) \right]
\end{aligned} \tag{56}
$$

$$
\frac{dB}{dZ} = \frac{\epsilon\beta B D^2}{4p^2 A^2} \frac{F\left(2+\frac{1}{p}, 1+\frac{1}{p}, 2+\frac{1}{p}; \frac{a-1}{2a}\right)}{F\left(\frac{1}{p}, \frac{1}{p}, \frac{1}{2}+\frac{1}{p}; \frac{a-1}{2a}\right)} \frac{B\left(1+\frac{1}{p}, 1\right)}{B\left(\frac{1}{p}, \frac{1}{2}\right)} \tag{57}
$$

where $F(\alpha, \beta; \gamma; z)$ is the Gauss' hypergeometric function defined as

$$
F(\alpha, \beta; \gamma; z) = \frac{\Gamma(\gamma)}{\Gamma(\alpha)\Gamma(\beta)} \sum_{n=0}^{\infty} \frac{\Gamma(\alpha+n)\Gamma(\beta+n)}{\Gamma(\gamma+n)} \frac{z^n}{n!} \tag{58}
$$

and $B(l, m)$ is the usual beta function. Thus, one can obtain using, (47)–(50)

$$
\frac{d^2(\Delta T)}{dZ^2} + \epsilon\beta G \frac{d(\Delta T)}{dZ} + F_2^{(1)} - F_2^{(2)} = 0 \tag{59}
$$

where $G > 0$ represents the coefficient of $-\epsilon\beta\Delta B$ in $d(\Delta B)/dZ = dB_1/dZ - dB_2/dZ$. Now equation (59) shows that there is a damping in the separation of solitons thus proving that there will be a suppression of the SSI in presence of the perturbation terms given by (2).

The initial conditions for these variables, corresponding to the initial waveform (7), are

$$
A = 1, \quad B = 0, \quad \Delta A_0 = 0, \quad \Delta B_0 = 0, \quad \Delta T_0 = T_0 \quad \& \quad \Delta\phi_0 = \phi_0
$$

## 4. NUMERICAL SIMULATIONS

The Mathematical set up, given by (56), (57) and (59) will be used to study various situations of the perturbed NLSE (2) to observe how SSI can be suppressed. For the fixed point of the dynamical system, given by (56) and (57), with $A = 1$

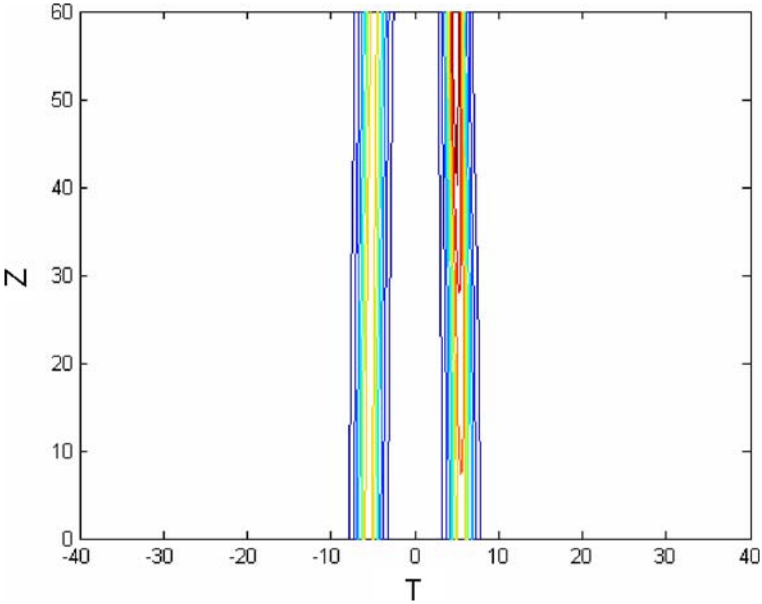


**Fig. 1.** m = 0, p = 1, $\delta = \sigma = 0.01$.

and $B = 0$, one recovers

$$\beta = \frac{\delta}{2^{\frac{m}{p}} a^{\frac{m}{p}}} \left(\frac{1+p}{2p^2}\right)^{\frac{1}{p}} \frac{F\left(\frac{m+1}{p}, \frac{m+1}{p}, \frac{m+1}{p}, +\frac{1}{2}; \frac{a-1}{2a}\right)}{F\left(2+\frac{1}{p}, \frac{1}{p}, \frac{3}{2}+\frac{1}{p}; \frac{a-1}{2a}\right)} \frac{B\left(\frac{m+1}{p}, \frac{3}{2}\right)}{B\left(\frac{1}{p}, \frac{3}{2}\right)}$$
$$+ \frac{\sigma a^{\frac{1}{p}}}{2^{\frac{p-1}{p}}} \left(\frac{1+p}{2p^2}\right)^{\frac{3}{2p}} \frac{1}{B\left(\frac{1}{p}, \frac{3}{2}\right) F\left(2+\frac{1}{p}, \frac{1}{p}, \frac{3}{2}+\frac{1}{p}; \frac{a-1}{2a}\right)}$$
$$\int_{-\infty}^{\infty} \frac{1}{(1+a\cosh\tau)^{\frac{1}{p}}} \left(\int_{-\infty}^{\tau} \frac{ds}{(1+a\cosh s)^{\frac{1}{p}}}\right) d\tau \tag{60}$$

In the following couple of numerical simulations, the choices $\epsilon = 0.1$ and $T_0 = 11$ were made while, the other parameters are chosen as follows:

1. In Figure 1, $m = 0,\ p = 1, \delta = \sigma = 0.01$ so that by (60), $\beta = 0.0496$.
2. In Figure 2, $m = 1,\ p = 1, \delta = \sigma = 0.01$ so that by (60), $\beta = 0.0895$.

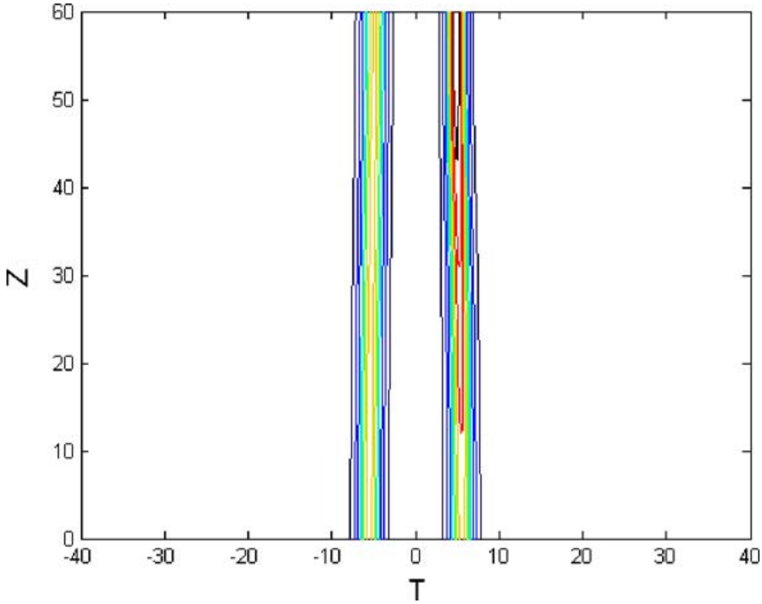


**Fig. 2.** m = 1, p = 1, $\delta = \sigma = 0.01$.

## 5. CONCLUSIONS

In this paper, the SSI of the NLSE, with dual-power law nonlinearity in presence of nonlinear gain, saturable amplifiers and filters are investigated. It is observed that the SSI can be suppressed in presence of these perturbation terms for various values of the degree of nonlinear gain. The QPT, due to these perturbation terms, was developed and the analytical reasoning of the suppression of the SSI was formulated.

Thus, in the applied soliton community two solitons can be injected into a single channel, close to one another and also suppress their mutual interaction so that performance enhancement can be acheived. This conclusion is based on numerical and analytical results due to the quasi-particle theory of SSI.

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