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If EM fields do not operate on each other, why do we need many modes and large gain bandwidth to generate short pulses?

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ABSTRACT

We recognize that the superposed light beams do not interact with each other to re-distribute their energy in space or time in the absence of interacting material dipoles. This platform requires that we re-visit the physical model behind the generation of pulses from the so-called "mode-locked" lasers. In the process, we come across the mathematical models behind formulating (i) the autocorrelation due to pulsed light and (ii) the group velocity of pulse propagation are based on the direct summation (integration) of non-causal infinite Fourier frequencies as if the EM waves can actually modify their energy distribution in the time domain. Accordingly, we show by modeling results and proposed experiments that time-frequency Fourier theorem can give rise to self-contradictory predictions, verifiable by simple laboratory experiments. Based on these results, we propose that we replace the paradigm of "interference of light" by "superposition effects due to light beams" as reported by the material dipoles of detectors and beam splitters.

[Key words: Mode-lock lasers, pulse autocorrelation, group velocity, interference, superposition effects of light.

1. INTRODUCTION

1.1. Traditional explanation for mode locking

In most text books [1] and journals [2] it is a standard practice to explain the mode locking process as the "interference" (summation) of a periodic array of CW frequencies (longitudinal modes) on which we impose techniques to assure that the arbitrary phases of all the modes become constant and remain so. Then their mathematical summation gives rise to an infinite train of temporal pulses with a single carrier frequency, which is the central frequency of the array of longitudinal modes.

$$E(\nu_0,t)=\sum_{n=0}^{N-1}e^{i2\pi(\nu_0+n\delta\nu)t+i\phi_c}\approx\frac{\sin N\pi\delta\nu t}{\sin\pi\delta\nu t}e^{i2\pi\nu_0 t+i\phi_c}\equiv Na(t)e^{i2\pi\nu_0 t+i\phi_c} \tag{1a}$$

We have assumed, for convenience, that the amplitudes of all the coherent frequencies are equal to unity and defined $a(t)$ as the normalized "carrier envelope" for the resultant modified field. The carrier frequency is ν_0 and is the mean frequency of the series. The traditional mode-locked optical intensity envelope, the peak normalized to unity, is now given by:

$$\begin{aligned} I(t)&=(1/N^2)\left|E(\nu_0,t)\right|^2=(1/N^2)[\sin^2 N\pi\delta\nu t/\sin^2\pi\delta\nu t]\\ &=\frac{1}{N^2}E^*E=\frac{1}{N}+\frac{2}{N^2}\sum_{p=1}^{N-1}(N-p)\cos[2\pi p\delta\nu t]\end{aligned} \tag{1b}$$

Unfortunately, Eq.1a contradicts the key assumption behind the success of more than a century old measurement technique known as Fourier transform spectroscopy (FTS) [3]. The assumption, validated by experiments, is that light beams containing different optical frequencies, do not "interfere" with each other irrespective of the temporal rate of fluctuations of their individual absolute phases (or temporal coherence length). The paradox becomes more perplexing if we literally take a CW gas laser (inhomogeneously broadened gain medium) of cavity length L running in N discrete longitudinal modes $2L=m\lambda_m=mc/\nu_m$ and detect the signal with a fast detector and sampling scope. We would find that the detector current envelope closely resembles the

mode lock envelope (as in the first expression of Eq.1b), even though the laser is not mode locked [4]. However, the shape will occasionally change as the relative phase differences between the modes slowly drifts around. But if the fast detector current is analyzed by an electronic spectrum analyzer (ESA), we will find a discrete set of sharp lines at frequencies $f = p\delta\nu = p(c/2L)$, which can be appreciated from the second term in the second expression of Eq.1b. These sharp lines are simply the beat between all the individual longitudinal modes. This of course, is the simplest version of light beating spectroscopy (LBS) or heterodyne spectroscopy. So, unlike FTS, both LBS and "mode locking" effectively allows "interference" between different optical frequencies. For LBS the "temporal coherence length" (or phase stability duration) does not need to be infinitely long, which is the key requirement for the "mode locking" process [Eq.1a]. How do we get out of this state of confusion, or paradoxical situation? That is the subject of this paper.

1.2 Absence of interaction or interference between light beams

The solution lies with a mundane observation that we have been trained to ignore for centuries. Well formed light beams do not operate on (interfere with) each other. They pass through each other without perturbing or re-distributing their temporal or spatial energy distributions [5]. Otherwise, we could not have consistently measured the unique Doppler shifts of light beams coming from stars billions of light years away and crossing through trillions of other beams (measuring the expanding universe!). There is no interference of light.

There are only superposition effects due to light as experienced by our material dipole detectors. The observed results are constrained by their quantum properties as to whether they can (or are allowed to) respond to all the superposed field-induced stimulations simultaneously. In other words, the superposition effects need to be represented by the summation of the allowed dipole stimulations, rather than the fields themselves. Mathematically, this is a very simple transition. We need to replace the fields $E(\nu,t)$ by the detector stimulations $d(t) = \chi_{(1)}E(\nu,t)$, which is the product of the linear susceptibility of the detecting molecules to dipolar stimulation and the stimulating EM field. We can never measure the so called "intensity" $E^*(t)E(t)$; we can only measure $d^*(t)d(t)$ as photo electric current, or photo chemical reaction, etc. Note that $\chi_{(1)}$ is a function of all the quantum constraints the detector has that includes the optical frequencies, the direction of polarizations, etc., it can respond to. In other words, Eq.1b should be re-written as the detector current $i(t)$ as:

$$\begin{aligned} i(t) &= |d(t)|^2 = (\chi_{(1)}^2/N^2)[\sin^2 N\pi\delta\nu t/\sin^2 \pi\delta\nu t] \\ &= d^*d = \frac{\chi_{(1)}^2}{N} + \frac{2\chi_{(1)}^2}{N^2}\sum_{p=1}^{N-1}(N-p)\cos[2\pi p\delta\nu t] \end{aligned} \tag{2}$$

The mathematical expressions of Eq.1b and Eq.2 are identical except for the factor $\chi_{(1)}^2$ that is a constant for a given detector. Thus we become oblivious to the profound mistake that we make by ignoring the physical process behind all optical phenomena. We "see" light only through the "eyes' of the detectors!

In general, all detecting material dipoles are susceptible to all the linear and non-linear polarizability, and in reality $d(t)$ should be expressed as quite a complex relation [6]:

$$\vec{d}(t) = \sum\nolimits_n \chi_{(n)}\vec{E}^n(\nu,t) \tag{3}$$

It is the linear and non-linear quantum properties of the intra-cavity saturable absorber or the Kerr medium and their temporal gate-like behavior that determines the generation of time domain intensity pulses out of laser cavities. It is not because the EM waves of the cavity modes interact with each other and then redistribute their energy in the time domain to create the "mode locked" pulses.

1.3. A new mean frequency out of many oscillating modes is not generated during mode locking process

Let us re-write Eq.1a in terms of linear dipole stimulation:

$$d(t) = \chi_{(1)}E(\nu_0,t) = \chi_{(1)}\sum_{n=0}^{N-1}e^{i2\pi(\nu_0+n\delta\nu)t+i\phi_c} \approx \chi_{(1)}\frac{\sin N\pi\delta\nu t}{\sin\pi\delta\nu t}e^{i2\pi\nu_0 t+i\phi_c} \equiv \chi_{(1)}Na(t)e^{i2\pi\nu_0 t+i\phi_c} \tag{4}$$

Do the dipoles of the saturable absorber really undulate and re-emit EM waves at the mean frequency ν_0 with a temporal amplitude envelope $a(t)$? In reality, the efficiency of the saturable absorber to become an open gate for

the transmission of the stored cavity energy is enhanced when the phases of all the randomly emitted frequencies match. It functions only as a temporal gate by being transparent after absorbing energy $d^*(t)d(t)$. The cavity feed back is enhanced as the 'gate' starts opening up through saturation due to increasing absorption of energy $d^*(t)d(t)$. The process becomes a feed-back loop to bring the spontaneous emissions in phase. But the saturable absorber cannot change the intrinsic quantum properties of the lasing molecules. The unperturbed intrinsic quantum properties of the lasing molecules dictate the strengths of the oscillating optical field and its frequency through the cavity properties (length and feed back). For inhomogeneously broadened gain medium, like He-Ne lasers, all the cavity allowed frequencies (modes) $2L = m\lambda_m = mc/\nu_m$ will be oscillating and their mode phases have to be "locked" for efficient absorption by the saturable absorber [7]. In fact, the spectrum of a mode locked He-Ne laser clearly shows all the discrete longitudinal modes, just as for the CW state. The total spectrum is not "transform limited", although each of the separate cavity mode frequencies show transform limited broadening due to intensity pulsing (unlike CW line width). For a homogeneously broadened gain medium, the natural tendency of the lasing medium is to oscillate only at the central frequency. This is where the gain has maximum value. All the excited atoms or molecules are capable of responding to the influence of the stimulating photons at this central frequency. So, for homogeneously broadened gain media, short pulse generation is not due to the summation of all the longitudinal modes of the cavity because they do not normally exist. Can the temporal modulation of the cavity feed-back signal make a homogeneously broadened gain medium into an inhomogeneous medium? Furthermore, the width of the spectral fringe from such perfectly mode locked lasers show transform limited behavior, which implies that the carrier frequency of the mode locked pulses is essentially a single frequency. We have carried out direct time-domain analysis of typical spectrometers' response to any pulse using its carrier frequency as the physical frequency (actual E-vector undulation) and have shown that the time-integrated spectral fringe width does correspond to *transform limited* spectrum. It is a convolution of the of the Fourier intensity spectrum of the temporal envelope and the CW single frequency response of the spectrometer [8]. If there were to exist multiple carrier frequencies in a pulse, the spectral fringes would be correspondingly broadened beyond "transform limited" response due to convolution of the CW-response with all these separate mode frequencies.

So far, we have not justified the reason behind the observed necessity for larger and larger gain band widths to obtain shorter and shorter pulses. Here, we are "hand-waving" by saying that a fast amplification of a short pulse requires the accessibility of a very large number of lasing atoms or molecules at the excited states with very short life times due to broadening of the energy level from a sharp line to a wider band. Accordingly, short pulse generation from a given volume and density of lasing atoms, the best choice would be a homogeneously broadened gain medium. But this argument only supports why we have been successful in making short pulses using homogeneously broadened lasing media in contrast to in-homogeneously broadened lasing media. We have another reason to believe that the so-called *transform-limited* "mode-locked" pulses are not due to superposition of multiple longitudinal modes. We have already argued, on the basis of (i) non-interference of light and (ii) our time-domain spectroscopic theory [5, 8], that a transform-limited pulse actually contains only one carrier frequency (or a very, very narrow band). This frequency is very likely to be the gain line center of the lasing medium, unless cavity-internal frequency selector moves it off-center at the cost of overall system gain. So, homogeneous broadening is definitely a necessary condition for all the lasing atoms to participate at the gain line center. But, is this sufficient? This opens up the need for a new hypothesis. The rate of stimulated emission probability of an excited atom is higher when its transition line width is broader. Let us re-phrase the statement. An isolated excited atom (say, in a gas) has lower probability (needs longer time) to release its excess energy to a stimulating field of proper frequency than when its intrinsic sharp line has been significantly broadened by embedding it in a solid or liquid medium. This hypothesis could be verified by comparative quantitative measurements of the amplification and the pulse width modification of the same short pulse passing through two drastically different gain media with the same gain line center (matching with the pulse carrier frequency), but with very different gain bandwidth. We suggest that one should choose δt such that $\delta\nu_1 \,\&\, \delta\nu_2 > (1/\delta t)$ for the two gain media to avoid being trapped by the circuitous argument that normally used using the time-bandwidth product limit $\delta\nu\delta t \geq 1$ from the Fourier theorem. To assure the narrowness of the pulse carrier frequency, it would be better to derive the short pulse from a well stabilized external cavity CW laser with the help of an external modulator ("chopper"). Obviously, the sharp rise and fall of the "imagined" envelope of the undulating EM field has physical impact on the temporal behavior of the dipolar stimulation induced on the excited atom. This subtle time varying stimulation of material dipoles should not be modeled by using Fourier frequencies of the pulse which eliminates time from the real physical process.

1.4. Limits of linear superposition of fields accepted by Fourier's time-frequency and Maxwell's wave equations

Explicit recognition that the superposition effects due to multiple light beams can become manifest only if some suitable detecting molecules can simultaneously sum the attempted stimulation by all the superposed beams has important consequences on several important mathematical relations tied to Fourier's time-frequency and Maxwell's wave equations.

The solutions for Maxwell's wave equation can be any simple sinusoid or a linear combination of them. But 'linear combination' implies a summation operation that cannot be carried out by the EM fields themselves. EM sinusoids activate the susceptibility of atoms and molecules to dipolar undulations through all possible linear and non-linear stimulations governed by rules of QM. The formalism of linear combination becomes physically manifest only for the QM-allowed sinusoids as the response of the detecting dipoles, but each multiplied by the characteristic first-order (linear) susceptibility factor. The rest of the sinusoids contribute to weak nonlinear effects governed by nonlinear susceptibilities. *Even for a very successful mathematical formalism, like Maxwell's wave equation, all of its mathematically acceptable rules cannot be taken for granted as representing nature's actual processes.* In general, superposed EM fields do not sum themselves into a new compound field.

Mathematically, the time-frequency Fourier theorem allows one to synthesize a pulse or to analyze a pulse as a 'linear superposition' of sinusoids. The 'linearity' of the summation applied to 'linear physical systems' is considered as the justification. Unfortunately, the pure mathematical logic of linearity cannot override the need for real physical processes in the absence of interaction between EM fields. Yet the success of Fourier theorem has been staggering, given that most material-based classical physical undulations (waves, pendulums, and so on) and QM systems can be represented by sinusoids. This is achieved by ad hoc customization of definitions for the Fourier conjugate variables according to the needs of the problem under consideration. We have just seen that 'mode locking,' or synthesizing a pulse out of the cavity 'modes' can be achieved only with the help of saturable absorbers or their equivalent. One cannot produce short pulse by simply superposing a set of laser beams with a periodic array of frequencies and highly stabilized with the best possible feed-back loops. A material medium with the right absorption bandwidth is essential. The time-frequency bandwidth limit $\delta\nu\delta t \geq 1$ is justified by the same theorem using 'decomposition' logic but a non-interacting linear system (classical spectrometers) cannot do the actual decomposition. Therefore, this 'bandwidth limit' cannot be a fundamental principle of nature, which we have demonstrated as spectral super resolution for amplitude modulated (AM) pulse using heterodyne spectroscopy [8a]. Appreciating this reality should open up the possibility of designing spectrometers with super resolution and a better understanding of natural line widths [8b].

1.5. Approach to modeling pulse amplification by a laser gain medium

Semiconductor laser amplifiers are playing major roles for many diverse applications. With the background information presented here related to non-interference of light that imposes limits on the applicability of Fourier's time-frequency relation, we would like to raise the questions regarding the best approach to model the propagation and amplification of laser pulses through some Fabry-Perot (FP) amplifier. Before we can do that we need to further develop the background on the proper definition and interpretation of "spectrum" and "coherence function" due to pulsed light, which are at the core of propagating light pulses through different media. For example, the autocorrelation (Wiener-Khintchine) theorem [9,10] tells us that the "spectrum" of a pulse is the Fourier transform of the measurable autocorrelation (fringe visibility, etc.) function due to the pulse. In our view, since Fourier frequencies are not physical, we would like to propose that only the actual carrier frequency of a light pulse should be considered as the optical spectrum, as opposed to the mathematical Fourier spectrum. We will show that explicit recognition of these points allows one to distinguish between the same autocorrelation function derived (i) either from a pulse with a single carrier frequency, (ii) or from a CW spectrum that is the same as the Fourier transform of the pulse envelope function. In the first case, the reduced fringe visibility (or autocorrelation) is due to the superposition of time-varying unequal amplitudes on the detector. In the second case, it is due to overlap of sinusoidal fringes whose "spatial" frequency varies due to physically different CW carrier frequencies contained in the beam under analysis. So, we propose to introduce a new concept. The first one is truly *temporal coherence* and well recognized in the literature [9,10] and the second one is *spectral coherence*. Unfortunately, the concept of *spectral coherence* has been folded into the concept of *temporal coherence* through the use of a corollary of the time-frequency Fourier theorem, known as the Autocorrelation or the Wiener-Khintchine theorem.

Section 2 reviews autocorrelation measurement by a standard Michelson interferometer and section 3 introduces autocorrelation measurements with a long fiver and an amplifier in one arm of a Michelson. After building up these backgrounds, section 4 considers the propagation of a pulse through a Fabry-Perot laser as an amplifier, where we introduce the paradox of classical dispersion theory based on propagating non-physical Fourier frequencies.

2. REVIEW OF AUTOCORRELATION MEASUREMENT BY A MICHELSON INTERFEROMETER

2.1 Developing fundamentals:

Consider a simple rectangular pulse of width ∂t_w with a carrier frequency ν. An electromagnetic signal is generally presented as:

$$a(t) = E(t)e^{i2\pi\nu t} \tag{5}$$

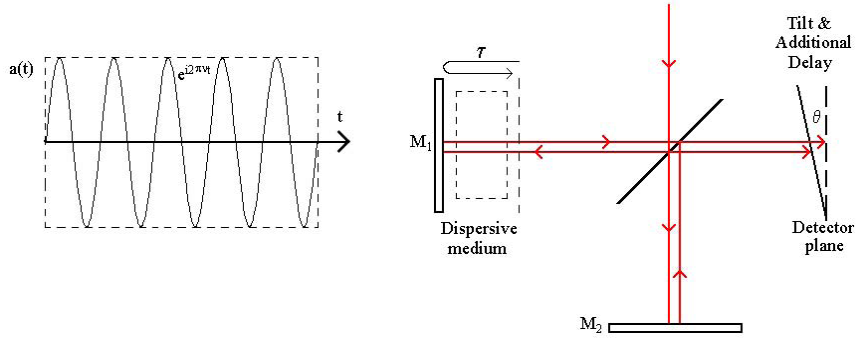


Figure 1. This figure depicts a time finite pulse with a single carrier frequency and a Michelson interferometer (MI) to measure the autocorrelation of the pulse through fringe visibility by superposing a replicated pair of pulses with variable delay. The MI can be used as a traditional free space device or one can insert a dispersive medium in one arm. The MI can also be set up in the "fringe mode" by tilting one of the mirrors to record fringe visibility variation due to a single pulse with a CCD camera.

If we send the pulse through a Michelson interferometer with a total relative path delay τ, a fast detector will register time varying intensity given by:

$$i(t,\tau) = \chi^2_{(1)} \,|\, a(t) + a(t+\tau)\,|^2 = \chi^2_{(1)}[\,|\, E(t)\,|^2 + |\, E(t+\tau)\,|^2 + 2E(t)E(t+\tau)\cos(2\pi\nu\tau)] \tag{6}$$

The time integrated record can be represented by:

$$\bar{i}\,(\tau) = 2[1+\gamma(\tau)\cos(2\pi\nu\tau)] \tag{7}$$

where

$$\gamma(\tau) = \chi^2_{(1)}\int E(t)\; E(t+\tau)\; dt \Big/ \chi^2_{(1)}\int |\, E(t)\,|^2\; dt \tag{8}$$

Notice that the normalization process makes $\gamma(\tau)$ independent of $\chi^2_{(1)}$ as encounters in standard text book expression [9]. By autocorrelation (or Wiener-Khintchine) theorem, $\gamma(\tau)$ and $\tilde{A}(\nu)$ form a Fourier transform pair, where the Fourier intensity spectrum $\tilde{A}(\nu)$ is the square modulus of $\tilde{a}(\nu)$. $a(t)$ and $\tilde{a}(\nu)$ again form a normalized Fourier transform pair [9]:

$$\gamma(\tau) = \int \tilde{A}(\nu)e^{i2\pi\nu\tau}\,\mathrm{d}\,\nu \tag{9}$$

$$a(t) = \int \tilde{a}(\nu)e^{i2\pi\nu t}\,\mathrm{d}\,\nu \tag{10}$$

For a rectangular function, $\tilde{A}$ (ν) is a normalized sinc^2 function and γ(τ) is a triangular function.

$$\tilde{A}\,(\nu) = \tilde{a}\,(\nu)|^2 \Big/ \int |\tilde{a}(\nu)|^2 \mathrm{d}\,\nu \tag{11}$$

Notice the differences in the Fourier transform conjugate variable pairs for Eq. 9 &10. For $\gamma(\tau) \Leftrightarrow \tilde{A}(\nu)$ the variables are (ν,τ), but for $a(t) \Leftrightarrow \tilde{a}(\nu)$ the variables are (ν,t). For $\gamma(\tau) \Leftrightarrow \tilde{A}(\nu)$ the integral is linear in the intensity and hence there is no interference between the different but "coherent" Fourier component frequencies of $\tilde{a}(\nu)$. For $a(t) \Leftrightarrow \tilde{a}(\nu)$ the Fourier integral is linear in amplitude and there is inherent assumption of "interference" between the Fourier component frequencies $\tilde{a}(\nu)$ to generate the resultant time finite pulse out of infinitely long coherent frequencies. Because of this contradiction, the roundabout relationship between

$a(t) \& \gamma(\tau)$ is purely mathematical. The Fourier relationship between $\gamma(\tau) \Leftrightarrow \tilde{A}(\nu)$ works for fringe visibility measurements only under the conditions that different optical frequencies do not interfere with each other. This assertion is valid only when the detector used is slow, as in traditional Fourier transform spectroscopy. In light beating (heterodyne) spectroscopy, one uses a fast detector where the superposition of different optical frequencies generates the difference (beat) frequencies.

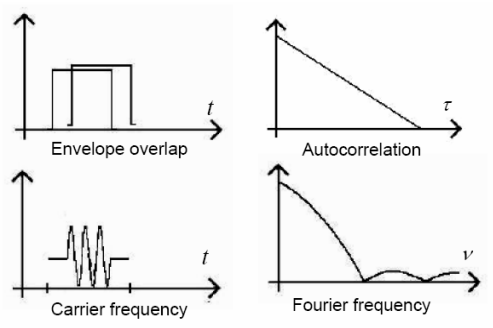


Figure 2a

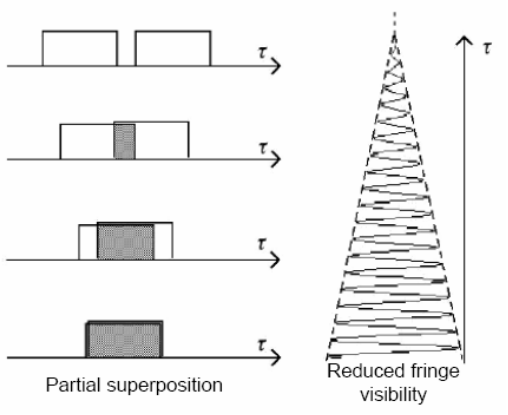


Figure 2b

Figure 2. These figures depict the effects of partial superposition between a pair of replicated rectangular pulses produced by a Michelson interferometer and the corresponding fringe visibility or autocorrelation curves.

Let us carry out the fringe visibility (autocorrelation) measurement by moving one of the Michelson mirrors [Fig 1b] slowly and recording the intensity with a fast detector. Notice that if the two pulses are partially superposed [Fig.2], then for different values of τ the fringe energy undulates between zero and maximum, giving rise to perfect (unit) visibility fringes for a finite domain of τ. But beyond the superposition region one will record intensity without fringes as long as the un-superposed pulses are there.

The situation can be appreciated from Fig.2b. The linear envelope of the superposition fringe intensities gives rise to the triangular γ(τ) already depicted in Fig 2a. The experimental generation of γ(τ) can be done in two ways. First, when the Michelson interferometer is in "filter mode", meaning, the incident pulsed beam, is collimated and the Michelson mirrors are perfectly parallel. In this mode one can get only one data point for each value of τ, requiring the experiment to be repeated many, many times by producing identical pulses. The other alternative is to use the Michelson in "fringe mode", implying the introduction of a tilt in one of the mirrors with one of its edges remaining at relatively zero path delay. Then the superposition effects due to two tilted but collimated pulses can be recorded on a photographic plate or a CCD-camera to time integrate the effect of the partially superposed pulses with variable delays along one spatial axis. The variable fringe visibility envelope is recorded in a single "shot", assuming of course that the Michelson mirror size and the tilt can accommodate the maximum delay τ_{max} that is greater than twice the pulse width $2\partial t_w$. If one is allowed to only examine the final experimental result (Fig.2) and the predictions Eq.7-9, one would be very comfortable with the power of the theory. But one who has carried out the actual experiment will recognize the care and attention required to accommodate the fact that the incident signal is a single pulse. Let us contrast this point by considering a hypothetical laser that can produce the CW spectrum given by Fig.2a or Eq.11. Michelson measurement will produce identical γ(τ) and fringe visibility given by Eq.11 even though the signal is always CW and the experimental difficulties are much simpler than the case when the signal is a single rectangular pulse.

We have already underscored that the integral of Eq.9 for $\gamma(\tau) \Leftrightarrow \tilde{A}(\nu)$ neglects the interference between the different Fourier components $\tilde{a}(\nu)$ whereas the pulse $a(t) \Leftrightarrow \tilde{a}(\nu)$ can be simulated only through interference (summation) of Fourier component frequencies with their amplitudes and phases built into Eq.10. The prediction of Eq.9 is valid only when we use a slow detector at the output of a Michelson in either filter or fringe mode. The response of boundary dielectric molecules of the passive (energy non-absorbing) beam splitter of the Michelson interferometer facilitates the superposition of the two beams. The boundary layer of the beam splitter responds to

different frequencies independent of each other even when they are superposed (combined) over the same region of the beam splitter. Further, detectors that are electrically slow cannot respond to the beat currents (difference frequencies) due to superposition of actual different frequencies. This last point can be appreciated by replacing the slow detector in this experiment with a high speed one. To illustrate this point with mathematical convenience, we will use a CW spectrum containing only two discrete lines, (ν_1 and ν_2), instead of the continuous spectrum of Fig. 2a. For CW light of unit amplitude Eq.6 can be re-written as:

$$i_{fast}(t,\tau) = \left|\chi_{(1)}a_1 e^{i2\pi\nu_1 t + i\phi_1} + \chi_{(1)}a_2 e^{i2\pi\nu_2 t + i\phi_2}\right|^2 = \chi^2_{(1)}\left[a_1^2 + a_2^2 + 2a_1a_2\cos\{2\pi(\nu_1-\nu_2)t + (\phi_1-\phi_2)\}\right] \quad (12)$$

The two amplitudes a_1 & a_2 transmitted by the beam splitter of the Michelson interferometer corresponding to the two frequencies ν_1 & ν_2 are given by:

$$a_1 = \sqrt{2}(1+\cos 2\pi\nu_1\tau)^{1/2} \quad \& \quad a_2 = \sqrt{2}(1+\cos 2\pi\nu_2\tau)^{1/2} \quad (13)$$

These are "reduced" amplitudes due to superposition effects on the Michelson beam splitter in filter mode. The filter mode is required for beat (heterodyne) detection. The Poynting vectors of the two fields belonging to the two different frequencies need to be collinear in addition to their wave fronts being similar in order for the detecting dipoles to respond to both the superposed fields. Otherwise, time varying energy transfer from the two superposed fields to the assembly of detecting molecules cannot take place. These are observed results from various heterodyne experiments and are now easy to appreciate when we accept that the observed superposition effects are due to restricted quantum mechanical behaviors of the detecting dipoles. If we use a slow detector, then the recorded intensity is (no "cross-talk" between two frequencies)

$$i_{slow}(\tau) = \chi^2_{(1)}(a_1^2 + a_2^2) = 2\chi^2_{(1)}\left[2 + \cos 2\pi\nu_1\tau + \cos 2\pi\nu_2\tau\right] = 4\chi^2_{(1)}\left[1 + \cos 2\pi\frac{\nu_1-\nu_2}{2}\tau\cos 2\pi\frac{\nu_1+\nu_2}{2}\tau\right] \quad (14)$$

This is superposition of the fringe intensities due to two independent frequencies. When the "DC" term is removed, the spatial fringes would appear to vary with a frequency $(\nu_1+\nu_2)/2$ with a modulation at a frequency $(\nu_1-\nu_2)/2$. The corresponding autocorrelation function (cf. Eq.7) is given by:

$$\gamma(\tau) = \cos 2\pi(\nu_1-\nu_2)\tau/2 \quad (15)$$

The same result can also be obtained by inverse Fourier transforming [Eq.9] the normalized symmetric spectral density function for this case:

$$\tilde{A}(\nu) = (1/2)\delta(\nu+\partial\nu) + (1/2)\delta(\nu-\partial\nu); \text{ where, } \partial\nu \equiv (\nu_1-\nu_2)/2 \quad (16)$$

Recall that we started with two CW frequencies at ν_1 & ν_2. Thus, the autocorrelation theorem in real world observation is only conditionally valid as in Eq.15 for slow detectors. For a fast detector (Eq.12), one cannot record any stationary autocorrelation factor. The photo detector current is always varying at the beat frequency (ν_1-ν_2). The undulatory beat current can completely vanish if either a_1 or a_2 or both become zero which is possible whenever either $2\nu_1\tau$ or $2\nu_2\tau$, or both assume exactly odd integral values [see Eq.13]. If only one of them assumes a non-integer value, one will have a DC current.

2.2 Distinguishing between temporal & spectral coherence: Autocorrelation with an infinite train of pulses.

Let us assume that with the help of an external chopper we create an infinite train of pulses from a CW single mode laser such that:

$$b(t) = a(t)\sum\nolimits_m \delta(t - m\partial t_s) \quad (17)$$

Where ∂t_s is the pulse separation and $a(t)$ is given by Eq.1 whose width is ∂t_w. Then the Fourier amplitude "spectrum" will be:

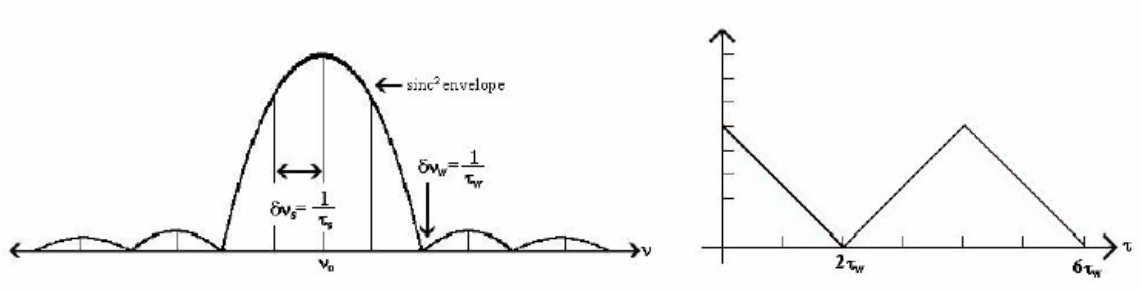
$$\tilde{b}(\nu) = \tilde{a}(\nu)\sum\nolimits_m \delta(\nu - m\partial\nu_s) \quad (18)$$

Or, the normalized Fourier intensity "spectrum" will be:

$$\tilde{B}(\nu) = \left|\tilde{b}(\nu)\right|^2 = \left(\tilde{A}(\nu)/m\right)\sum\nolimits_m \delta(\nu - m\partial\nu_s) \quad (19)$$

If this pulse train is sent through a Michelson with changing relative delays τ, one would record $\gamma(\tau)$ as periodically repeating the result of Fig 2, as shown in Fig.3b. This is easy to see from Eq.9 and Eq.19. Since the carrier frequency is the same in all the pulses and the pulses are derived from a stabilized CW laser with the same carrier

frequency, all delayed pulses will produce superposition fringes. Fig.4 shows the repetitive partial superposition between the two replicated and delayed train of rectangular pulses.



(a) (b)

Figure 3. Fourier spectral intensity function (a) and the corresponding repetitive triangular autocorrelation function (b) due to an infinite train of rectangular pulses with 50% duty cycles with the same fixed carrier frequency ν_0 .

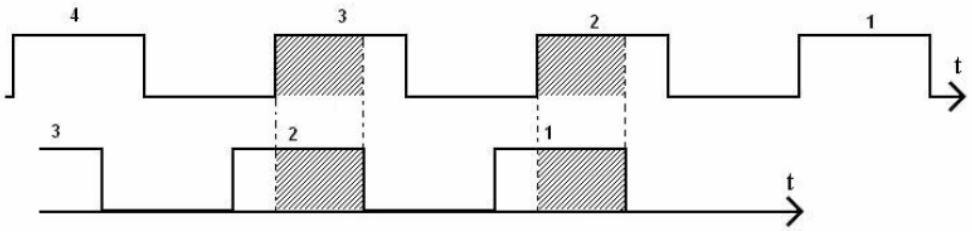


Figure 4. This figure depicts the repetitive partial superposition between the two replicated and delayed train of rectangular pulses. Thus, the measurement process periodically repeats the same autocorrelation function measured with a single pulse of same shape.

Since $\tilde{B}(\nu)$ is periodic in $A(\tilde{\nu})$, the corresponding autocorrelation function $\gamma(\tau)$, the Fourier transform of Eq.19 is also periodic. Again we can distinguish between $\gamma(\tau)$ derived from measurement reality depicted by time integration as in Eq.8 and the Fourier transform of the normalized spectrum $\tilde{B}(\nu)$ [Eq.19]. If we literally start with a CW spectrum source as shown in Fig.4a, we will measure the same $\gamma(\tau)$ as for the infinite train of pulses [Fig.5], even though the situations are physically very different. We ignore the time domain information if we focus only on the time integrated $\gamma(\tau)$. Note that the two cases discussed here for a single and an infinite trains of pulses, the carrier frequency has a single value. There is no spread of actual optical frequency simply due to chopped amplitude. We have demonstrated this in a separate heterodyne experiment [8].

3. AUTO CORRELATION WITH AN AMPLIFIER IN ONE ARM OF THE MICHELSON

In the previous section, the two Michelson interferometer (MI) arms were in free space and the delay was frequency insensitive. In this section, we introduce a dispersive medium, say, a semiconductor amplifier, as shown in Fig.1b. We can think of practical situations where a very weak pulse has to be amplified before it is used for some other sophisticated applications where its physical frequency content is necessary critical information, e.g., WDM communication through very long fiber cable.

3.1 Single pulse case

Let us assume that the refractive index of the medium is n at the physical carrier frequency ν of the pulse. So, if we experiment with a CW light, then the extra relative delay in the MI arm due to the medium of length l would be:

$$\tau' = \tau + (n-1)2l/c \tag{20}$$

τ also includes the space delay in the same arm. However, the velocity of a pulse center is given traditionally by

group index n_g, given by [12]

$$n_g = n - \lambda(dn/d\lambda) = n + \nu(dn/d\nu) \tag{21}$$

So, if the Michelson free space delay is τ and we insert a medium of length l and send a pulse, the relative delay τ' that it will experience is

$$\tau' = \tau + (n_g - 1)2l/c \tag{22}$$

Everybody would recognize that theoretically we cannot just replace τ by τ' in Eq.7 & 8 to find $\gamma'(\tau')$ from measured fringes $i'(\tau')$. The autocorrelation shape $\gamma'(\tau')$ for pulses using dispersive delay will be different from $\gamma(\tau)$ for non-dispersive delays since the pulse shape would be distorted due to "dispersion", typically represented by differential phase shifts experienced by the different Fourier component frequencies. But, we have already seen that Wiener-Khintchine theorem works even though the Fourier frequencies are not physical. We also know that the pulses are broadened by dispersive media in the real world. Before we can search for the deeper physics behind pulse broadening, we need to assure ourselves that the Fourier approach is not the only possible approach to modeling pulse dispersion even though its prediction does correspond closely to the measured results in many circumstances. However, serious contradictions (e.g. superluminal group velocity) are already being encountered by many publications in the field of slow/fast light [12] where people are routinely propagating the Fourier transformed frequencies for a given pulse. We are trying to establish, through a series of mathematical and experimental contradictions that the very approach of using superposition of Fourier frequencies cannot properly model real physical situations simply because EM waves do not sum themselves as is being assumed by the Fourier theorem.

Clearly $\gamma'(\tau')$ with dispersive media will not be equal $\gamma(\tau)$ without dispersive medium. Our numerical computations indicate that for a short length of dispersive medium we will not be able to discern the differences between Fourier predictions (Eq.9 & Eq.22). But we know that there is a built in contradiction because Eq.9 neglects "interference" between different Fourier components, whereas Eq.10 and the derivation of Eq.22 accepts the field superposition effects as physically real. We need a very long-path dispersive medium (large τ') and hence we must deal with an infinite train of pulses to introduce pulse superposition with very large delays and dispersion, because a very long pulse will not display measurable dispersion effects.

3.2 Train of pulses

Let us insert a very long spool of single mode fiber into one arm of the Michelson with the help of appropriate optics. Let us now apply dispersive delay a la Fourier on the Fourier frequencies depicted in Fig.3a. Modeling tells us that because of rapid decay of the sinc envelope, one can demonstrate the essential effects with about thirteen central Fourier frequencies due to a 50% duty cycle rectangular pulses (from Eq.18).

$$\tilde{b}(\nu) = \sum\nolimits_m b_m e^{i2\pi\nu_m(t+\tau_m)}$$

$$\tilde{b}(\nu) = \sum\nolimits_m b_m e^{i2\pi\nu_m t} = b_0 e^{i2\pi\nu_0 t} + b_{\pm 1} e^{i2\pi(\nu_0 \pm \partial\nu)t} + b_{\pm 3} e^{i2\pi(\nu_0 \pm 3\partial\nu)t} + b_{\pm 5} e^{i2\pi(\nu_0 \pm 5\partial\nu)t} + \ldots \tag{23}$$

τ_m is the total round trip through the fiber inside the MI; the free space delay is balanced out by the other MI arm. The values of $b_{\pm m}$'s are secured through normalized spectral amplitude envelope sinc, and not sinc^2, as shown in Fig.3a.

Let us now assume that we have found the unique roundtrip length of the fiber L_0 such that all the thirteen Fourier frequencies, chosen for this computation, undergo modulo $|2\pi|$ or an integral number q_m times 2π phase delay (see Fig.4a & d):

$$2\pi\nu_m\tau_m = 2\pi\nu_m(n_m 2L_0/c) = 2\pi q_m \tag{24}$$

When this "integer condition" is met, all the thirteen frequencies of Eq.23 will emerge from the MI in phase and the regenerated pulse shape due to superposition will not show any dispersive effects at all. In other words, fringe visibility measurements around L_0 or its integral multiple length (albeit superposition of two different delayed pulses) will give autocorrelation values very similar to Fig.2, implying $\gamma'(\tau') = \gamma(\tau)$ (for convenience measurement, we have assumed that the free-space delay differential in the Michelson has also been adjusted to modulo $|2\pi|$ delay for the central carrier frequency ν_o).

Now, let us look at the results of modeling. For propagation delays L_0 equal to 912km and 1824km, the pulses appear undistorted (Fig.4d) compared to the input pulses (Fig.4a). For a propagation delay of L_d equal to 600km, the width of the main pulse has reduced by almost a factor of three (Fig.4c). Is it pulse compression? For the case of L_d, the numerical values for $\nu_m \tau_m$ are appreciably different from integer. We have chosen this case to dramatize the contradiction we face when we use mathematical Fourier frequencies as if they are real physical.

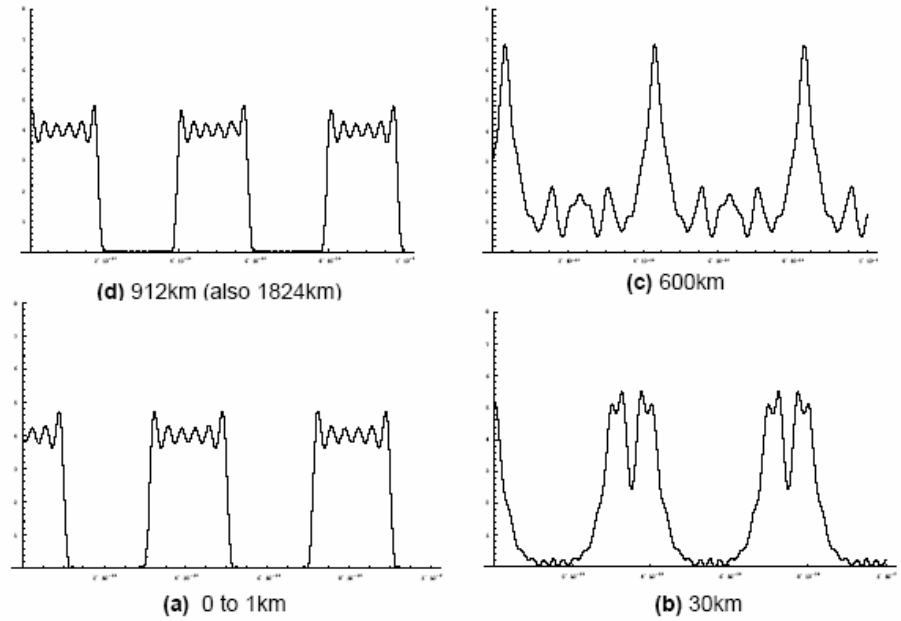


Figure 4. Pulse dispersion effects through a long fiber using Fourier frequencies of an infinite train of rectangular pulses (200ps with 50% duty cycle). Plots are from the square modulus of the Eq.23 for different lengths of fiber. (a) Input pulses (remaining undulations are due to modeling that uses only 13 of the large number of Fourier frequencies with appreciable energy). (b) Pulse distortion at 30km; the pulse is almost split into two. (c) Pulse distortion at 600km; notice that the main pulse is now much narrower than the input pulse! The peak energy is distributed in the secondary pulses. (d) Restoration of the original pulse shape after long propagation. Thus, Fourier theorem predicts dispersion free propagation for certain specific lengths of propagation!

We believe that the predicted results of Fig.4 are not achievable in practice with AM signals, but can be obtained with a CW source that actually contains the carrier frequencies given by Eq.23. We know from CW measurements that the light velocity in a medium is a function of the actual frequency. This is not in question. We are trying to explore the physical origin of pulse broadening that should not be assigned to differential propagation delay due to mathematical Fourier frequencies that are not physically present in the pulse. The silica molecules on the fiber respond to light in less than a femto second. They cannot anticipate the shape and repetition rate of an incident signal, nor do these molecules have the capability to carry out Fourier transformation of a very long duration pulse train. We believe the pulse broadening is due to "time diffraction" that we have modeled in another paper [13]. See Fig.5 below. For the simplest case of free space propagation, we have derived the origin of pulse broadening using Huygens-Fresnel principle (HF), which is directly proportional to the diffraction order. A pulse of width δt will be broadened approximately to:

$$\delta t_s = \partial t + \tau_{g-m}; \quad \tau_{g-m} = N(m\lambda / c) = mN / \nu \tag{25}$$

All spectrometers, including gratings, have a time constant, which we have defined [14, 8] as the resolving power (number) divided by the carrier frequency of the light. In our case here τ_{g-m} is the time constant for the grating, where N is the total number of grating lines and m is the diffraction order. In contrast, a very high finesse Fabry-Perot (FP) can provide a very large temporal delay [14].

$$\tau_{FP} = mN / \nu; \quad N = \pi\sqrt{R}/(1-R); \quad m = 2nd / \lambda \tag{26}$$

This is very similar to Eq.25, but now N is the finesses of the FP and d is the thickness of the FP plate of refractive index n with reflective coatings R on both sides. Large R can give large N and hence large FP spectrometer time

constant. For a free-space grating, the diffractive spreading of a pulse δt_s [Eq.25] is much easier to visualize [Fig.5] than for an FP etalon with material dispersion. Waveguides and fibers introduce complex temporal diffraction delays (due to effective index generated by core and cladding and boundary conditions) compared to a free space grating. So, the direct time domain propagation of a pulse through a waveguide is quite complex. We have used FDTD method to find the time integrated response of an AWG (array wave guide grating), as a simplified example of our concept [see Fig.5 in Ref.13a].

If one really carries out the above experiment with a CW spectrum equivalent to the Fourier spectrum of Fig.4a or its close equivalent using a high quality frequency stabilized (to 1kHz) gas laser (like He-Ne), one would actually be able to measure such autocorrelation curves equivalent to the ones predicted above.

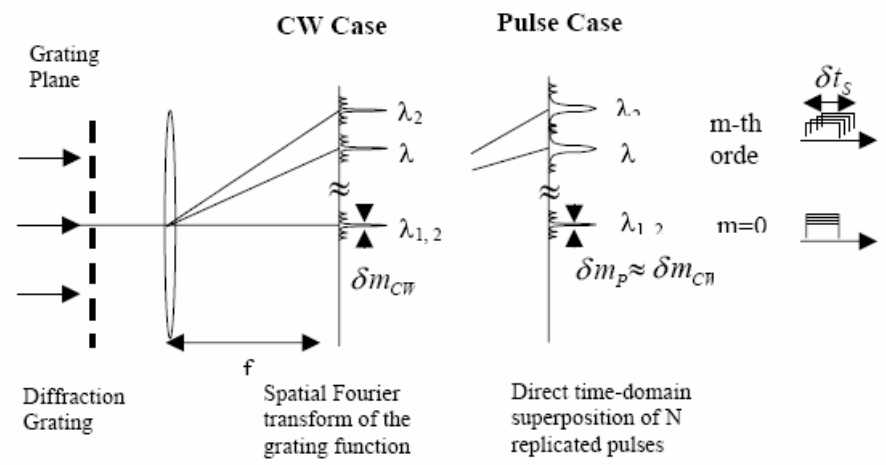


Figure 5. Pulse broadening by a free space grating is due to superposition of N replicated pulses produced by the N-slits and then superposed in the far-field, or by the optical system in a spectrometer with a periodic delay $\tau_m = m\lambda / c = m / \nu$ at the m-th order diffraction. The results are derived by using direct time-domain propagation of the pulse using Huygens-Fresnel diffraction principle. For a free-space grating, the zero order propagation is free of apparent "dispersion". [From ref.13].

4. PULSE PROPAGATION THROUGH A SOLID FABRY-PEROT

4.1 Passive high finesse Fabry-Perot of a single pulse

For a free-space FP, the spectral fringe width broadened due to finite pulse $a(t)$ is given by [8, 14.]:

$$I_{FP-pls.}(\nu,\tau) = I_{FP-CW}(\nu,\tau) \otimes \tilde{A}(\nu); \quad \tau = 2d/c; \; \Delta\nu_{fsr} = c/2d = 1/\tau \tag{27}$$

But for a solid etalon, this should change for pulsed light because we need to use the group velocity via n_g.

$$I_{FP-pls.}(\nu,\tau') = I_{FP-CW}(\nu,\tau') \otimes \tilde{A}(\nu); \quad \tau' = 2n_g d/c; \; \Delta\nu_{fsr\text{-}pls.} = 1/\tau' \tag{28}$$

In general, n_g (Eq.21) is not particularly sensitive to pulse width or the carrier frequencies except around resonance frequencies of the material. So, one way to find the validity of n_g is to measure differential free-spectral range for a light signal of the same frequency under CW and pulsed conditions (Fig.6). By classical dispersion theory, the CW and pulse velocities are (c/n) & (c/n_g). The free spectral range through a solid FP etalon for the same carrier frequency should be different:

$$\partial(\Delta\nu_{fsr}) = \Delta\nu_{fsr-CW} - \Delta\nu_{fsr-pls.} = \frac{c}{n2d} - \frac{c}{n_g 2d} = \frac{c}{n2d}(1 - n/n_g) \tag{29}$$

We have mathematically shown for a free-space FP [8, 14] that for any pulse whose width is longer than the etalon time constant $\tau_{FP} = mN/\nu$ [Eq.26], the FP response curve becomes identical to the CW response function. So, we have no reason to accept the idea that a long pulse (relative to the FP thickness and finesse) will experience a

velocity (c/n_g) rather than (c/n) when the carrier frequency is same even for a solid FP. But to dispel the doubt, we are in the process of carrying out the following experiment in our lab using a solid etalon and by measuring the free spectral range with CW and pulsed light with the same carrier frequency, making sure the pulse length is definitely shorter than the etalon time constant.

Let us consider a 1cm silica etalon with measured finesse N=150, which gives τ_{FP}=10.444ns for CW case. If we further consider $\lambda = 1.55\mu$, $n = 1.444388$, $(dn/d\lambda) = -0.01189888\mu^{-1}$ and $(1 - n/n_g) = 0.0126078$ from Eq.21 and ref.12. Then,

$$\partial(\Delta\nu_{fsr}) = 130.932MHz \ \ (d = 1cm) \tag{30}$$

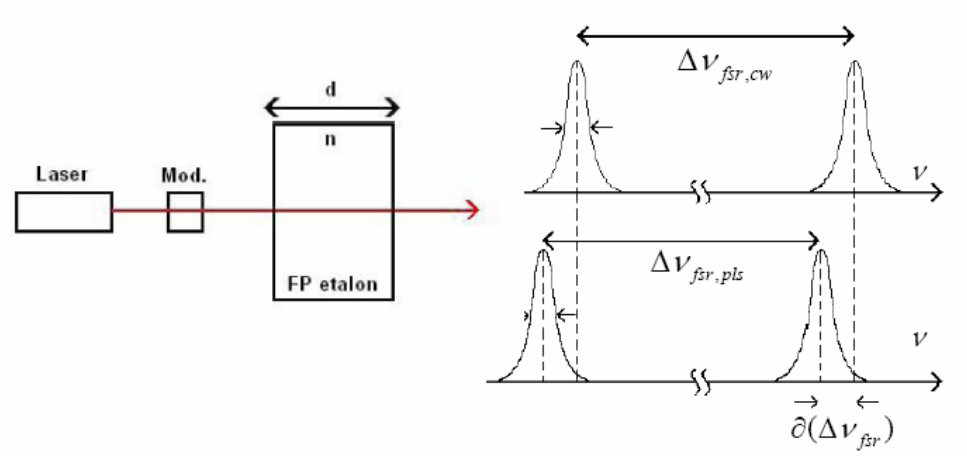


Figure 6. Measurement of the free spectral range of a solid Fabry-Perot etalon (or an amplifier) using CW and pulsed light to test whether carrier frequency or the Fourier frequencies determine the peak transmissions through a high finesses Fabry-Perot.

Thus, using a precision tunable CW laser around 1550nm, an external amplitude modulator and a high finesse silica FP etalon, one can measure the differential free spectral range using CW and single pulses to test the validity of classical dispersion theory based on propagating Fourier frequencies. For the convenience of precision measurement, one should note that the pulse length should be much larger than the single-pass delay for superposition effects to take place and to generate high finesse fringes. FP fringes broaden as the pulse width decreases towards the single-pass delay; again that is not due to Fourier frequencies [8, 14]. But, since the classical dispersion theory is blind to the pulse width, the reader should be happy to use a pulse width longer than the FP time constant τ_{FP}. The numerical case sighted above gives an idea that it is actually verifiable with traditional laboratory equipment available these days. However, it is critical to choose the precision of the tunable laser that is capable of achieving what is dictated by Eq.30 by manipulating the etalon thickness and its finesse. Our prediction is that we will find $\partial(\Delta\nu_{fsr})$ to be close to zero within our measurement precision.

4.2 A single pulse through an active FP laser

Consider now the FP etalon to be a very narrow line width FP diode laser. They are "strongly" homogeneously broadened lasers and runs at a single mode with line width easily less than 100MHz when the facet coatings are highly reflective.

We want to measure the gain of an external input frequency at CW and pulsed state. People have demonstrated that diode lasers can be mode locked at two different frequencies by proper external cavity arrangements. We use this cue to amplify a frequency ν_{input} which is several cavity modes apart from the gain line center naturally chosen by the laser, ν_{laser}. A tunable second laser is used to transmit and measure the output power by carefully tuning the ν_{input} to match up with the FP laser's internal FP resonances given by $\delta\nu = c/2nL$. Once this is validated with the external tunable laser line, it could be modulated and checked for any change in the input frequency for maximum amplified output power, because, by classical dispersion theory the resonance frequencies should be different:

$$2dn = m\lambda_{m,cw} = mc/\nu_{m,cw}; \ \ 2dn_g = m\lambda_{m,pls} = mc/\nu_{m,pls} \tag{31}$$

The laser mode spacing and free spectral range for a solid FP are identical in mathematical expression and hence the

differential value for the frequency can be expressed by copying Eq.29:

$$\partial \nu_{res} = \frac{c}{n2d}(1 - n/n_g) \tag{32}$$

Our quick calculation implies that $\partial \nu_{res}$ would be in the domain of many GHz and easily measurable if n & n_g were really different rather than being n for both the CW and pulse cases, but determined by the carrier frequency, as long as the pulse contains many optical cycles.

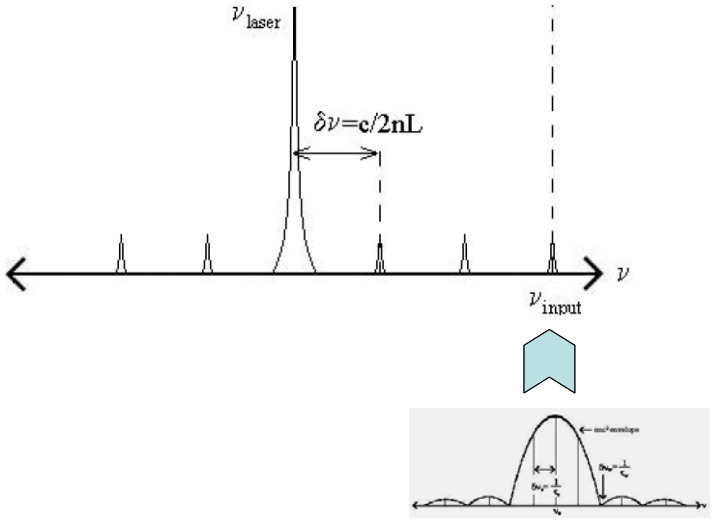


Figure 7. Determination of Fabry-Perot amplifier resonance frequency for pulsed and CW input light to validate whether Fourier frequency driven group velocity or the carrier frequency driven CW index determine the Fabry-Perot resonances in an active laser gain medium. If amplitude modulation really generates Fourier frequencies, then all of these new frequencies would not resonate with the amplifying laser's satellite mode, as depicted above.

As mentioned earlier, pulses generated by an external amplitude modulator from a single frequency CW laser contain only the original laser frequency, what we call the carrier frequency. Our point is that we should propagate the carrier frequency in the actual time domain, rather than the non-causal Fourier frequencies existing over all time. Note that, using heterodyne techniques, one can also compare the actual line width (distribution of actual carrier frequency) of the amplified signal for both CW and pulse experiments. We have measured the carrier frequency of an amplitude-modulated laser directly in an experiment [8, 15] and found no Fourier frequencies.

The case presented here is essentially to question the validity of classical dispersion theory based on the Fourier frequencies of a pulse. However, it can also be taken as a guide for using n instead of n_g in modeling, say, the amplification of a weak communication pulse train.

4.3 A train of pulses through an active FP

Let us now consider the same experiment but with an infinite train of rectangular pulses rather than a single pulse. Further, we assume that we can vary the duty cycle. Then the Fourier spectrum will effectively be a finite set frequency like those shown in Fig3a. Again precise measurement should be carried out to find any changes in the differential free-spectral range (Eq. 28) under CW and pulse train conditions. But, a more interesting experiment would be to measure the change in the gain of the output pulse train as the pulse duty cycle is varied. Since this will change the spacing between Fourier frequencies centering on the FP-laser mode coincident with the carrier frequency ν_{input}, amplification of the output pulses should vary with the duty cycle. However, our prediction is that there would not be any change in the amplification (by avoiding saturation, etc.) of the pulses since the carrier frequency ν_{input} remains fixed and aligned with the FP laser side mode. One should also carry out separate heterodyne measurements to check the presence of Fourier frequencies given by Fourier theorem for amplitude modulated light [8].

5. DISCUSSIONS

Our key platform has been the explicit recognition that light beams by themselves do not "interfere" with each other. We see light always through the "eyes" of the material detectors that must first respond to the total resultant dipolar stimulation(s) induced by the impinged EM field(s). The "superposition effects" due to light [5, 15] are "colored" by the detectors as they respond to the stimulation by the light beam(s) using their characteristic and unique (limiting) quantum (response) properties.

The Fourier theorem, by implying that light beams can sum by themselves to create a new resultant EM field, effectively ignores quantum mechanical rules for actual detection processes behind superposed light beams. So we must carefully re-visit all possible physical conclusions that we derive using the various aspects of the time-frequency Fourier theorem. From galactic systems to single biological cellular system, all are evolving through interactions between elementary particles, atoms and molecules. They all have unique characteristic life times and there are specific time constants for all the specific interactions. So, just for the convenience of computations, we can not arbitrarily get rid of the time from a finite duration pulse and replace it with time-free Fourier frequencies. It is bound to introduce unnecessary non-causal and/or contradictory results, as we have shown in this and other papers [13, 16, 17]. Other fields like slow/fast light are also encountering many self-contradictory and non-causal results like superluminal group velocity [12]. Even the quantum mechanical indeterminacy relation is only a corollary of a Fourier transform relation, mathematically strengthened by Schwartz's inequality theorem [18].

We started with the question: If EM fields do not operate on each other, why do we need many modes and large gain bandwidth to generate short pulses? We have responded by arguing that, since the modes by themselves cannot produce pulses, it has to be the quantum properties of the lasing molecules and the saturable absorber medium that do so. We believe that such re-thinking will help us model laser pulse generation, propagation and amplification more precisely by relying on material properties. Then we focused on bringing a series of simple paradoxes out in the open that have always been there behind the use of time-frequency Fourier theorem and Wiener-Khintchine (autocorrelation) theorem. This has been done through a series of modeling tied to proposed experiments that can be carried out in any reasonably equipped laboratory for validation of our assertions. Since the light signals are always space and time finite (even a CW laser has to be turned on and off), we suggest that we should re-derive all the optical results by direct time-domain propagation of the pulse through the physical system under consideration using the Huygens-Fresnel principle. Actual E-vector undulation frequency (or frequencies), what we call carrier frequencies, contained in the pulse should be treated as the physical spectrum. Such E-vector undulations (carrier frequency) of ultra short pulses have recently been measured [19]. Such an approach opens up better understanding of the actual light-matter interaction processes. It will also open up possibilities for new innovations leading to better optical instruments.

The purpose and the first priority of a hypothesis to model nature should be based on actual observations under diverse conditions. Secondly, it should be an aid to our imagination (and visualization) of the real physical process occurring in nature [20, 15]. If the starting paradigm (hypothesis) is not anchored to some real physical process of nature, it is eventually bound to give rise to paradoxical results. We are aware that any time we encounter one paradox between a well established theory and an experiments, it is an occasion for us to re-visit the logical congruence between the specific derivation and the particular experiment. But, when we encounter a good number of paradoxes over a widely differing derivations and experiments, as we have outlined in this and other papers [4, 5, 8, 13-18], then it is time to re-visit the relevant paradigm itself instead of developing separate explanations for individual experiments. In our case, the common identifiable paradigm is the assumption that superposed light beams can generate a new resultant field by themselves. So, we propose that the paradigm of "*interference of light*" should be replaced by the paradigm of "*superposition effects due to light*" (as experienced and recorded by detectors).

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