

Reality of superposition principle and autocorrelation function for short pulses

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ABSTRACT

We “see” light only when some material detectors (dipoles) respond to the incident EM field. EM fields do not operate on each other to make themselves visible to us. Superposition of multiple fields becomes manifest only when the intrinsic properties of these dipoles allow them to respond to all the superposed fields simultaneously and thereby summing the effects of all the fields. Accordingly, depending upon the different intrinsic properties of the detectors and the physical conditions of measurements (integration times, etc.) the manifestation of the “coherence” properties for the same set of superposed fields could be different. It is then prudent to represent the autocorrelation function for superposed fields in terms of the dipole undulation of the detectors rather than the fields themselves. Then the physics of the detectors and the measurement conditions automatically becomes an inherent part of the discussion on coherence. We illustrate our premise by presenting the analysis to understand the behavior of beam splitters, two-beam interferometers and an N-beam grating “interferometer” in terms of the autocorrelation functions due to a short pulse as would be experienced by the material dipoles of the beam splitters and detectors. Our approach reveals that superposition effects to become manifest the multiple fields must be physically superposed simultaneously on the detecting dipoles and hence the process is causal.

Key Words: Short pulse autocorrelation function, Coherence function; Interference of light, Non-interference of light; Two-beam interferometers, Grating fringes with a short pulse.

1. INTRODUCTION

1.1. Justification of the paper, the sustained paradoxes.

The timeliness of the paper can be appreciated by the fact that Glauber has shared the 2005 Physics Nobel Prize for developing the theory of quantum coherence. The engineering importance of the paper derives from our belief that innovative ideas for new applications of ultra short light pulses may be generated our community if we can *bring the conceptual congruence between the engineering and the physics definition for the same photon that constitute our light pulses*. The many engineers tacitly assume the photons to be contained within the volume of the pulse and hence they are “local” in space and time because the measurability of the energy of the pulses is determined by the classical optical arrangements (like interferometers, laser cavities, optical modulators, photo detectors, etc.). The measured results (“reports”) by these instruments and detectors help us define (imagine) all the spatial and temporal characteristics of the light pulses, which are always well defined in space-volume and time-duration. In contrast, the physicists tend to define photons as the Fourier monochromatic modes of the vacuum [1], which, by virtue of the time-frequency Fourier theorem, are necessarily “non-local”. We now add to this physics model the traditional hypothesis that single photo-electron “clicks” are exclusively due to the absorption of “one at a time”, indivisible, single photons. This enforces one, specifically while carrying out interference and diffraction experiments at extremely low light levels, to assign several non-causal properties to the photons, like (i) “self-interference”, (ii) “delayed choice”, (iii) “teleportation”, etc.[2]. Thus, we have some paradoxically divergent

concepts apparently held by many engineers and physicists. Recent literature shows that attempts are being made to bring some convergence [2, 3, 4]

1.2. Reality of superposition principle

For centuries we have been loosely using the phrase “interference of light” even though we know that electromagnetic fields (EM) do not operate on each other in the absence of mediating material media, which respond to EM field vectors as dipoles. In free space or in a non-interacting medium, the superposed fields can propagate uninfluenced by each other even while occupying the same physical volume, whether the beams are collinear or just crossing through each other. Other wise the visual world would have always appeared as scintillating speckles and the WDM fiber optics communication would not have worked. This is true whether the superposed light beams (i) are generated by a “coherent” laser source or an “incoherent” thermal source, (ii) contained distinctly different E-vector oscillation frequencies, (iii) have E-vectors that are oscillating parallel or orthogonal to each other. Light is observable (visible) to us only through the energetic transformation experienced by some polarizable detecting atoms or molecules. The effects of superposition of more than one EM fields in a physical domain can become manifest only when the observing detector is capable of responding to all the E-vectors simultaneously and thereby summing and displaying the superposition effect. Thus, what we call “interference of light” is actually the vectorial summing capability or incapability of the detecting molecules and the correspondingly proportionate energy absorption from the fields by them. As a result, the spatial and/or temporal variations in the number of transitions of elementary detectors within the domain of superposed fields appear to us as “interference” fringes. Consequently, registrations of “interference” fringes require real physical superposition (presence) of multiple fields on the detecting molecules that can be simultaneously experienced by them [5]. Further, the response times of most of our photo detectors (silver halide molecules, photo conductors, etc.) are mostly femto seconds or shorter. We also know that all light beams have finite time duration; they are fundamentally pulsed. Even a laser has to be turned on and off for an experiment. If the pulses carry multiple carrier frequencies and phases, or these parameters are chirping, the spatial and temporal location of the registered “fringes” are different and the integrated result appears as reduced visibility fringes. This is not intrinsic “incoherence” of the EM fields; they have successfully delivered their energy in sub-femto second time periods. Thus the “coherence time” should be defined with reference to integration (or exposure) time of the detecting system and in terms of the degree of overlap of the superposed pulses, and of course, frequency and phase distributions and their chirping of the EM fields would be essential input parameters. Pulses with extremely steep (femto second domain) phase and frequency chirping will require deeper investigation of the time required for light induced measurable transitions, especially since EM fields deliver energy while on their eternal highest velocity existence.

Since light does not “interfere” with light (QM defines them as Bosons), it is hard to create a physical model for the process behind the observed superposition effects that are congruent with the popular belief that an indivisible single photon can “interfere” with itself. Because of the inherent quantum nature of electrons and their binding energies in different detector molecules, all photo chemical and photo electric detectors are forced to produce discrete set of “clicks”. This is not a definitive proof that multiple superposed wave packets (photons) cannot induce excitation and cannot collaboratively share energy to bring about transitions of the detecting dipoles [5]. In fact, in a series of very careful experiments, Panarella [6] has found that the intensity of the light beam must be equivalent to four photons at a minimum to produce interference effects (diffraction rings by a pinhole). Since interference effects and coherence functions are almost synonymous in current literature [7], we present arguments that the actual coherence functions should be interpreted, not as auto- or cross-correlation between different superposed fields, but as auto- or cross-correlation between the different dipole undulations induced by the superposed fields. For this paper, we do not need to use any quantum mechanics. Our arguments will be based on easy, semi-classical formalism since Sudarshan has already established (“Optical Theorem”) that Glauber’s quantum coherence functions are equivalent to Wolf’s classical coherence functions [7]. Besides, Jaynes, Lamb, Scully, etc. have consistently shown that semi-classical formulation explains photo induced reactions very well [8, 9, 10]. We will treat light as classical wave packet with a uniquely defined carrier frequency and an amplitude envelope, $\bar{a}(t)\exp[i2\pi\nu t]$. Section 2 is devoted to explaining two-beam superposition effects in terms of second order temporal coherence function. We will underscore the critical role played by the beam-splitters and the relative phase shift “ π ” they introduce on “external reflections” (reflection from a denser boundary back into the lighter medium) [11], which is essential to conservation of energy. Section 3 will model the response of an N-slit grating in terms of summation of all possible combination of second order temporal coherence functions between the N beams produced by the grating.

2. TWO-BEAM SUPERPOSITION

2.1. Re-direction of field energy by a passive dielectric boundary

All two-beam interferometers have one or more beam splitters in some form or other. Accordingly we need to understand the specific roles it play while interacting with electromagnetic fields incident on it from opposite directions. Let us consider the Fig.1(a) where two well defined beams, which are coherent to each other, are intersecting in free space (ignoring the photographic plate in their junction for a moment). Neither of the beams will experience any change in their spatial energy distributions in the absence of any photo-refractive material. Both classical and quantum physics accept that EM fields do not operate on each other in the absence of interacting material medium. But when a photo detector array or a photographic plate is inserted within the volume of superposition of the well defined beams, one can register the fringes parallel to the bisector of the angle between the beams. One should appreciate that the beams themselves do not re-distribute their energy spatially and re-constitute themselves to the original form before emerging out of the volume of superposition. Light beams carry energy at the uniform speed as per the Poynting vector. The local detecting material particles (detecting dipoles) attempt to respond to all the fields simultaneously present on it when allowed by quantum mechanical rules. The regions where the two E-vectors try to induce undulations in the opposite directions, the dipoles fail to get excited and the region registers a dark spot (fringe); it is not because the E-field energy did not pass through that region [12, 4b].

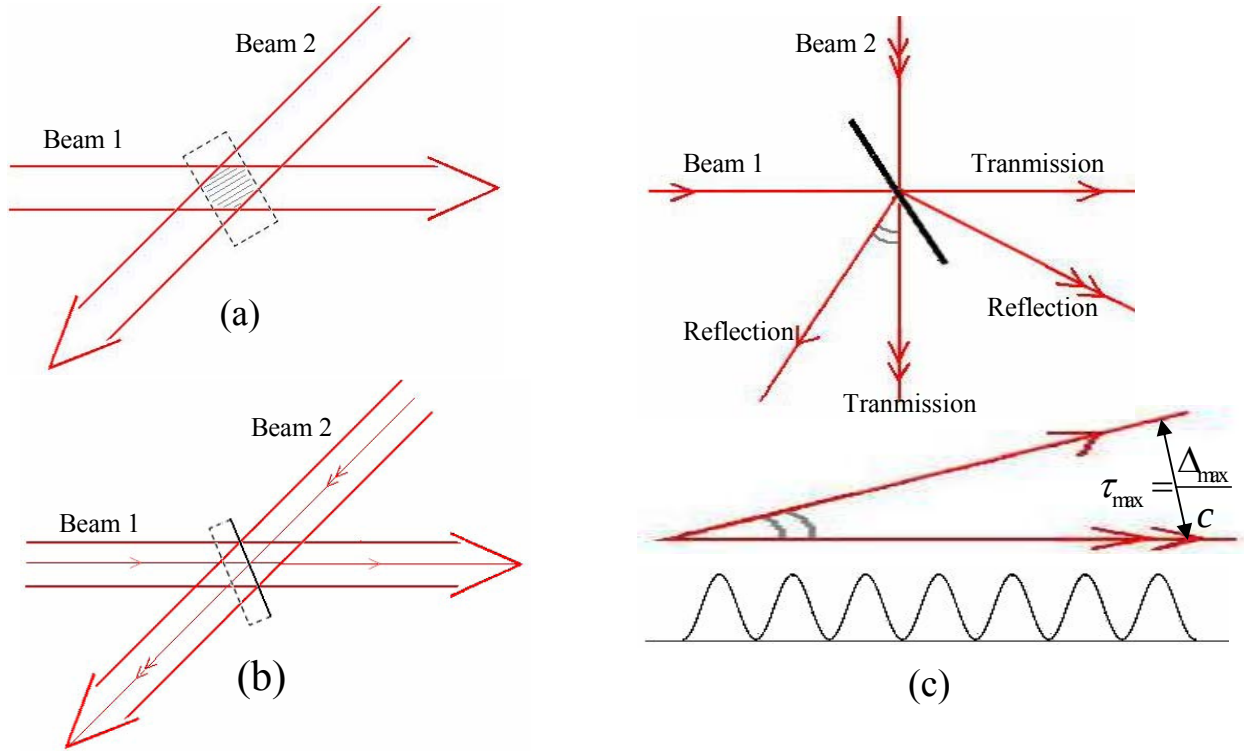


Figure 1. Appreciating non-interference of light in the absence of interacting material medium. In (a) we see two well defined beams cross each other without any mutual influence until we put a photographic plate or a photo detector array when we observe spatial fringes. (b) shows that a dielectric boundary can re-direct the energy of one beam into the other. Under right condition of real physical superposition of two beams from opposite directions, the beam splitter can become a 100% transmitter or a 100% reflector provided the Poynting vectors for the reflected and transited beam pairs are perfectly co-linear. (c) shows the situation when the Poynting vectors make an angle with each other. Then the beam splitter cannot impose any superposition effects directly on the beam energy. The two tilted wave fronts propagate out without influencing each other unless one inserts a detector array within the superposed beams that reveals spatial fringes, as in (a).

Let us now insert a passive dielectric boundary, like a common beam splitter, within the volume of superposition as shown in Fig.1(b). All of our practical interferometers have one or more beam splitters that help divide the incident beam(s) and recombine them to produce the superposition effects. Let us review the process according to classical physics. We assume that the two beams have been produced from the same single frequency laser but with a relative delay of τ between them. When the Poynting vectors are perfectly collinear, the energy of

the beams can be redirected by this passive beam splitter either partially or completely depending upon the strengths of the incident amplitudes, the π phase shift between the “external” and the “internal” reflections [11] and the relative propagational phase delays. First we are considering the case of co-linear Poynting vectors. The right-going and down-going amplitudes are given by Eq.1. Just to appreciate the role of the dielectric boundary, we are considering the incidence of two light pulses of same frequency and very long duration (effectively CW) of amplitudes $\vec{d}_1$ and $\vec{d}_2$. We are also assuming that both the polarizations are parallel to each other and perpendicular to the plane of incidence to avoid polarization dependence of r and t . The corresponding transmitted and reflected amplitudes are designated by $\vec{d}_1$ and $\vec{d}_2$. The vector symbol $\vec{d}$ has been used to underscore the role of the material dipoles, which is equal to $\vec{d}$ within a non-absorbing medium. The general case for pulsed light can be taken care of by making $\vec{d}_1(t-t_1)$ and $\vec{d}_2(t-t_2)$ time dependent:

$$\begin{aligned}\vec{d}_{\text{down}}(t, \nu, \tau) &= \vec{d}_1 \cdot r \cdot e^{i2\pi\nu(t-t_1)} + \vec{d}_2 \cdot t \cdot e^{i2\pi\nu(t-t_2)} \\ \vec{d}_{\text{right}}(t, \nu, \tau) &= \vec{d}_1 \cdot t \cdot e^{i2\pi\nu(t-t_1)} + \vec{d}_2 \cdot r \cdot e^{i\pi} \cdot e^{i2\pi\nu(t-t_2)}\end{aligned}\quad (1)$$

The corresponding intensities are:

$$\begin{aligned}I_{\text{down}}(\nu, \tau) &= |d_{\text{down}}|^2 = d_1^2 r^2 + d_2^2 t^2 + 2\vec{d}_1 \cdot \vec{d}_2 tr \cos 2\pi\nu\tau \\ I_{\text{right}}(\nu, \tau) &= |d_{\text{right}}|^2 = d_1^2 t^2 + d_2^2 r^2 - 2\vec{d}_1 \cdot \vec{d}_2 tr \cos 2\pi\nu\tau\end{aligned}\quad (2)$$

And the general conservation of intensity (energy) can be easily found from the Eq.2:

$$I_{\text{total}} = I_{\text{right}} + I_{\text{down}} = d_1^2(T + R) + d_2^2(T + R) = d_1^2 + d_2^2 \quad (3)$$

One can now appreciate the possibility of converting a glass plate or a 50% beam splitter into a 100% reflector or transmitter by manipulating the reflectivity and/or the two amplitudes. Let us consider the particular case when all the energy is directed to the lower beam, “robbing” from the right beam. We need the incident amplitudes to be equal $d_1 = d_2 = d$ and also $R = r^2 = T = t^2 = 0.5$. Then,

$$\begin{aligned}I_{\text{down}}(\nu, \tau) &= d^2(t + r)^2 = 2d^2 \\ I_{\text{right}}(\nu, \tau) &= d^2(t - r)^2 = 0\end{aligned}\quad (4)$$

Mathematically these are trivial relations from undergraduate optics. But, the subtle point that are intended here are as follows. The energy conservations in Eqns.3 and 4 are preserved only because of the π phase shift (Eq.1) introduced by the “external reflection”. The 100% energy re-direction of Eq.4 is also a byproduct of this π phase shift. But the most important point is that since light does not interfere with light, the role of the dipoles of the beam splitter (albeit passive) is physically undeniable. If we block off, say, the Beam 2 [Fig.1(b)], then the energy of the Beam1 will be simply split according to traditional beam splitter rule into right-going $d_1^2 T$ and down-going $d_1^2 R$ and the corresponding re-radiated amplitudes will be $d_1 \sqrt{T}$ and $d_1 \sqrt{R}$. If we now suddenly and very briefly open the shutter from the Beam2 again under the conditions of Eq.4, we will be able to turn-off the right-going beam briefly while all the energy will be sent along the down-going beam. It is obvious that the dipoles of the passive beam splitter collectively participate in re-directing and re-radiating the field energy based on the phase conditions and the Poynting vectors of the incident beams. Classical physics has been fully aware of this through the Maxwell’s equations [12]. The simultaneous presence of real EM field from both the directions on the beam splitter boundary is essential to generate the energy re-direction capability of a passive dielectric boundary (for the “interference” effects to become manifest). That even (non-absorbing) passive material dipoles actually dictates re-radiation of EM field is well known in classical physics from the explanation for the Brewster angle. At this angle ($\tan \theta_B = n$), the reflection of a beam becomes zero when the incident state of polarization is parallel to the plane of incidence because the E-vector induced undulation inside the medium (refracted direction) becomes parallel to the direction of reflection and dipoles cannot radiate along its axis of undulation [11].

When the Poynting vectors of the transmitted and reflected beam pairs are not co-linear, the beam splitter will simply produce two pairs of four independent beams with a finite tilt angle between them as shown in Fig.1c. As before, the strengths of the beams will be dictated by the same amplitude transmittance T and reflectance R of the beam splitter, but the beams’ energy transmittance and reflectance will not be modified by this dielectric boundary

[11]. However, detector arrays in both directions will register spatially distributed fringes as long as the tilted but finite size beams remain physically overlapped. The down-going beam pair will have amplitudes $(d_1\sqrt{R}, d_2\sqrt{T})$ and the right going pairs will be $(d_1\sqrt{T}, d_2\sqrt{R})$. The sum total energy of the four beams, of course, is conserved as $(d_1^2 + d_2^2)$ since $T + R = 1$.

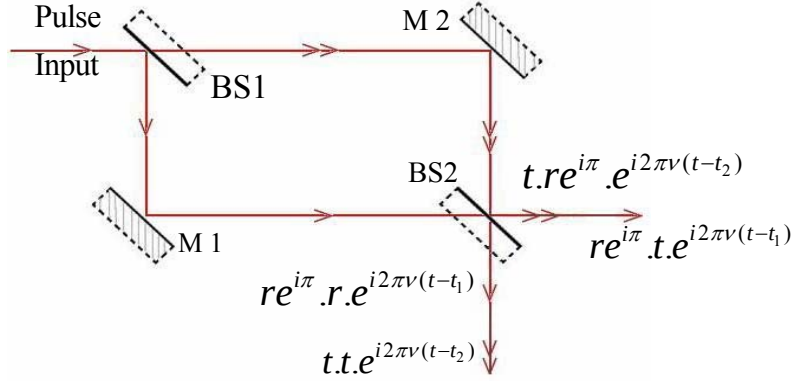


Figure2. A Mach-Zehnder interferometer (MZ) consists of two mirrors (M1 & M2) and two beam splitters (BS1 & BS2). The beam splitters introduce π phase shift at all “external” reflections. This is critical to energy conservation in re-directing energy from one beam to the other by the beam splitter’s boundary molecules.

2.2. Mach-Zehnder Interferometer

Let us now introduce an identical pair of such beam splitters in the popular Mach-Zehnder interferometer (MZ) as shown in Fig.2. We will now consider an input pulse $\vec{d}(t)\exp[i2\pi\nu t]$ of width δt . The down-going and the right-going amplitudes are given by:

$$\begin{aligned}\vec{d}_{\text{down}}(t, \nu, \tau) &= \vec{d}(t-t_1) \cdot re^{i\pi} \cdot re^{i2\pi\nu(t-t_1)} + \vec{d}(t-t_2) t^2 e^{i2\pi\nu(t-t_2)} \\ &= e^{i2\pi\nu(t-t_1)+i\pi} [R\vec{d}(t-t_1) - T\vec{d}(t-t_2)e^{i2\pi\nu\tau}], \\ \vec{d}_{\text{right}}(t, \nu, \tau) &= e^{i2\pi\nu(t-t_1)+i\pi} tr[\vec{d}(t-t_1) + \vec{d}(t-t_2)e^{i2\pi\nu\tau}]; \text{ where, } \tau \equiv (t_1 - t_2)\end{aligned}\quad (5)$$

Again, the re-radiated field amplitudes are denoted by d to underscore the role of the collective dipole undulations of the beam splitter. Then the re-radiated time varying intensities are:

$$\begin{aligned}I_{\text{down}}(t, \nu, \tau) &= |\vec{d}_{\text{down}}|^2 = [R^2 d^2(t-t_1) + T^2 d^2(t-t_2) - 2RT\vec{d}(t-t_1) \cdot \vec{d}(t-t_2) \cos 2\pi\nu\tau] \\ I_{\text{right}}(t, \nu, \tau) &= |\vec{d}_{\text{right}}|^2 = TR[d^2(t-t_1) + d^2(t-t_2) + 2\vec{d}(t-t_1) \cdot \vec{d}(t-t_2) \cos 2\pi\nu\tau]\end{aligned}\quad (6)$$

If we now use an energy integrating detector, the total recorded energy in the two directions will be:

$$\begin{aligned}E_{\text{down}}(\nu, \tau) &= \int_0^{T>2\delta t} I_{\text{down}} dt = E_{\text{pls}}[T^2 + R^2 - \gamma_{\text{BS}}(\tau)2TR \cos 2\pi\nu\tau] \\ E_{\text{right}}(\nu, \tau) &= \int_0^{T>2\delta t} I_{\text{right}} dt = 2TRE_{\text{pls}}[1 + \gamma_{\text{BS}}(\tau) \cos 2\pi\nu\tau]\end{aligned}\quad (7)$$

Since the dipole undulation follows the same frequency (some times with phase shifts), the normalized mathematical structure of the autocorrelation function of the EM fields within the material medium $\gamma_{\text{BS}}(\tau)$ and in free space $\gamma_{\text{field}}(\tau)$ are identical, as indicated below:

$$\begin{aligned}\gamma_{\text{BS}}(\tau) &= \int_0^{T>2\delta t} \vec{d}(t-t_1) \cdot \vec{d}(t-t_2) dt / \int_0^{T>2\delta t} d^2(t) dt \equiv \gamma_{\text{field}}(\tau) \\ E_{\text{pls}} &= \int_0^{T>2\delta t} a^2(t) dt\end{aligned}\quad (8)$$

One can demonstrate the conservation of energy for all values of τ , using $(T + R) = 1$:

$$E_{\text{total}} = E_{\text{down}}(\nu, \tau) + E_{\text{right}}(\nu, \tau) = E_{\text{pls}}(T + R)^2 = E_{\text{pls}}\quad (9)$$

Again, the key point to appreciate is that the E-fields must be present simultaneously on the beam splitter boundary from both the sides of the out put beam splitter of MZ for the superposition effect to physically materialize. We have also added a subscript *BS* to the normalized autocorrelation function $\gamma_{BS}(\tau)$ to underscore that it represents the behavior of the material undulations of the beam splitter dipoles induced by the EM fields rather than correlation of the pure free-space EM fields by themselves. This definition of $\gamma_{BS}(\tau)$ automatically assures that orthogonally polarized light cannot produce superposition effects because their dot product makes the cross-term (superposition term) of Eqns.6 & 7 become zero. The physical implication is that the same set of material dipoles cannot undulate in two orthogonal directions at the same time. We do not need to use an in-congruent statement like “orthogonally polarized light do not interfere” because light does not interfere anyway. Note also that if we set $\tau > 2\delta t$, then $\gamma_{BS}(\tau) = 0$ and again cross-term drop; but this time because the overlap product of $\vec{d}(t-t_1) \cdot \vec{d}(t-t_2)$ is zero for lack of overlap. While this situation can be “explained away” by saying that the “coherence length” of the pulse is shorter than the relative path delay, a more direct explanation is that the “superposition effect” is non-existent simply because physically the two pulses were not simultaneously acting on the molecules of the beam splitter boundary from the two sides; there were no real physical superposition of the two pulses. When $\delta t \gg \tau$, or when $\tau \rightarrow 0$, $\gamma_{BS}(\tau) \rightarrow 1$, then it is possible to re-direct most of the energies of the two pulses coming from the two directions into one, which can be derived from Eq.7 by setting a value for τ such that $\nu\tau = n$, or $[n + \frac{1}{2}]$, n being an integer. Let us choose the case of $\nu\tau = [n + \frac{1}{2}]$:

$$\begin{aligned} E_{down}(\nu, \tau) &= 2R^2 E_{pls} [T + R]^2 = E_{pls} \\ E_{right}(\nu, \tau) &= 2R^2 E_{pls} [1 - 1.1] = 0 \end{aligned} \quad (10)$$

Proper co-linearity of the Poynting vectors is essential for the Eqns. 10 to hold as discussed in section 2.1. When the angle between the Poynting vectors is non-zero, one generates non-interfering but superposed beams that will generate spatial fringes if only a detector array is placed after the beam splitter.

We will now consider the case of a low photo electron counting situation by severely reducing the input intensity under the condition of co-linear Poynting vectors. Let us assume that (i) the input pulse is a rectangle of width δt , (ii) the relative delay between the two MZ arms is $\tau \cong \delta t / 2$ giving us $\gamma_{BS}(\tau) = 0.5$ such that $\nu\tau = n$, an integer implying maximum energy re-direction by the beam splitter and (iii) $T = R = 0.5$. Under these specific assumptions the down-going and right-going energies can be derived from Eq.7:

$$E_{down}(\tau = \delta t / 2) = 0.25 E_{pls}; \quad E_{right}(\tau = \delta t / 2) = 0.75 E_{pls} \quad (11)$$

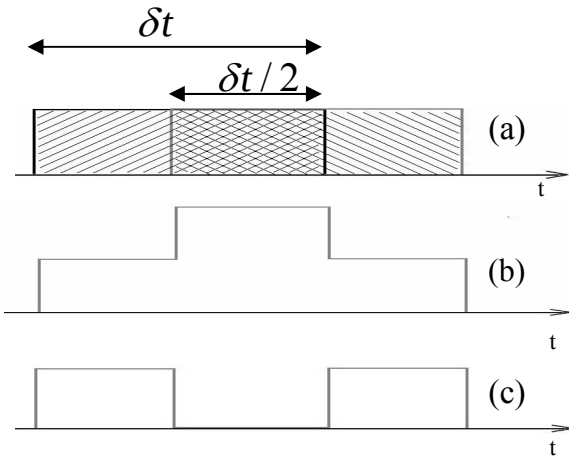


Figure 3. Two different types of out put pulses from an MZ due to a single incident rectangular pulse. (a) shows the physical superposition of the two pulses on BS2 of Fig.2 from two opposite sides. (b) shows output at in-phase condition with more energy in the middle of the stretched pulse. (c) shows the output for out-of-phase condition and it appears as two separate pulses with a “null” in the middle. If photons are indivisible packets of energy and the input pulse contained only one photon, then we loose all these elegant time evolving superposition effects that are predicted by semi-classical model and can be validated when the pulse contains a very large number of photons.

Now consider the case where the total energy of the input pulse has been reduced to a value that is equivalent to a total of n -photo electron emission. Then the integrating detector on the right output of the MZ should register a total of $0.75n$ counts and that on the downward beam will register $0.25n$ counts, as per Eq.11 assuming a lossless MZ system. But if we use a pair of fast pico second detectors to be able to follow the temporal evolution of the transmitted nano second pulses generated by the MZ system [Eq.6], we should be able to measure the following

interesting effects [Fig.3]. (i) The output pulse duration has been stretched from δt by the MZ to $(1.5\delta t)$. (ii) The RHS detector receives photons during the entire period of $(1.5\delta t)$, but its count rate surges up during the central $(\delta t/2)$ period, while (iii) the detector on downward output will experience zero photo count during the same central $(\delta t/2)$ period. Classical optics does not have any problem explaining any of these variations in the rate of photo electron counts. Let us now reduce the energy of the input pulse equivalent to only one single photon $h\nu$ with the help of some rigorous detection system placed before the input to the MZ. What would we observe under this situation? Combining the prevailing notion that photons are indivisible energy packets (“it goes only one way”) with Panarella’s experimental conclusions [6] that a minimum of four photons are required to create superposition effects (in our case, on the MZ output beam splitter), we should not be able register any of the classical prediction of time dependent intensity variations given by Eq.6 and shown in Fig.3. If we now increase the energy of the input pulse equivalent to, say, 5 photo counts, will we be able to validate the photo count distribution predicted by Fig.3? We, of course, believe that photons are divisible and diffractable wave packets as they propagate and evolve. They can collectively share their energy to an atom or a molecule when superposed on them and provide the required total amount of energy for the quantum transition $\Delta E = h\nu$ as long as all the wave packets have the same carrier frequency ν [5]. Based on Panarella’s observation [6], we want to underscore that a simple slow down of the rate of photo electron emission to n-electrons per second should not be interpreted as the conclusive evidence that the actual EM field constitutes the propagation of exactly $(n \cdot h\nu)$ photons per second.

2.3. Michelson’s Fourier transform spectrometer

Michelson’s Fourier transform spectrometer (FTS) very elegantly creates delayed beams with the help of a single beam splitter and a pair of mirrors, one of which can be continuously moved to introduce variable delay τ between the two interfering beams [13]. As for the case of MZ, FTS can also be set to produce spatial fringe mode with appropriate tilt in the mirrors by making a small angle between the Poynting vectors of the two superposed beams. This way the total energy is split between returning and outgoing beams equally. The classic “etendue-advantage” (light gathering power) [14] is derived when the Poynting vectors are perfectly co-linear as explained earlier. The interferometer sketch is shown in Fig.4.

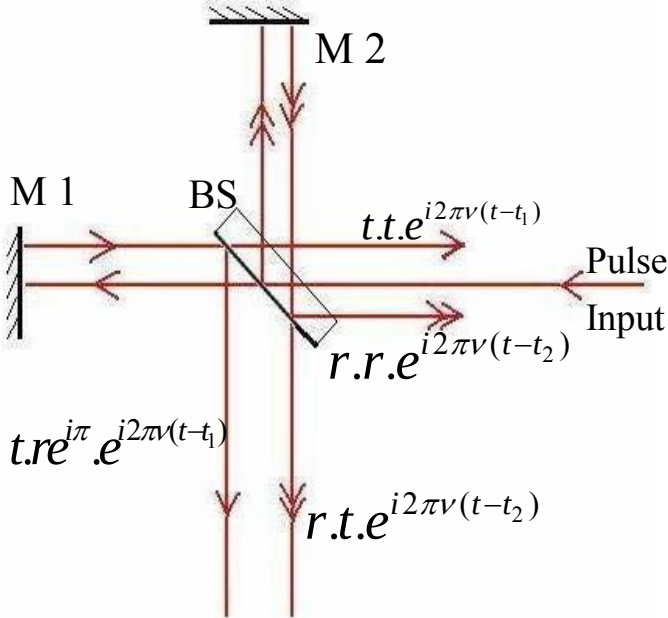


Figure 4. The role of a dielectric beam splitter (BS) in a Michelson interferometer, which replicates one incident pulse into two with a controllable delay with the aid of two more mirrors, M1 & M2. The phase shift π for “external” reflection is critical to conserve energy when two beams are superposed from the opposite directions. The role of the molecules of the dielectric boundary is to re-direct energy from one beam to the other when the beam Poynting vectors are collinear.

The downward beams are the useful beams for this interferometer and the right-going beam is a system induced waste. As before, the down-going and the right-going re-radiated amplitudes by the beam splitter are given by:

$$\begin{aligned}
d_{\text{down}}(t, \nu, \tau) &= d(t-t_1).t.re^{i\pi}.e^{i2\pi\nu(t-t_1)} + d(t-t_2).r.t.e^{i2\pi\nu(t-t_2)} \\
&= tre^{i2\pi\nu(t-t_1)+i\pi}[d(t-t_1)-d(t-t_2)e^{i2\pi\nu\tau}] \\
d_{\text{right}}(t, \nu, \tau) &= d(t-t_1).t.t.e^{i2\pi\nu(t-t_1)} + d(t-t_2).r.r.e^{i2\pi\nu(t-t_2)} \\
&= e^{i2\pi\nu(t-t_1)}[Td(t-t_1)+Rd(t-t_2)e^{i2\pi\nu\tau}]; \text{ where } \tau \equiv (t_1-t_2).
\end{aligned} \tag{12}$$

Then the re-radiated time varying intensities are:

$$\begin{aligned}
I_{\text{down}}(t, \nu, \tau) &= |d_{\text{down}}|^2 = TR[d^2(t-t_1) + d^2(t-t_2) - 2d(t-t_1)d(t-t_2)\cos 2\pi\nu\tau] \\
I_{\text{right}}(t, \nu, \tau) &= |d_{\text{right}}|^2 = T^2d^2(t-t_1) + R^2d^2(t-t_2) + 2TRd(t-t_1)d(t-t_2)\cos 2\pi\nu\tau
\end{aligned} \tag{13}$$

As before, energy integrating detectors will register the following total energies in the two directions:

$$\begin{aligned}
E_{\text{down}}(\nu, \tau) &= \int_0^{T>2\delta t} I_{\text{down}} dt = 2TRE_{\text{pls}}[1 - \gamma_{\text{BS}}(\tau)\cos 2\pi\nu\tau] \\
E_{\text{right}}(\nu, \tau) &= \int_0^{T>2\delta t} I_{\text{right}} dt = E_{\text{pls}}[T^2 + R^2 + 2TR\gamma_{\text{BS}}(\tau)\cos 2\pi\nu\tau]
\end{aligned} \tag{14}$$

As before, the pulse autocorrelation $\gamma_{\text{BS}}(\tau)$, as experienced by the beam splitter medium, and the input pulse energy E_{pls} , are given by the pair of Eq.8. The conservation of energy is preserved since $T + R = 1$:

$$E_{\text{total}} = E_{\text{down}} + E_{\text{right}} = E_{\text{pls}}[T^2 + R^2 + 2TR] = E_{\text{pls}}[T + R]^2 = E_{\text{pls}} \tag{15}$$

Let us now discuss the situation when the Poynting vectors are not co-linear for the downward and the right-going beam pairs. The emergent plane parallel beams are now at an angle with each other as depicted in Fig.1c. Since the beam splitter is no longer introducing variable energy re-direction with τ (we are not moving one of the mirrors parallel to itself), the spatially distributed “interference” fringes will become visible only when we insert a detector array within the volume of the superposed outgoing beams. We are now considering only the downward pair of beams. The variation of the sum of the dipolar amplitudes as would be experienced by a detector array along the spatial X-axis is still given by an equation similar to that in Eq.11, but we should acknowledge that the dipole stimulation of the detector $\vec{d}(t)$ that precedes energy absorption is given by the linear polarizability relation $\vec{d}(t) = \chi_1 \vec{a}(t)$:

$$\vec{d}_x(t, \nu, \tau) = tre^{i2\pi\nu(t-t_1)}[\vec{d}(t-t_1) - \vec{d}(t-t_2)e^{i2\pi\nu\tau}] \tag{16}$$

$$\tau = \Delta / c = (x \tan \theta) / c; \quad \theta \text{ small.} \tag{17}$$

The suffix x has been added to $\vec{d}$ to underscore that we are now dealing with spatial fringes as registered by the detector array. The time varying intensity at different spatial points is now given by:

$$I_x(t, \nu, \tau) = |a_x|^2 = TR[d^2(t-t_1) + d^2(t-t_2) - 2\vec{d}(t-t_1) \cdot \vec{d}(t-t_2)\cos 2\pi\nu\tau] \tag{18}$$

The time integrated fringe pattern and the detector autocorrelation functions are:

$$F_x(\nu, \tau) = \int_0^{T>2\delta t} I_{\text{down}} dt = 2TRE_{\text{pls}}[1 - \gamma_{\text{det}}(\tau)\cos 2\pi\nu\tau] \tag{19}$$

$$\gamma_{\text{det}}(\tau) = \int_0^{T>2\delta t} \chi_1 \vec{a}(t-t_1) \cdot \chi_1 \vec{a}(t-t_2) dt / \int_0^{T>2\delta t} \chi_1^2 a^2(t) dt \equiv \gamma_{\text{field}}(\tau) \tag{20}$$

Notice again the identical mathematical expression for $\gamma_{\text{det}}(\tau)$ and $\gamma_{\text{field}}(\tau)$. Michelson’s classic visibility function V is the autocorrelation of the pulse, derived from Eq.18:

$$V(\nu, \tau) = (E_{x,\text{max}} - E_{x,\text{min}}) / (E_{x,\text{max}} + E_{x,\text{min}}) = \gamma_{\text{det}}(\tau) \tag{21}$$

We have kept χ_1 explicitly in Eq.20 to underscore the physically different processes that are giving rise to the time integrated “interference” terms under different conditions of co-linear and non-co-linear Poynting vectors for the superposed beams. The detectors can absorb energy whether the Poynting vectors of the superposed beams are at an

angle or co-linear, but the passive beam splitter cannot redirect energy unless the Poynting vectors are co-linear. We have neglected the effects of higher order polarizability $\chi_{(n)} a^n(t)$ assuming that they are very weak, which may not be true for all possible detecting molecules.

For spectrometric analysis, it is customary to extract the oscillatory part of the fringe data from the time integrated fringe pattern of Eq.19 [13, 14, 11c]:

$$F_{osc}(\nu, \tau) = C \gamma_{det}(\tau) \cos 2\pi\nu\tau; \quad C \equiv 2TRE_{pls} \quad (22)$$

If $\tau_{max} = \Delta_{max} / c > 2\delta t$ [see Fig.1c?] then the interferogram would record complete fringe information required by Eq.20 [see Fig.5]. The carrier frequency ν of this ideal single pulse can be determined by measuring τ_{max} and counting the total number of fringes $\nu\tau_{max} = N$.

Let us now consider the case of light coming from an atomic discharge lamp, say Cd-red line, as it is spectrally “single line” [13]. Let us assume that all the atoms produce identical and approximately same shaped pulses $\bar{a}(t) \exp[i2\pi\nu t]$ but their carrier frequencies are broadened by the Doppler shift due to Maxwellian velocity

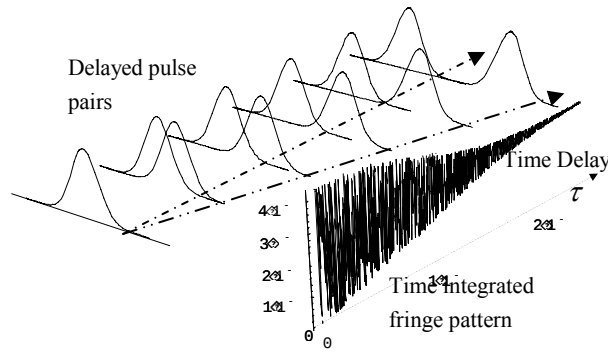


Figure 5. A pictorial explanation of the degrading fringe visibility due to overlap of unequal amplitudes of the superposed pulses [15d]. One should not interpret this degradation as due the presence of different optical frequencies given by the square modulus of the Fourier transform of the incident pulse amplitude. The amplitude envelope of the pulse is shown to be a Gaussian. The dense “hash” represents the rapid spatial variation of the fringes represented by Eq.18, and the envelope represents the autocorrelation function of Eq.20 registered by the detector array.

distribution. We assume that $D(\nu)$ is the spectral intensity distribution function due to this Doppler broadening. Let us now ask the question: Is it possible to determine the shape of the wave packet $a(t)$, the photons that atoms emit, by classical formulation of spectroscopy? Since our key premise is non-interference of light beams, all these atomic pulse will create their own intensity fringes without mutual “interference”. Mathematically speaking, there are no superposition cross-terms between the different frequencies. The resultant intensity fringe function can be derived by integrating Eq.18 for all $D(\nu)$. Alternately, the corresponding oscillatory component of the time integrated fringes can be derived by integration Eq.21 for all $D(\nu)$:

$$F_{osc, discharge}(\nu, \tau) = C \gamma_{det}(\tau) \int_{\nu_{min}}^{\nu_{max}} D(\nu) \cos 2\pi\nu\tau \, d\nu = C \gamma_{det}(\tau) \tilde{D}(\tau) \quad (23)$$

Now we can explore the possibility of determining the shape of the wave packets emitted by atoms by spontaneous emission in a discharge lamp:

$$\left[\int_0^{T>2\delta t} a(t-t_1)a(t-t_2)dt \right] / \left[\int_0^{T>2\delta t} a^2(t)dt \right] \equiv \gamma_{det}(\tau) = \left[F_{osc, discharge}(\nu, \tau) / C \tilde{D}(\tau) \right] \quad (24)$$

Eq.23 clearly opens up the possibility of determining $a(t)$, the envelope of a spontaneously emitted photon, since $F_{osc, discharge}(\nu, \tau)$ can be extracted from the recorded fringes, C represents known parameters of the interferometer and $\tilde{D}(\tau)$ can be analytically computed since $D(\nu)$ has been derived analytically and also validated by various other laboratory methods. To our knowledge, the above prescription for deriving the shape of the wave packet for spontaneously emitted photon is new.

It is worth noting that when the incident beam contains more than one frequency and the optical set up is arranged for making the Poynting vectors co-linear and co-directional for each pair of the beams generated by the beam splitter, it re-directs the energy corresponding to each frequency separately. There is no cross-talk or superposition terms between different frequencies. Otherwise, Fourier transform spectroscopy would have never worked. In fact, Michelson’s FTS is the key experimental support behind the concept of “non-interference” or

“incoherence” of different frequencies. But, Forrester et al [16] demonstrated that beat frequencies due to the superposition of light beams containing different frequencies are detectable with fast electronics and photo conductive detectors. This paradox is resolved with our hypothesis that superposition effects are always displayed by material dipoles whenever their intrinsic quantum properties allow them to simultaneously respond to all the superposed fields. Light beams by themselves never interact with each other. In other words, “non-interference” of light is the general behavior of light under all conditions. It is easy to appreciate that dipoles of photo conductive detectors with broad valance and conduction bands, can simultaneously respond to multiple frequencies allowed by the band-gap and hence undergo time varying excitation and the rate of transfer of electrons to the conduction band becomes oscillatory, which is the beat frequency [17]. Then we are also forced to conclude that collectively the dipoles of a dielectric boundary are capable of re-directing the energies of beams that are incident co-linearly from the opposite directions, but only for the same frequency E-fields. The stimulations due to different frequencies propagate un-perturbed by each other just as the beams due to same frequency but whose Poynting vectors are non-co-linear.

3. MULTIPLE-BEAM SUPERPOSITIONS

Gratings and Fabry-Perot interferometers are the two generic examples for which one can make analysis of multiple-beam superposition by using our method of real time propagation of a pulses replicated by these instruments [15]. We would like to briefly summarize the case of an N-slit grating and express its “spectral fringe” formation in terms of summation of a number of two-beam autocorrelation functions, which represent all possible combination of two-beam superposition out of all the N-beams produced by the N-slit grating. Fig.6 gives the simple physical picture for grating as a pulse replicator – we see N pulses delayed by the periodic step delay τ , where $\tau_0 \equiv N\tau$ is defined as the spectrometer time constant [15a]. When the appropriate optics helps them get superposed on a detector array, we get the traditional grating fringes.

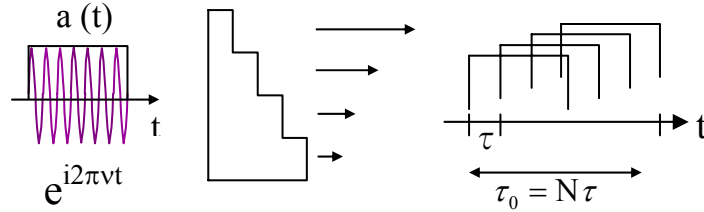


Figure 6. A grating replicates an incident pulse into N new equal amplitude pulses with a periodic delay τ . When this train of pulses are superposed on a detector array, one can register the grating fringes. The out put pulse from all multiple beam spectrometers are stretched and have a characteristic time constant $N\tau$.

The train of out put amplitudes is denoted by:

$$\vec{d}_{out}(t, \nu) = (1/N) \sum_{n=0}^{N-1} \vec{d}(t - n\tau) \cdot e^{i2\pi\nu(t - n\tau)} \quad (25)$$

The time varying intensity represents a stretched pulse by approximately an amount of $\tau_0 \equiv N\tau$:

$$I'_{pls}(t, \nu, \tau) = |\vec{d}_{out}(t)|^2 = \sum_{n=m=0}^{N-1} \frac{1}{N^2} \cdot a^2(t - n\tau) + 2 \sum_{n \neq m} \frac{1}{N^2} \cdot \vec{d}(t - n\tau) \cdot \vec{d}(t - m\tau) \cdot \cos[2\pi(n - m)\nu\tau] \quad (26)$$

Then the time integrated fringe intensity is given by:

$$I_{pls}(\nu, \tau) = \frac{1}{N} + \frac{2}{N^2} \sum_{p=1}^{N-1} (N - p) \gamma_{det}(p\tau) \cos[2\pi p\nu\tau] \quad (27)$$

Where the auto correlation function is defined as:

$$\gamma_{det}(p\tau) \equiv \gamma_{nm}(\tau) \equiv \gamma(|n - m|\tau) = \frac{\int \vec{d}(t - n\tau) \cdot \vec{d}(t - m\tau) dt}{\int d^2(t) dt} \quad (28)$$

If the incident pulse is longer than $\tau_0 \equiv N\tau$, then

$$\lim_{\delta t \rightarrow \tau_0 = N\tau} Lt. [\gamma_{det}(p\tau)] \rightarrow 1 \quad (29)$$

and we can recover the traditional CW grating formula without even the need for real CW light:

$$\lim_{\delta t \rightarrow \tau_0 = N\tau} [I_{pls}(\nu, \tau)] = I_{cw}(\nu, \tau) = \frac{1}{N} + \frac{2}{N^2} \sum_{p=1}^{N-1} (N-p) \cos[2\pi p\nu\tau] \equiv \frac{1}{N^2} \frac{\sin^2 \pi N\nu\tau}{\sin^2 \pi \nu\tau} \quad (30)$$

The strength of our approach is that we have started with a generalized pulse and derived the expression for the grating spectral fringes [Eq.26] in terms of a series of two-beam autocorrelation functions as would be experienced by a detector array. Further, in the limit of long pulse length, $\delta t \gg R\lambda / c = mN\lambda / c = N\tau \equiv \tau_0$, where R is the classical resolving power of the grating, we recover the traditional CW formulation as a special case. The details of the evolution of these concepts and derivations for both gratings and a Fabry-Perots can be found in ref. 15.

4. CONCLUSIONS

We have underscored the need to re-visit the phenomenon of “interference” in view of “non-interference of light” in general, and not just for orthogonal polarizations and different frequencies. Since superposition effects are displayed by the summing capability of material dipoles of the induced dipole undulations by the multiple superposed fields on them, we have developed the generic expressions for two-beam and multiple-beam superposition effects in terms of autocorrelation functions for the relevant material dipole under the illumination by a generic light pulse. The (i) time-domain and (ii) material-dipole driven approach forces us to consider the roles of different physical properties of different materials in studying superposition effects, which have been missing in the traditional coherence theory. In the process we have found that the current model of interference by single photon with its indivisible packet of energy does not conform to the required reality of the simultaneous superposition of multiple signals on the same detecting molecule. We have also found that the acceptance of the semi-classical model of atomic emissions as classical wave packets allows one to extract the shape of the envelope of this wave packet by inverting the corresponding autocorrelation function which is easily measurable by a two-beam interferometer. Our time-domain approach for a grating spectrometer reveals that we have obtained a very generalized expression for the spectral fringes due to a short pulse that naturally evolves into the traditional CW expression when the pulse length simply exceeds the spectrometer time constant, a property that classical spectrometry have not yet formally recognized.

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