The Physical Reality of Applied Quantum Optics

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Physics versus Mathematics

By

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CHAPTER ONE

INTRODUCTION

Over the past 60 years, the field of Quantum Optics has seen technological developments in various attempts to develop sources emitting a single-photon per radiation mode, as well as corresponding photodetectors for resolving very low numbers of photons (e.g., Hepp *et al*. [1]; Lodahl *et al*. [2]).

Quantum Optics (e.g., Garrison and Chiao [3]) is predicated on the superposition and entanglement of apparently single-photon number states. The processing of such quantum states by means of beam splitters is supposed to enhance the performance and sensitivity of various information assessing operations. The theoretical model is based on the mixed quantum states of an ensemble of measurements, which provides time-independent distributions for the numerical evaluation of probability amplitudes associated with alternative propagation pathways of one single-photon.

The detection and measurement of correlated photons and their degrees of freedom constitute the experimental demonstration of concepts in Quantum Optics. The three stages of generation, propagation, and detection impact, in various ways, on the properties of measured outcomes.

An early application of Quantum Optics occurred in the experiments intended to prove quantum nonlocality by generating a pair of polarised, highly correlated (or entangled) photons, with an optically nonlinear crystal (e.g., Garrison and Chiao [3]). One photon is sent in one direction and the other in the opposite direction. Remote and independent measurements would appear to be correlated, leading to the concept of quantum nonlocality. Surprisingly, though, the same correlation function can be derived without entanglement, by using single and independent quantum polarisation states, or qubits (Vatarescu [4-5]).

In an article published in August 2014 by Tipler [6] and titled "The quantum nonlocality does not exist", the author detailed physical arguments for the statistical nature of the experimental results contradicting the concept of quantum nonlocality. While initiating the analysis with ensemble entangled states, Tipler points out that quantum wave functions need not collapse into a specific state upon measurement. Rather, the overall wave function continues to evolve and branches out into one of four possible options, thereby generating an ensemble distribution of measured values.

A recent article in the Physical Review A authored by R. B. Griffiths [7], appears to be the first editorial exception to the uncompromising protection of the concept of quantum nonlocality. Additional analyses disproving the physical validity of the concept of quantum nonlocality have been published recently in other journals, e.g. (Boughn [8]; Khrennikov [9]; Kupczynski [10]).

It is noteworthy that a large body of analytic rebuttals of the concept of quantum nonlocality has been continually ignored in tens of articles which are published every year on this subject. These continual omissions, in the legacy journals of professional literature, of challenging and physically meaningful interpretations of the experimental results can only be an indication of the resistance organised by vested interests. As yet, not one single article has reported any evidence – at the level of pure quantum states of a single measurement – of a correlated or entangled collapse of the quantum wave function at one location as a result of a measurement carried out at another remote location. The global, mixed quantum states which are claimed to generate those quantum correlations are distributions of ensemble measurement possessing no dependence of the time and location of the measured observables. Despite the physical impossibility of a photon to maintain its polarisation state or even survive propagation through a dielectric medium because of the quantum Rayleigh scattering (Louisell [11]; Marcuse [12]), the concept of quantum nonlocality is still believed to provide an untapped resource for some future applications.

Additionally, in 2015, experimental results presented by Qian *et al.* [13] in the classical regime of large numbers of photons, found strong correlations, known as concurrences, between variables of polarised light. These results broadened the concept of correlation between observable values of quantum variables, raising questions about the quantum signature of such correlations.

These analytic and experimental results prompted this author to scrutinise the physical processes and interactions involved in the experimental setups, and which have continually been ignored and overlooked in major professional journals. Although detrimental to scientific development, the editorial policy of many journals would aim to preserve the status quo of physical understanding.

Another application of Quantum Optics has to do with the combination of a single-photon source, a beam splitter and simultaneous detections with two separate photodetectors. This combination is commonly described by means of a time-independent mixed state of correlation measurements (Garrison and Chiao [3]; Mandel [14]).

However, a few questions have been overlooked: 1) How can a single-single propagate in a straight line in a homogenous dielectric medium given the multitude of electric dipoles it encounters? 2) How can a probability amplitude, instead of an optical field associated with a photon throughout its propagation, trigger or activate a photodetector? 3) How can a single-photon state of an ensemble of measurements, at a given time, interfere with another state that is physically absent at the same time? 4) How can the physical duration of a monochromatic photon be described mathematically by a polychromatic Fourier wave packet of the ensemble of measurements, given that only one Fourier spectral line would be measured at any given time?

Technological advances in material fabrication for integrated photonic devices and circuits will need to be matched by improved physical understandings of light-matter interactions. These will facilitate the design and operations of functional devices such as phase-sensitive amplifications of photons, sub-Poissonian sources of photons, low-power phase-sensitive switches and modulators, etc. and will have the potential to open up new applications in optical communication relying on easy to control quantum interactions.

The tenets of Quantum Optics consist of single photons propagating in a straight line in a dielectric medium, the interference patterns of probability amplitudes based on ensemble-evaluated mixed states of photons, and the use of beam splitters as entangling devices. Nevertheless, any probability amplitude of a quantum event should be evaluated from wave functions that reflect the physical reality. Discarding temporal information – which becomes a lack of information – about the propagation pathway of a single photon does not create a physical effect; it can only mask or obscure the existence of physical interactions.

This book adopts a physical approach. As photons propagate through a dielectric medium, the quantum Rayleigh spontaneous emission replaces entangled photons with independent ones in homogeneous dielectric media where single photons cannot propagate in a straight line. Pure quantum states of wavefronts of independent groups of photons deliver the intrinsic field profile associated with a photonic wavefront and the correct expectation values for its number of photons, its complex optical field, and phase quadratures. The spatial distribution of a photon, both longitudinally and laterally, is found from the quantised Maxwell equations in the context of a Wigner-type monochromatic time-varying spectral component. These photonic properties enable a direct analysis of various beam splitters and interferometric filters, leading to generalized expressions for the correlation functions characterizing counting of coincident numbers of photons for the fourth-order field interference.

1.1 A Historical Perspective

The need to detect and analyse very week optical signals from distant stars led in the 1950s to the method of intensity interferometry in a bid to overcome the sensitivities associated with optical fields' interference. The Hanbury Brown and Twiss experiment [15] of 1956 with a mercury arc lamp as the optical source of radiation, measured "correlation between photons in coherent light rays". In that experiment, the light was produced by many different atoms, and one spontaneously emitted photon would have been slightly amplified on its way out. Thus, some photons might arrive in pairs at the half-silvered mirror (or beam splitter), which may explain the results showing a correlation between pairs of photon counts in terms of statistical distributions of bosons (Purcell [16]).

In 1961, Fano developed a theory of two-photon interference involving two emitting atoms and two photo-detecting atoms [17]. Employing generic transition or interaction matrix elements, the theory leads to "… a cosine function of both space and time" of the joint one-photon absorption by each of the two detectors.

The next stage of the theoretical development in 1963, saw Glauber specify the probability amplitude of photon detection in terms of the annihilation operator and input quantum states [18-19]. Given that the photon creation operator and the annihilation operator are the adjoint of each other, the detection probability became identical to the expectation value of the number operator. This enabled the use of number states defined as the eigenstates of the free-space Hamiltonian operator of optical fields. An equivalence with classical optical fields was derived in the form of an ensemble eigenstate of the annihilation operator, and which became known as the coherent states of light. While the removal of one photon may not have any consequence for a very large number of photons as is the case in the classical regime, in the quantum regime of a few photons per radiation mode, ironically, the coherent state is impractical; additionally, the coherent states fail to deliver the eigenvalues of phase quadratures (Carruthers and Nieto [20]).

The Glauber theory of photon detection and correlation (or coherence) is based on ensemble distributions of photons However, the corresponding quantum states lack information about time-dependent, instantaneous measurements or interactions which require the use of pure quantum states that are allowed to be time-dependent. This shortcoming of quantum evaluations was pointed out by Mandel and Wolf in 1965 [21, Section 7.3] where one finds the following statement:

> "We have already shown in Sec. 3.2 that a description of ordinary interference effects may readily be given in terms of the quantized field, and that it follows the classical treatment fairly closely. It might therefore be thought that the transient superposition effects discussed in the last two sections can also be described quantum mechanically, in a closely parallel manner. However, here we come up against the basic feature that quantum mechanics is always concerned with expectation values of observables, whereas the calculation of expectation values was deliberately avoided in the simple treatment leading to Eqs. (7.6) and (7.10)."

This shortcoming of the quantum approach was ignored in the developments of following decades. This book aims to rectify this deficiency, prompted by the very fact that experimental results are measured one value at a time, and the final resultant distribution of the ensemble of measurements is time-independent once the experiment is complete.

Equally, the Glauber theory would have properties of instantaneous photons determined by the ensemble distribution to which they belong. For the number states, the photons would have no optical field, while for the coherent states the optical field arises from the overall superposition of a very large, if not infinite, number of photons. Yet, photons interact with dipoles and are detected at a given location and a particular time.

Another branch of Quantum Optics was initiated in 1965 by Jaynes and Cummings with an article analysing the interactions between single photons and atoms placed in resonant cavities [22]. In this case the photonic state is time-dependent, consisting of two consecutive number states.

Over the next two decades, 1965 to 1985, laser sources were used to prove interference between independent radiation modes of various numbers of photons. A review of possible single-photon interference patterns was presented by Walls in 1977, and space and time overlaps were included [23]. As the photon is the lowest amount of indivisible energy carried by an electromagnetic field, it can only be detected at one of two photodetectors located in the two alternative pathways. By contrast, the ensemble distribution of measurements would attach a non-zero value to both possibilities of the transition matrix evaluated with a pathway-entangled state of one photon. It is this discrepancy between the single measurement and the overall distributions that gives rise to a mathematical single-photon interference pattern in the context of a large number of measurements. The corresponding experimental results seemingly supporting quantum interference of probability amplitudes, are interpreted on the basis of there being only one photon emitted by the optical source for each measurement.

In a 1986 article (Grangier *et al*. [24]), an optical source based on Calcium atoms was reported to yield an anti-correlation parameter as low as 0.18 which the authors classified as practically a single-photon source. The apparent quantum interference patterns of probability amplitudes obtained with a Mach-Zehnder configuration containing two beam splitters are interpreted as a clear evidence of single-photon interference. However, as one single photon would be scattered randomly in a quantum Rayleigh interaction with electric dipoles in the dielectric medium of the beam splitter, the possibility does exist that the interference was, in fact, created by the 18% of the groups of photons emitted spontaneously and slightly amplified on their way out, and which split at the first beam splitter and recombined at the second one, on their way to the same photodetector.

Over the next 15 years, to the turn of the century, optically nonlinear parametric crystals occupied centre stage as the most practical source of allegedly single photons. Pairs of spontaneously emitted photons are generated simultaneously by the interaction of an optical pump with nonlinear crystals. The optical frequencies and wavevectors of the two emitted photons obey conservation laws resulting in a high degree of correlation between the same degrees of freedom, and known as entanglement of photon states. From a physical perspective, measuring one photon's characteristic values would indicate the pair photon's values of frequency, wavevector, and, given the anisotropic polarisation or birefringence of the crystal, the polarisation of the other photon.

Although the optical pump pulses exciting the nonlinear crystals have relatively low levels of power, parametric amplification cannot be prevented resulting in a few, rather than one, photons per temporally discrete group. Therefore, the experimental results presented in the 1999 review article by Mandel [14] can be explained without quantum interference of probability amplitudes, and, in so doing, taking the counterintuitive element out of the picture [25]. This is, particularly, the case for the Hong-Ou-Mandel dip associated with a reduced counting of photon coincidences between the output modes of a beam splitter mixing two input synchronised and identical streams of single photons.

Similarly, the claim of remotely collapsing a wave function is highly questionable (Fuwa *et al*. [26]). A single photon propagating through a beam splitter would be deflected from its planned pathway by quantum Rayleigh scattering. Furthermore, the maximum likelihood method of numerically reconstructing a quantum state from raw data "aims to find, among the variety of all possible density matrices, the one that maximizes the probability of obtaining the given experimental data set and is physically plausible" (Lvovsky and Raymer [27]). From the experimental point of view, the observer B's "photoreceivers do not have to be efficient, and he can post-select on finding his system in a particular subspace" [26]. The maximum likelihood method of reconstruction requires a target state, and the sign parameter *s* is delivered from observer A for the reconstruction of the quantum state by observer B. Therefore, the reconstruction is not independent. Indeed, the quantum Rayleigh scattering would deflect any single-photon crossing the beam splitter, and bearing in mind the slight parametric amplification inside the source, the two observers share the same photon phase from the same group of photons split at the beam splitters. There was no direct link at the level of a pure quantum state of a single measurement between the two observers, if only, because with only one photon in the experimental setup at any given time, only one detector can be triggered, whether or not the photon is entangled. Additional experiments (Ringbauer *et al*. [28, p. 4]) attempting to identify a cause and effect for quantum nonlocality between remote photodetectors found that "a direct causal influence from one outcome to the other can therefore not explain quantum correlations ".

This century, the next stage in the quest for practical single-photon sources involved semiconductor quantum dots placed inside dielectric micro-cavities (e.g., Hepp *et al*. [1]; Lodahl *et al*. [2]); Senellart *et al*. [29]). It is pointed out in [30] that a quantum dot "emits a cascade of photons and a single photon is obtained only through *spectral filtering* of one emission line". High-finesse optical cavities incorporated in a measurement setup distort the temporally regular sequence of single photons because of multiple internal reflections. The emerging stream may contain groups of a few temporally overlapping photons, e.g. five, which may be unevenly split by a beam splitter and reduced in number through quantum Rayleigh spontaneous emission, so as to generate no coincidence for a zero delay-time, in a Hanbury Brown and Twiss measurement. Obviously, the beam splitter can precede the interference filter, in which case quantum Rayleigh stimulated emission can cause two photons from different radiation modes to interact with the same dipole so that one of the photons is coupled into the other radiation mode.

A quantum dot placed in a high finesse micro-cavity of a few-wavelengths long and excited with a picosecond pulse, can emit a photon spontaneously and be re-excited within the duration of the same pulse. If the photon was reflected towards the quantum dot, stimulated emission may occur due to the small dimensions of the micro-cavity. This will result in two, or more, photons leaving the emitter simultaneously, as well as a reduced lifetime of the excited state of the quantum dot, manifesting itself as a higher decay rate overshadowing the Purcell effect.

The conventional description of a photon as an ensemble wave packet composed of Fourier components - which can only exist individually one at a time - leads to counterintuitive explanations. A physically meaningful description of a photon can be identified as a monochromatic Wigner-type spectral component which varies with time. Furthermore, any photon-dipole interactions occurring during the propagation is completely ignored in the professional literature, with the propagation phase being attached to the optical field operator, as the number states carry no optical field.

Equally, in the professional literature, an interference term for a single-photon is provided, mathematically, by creating a pathway-entangled quantum wave function for the photon’s propagation (Garrison and Chiao [3]; Walls [23]). Thus, a mysterious quantum effect appears as a result of an ensemble distribution, even though each individual measurement described by a pure quantum state, i.e., only one component of the ensemble, fails to generate that particular effect. This contradiction leads to the question of whether or not the optical source emits only one photon. Additionally, rather puzzlingly, the operations of beam splitters and interference filters are modelled in terms of continuous waves, ignoring the temporally discrete nature of the stream of photons.

The probability amplitude approach to photonic quantum interference leads to physical contradictions and counterintuitive conclusions which are held up as evidence of non-classical features. These are, however eliminated and physically explained by identifying the intrinsic field of photonic wavefronts, as explained throughout this book.

1.2 An Outline of This Book

Four major elements underpin the purpose of this book. The first two elements are linked to the presence, in a homogeneous dielectric medium, of the quantum Rayleigh conversion of photons. As a result, one photon cannot propagate in a straight line and, initially, entangled photons are annihilated and replaced with independent photons. The other two elements arise from employing a time-dependent pure quantum state to deliver the measured values of photonic degrees of freedom. As a consequence, the quantization of the optical field is derived without any equivalence to quantum harmonic oscillators, and the photon is identified as an energy excitation characterised by a Wigner-type or mixed time-frequency representation of a monochromatic signal pulse.

With a view to identifying and probing a possible boundary between the quantum and classical regimes of optics, Chapter 2 of this book headlined “The Quantum Rayleigh Coupling of Optical Waves” describes the functional roles of the quantum Rayleigh emissions of photons and the resultant classical manifestations.

The conventional interpretation assigns no optical field to photon number states which are the number eigenstates of the electromagnetic field. Any propagation effect is attached to the field operators and the absorption of a particle-like photon requires a transition between two consecutive number states. An optical field, known as a coherent state, is generated by an infinite superposition of number states under the condition of an ensemble eigenstate of the annihilation operator. Nevertheless, a photonic wavefront interacts with electric dipoles instantaneously, and the need arises for an intrinsic and instantaneous optical field for any number of photons, regardless of the overall distribution of the optical beam. A physical solution to this problem is presented in Chapter 3 which identifies dynamic and coherent number states under the headline of “The Intrinsic Optical Field of Photons”

“Photonic quantum noise reduction” can be implemented by means of parametric processes. Their scrutiny reveals common features such as phase dependent gain coefficients accompanied by a phase pulling effect as described in Chapter 4. These physical mechanisms can generate sub-Poissonian distributions of photons through a saturation-like effect and using only integrated photonic circuits.

In Chapter 5, “The Quantum Regime Operation of Dielectric Devices” is analysed in the light of evidence that emerged from the previous chapters, leading to different physical processes for various types of beam splitters, and to the temporal role that interferometric filters play in altering the original time sequence of a beam of single photons. As a consequence, individual measurements reveal the physical processes involved in creating interference patterns and are represented by pure quantum states which are dependent on the position and time of the measurements. The ensemble statistical distribution ensues from repeated measurements.

“Photonic coincidences and correlations“ is the headline of Chapter 6 which identifies temporal and spatial properties and aspects of an individual measurement, as well as photon-dipole interactions by means of the pure dynamic and coherent number states derived in Chapter 3. The use of pure, dynamic and coherent quantum states enables us to demystify the quantum nonlocality alleged to create coincidences between two separate photocurrents generated by mixing single photons across a beam splitter.

"Quantum Rayleigh Annihilation of Entangled Photons" of Chapter 7 specifies the Hamiltonian of interaction between the electric dipoles in a dielectric medium and the optical field of photons propagating through that medium. As a result, one single photon is absorbed and spontaneously emitted in a random direction with a random state of polarisation. Yet, the remote correlations between these individual photons lead to the same correlation functions producing the same Bell-type outcomes as the absorbed entangled photons. The statistical character of the "quantum nonlocality" outcomes is reinforced by the possibility of obtaining the same correlations with two separate sources placed in the vicinity of the respective photodetectors.

In Chapter 8, recent experimental and theoretical developments are explained in the context of this textbook's formalism, pointing to a blurred boundary between quantum and classical regimes, which is borne out by the analyses of the previous chapters. A paradigm shift in the interpretation of experimental outcomes of Quantum Optics is highly necessary, being based on the following physical processes and elements:

1. The quantum Rayleigh spontaneous and stimulated emissions;
2. The unavoidable parametric amplification of spontaneous emission, and the formation of groups of monochromatic photons in a high finesse cavity incorporating a quantum dot;
3. Self-contained quantisation of the optical field without any harmonic oscillators leading to the dynamic and coherent number states;
4. The intrinsic optical field of photons and their localised spatial distributions;
5. The description of instantaneous and localised photon-dipole interactions by means of pure, dynamic and coherent number states;
6. The quantum evolution and predictions being described by the Ehrenfest theorem, for any level of optical field excitation, to evaluate the expectation value of an operator in the context of a given set of pure wavefunctions.
7. Identifying quantum phenomena at the level of single events and measurements with a space- and time-dependence, leading to quantum locality and realism.

Overall, there are no quantum optic "miracles" once the physically present effects are correctly identified.

1.3 Remarks

The smooth transition from the quantum regime of one or a few photons to the classical one of a large number of photons is due to the adequate choice of a wavefunction in the form of the *pure*, *dynamic and coherent* number states derived in Chapter 3. These states deliver the correct number of photons carried by a radiation mode, its field amplitude and phase quadratures. The equations of motion for the evolution of these variables are derived in Chapter 3 and applied in the following Chapters for any levels of mode excitation and photon-dipole interactions.

Measurements of instantaneous wave fronts are described by dynamic and coherent number states which erase any quantum-classical boundary. While the intrinsic optical field of photons is critical for time-correlations at the level of one or a few photons, its importance could still be significant at high levels of photon numbers for a monochromatic group of photons as the optical field does not have a Fourier spectrum but is represented by a mixed time-frequency structure of the form S (ω, t). That is, the monochromatic spectral component exhibits a time-varying amplitude which is not related to an optical Fourier spectrum corresponding to time-independent spectral components.

Entangled photons are scattered by the quantum Rayleigh spontaneous emission but remote correlations of measured states of polarisations are still reproduced by single and independent photons.

The quantum Rayleigh photon-dipole interaction may involve two photons colliding at a dipole, with a possible outcome being the transfer of the excitation from one radiation mode to the other. This process may take place inside the dielectric medium of an interference filter or beam splitter, thereby creating groups of monochromatic photons from initially independent photons bouncing back and forth inside a resonant cavity. This quantum Rayleigh coupling of photons may explain the Hong-Ou-Mandel dip as one photon carried by one radiation mode may be captured by another photon associated with a second radiation mode inside a cavity.

Equally, a remote wavefunction state preparation through a detection of one of the entangled photons is practically impossible because the photodetector's excitation is triggered by an energy level of its structure rather than a quadrature state of a quantum harmonic oscillator.

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CHAPTER TWO

THE QUANTUM RAYLEIGH COUPLING OF OPTICAL WAVES

Most activities in Quantum Optics aim to generate and manipulate one photon per radiation mode. It is assumed that once generated, a single photon will propagate unimpeded through a dielectric medium despite encountering a large number of electric dipoles. However, as a result of photon-dipole interactions, the process of quantum Rayleigh scatterings or emissions comes into play involving the absorption of one photon and the spontaneous emission of another photon of the same energy or frequency. At the photodetection stage, the temporally discrete electronic signals are assumed to be triggered by one single photon despite the possibility of a few photons arriving simultaneously and triggering a similar signal.

From a physical perspective, in dielectric devices such as optical fibres and integrated optic waveguides, beam splitters, interference filters, polarisation controllers, etc., a single photon cannot propagate in a straight line, thereby raising significant questions about the validity of the conventional model or interpretation of Quantum Optics experimental outcomes.

Despite having been well documented before the 1980s, the quantum Rayleigh spontaneous and stimulate emissions were totally ignored in any explanation of optical mode coupling devices such as directional waveguide couplers. Two optical beams of the same frequency but different wave vectors propagating through a homogeneous dielectric medium can exchange photons with each other and/or undergo mutually induced phase shifts as a result of stimulated Rayleigh emission underpinning the coupling term of the Poynting theorem. Quadrature fields of the same optical wave exchange power as they propagate through a homogeneous and linear dielectric medium. Consequently, coupling of photons between two optical waveguides takes place in the shared cladding region.

The widely used coupled-mode theory describing optical power coupling between two adjacent waveguides – introduced in the 1960s – relies on a perturbation of the cladding as a physical mechanism. However, the gradient of the perturbed dielectric constant gives rise to a randomly scattering term in the comprehensive wave equation. This approach is inconsistent with directional coupling of photons and leads to physically impossible outcomes raising questions about the validity of its application as in recently published articles in the *IEEE J. Quantum Electron*., vol. 54, 2018 (no. 1, article 6300206, and no. 2, article 6800207).

A physically meaningful and correct identification of processes underpinning the description of propagation and coupling of photons in dielectric media is crucial in order to open up new practical ways of designing, fabricating, and operating integrated photonic devices. The quantum Rayleigh conversion of photons (QRCP), provides a meaningful explanation for the operation of the optical directional couplers, in contrast to the perturbation approach based on the fictional splitting of the dielectric constant $\varepsilon (x, y, z) = \varepsilon_b + \Delta \varepsilon$ into a uniform background ε_b and a perturbation $\Delta \varepsilon (x, y, z)$ which is supposed to generate a coupling polarisation $\Delta \boldsymbol{P} = \Delta \varepsilon \ \boldsymbol{E}$ induced by an overlapping optical field. No explanation has been provided, at least in this context, as to how the optical field can discriminate, physically, between the total local value of the optical susceptibility and the added perturbation to the refractive index, as opposed to the mathematical splitting of the permittivity.

2.1 Coupled-Wave Interactions in a Homogeneous Dielectric Medium

Optical parametric processes are well established for the second- and third- order susceptibilities (Shen [1]; Boyd [2]) and are characterized by conservation of the total energy of photons before and after the interactions. The parametric gain displays a strong dependence on the relative phase between the pump and signal waves. A quantum feature of a parametric process of photon conversion is the amplification of spontaneously emitted photons to generate another optical wave (Vatarescu [3]; Inoue and Mukai [4-5]). These interactions can take place in both homogeneous and inhomogeneous dielectric media.

Similarly, an optically linear parametric (OLP) interaction consists of an electric dipole absorbing one photon and emitting one photon of the same frequency or energy, e.g. the elastic or Rayleigh scatterings (Marcuse [6]; Louisell [7]). The corresponding Hamiltonian of interaction and the Heisenberg equation of motion for the photon annihilation and creation operators are presented below in this Chapter.

This quantum process is localized, can take place in both homogeneous and inhomogeneous dielectric media, and is described macroscopically, or classically, by means of the interaction term included in the Poynting theorem of the flow of energy. The resultant combination of in-quadrature waves is similar to the quadrature states of light (Mandel and Wolf [8], Ch. 21) exchanging photons through parametric interactions.

The classical Rayleigh scattering which is attributed to local perturbations or fluctuations in the dielectric constant is linked to optical power losses in optical fibers.

Quantum electronically, an oscillating electric dipole polarisation can be the source of spontaneous and stimulated emissions of photons [6-7]. Consequently, coupling of photons between two arbitrary waves interacting simultaneously with the same dielectric medium can take place with one beam of photons exciting the electric dipole polarisation and the other beam de-exciting the dipoles and gaining power through stimulated emission. The direction of coupling will depend on the relative phase between the waves. These physical processes require that Maxwell's *curl* ***H*** equation of *each* wave be driven by the *total* electric dipole polarisation available in the medium.

A practical device based on optical power coupling is the two-waveguide optical directional coupler. The conventional electromagnetic coupled- mode theory, e.g., (Huang [9]; Yishen,*et al*. [10]; Huang and Mu [11]; Marcatili [12]) links the optical power coupling between two single-mode waveguides to the perturbation of the permittivity of the optical waveguide cladding and the unperturbed evanescent modal fields that existed in the absence of the second waveguide. But the tail end of these fields no longer exists physically, having been disturbed by the introduction of the other waveguide. In reference [12], two sets of wave equations are mixed up to generate an "interaction" between the guided modes of the individual waveguides. One set of equations involves the normal, even and odd, modes of the coupler and the other set involves the modes of the individual waveguides. But no physical effect underpins this mathematical technique. Equally, the incoming guided mode of one waveguide is *instantly* converted, at the input to the coupler, into a superposition of the normal modes. But no explanation is provided as to how the propagation constant of the incoming photons is converted into the propagation constants of the normal modes.

Additionally, the approach based on the normal modes [9], [12] of the two-waveguide structure fails to explain how the incoming guided wave is converted instantly, at the input, into the orthogonal even and odd modes or why the same converting process does not occur between the two single mode waveguides in the cladding region they share. Equally, a group of photons cannot simultaneously cross from one waveguide into the other while propagating in the same normal mode *without converting* the propagation constant. Furthermore, the comprehensive wave equations incorporating the gradient of the dielectric constant would cause the tail end of the evanescent field of one waveguide to be scattered by the existence of the second waveguide.

It is the effect of stimulated emission which is capable of causing photons to propagate with the same wave vector as the stimulating field [6-7] instead of being spontaneously emitted or classically scattered. The gain providing stimulated emission is always the source of spontaneous emission. The spontaneous emission, which is a feature of the quantum wave-dipole interactions, is dependent on the optical susceptibility and the level of the optical pump and is distinct from the zero-point fluctuations of the electromagnetic field, as presented below in this Chapter.

The phase-dependent parametric gain in optical fibers for the third-order susceptibility involving four-photon mixing interactions was analyzed [3] and demonstrated experimentally [4-5]. Conceptually, the following analysis adapts reference [3] to an *optically linear* medium.

In this Chapter, a physically meaningful framework is developed – Section 2.1.1 – for the optically linear parametric (OLP) coupling of photons by making it consistent with the quantum effects of spontaneous and stimulated emissions associated with the linear parametric gain. The optical wave propagates forwards through stimulated emission as a result of the optimal phase-dependent gain. This propagation involves a cascade of photonic conversions between quadrature waves in dielectric media, emerging from the Poynting theorem of the flow of energy. The refractive index results from local exchanges of energy between the optical field and the electric dipole polarisation.

In the case of an optical directional coupler – Section 2.1.2 – composed of two waveguides, the extinction of the wave launched into one waveguide gives rise to in-quadrature waves, one in each waveguide. The related coupling coefficients are determined from the total value of the local susceptibility at every point in space where any two waves overlap. No approximations are made in the derivation of results and no assumptions are needed. Previous experimental results are reassessed and physical aspects of optically linear parametric interactions (OLP) are outlined.

2.1.1 Inter-quadrature Coupling through Optically Linear Parametric Interactions

An optical collimated beam, or a travelling radiation mode, characterized by a field amplitude E_o, an initial phase φ, an angular frequency ω, a wave vector $\boldsymbol{k}$ and a field polarisation unit $\boldsymbol{u}$, is represented, at time t and distance $\boldsymbol{r}$= (x, y, z) from origin, by the relations:

$$\boldsymbol{E}\,(\boldsymbol{k},\varphi) = E_o\,(z)\,f(\boldsymbol{r})e^{-i\varphi}\,e^{i(\omega t - \boldsymbol{k}\cdot\boldsymbol{r})}\,\boldsymbol{u} \tag{2.1a}$$

$$\iint f^2\,(x,y,z)dx\,dy = 1 \tag{2.1b}$$

$$\mathrm{P}\,(z) = 0.5\,\varepsilon_o n\,c\,E_o^2\,(z) \tag{2.1c}$$

The field distribution $\boldsymbol{E}$ is given in terms of the peak amplitude E_o and the spatial distribution $f\,(x, y, z)$ which has units of m^{-1} and is normalized across any (x, y)-plane in eq. (2.1*b*) so that P(z) represents the total average power crossing that surface. Additionally, ε_o is the permittivity of free space, n is the refractive index of the medium, and c is the speed of light in vacuum.

The dipole polarisations involved in the interaction are linear, i.e.

$$\boldsymbol{P} = \varepsilon_o\,\chi\,\boldsymbol{E}$$

where χ (x, y, z) is the susceptibility of the medium, and in terms of the photon frequency notation,

$$\boldsymbol{P}(\omega) = \varepsilon_o\,\chi\,(\omega\,;\,\omega)\;\boldsymbol{E}\,(\omega)$$

indicating a parametric dipole polarisation with one photon being absorbed and one photon being emitted, both having the same energy [1-2].

The conventional approach in electromagnetic theory is to insert the total field $\boldsymbol{E}_{tot} = \boldsymbol{E}_1 + \boldsymbol{E}_2$ in the Maxwell equations and then select terms of identical indices on both sides. This approach, however, leads to two difficulties: 1) the need to single out one wave by using an orthogonality condition; and 2) it disregards the quantum process of two optical waves of the same frequency but different wave vectors interacting with each other in an optically linear dielectric medium (as described below). As any wave propagating through the dielectric medium will interact with every electric dipole polarisation $\boldsymbol{P}$ oscillating at the same frequency, the Maxwell *curl* $\boldsymbol{H}$ equation (Ampère's law) needs to be driven by the total electric dipole polarisation or the macroscopic dipole polarisation, leading to an additional term:

$$\nabla \times \boldsymbol{H}_1 = \frac{\partial}{\partial t}\,(\boldsymbol{E}_1 + \boldsymbol{P}_1) + \frac{\partial}{\partial t}\,(\boldsymbol{P}_2) \tag{2.2}$$

And a similar equation applies, with the indices interchanged, for the magnetic field $\boldsymbol{H}_2$.

As the optical beam crosses an (x, y) - plane interface from free-space into a dielectric material, the wave equations for its electric $\boldsymbol{E}_1$ and magnetic

$\boldsymbol{H}_1$ fields describing its propagation are derived from the Maxwell equations (Orfanidis [13, Ch. 14]) as

$$\nabla^2 \boldsymbol{E_1} + (\,k_o n)^2\, \boldsymbol{E_1} = \nabla\left(\boldsymbol{E_1} \cdot \frac{\nabla\varepsilon}{\varepsilon}\right) - \omega^2\, \mu_0\, \boldsymbol{P_2} \qquad (2.3a)$$

$$\nabla^2 \boldsymbol{H_1} + (\,k_o n)^2\, \boldsymbol{H_1} = -i\,\omega\,\nabla\,\varepsilon\, \times \boldsymbol{E_1} - \,i\,\omega\,\nabla \times\, \boldsymbol{P_2} \qquad (2.3b)$$

where $k_o = \omega / c$, $\varepsilon / \varepsilon_o = n^2 = 1 + \chi$ is the dielectric constant, and the time derivative operator was replaced with $\partial \boldsymbol{E}$, (or $\boldsymbol{H}$) / $\partial t = i\, \omega\, \boldsymbol{E}$, (or $\boldsymbol{H}$). The polarisation density $\boldsymbol{P}_1$ was included on the left-hand side of eqs. (2.3) and a second source polarisation $\boldsymbol{P}_2$ was added in the *curl* $\boldsymbol{H}$ Maxwell equation in order to point out the various quantum operations that these polarisations carry out simultaneously. Equivalent equations apply to $\boldsymbol{E}_2$ with the indices interchanged.

The initial direction of propagation of the refracted wave from the boundary between the free-space and the medium is determined from Snell's law which is indicative of the conservation of wave momentum in the (*x, y*) - plane. The refracted wave is generated by a boundary layer of source terms incorporating the gradient $\nabla\varepsilon$ and $\boldsymbol{E}_1$ in eqs. (2.3). Inside the medium, for $\nabla\varepsilon \neq 0$, scattering of the optical wave takes place – see the right-hand side of eqs. (2.3) – unless the field propagates along an optical waveguide, satisfying the boundary conditions for a guided mode. The polarisation $\boldsymbol{P}_2$ can radiate into mode $\boldsymbol{k}_1$ and a mutual interaction emerges from the Poynting theorem.

The differential, local and temporal, Poynting theorem [13, Ch. 1] of the optical flow of energy has the following form, with the asterisk denoting the complex conjugate of the variable:

$$\nabla \cdot \boldsymbol{\mathcal{P}_1} = -i\,\omega\, \boldsymbol{D}_1 \cdot \boldsymbol{E}_1^* - i\,\omega\, \boldsymbol{B}_1 \cdot \boldsymbol{H}_1^* - i\,\omega\, \boldsymbol{P}_2 \cdot \boldsymbol{E}_1^* \qquad (2.4)$$

where $\boldsymbol{\mathcal{P}_1} = \boldsymbol{E_1} \times \boldsymbol{H}_1^*$ is the Poynting vector parallel to the wave vector $\boldsymbol{k}_1$, and the vectors $\boldsymbol{k}$, $\boldsymbol{E}$ and $\boldsymbol{H}$ are perpendicular to each other for the same radiation mode. We align $\boldsymbol{k}_1$ to be parallel to the z-axis in a bulk medium. The constitutive relations are: $\boldsymbol{D} = \varepsilon\, \boldsymbol{E}$ and $\boldsymbol{B} = \mu\, \boldsymbol{H}$. Making use of the identity $1/c^2 = \varepsilon_o\, \mu_o$ and, for a radiation field $\varepsilon E^2 = \mu H^2$, we obtain from the real part of eq. (2.4) the longitudinal rate of change of the optical intensity E_j^2 at a point (*x, y, z*), and from the entire complex equation, the rate of change of the field $E(\boldsymbol{k}_j) = E_j \exp(-i\,\varphi)$ of eq.(2.1a) after setting $E_j = E_{oj}\, f_j\,(x, y, z)$, with j = 1 or 2 :

$$\frac{d\,E_1^2}{dz} = -2\,\gamma E_1\,E_2\,sin\,\theta_{21} \tag{2.5a}$$

$$\frac{d\,E\,(\boldsymbol{k}_1)}{dz} = -i\,\left(k_0\,n + \gamma\,\frac{E_2}{E_1}cos\,\theta_{21}\right)\,E_1 - \gamma E_2\,sin\,\theta_{21} \tag{2.5b}$$

$$\frac{d\,\theta_{21}}{dz} = (\,\boldsymbol{k}_2 - \boldsymbol{k}_1)\cdot\,\boldsymbol{u}_z + \gamma\left(\frac{E_1}{E_2} - \frac{E_2}{E_1}\right)cos\,\theta_{21} \tag{2.5c}$$

$$\frac{d}{dz}\varphi_1 = \gamma\frac{E_2}{E_1}\,cos\,\theta_{21} \tag{2.5d}$$

$$\theta_{21} = (\,\boldsymbol{k}_2 - \boldsymbol{k}_1)\cdot\boldsymbol{z} + \varphi_2 - \varphi_1 \tag{2.5e}$$

where $\gamma = k_o\,\chi\,\boldsymbol{u}_1\cdot\,\boldsymbol{u}_2\,/\,(2\,n)$ is the local coupling coefficient, $\boldsymbol{k}_1 = k_o\,n\,\boldsymbol{u}_z$ is the wavevector, with $\boldsymbol{u}_z$ being a unit vector in the z-direction, and θ_{21} is the phase difference between the two fields. Similar equations to (2.5) hold for E_2 with the indices interchanged.

The refractive index n in eq. (2.5*b*) arises from the first two terms on the right-hand side of eq. (2.4) and is the result of local and instantaneous exchanges of energy between the optical field and the dielectric medium. It corresponds to the self-coupling term with $\boldsymbol{P}_1$ from $\boldsymbol{D}_1 = \boldsymbol{E}_1 + \boldsymbol{P}_1$ replacing $\boldsymbol{P}_2$ in eq. (2.4) and setting $\theta = 0$ in eqs. (2.5). The second term on the right-hand side of eq. (2.5*b*) emerges from the last term of eq. (2.4), indicating a possible phase shift brought about by a mutual interaction between the optical waves E_2 and E_1.

As indicated by the last term of eq. (2.4) or by eq. (2.5*a*), for $E_1 \neq 0$, coupling of power can take place in a homogeneous medium, i.e. where $\nabla\varepsilon = 0$. This effect is not identified by the wave equations (2.3). For a real value χ , the coupling term in eq. (2.4) conserves the number and energy of the photons involved in the process of stimulated emission. A non-vanishing stimulating field of a particular radiation mode $\boldsymbol{k}_1$ is obtained from spontaneous emission or scattering of the incoming wave, e.g. elastic or Rayleigh scatterings. An expression evaluating the amount of spontaneous emission is outlined at the end of this Chapter.

The maximum parametric gain is found from eqs. (2.5) to occur in the same direction as that of the pump wave E_2 , i.e. $\boldsymbol{k}_1 = \boldsymbol{k}_2$, and for $\theta = -\pi\,/\,2$. Groups of photons spontaneously emitted by $\boldsymbol{P}_2$ and identified as E_1 will have their arbitrary phase changed rapidly by the interaction as pointed out by eq. (2.5*c*), and the relative phase will become locked at $\theta = -\pi\,/\,2$. As a result, an optical field E_1 will appear and the outgoing field will consist of an amplitude modulation, taking the form:

$$E_{out} = (E_2 + i\, E_1)\, e^{i\,(\omega\, t - \boldsymbol{k}\cdot \boldsymbol{r})} =$$

$$= E_o[\, cos\, (\gamma\, z) + \, i\, sin\, (\gamma\, z)]\, e^{i\,(\omega\, t - \boldsymbol{k}\cdot \boldsymbol{r})} \qquad (2.6a)$$

$$ReE_{out} = E_o\, [\, cos\, \ (\gamma\, z)\ \ cos(\omega\, t - \boldsymbol{k}\cdot\boldsymbol{r})\ -$$

$$- \, sin\, (\gamma\, z)\ \ sin\ \ (\omega\, t - \boldsymbol{k}\cdot\boldsymbol{r})] \qquad (2.6b)$$

which has the appearance of a phase modulation. The real part of E_{out} reveals two quadrature waves [8] exchanging power as they propagate alongside each other in the z- direction, i.e. $\boldsymbol{k}\cdot\boldsymbol{r} = k_z \cdot z$. Each quadrature field rotates in the (x, y) plane with its amplitude varying periodically.

2.1.2 The operation of optical waveguide directional couplers

Bearing in mind the physical elements of the quantum Rayleigh emissions as outlined in Section 2.1.1 above, and relying on the quadrature states derived in the previous Section, we turn our attention to the case of an optical directional coupler composed of two single-mode waveguides. After identifying the waveguides by the letters *a* and *b*, we define the normalized fields $e = E\, /\, E_o$ (where E_o^2 corresponds to the normalizing input power) as

$$e_{\,a} = \left(p_a\, e^{-i\varphi_{pa}} + \ q_a\, e^{-i\varphi_{qa}}\,\right) f_a\, (\boldsymbol{r})\, e^{i\,(\omega\, t - \boldsymbol{k}_a\cdot \boldsymbol{r})} \qquad (2.7a)$$

$$e_{\,b} = \left(p_b\, e^{-i\varphi_{pb}} + \ q_b\, e^{-i\varphi_{qb}}\,\right) f_b\, (\boldsymbol{r})\, e^{i\,(\omega\, t - \boldsymbol{k}_b\cdot \boldsymbol{r})} \qquad (2.7b)$$

where the subscripts p and q correspond, respectively, to the initial quadrature phases of $\varphi = 0$ and $-\pi/2$, and amplitudes $p_{a\,(b)}$ and $q_{a\,(b)}$.

A mathematical solution to eqs. (2.5) can be derived by means of elliptic functions for the total field phasors of e_a and e_b as defined in eqs. (2.1) but its complexity obscures physical features of the optically linear parametric interactions. Such characteristics are outlined in the remainder of this Section.

As the waves propagate along a directional coupler – illustrated in cross-section in Fig. 2.1 – changes in the evanescent fields spread across the entire modal field through the boundary conditions. For an input launched into one of the waveguides, the generated in-quadrature wave appears in both waveguides – see Fig. 2.2 – from amplified spontaneously emitted photons captured by each modal field. The evanescent fields of the two waveguides overlap in the cladding region they share – see Fig. 2.1 - bringing about coupling interactions between the two modal fields through the quantum Rayleigh field-dipole effects.

The tail of the evanescent field of either waveguide is scattered by the gradient of the local dielectric constant created by the core-cladding boundary of the other waveguide, as indicated by the source terms of the wave equations (2.3) containing $\nabla\varepsilon$. This scattering will lead to a loss factor being added to eqs. (2.4) and (2.5) and the scattered photons could become seed photons to be amplified in the other waveguide.

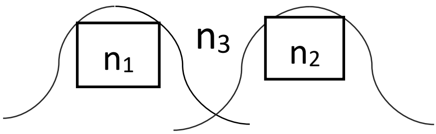

Fig. 2.1 The Rayleigh induced coupling of photons takes place in the cladding between the two waveguides with refractive index n_3

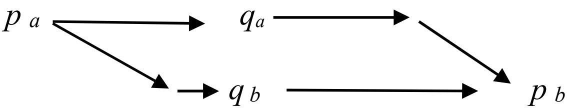

Fig. 2.2 A diagram of the longitudinal stages of the comprehensive coupling effects between the two waveguides of a directional coupler.

These effects are illustrated in (Marcatili *et al*. [14, Fig. 1(c)]; Syms and Peall [15]) where the measured optical power is coupled into the output waveguide through the cladding region, predominantly from the wave emerging from the terminated input waveguide. A small fraction of power associated with the evanescent tail of the input guided wave may directly excite the outgoing modal core field. These interactions are also present in Fig. 1(a) and (b) of [14] where two coupling stages can be identified: the conventional two-waveguide coupler and the coupling between the mode of the waveguide continuing to the output and the optical field propagating alongside the waveguide and originating from the terminated waveguide.

The exchange of power between any two modal fields is evaluated from eq. (2.4) integrated over the (*x, y*) - plane. Eqs. (2.5) are modified by substituting $P^{1/2}$ from eq. (2.1c) for the fields *E*, by replacing $\boldsymbol{k}$ with the propagation constant $\beta = \boldsymbol{k} \cdot \boldsymbol{u_z}$ and γ with this coupling coefficient: