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Gerd Marowsky *Editor*

Planar Waveguides and other Confined Geometries

Theory, Technology, Production, and Novel Applications

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Gerd Marowsky
Laser-Laboratorium Göttingen e.V.
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Preface

The aim of this book is to provide an actual collection of highlights in the technology, production, and novel applications of waveguides and other confined geometries, in particular optical fibers.

It is well known that waveguide phenomena are nowadays a textbook subject and some energy is needed to prove that waveguide problems are still of great scientific and commercial interest.

Any textbook approach should go back to the books by Theodor Tamir,[1] who always considered waveguiding as the heart of any integrated optics. It is very interesting to study the details of his first book on this subject in as early as 1975. There he dated the beginning of integrated optics to the late 1960s and mentioned as primary applications signal processing and optical communications. It is striking to note that he did not mention at this time any analytical applications (lab-on-chip) or sensors. However, he already did mention the possible chances of novel optical phenomena opened *by the concentration of high intensity fields in thin films and effects due to attendant enhancement of nonlinear and/or active properties of such films*. It should be emphasized that at least two contributions in this book deal with nonlinear phenomena in waveguiding structures and potential applications.

Contributions came from Europe, Canada, and Israel. The manuscripts have been arranged as follows:

- Theoretical considerations for semi-infinite planar waveguides. Unusual, novel ways of efficient frequency conversion in structured fibers and nanoplasmonic waveguides
- Five contributions on laser writing of waveguiding structures in various geometries and materials
- Channel formation in cluster targets as an example of self-organization of waveguiding structures

[1] Still available is the second corrected and updated edition of the book "Integrated Optics," T. Tamir editor, Springer-Verlag, Berlin, Heidelberg, New York, 1979.

- An example of a novel application of Bragg grating technology using optical fibers for 3D shape sensing

The contributions are beyond the normal textbook style, representing cutting-edge research results. A short preview/summary of some of the selected contributions is given below:

In Chap. 1 of this book radiation emerging from an open-ended waveguide is studied in theory. The idealized case of a planar waveguide open to an infinite half space is considered. An iterative scheme is constructed in analogy to the exact solution of reflection from a step-potential in one-dimensional quantum mechanics and Fresnel reflection of electromagnetic radiation from a half space. For the planar waveguide, the first step in the iterative scheme is evaluated. This yields an approximation to the radiation pattern emerging from the waveguide, as well as the amplitude of the wave reflected back into the waveguide.

Another selected example, Chap. 2, considers second-harmonic generation in a circular cylindrical waveguide with embedded periodically arranged tubelets of high nonlinear susceptibility. The configuration is similar to usual quasi-phase-matched arrangements, but places the nonlinear polarization sheet at intensity peaks of the guided fundamental wave. This leads (yet only in theory) under optimized conditions to high conversion efficiencies and hence pump depletion has to be considered in the framework of mode-coupling theory.

A major advance on femtosecond laser structuring inside optical fibers in Chap. 4 presents powerful new directions for inscribing optical, microfluidic, and mechanical structure devices in the fiber cladding that efficiently interconnect with the fiber core waveguide. This all-fiber platform circumvents tedious packaging and assembly process steps to enable highly functional photonic microsystems and lab-in-fiber devices to form into a compact and flexible optical fiber system that greatly enhances the value of today's widely deployed fiber cables.

The modification of optical properties of dielectrics has led to the fabrication of waveguides in a variety of glasses and crystals as shown in Chap. 6. UV photosensitivity of germanium-doped glass led to a new technology of fiber Bragg gratings; now femtosecond lasers have allowed yet another aspect of laser-material interaction to be exploited, bordering in the twilight zone between optical damage and plasma formation. This not fully understood process has also led to the production of 3D waveguides and holds the promise of easing the process of waveguide fabrication in all sorts of transparent media, provided the optical loss can be tamed. At the other end of the spectrum, the brute force of carbon dioxide lasers has also been used to cut glass to form delicate, yet complex, low-loss optical devices for fast prototyping and integration. These technologies are reviewed by examining the induction of refractive index changes and waveguides in several different materials.

An unusual case of waveguide generation is discussed in Chap. 9: channel formation due to focusing of high-power laser radiation in a self-generated rare gas plasma. The discovery of a stable highly nonlinear mechanism that can compress power to a density exceeding thermonuclear values led to a new era of power

compression. The associated steps, increasing by a factor of 10^{10} in power density, mark different stages of technological breakthroughs.

Chapter 10 discusses a novel method for direct femtosecond laser-based writing of Bragg gratings and micro-waveguides not only inside the core but also in the cladding of a standard telecom fiber. This enables an optical 3D shape monitoring system using just a single one-core optical fiber. For a 1.5 m long fiber, the tip position can be tracked with an accuracy of less than 1 mm for the x,y,z-coordinates. Applications of this new technology range from medical tracking systems to long-term inspection of boreholes and drilling wells in the oil and gas industry.

Göttingen, Germany
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Gerd Marowsky

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The editor took the liberty to make appropriate changes in some of the manuscripts for better readability of the book.

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Chapter 1
Radiation from a Semi-infinite Unflanged Planar Dielectric Waveguide

B.U. Felderhof

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1.1 Introduction

In a classic paper Levine and Schwinger [1] studied the radiation of sound from an unflanged circular pipe. Later they extended their theory to electromagnetic radiation [2]. Their work constituted the first major advance in the theory of diffraction after Sommerfeld's exact solution of the problem of plane wave diffraction by an ideally conducting half plane [3]. In the theory of diffraction of sound, microwaves, or light, the radiation is assumed to propagate in uniform space with reflection by rigid objects or idealized boundaries. For a lucid introduction to the theory of diffraction we refer to Sommerfeld's lecture notes [4]. The early theory of diffraction was reviewed by Bouwkamp [5]. Later developments are discussed by Born and Wolf [6]. A brief review of the principles and applications of open-ended waveguides with idealized walls was presented by Gardiol [7].

The invention of the dielectric waveguide by Hondros and Debye [8] has led to the development of optical fibers and the subsequent advances in telecommunication. In the theoretical determination of running wave solutions to Maxwell's equations in a spatially inhomogeneous medium the radiation is assumed to propagate in a guiding structure of infinite length. The problem of emergence of radiation from a semi-infinite waveguide into a half-space is of obvious technical interest. In the case of sound the analysis is based on the exact solution of Levine and Schwinger [1,9–12]. For a dielectric waveguide the observation of the emerging radiation can be used as a tool to study the nature of the driving incident wave [13].

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It is advantageous to simplify the theoretical analysis by the use of planar symmetry. The theory of Levine and Schwinger was extended to planar geometry by Heins [14,15]. In the following we study radiation emerging from a semi-infinite planar dielectric waveguide. As an intermediate step the wave function in the exit plane must be calculated. In this case the integral equation technique of Schwinger [16] cannot be implemented, because of the complicated nature of the integral kernel. We evaluate the emitted radiation and the coefficients of reflection back into the waveguide approximately in a first step of an iterative scheme.

We show in two appendices that the iterative scheme converges to the exact solution in the related problems of reflection by a step potential in one-dimensional quantum mechanics and of Fresnel reflection of an electromagnetic plane wave by a half-space. For the planar dielectric waveguide it does not seem practically possible to go beyond the first step of the iterative scheme.

In a numerical example we study a planar waveguide consisting of a slab of uniform dielectric constant, bounded on both sides by a medium with a smaller dielectric constant [17–19]. In the case studied the first step of the iterative scheme leads to a modification of the wave function at the exit plane which is relatively small in comparison with the incident wave. Hence we may expect that the calculation provides a reasonable approximation to the exact solution.

1.2 Planar Open End Geometry

We employ Cartesian coordinates (x, y, z) and consider a planar waveguide in the half-space $z < 0$ with stratified dielectric constant $\varepsilon(x)$ and uniform magnetic permeability μ_1. In the half-space $z > 0$ the dielectric constant is uniform with value ε' and the magnetic permeability is μ_1. We consider solutions of Maxwell's equations which do not depend on the coordinate y and depend on time t through a factor $\exp(-i\omega t)$. Waves traveling to the right in the left half-space will be partly reflected at the plane $z = 0$ and partly transmitted into the right half-space. The solutions of Maxwell's equations may be decomposed according to two polarizations. For TE polarization the components E_x, E_z, and H_y vanish, and the equations may be combined into the single equation

$$\frac{\partial^2 E_y}{\partial x^2} + \frac{\partial^2 E_y}{\partial z^2} + \varepsilon\mu_1\omega^2 E_y = 0 \qquad \text{(TE)}. \tag{1.1}$$

We have used SI units. For TM polarization the components E_y, H_x, and H_z vanish, and the equations may be combined into the single equation

$$\frac{\partial^2 H_y}{\partial x^2} + \frac{\partial^2 H_y}{\partial z^2} - \frac{1}{\varepsilon}\frac{d\varepsilon}{dx}\frac{\partial H_y}{\partial x} + \varepsilon\mu_1\omega^2 H_y = 0 \qquad \text{(TM)}. \tag{1.2}$$

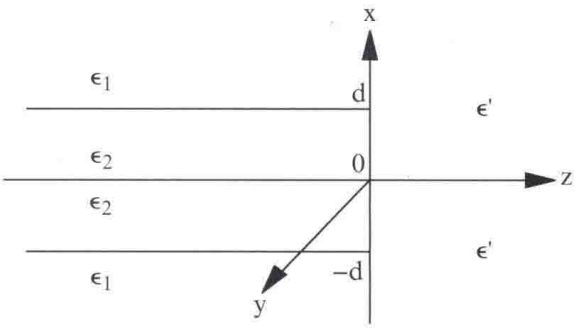

Fig. 1.1 Geometry of the planar waveguide

We assume that the profile $\varepsilon(x)$ is symmetric, $\varepsilon(-x) = \varepsilon(x)$, and has a simple form with $\varepsilon(x)$ increasing monotonically for $x < 0$ from value ε_1 to a maximum value ε_2 at $x = 0$. An example of the geometry under consideration is shown in Fig. 1.1. In the example the dielectric constant in the left half-space equals a constant ε_2 for $-d < x < d$ and a constant $\varepsilon_1 < \varepsilon_2$ for $x < -d$ and $x > d$.

For definiteness we consider only TM polarization. It is convenient to denote the magnetic field component $H_y(x, z)$ for $z < 0$ as $u(x, z)$ and for $z > 0$ as $v(x, z)$. The continuity conditions at $z = 0$ are

$$u(x, 0-) = v(x, 0+), \qquad \frac{1}{\varepsilon(x)} \frac{\partial u(x, z)}{\partial z}\bigg|_{z=0-} = \frac{1}{\varepsilon'} \frac{\partial v(x, z)}{\partial z}\bigg|_{z=0+}. \tag{1.3}$$

We consider a solution $u_{0n}(x, z)$ of Eq. (1.2) given by a guided mode solution

$$u_{0n}(x, z) = \psi_n(x) \exp(i p_n z), \tag{1.4}$$

where $\psi_n(x)$ is the guided mode wave function and p_n the guided mode wavenumber. We assume $p_n > 0$, so that the wave $u_{0n}(x, z) \exp(-i\omega t)$ is traveling to the right. The complete solution takes the form

$$u_n(x, z) = u_{0n}(x, z) + u_{1n}(x, z), \qquad v_n(x, z), \tag{1.5}$$

where $u_{1n}(x, z)$ and $v_n(x, z)$ must be determined such that the continuity conditions Eq. (1.3) are satisfied. The function $u_{1n}(x, z)$ describes the reflected wave, and $v_n(x, z)$ describes the wave radiated into the right-hand half-space.

Since the right-hand half-space is uniform the solution $v_n(x, z)$ takes a simple form and can be expressed as

$$v_n(x, z) = \int_{-\infty}^{\infty} F_n(q) \exp(i q x + i \sqrt{\varepsilon' \mu_1 \omega^2 - q^2}\, z)\, dq. \tag{1.6}$$

The contribution from the interval $-\sqrt{\varepsilon' \mu_1}|\omega| < q < \sqrt{\varepsilon' \mu_1}|\omega|$ corresponds to waves traveling to the right; the contribution from $|q| > \sqrt{\varepsilon' \mu_1}|\omega|$ corresponds to evanescent waves.

Similarly the solution $u_{1n}(x,z)$ in the left half-space can be expressed as

$$u_{1n}(x,z) = \sum_{m=0}^{n_m-1} R_{mn}\psi_m(x)\exp(-ip_m z) + \int_0^\infty R_n(q)\psi(q,x)\exp(-i\sqrt{\varepsilon_1\mu_1\omega^2 - q^2}\, z)\, dq, \tag{1.7}$$

where the sum corresponds to guided waves traveling to the left, with n_m the number of such guided modes possible at the given frequency ω, and the integral corresponds to waves radiating towards the left. We require that the mode solutions are normalized such that [20]

$$\int_{-\infty}^{\infty} \frac{\psi_m^*(x)\psi_n(x)}{\varepsilon(x)}\, dx = \delta_{mn}, \qquad \int_{-\infty}^{\infty} \frac{\psi_m^*(x)\psi(q,x)}{\varepsilon(x)}\, dx = 0,$$
$$\int_{-\infty}^{\infty} \frac{\psi^*(q,x)\psi(q',x)}{\varepsilon(x)}\, dx = \delta(q-q'). \tag{1.8}$$

The guided mode solutions $\{\psi_m(x)\}$ can be taken to be real. Orthogonality follows from Eq. (1.2). We show in the next section how the functions $u_{1n}(x,z)$ and $v_n(x,z)$ may in principle be evaluated from an iterative scheme. The coefficients $\{R_{mn}\}$ and the amplitude function $R_n(q)$ also follow from the scheme.

1.3 Iterative Scheme

The iterative scheme is based on successive approximations to the scattering solution. Thus we write the exact solution as infinite sums

$$u_n(x,z) = \sum_{j=0}^{\infty} u_n^{(j)}(x,z), \qquad v_n(x,z) = \sum_{j=0}^{\infty} v_n^{(j)}(x,z), \tag{1.9}$$

with the terms $u_n^{(j+1)}(x,z)$, $v_n^{(j+1)}(x,z)$ determined from the previous $u_n^{(j)}(x,z)$, $v_n^{(j)}(x,z)$. In zeroth approximation we identify $u_n^{(0)}(x,z)$ with the incident wave,

$$u_n^{(0)}(x,z) = \psi_n(x)\exp(ip_n z). \tag{1.10}$$

The corresponding $v_n^{(0)}(x,z)$ will be determined by continuity at the exit plane $z=0$. From Eq. (1.10) we have $u_n^{(0)}(x,0-) = \psi_n(x)$. This has the Fourier transform

$$\phi_n(q) = \frac{1}{2\pi}\int_{-\infty}^{\infty} \psi_n(x)\exp(-iqx)\, dx. \tag{1.11}$$

Using continuity of the wave function at $z = 0$ and the expression Eq. (1.6) we find correspondingly

$$v_n^{(0)}(x, z) = \int_{-\infty}^{\infty} \phi_n(q) \exp(iqx + i\sqrt{\varepsilon' \mu_1 \omega^2 - q^2}\, z)\, dq, \tag{1.12}$$

so that in zeroth approximation $F_n^{(0)}(q) = \phi_n(q)$. Clearly the zeroth approximation does not satisfy the second continuity equation in Eq. (1.3), and we must take care of this in the next approximation.

For the difference of terms in Eq. (1.3) we find

$$\rho_n^{(0)}(x) = \frac{-i}{\varepsilon(x)} p_n \psi_n(x) + \frac{i}{\varepsilon'} \int_{-\infty}^{\infty} \sqrt{\varepsilon' \mu_1 \omega^2 - q^2}\, \phi_n(q) \exp(iqx)\, dq. \tag{1.13}$$

By symmetry $\rho_n^{(0)}(x)$ is symmetric in x for n even and antisymmetric in x for n odd.

The next approximation $u_n^{(1)}(x, z)$ can be found by comparison with the solution of the problem where the profile $\varepsilon(x)$ extends over all space and radiation is generated by a source $\varepsilon(x)\rho(x)\delta(z)$ with a Sommerfeld radiation condition, so that radiation is emitted to the right for $z > 0$ and to the left for $z < 0$. This antenna solution can be expressed as

$$u_A(x, z) = \int_{-\infty}^{\infty} K(x, x', z)\rho(x')\, dx', \tag{1.14}$$

with kernel $K(x, x', z)$. The latter can be calculated from the Fourier decomposition

$$\delta(z) = \frac{1}{2\pi} \int_{-\infty}^{\infty} e^{ipz}\, dp, \tag{1.15}$$

in terms of the integral

$$K(x, x', z) = \frac{1}{2\pi} \int_{-\infty}^{\infty} G(x, x', p)\, e^{ipz}\, dp, \tag{1.16}$$

with the prescription that the path of integration in the complex p plane runs just above the negative real axis and just below the positive real axis. The Green function $G(x, x', p)$ can be found from the solution of the one-dimensional wave equation,

$$\frac{d^2G}{dx^2} - \frac{1}{\varepsilon}\frac{d\varepsilon}{dx}\frac{dG}{dx} + (\varepsilon\mu_1\omega^2 - p^2)G = \varepsilon(x)\delta(x - x'). \tag{1.17}$$

The solution takes the form [20]

$$G(x, x', p) = \sqrt{\varepsilon(x)}\frac{f_2(x_<, p) f_3(x_>, p)}{\Delta(f_2, f_3, p)}\sqrt{\varepsilon(x')}, \tag{1.18}$$

where $x_<(x_>)$ is the smaller (larger) of x and x' and the remaining quantities will be specified in the next section. The Green function satisfies the symmetry properties

$$G(x, x', -p) = G(x, x', p), \qquad G(-x, -x', p) = G(x, x', p), \tag{1.19}$$

and the reciprocity relation

$$G(x, x', p) = G(x', x, p). \tag{1.20}$$

Consequently the kernel $K(x, x', z)$ has the properties

$$K(x, x', -z) = K(x, x', z), \qquad K(-x, -x', z) = K(x, x', z), \tag{1.21}$$

as well as

$$K(x, x', z) = K(x', x, z). \tag{1.22}$$

The function $u_n^{(1)}(x, z)$ is now identified as

$$u_n^{(1)}(x, z) = -\int_{-\infty}^{\infty} K(x, x', z)\rho_n^{(0)}(x')\, dx'. \tag{1.23}$$

The minus sign is needed to provide near cancellation of the source density between the zeroth and first order solutions, $\rho_n^{(0)}(x) + \rho_n^{(1)}(x) \approx 0$. We find the first order function $F_n^{(1)}(q)$ by Fourier transform from the value at $z = 0$ in the form

$$F_n^{(1)}(q) = \frac{1}{2\pi}\int_{-\infty}^{\infty} u_n^{(1)}(x, 0)e^{-iqx}\, dx. \tag{1.24}$$

The corresponding function $v_n^{(1)}(x, z)$ is found from Eq. (1.6). The first order source density $\rho_n^{(1)}(x)$ is found to be

$$\rho_n^{(1)}(x) = \frac{-1}{\varepsilon(x)}\frac{\partial u_n^{(1)}(x, z)}{\partial z}\bigg|_{z=0} + \frac{i}{\varepsilon'}\int_{-\infty}^{\infty} \sqrt{\varepsilon'\mu_1\omega^2 - q^2}\, F_n^{(1)}(q)\exp(iqx)\, dq. \tag{1.25}$$

In principle the first order function $u_n^{(1)}(x, 0)$ in the exit plane $z = 0$ may be regarded as the result of a linear operator $\mathcal{R}^{(1)}$ acting on the state $\psi_n(x)$ given by the incident wave. The iterated solution then corresponds to the action with the operator $\mathcal{R} = \mathcal{R}^{(1)}(\mathcal{I} - \mathcal{R}^{(1)})^{-1}$, where $\mathcal{I}$ is the identity operator. In order j of the geometric series corresponding to the operator $\mathcal{R}$ the wave functions $u_{1n}^{(j)}(x, z)$ and $v_n^{(j)}(x, z)$ in the left and right half-space can be found by completing the function $u_{1n}^{(j)}(x, 0) = v_{1n}^{(j)}(x, 0)$ in the exit plane by left and right running waves, respectively.

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