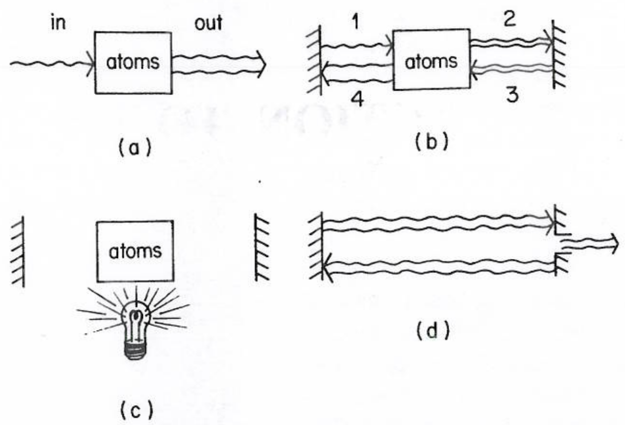

Figure 1.1 Basic elements of a laser.

into the amplifying medium to the extraction of light from the cavity is an opportunity for energy loss and entropy gain. One can say that the success of masers and lasers came only after physicists learned how atoms could be operated efficiently as thermodynamic engines.

One of the challenges in understanding the behavior of atoms in cavities arises from the strong feedback deliberately imposed by the cavity designer. This feedback means that a small input can be amplified in a straightforward way by the atoms, but not indefinitely. Simple amplification occurs only until the light field in the cavity is strong enough to affect the behavior of the atoms. Then the strength of the light as it acts on the amplifying atoms must be taken into account in determining the strength of the light itself. This sounds like circular reasoning, and in a sense it is. The responses of the light and the atoms to each other can become so strongly interconnected that they cannot be determined independently but only self-consistently. Strong feedback also means that small perturbations can be rapidly magnified. Thus it is accurate to anticipate that lasers are potentially highly

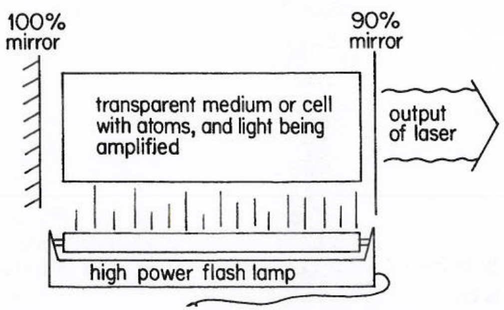

Figure 1.2 A complete laser system, showing elements responsible for energy input, amplification, and output.

erratic and unstable devices. In fact, lasers can provide dramatic exhibitions of truly chaotic behavior, and are the objects of fundamental study for this reason.

For our purposes lasers are principally interesting, however, when they operate stably, with well-determined output intensity and frequency as well as spatial mode structure. The self-consistent interaction of light and atoms is important for these properties, and we will have to be concerned with concepts such as gain, loss, threshold, steady state, saturation, mode structure, frequency pulling, and linewidth.

In the next several sections we sketch properties of laser light, discuss modes in cavities, and give a theory of laser action. This theory is not really correct, but it is realistic within its own domain, and has so many familiar features that it may be said to be ''obvious.'' It is also significant to observe what is not explained by this theory, and to observe the ways in which it is not fundamental but only empirical. These gaps and missing elements are an indication that the remaining 17 chapters of the book may also be necessary.

1.2 LASERS AND LASER LIGHT

Many of the properties of laser light are special or extreme in one way or another. In this section we provide a brief overview of these properties, contrasting them with the properties of light from more ordinary sources when possible.

Wavelength

Laser light is available in all colors from red to violet, and also far outside these conventional limits of the optical spectrum, as shown in Figure 1.3. Over a wide

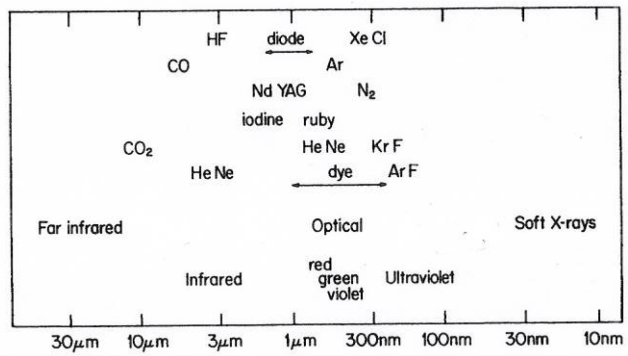

Figure 1.3 Near-optical regions of the electromagnetic spectrum, with types of lasers located approximately at their operating wavelength or wavelengths.

portion of the available range laser light is "tunable." This means that some lasers have the property of emitting light at any wavelength chosen within a range of wavelengths. Tunability is primarily a property of dye lasers. The longest laser wavelength can be taken to be in the far infrared, in the neighborhood of 100–500 μm. Devices producing coherent light at much longer wavelengths by the "maser–laser principle" are usually thought of as masers. The search for lasers with ever shorter wavelengths is probably endless. Coherent stimulated emission in the XUV or soft X-ray region (10–15 nm) has been reported. Appreciably shorter wavelengths, those characteristic of gamma rays, for example, may be quite difficult to reach.

Photon Energy

The energy of a laser photon is not different from the energy of an "ordinary" light photon of the same wavelength. A green–yellow photon, roughly in the middle of the optical spectrum, has an energy of about 2.5 eV (electron volts). This is the same as about 4×10^{-19} J (joules) $= 4 \times 10^{-12}$ ergs. It should be clear that electron volts are a much more convenient unit for laser photon energy than Joules or ergs. From the infrared to the X-ray region photon energies vary from about 0.01 eV to about 100 eV. For contrast, at room temperature the thermal unit of energy is $kT \approx \frac{1}{40}$ eV $= 0.025$ eV. This is two orders of magnitude smaller than the typical optical photon energy just mentioned, and as a consequence thermal excitation plays only a very small role in the physics of nearly all lasers.

Directionality

The output of a laser can consist of nearly ideal plane wave fronts. Only diffraction imposes a lower limit on the angular spread of a laser beam. The wavelength λ and the area A of the laser output aperture determine the order of magnitude of the beams's solid angle ($\Delta\Omega$) and vertex angle ($\Delta\theta$) of divergence (Figure 1.4) through the relation

$$\Delta\Omega \approx \frac{\lambda^2}{A} \approx (\Delta\theta)^2 \tag{1.2.1}$$

This represents a very small angular spread indeed if λ is in the optical range, say

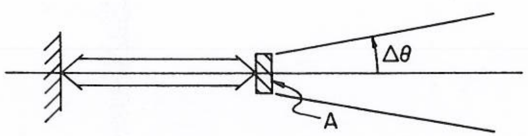

Figure 1.4 Sketch of a laser cavity showing angular beam divergence $\Delta\theta$ at the output

500 nm, and A is macroscopic, say $(5\text{ mm})^2$. In this example we compute $\Delta\Omega \approx (500)^2 \times 10^{-18}\text{ m}^2/(5^2 \times 10^{-6}\text{ m}^2) = 10^{-8}$ steradians.

Monochromaticity

It is well known that lasers produce very pure colors. If they could produce exactly one wavelength, laser light would be fully monochromatic. This is not possible, in principle as well as for practical reasons. We will designate by $\Delta\lambda$ the range of wavelengths included in a laser beam of main wavelength λ. Similarly, the associated range of frequencies will be designated by $\Delta\nu$, the bandwidth. In the optical region of the spectrum we can take $\nu \approx 5 \times 10^{14}$ Hz (hertz, i.e., cycles per second). The bandwidth of sunlight is very broad, more than 10^{14} Hz. Of course filtered sunlight is a different matter, and with sufficiently good filters $\Delta\nu$ could be reduced a great deal. However, the cost in lost intensity would usually be prohibitive. (See subsection on spectral brightness below.) For lasers, a very low value of $\Delta\nu$ is 1 Hz, while a bandwidth around 100 Hz is spectroscopically practical in some cases (Figure 1.5). For $\Delta\nu \approx 100$ Hz the relative spectral purity of a laser beam is quite impressive: $\Delta\nu/\nu \approx 100/(5 \times 10^{14}) = 2 \times 10^{-13}$. This exceeds the spectral purity (Q factor) achievable in conventional mechanical and electrical resonators by many orders of magnitude.

Coherence Time

The existence of a finite bandwidth $\Delta\nu$ means that the different frequencies present in a laser beam can eventually get out of phase with each other. The time required for two oscillations differing in frequency by $\Delta\nu$ to get out of phase by a full cycle is obviously $1/\Delta\nu$. After this amount of time the different frequency components in the beam can begin to interfere destructively, and the beam loses "coherence." Thus $\Delta\tau = 1/\Delta\nu$ is called the beam's coherence time. This is a general definition, not restricted to laser light, but the extremely small values possible for $\Delta\nu$ in laser light make the coherence times of laser light extraordinarily long.

For example, even a "broadband" laser with $\Delta\nu \approx 1$ MHz has the coherence time $\Delta\tau \approx 1$ μsec. This is enormously longer than most "typical" atomic fluorescence lifetimes, which are measured in nanoseconds (10^{-9} sec). Thus even lasers that are not close to the limit of spectral purity are nevertheless effectively 100% pure on the relevant spectroscopic time scale. By way of contrast, sunlight

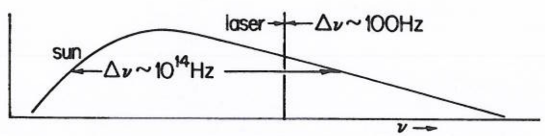

Figure 1.5 Spectral emission bands of the sun and of a representative laser, to indicate the

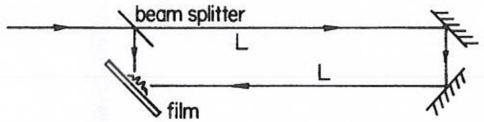

Figure 1.6 Sketch of a two-beam interferometer showing interference fringes obtained at the recording plane if the coherence length of the light is great enough.

has a bandwidth $\Delta\nu$ almost as great as its central frequency (yellow light, $\nu \approx 5 \times 10^{14}$ Hz). Thus for sunlight the coherence time is $\Delta\tau \approx 2 \times 10^{-15}$ sec, so short that unfiltered sunlight cannot be considered temporally coherent at all.

Coherence Length

The speed of light is so great that a light beam can travel a very great distance within even a short coherence time. For example, within $\Delta\tau \approx 1$ μsec light travels $\Delta z \approx (3 \times 10^8$ m/sec$) \times (1$ μsec$) \approx 300$ meters. The distance $\Delta z = c\,\Delta\tau$ is called the beam's coherence length. Only portions of the same beam that are separated by less than Δz are capable of interfering constructively with each other. No fringes will be recorded by the film in Figure 1.6, for example, unless $2L < c\,\Delta\tau = \Delta z$.

Spectral Brightness

A light beam from a finite source can be characterized by its beam divergence $\Delta\Omega$, source size (usually surface area A), bandwidth $\Delta\nu$, and spectral power density $P(\nu)$ (watts per hertz of bandwidth). From these parameters it is useful to determine the *spectral brightness* β_ν of the source, which is defined (Figure 1.7) to be the power flow per unit area, unit bandwidth, and steradian, namely $\beta_\nu = P_\nu / A\,\Delta\Omega\,\Delta\nu$. Notice that $P_\nu / A\,\Delta\nu$ is the spectral intensity, so β_ν can also be thought of as the spectral intensity per steradian.

For an ordinary *nonlaser optical source*, brightness can be estimated directly from the blackbody formula for $\rho(\nu)$, the spectral energy density (J/cm^3-Hz):

$$\rho(\nu) = \frac{8\pi\nu^2}{c^3} \frac{h\nu}{e^{h\nu/kT} - 1} \tag{1.2.2}$$

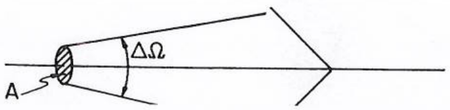

Figure 1.7 Geometrical construction showing the source area and emission solid angle appropriate to a discussion of spec-

The spectral intensity (watts/cm^2-Hz) is thus $c\rho$, and $c\rho/\Delta\Omega$ is the desired spectral intensity per steradian. Taking $\Delta\Omega = 4\pi$ for a blackbody, we have

$$\beta_\nu = \frac{2\nu^2}{c^2} \frac{h\nu}{e^{h\nu/kT} - 1} \tag{1.2.3}$$

The temperature of the sun is about $T = 5800\text{ K} \approx 20 \times (300\text{ K})$. Since the main solar emission is in the yellow portion of the spectrum, we can take $h\nu \approx 2.5$ eV. We recall that $kT \approx \frac{1}{40}$ eV for $T = 300$ K, so $h\nu/kT \approx 5$, giving $e^{h\nu/kT} \approx 150$ and finally

$$\beta_\nu \approx 1.5 \times 10^{-12}\ \text{W/cm}^2\text{-sr-Hz} \qquad \text{(sun)} \tag{1.2.4}$$

Several different estimates can be made for laser radiation, depending on the type of laser considered. Consider first a *low power He-Ne laser*. A power level of 1 mW is normal, with a bandwidth of around 10^4 Hz. From (1.2.1) we see that the product of beam cross-sectional area and solid angle is just λ^2, which for He–Ne light is $\lambda^2 \approx (6238 \times 10^{-8}\text{ cm})^2 \approx 3.89 \times 10^{-9}\text{ cm}^2$. Combining these, we find

$$\beta_\nu \approx 25\ \text{W/cm}^2\text{-sr-Hz} \qquad \text{(He–Ne laser)} \tag{1.2.5}$$

Another common laser is the *mode-locked neodymium-glass laser*, which can easily reach power levels around 10^4 MW. The bandwidth of such a laser is limited by the pulse duration $\tau_p \approx 30$ psec (30×10^{-12} sec), as follows. Since the laser's coherence time $\Delta\tau$ is equal to τ_p at most, its bandwidth is certainly greater than $1/\tau_p \approx 3.3 \times 10^{10}\text{ sec}^{-1}$. We convert from radians per second to cycles per second by dividing by 2π and get $\Delta\nu \approx 5 \times 10^9$ Hz. The wavelength of a Nd:glass laser is 1.06 μm, so $\lambda^2 \approx 10^{-8}\text{ cm}^2$. The result of combining these, again using $A\,\Delta\Omega = \lambda^2$, is

$$\beta_\nu \approx 2 \times 10^8\ \text{W/cm}^2\text{-sr-Hz} \qquad \text{(Nd:glass laser)} \tag{1.2.6}$$

Recent developments have led to lasers with powers at the terawatt level (10^{12} W), so β_ν can be even several orders of magnitude larger.

It is clear that in terms of brightness there is practically no comparison possible between lasers and thermal light. Twenty orders of magnitude in brightness separate our sun from a mode-locked laser. This raises an interesting question of principle. Let us imagine a thermal light source filtered and collimated to the bandwidth and directionality of a He–Ne laser, and the He–Ne laser attenuated to the brightness level of the thermal light. The question is: could the two light beams with equal brightness, beam divergence, and bandwidth be distinguished in any

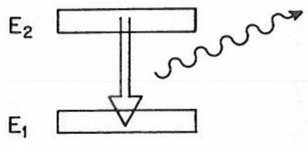

Figure 1.8 Photon emission accompanying a quantum jump from level 2 to level 1.

tions in the light beam. These fluctuations can reflect the quantum nature of the light source, and are detected by photon counting.

Active Medium

The materials that can be used as the active medium of a laser are so varied that a listing is hardly possible. Gases, liquids, and solids of every sort have been made to *lase* (a verb contributed to science by the laser). The origin of laser photons, as shown in Figure 1.8, is most often in a transition between discrete upper and lower energy states in the medium, regardless of its state of matter. He–Ne, ruby, CO_2, and dye lasers are familiar examples, but exceptions are easily found: the excimer laser has an unbound lower state, the semiconductor diode laser depends on transitions between electron bands rather than discrete states, and understanding the free-electron laser does not require quantum states at all.

Type of Laser Cavity

All laser cavities share two characteristics that complement each other: (1) they are basically linear devices with one relatively long optical axis, and (2) the sides parallel to this axis are open, rather than closed by reflecting material as in a microwave cavity. There is no single best shape implied by these criteria, and in the case of ring lasers the long axis actually bends and closes on itself (Figure 1.9). Despite what may seem obvious, it is not always best to design a cavity with the lowest loss. In the case of Q switching an extra loss is temporarily introduced into the cavity for the laser to overcome, and very high-power lasers sometimes use mirrors that are deliberately designed to deflect light out of the cavity rather than contain it.

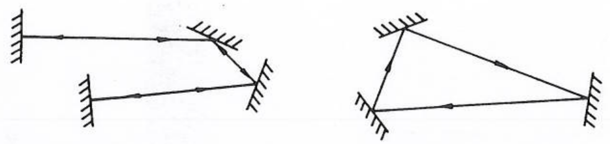

Figure 1.9 Two collections of mirrors making laser cavities, showing standing-wave and traveling-wave (ring) configurations on the left and right, respectively.

Applications of Lasers

There is apparently no end of possible applications of lasers. Many of the uses of lasers are well known by now to most people, such as for several kinds of surgical procedures, for holography, in ultrasensitive gyroscopes, to provide straight lines for surveying, in supermarket checkout scanners and compact disc players, for welding, drilling, and scribing, in compact death-ray pistols, and so on. (The sophisticated student knows, even before reading this book, that some "well-known" applications have never been realized outside the movie theater.)

1.3 LIGHT IN CAVITIES

In laser technology the terms *cavity* and *resonator* are used interchangeably. The theory and design of the cavity are important enough for us to devote all of Chapter 14 to them later in the book. In this section we will consider only a simplified theory of resonators, a theory that is certain to be at least partly familiar to most readers. This simplification allows us to introduce the concept of cavity modes, and to infer certain features of cavity modes that remain valid in more general circumstances. We also describe the great advantage of open, rather than closed, cavities for optical radiation.

We will consider only the case of a rectangular "empty cavity" containing radiation but no matter, as sketched in Figure 1.10. The assumption that there is radiation but no matter inside the cavity is obviously an approximation if the cavity is part of a working laser. This approximation is used frequently in laser theory, and it is accurate enough for many purposes because laser media are usually only sparsely filled with "active" atoms or molecules.

In Chapter 2 full solutions for the electric field in the cavity are given. For example, the z dependence of the x component of the field takes the form

$$E_x(z) = A \sin k_z z \tag{1.3.1}$$

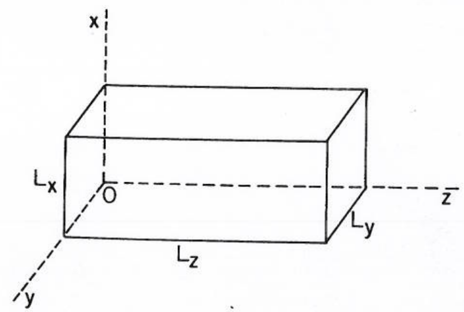

Figure 1.10 A rectangular cavity with side lengths L_x, L_y, L_z.

However, here we are interested only in the simplest features of the cavity field, and these can be obtained easily by physical reasoning.

The electric field should vanish at the walls of the cavity. It will do so if we fit exactly an integer number of half wavelengths into the cavity along each of its axes. This means, for example, that $L = n(\lambda/2)$ along the z axis, where $n = 1, 2, \ldots$, is a positive integer, and L is the cavity length. If we use the relation between wave vector and wavelength, $k = 2\pi/\lambda$, this is the same as

$$k_z = \frac{\pi}{L} n \tag{1.3.2}$$

for the z component of the wave vector. By substitution into the solution (1.3.1) we see that (1.3.2) is sufficient to guarantee that the required boundary condition is met, i.e., that $E_x(z) = 0$ for both $z = 0$ and $z = L$. Similar conditions apply to the x and y dependences, and to the other components of the wave vector, when we fit integer numbers of half waves along the x and y axes:

$$k_x = \frac{\pi}{L} l, \qquad l = 1, 2, \ldots \tag{1.3.3}$$

$$k_y = \frac{\pi}{L} m, \qquad m = 1, 2, \ldots \tag{1.3.4}$$

Note that we have taken the same length L for the x and y dimensions as well.

We can combine these three components of the wave vector to get an expression for the mode frequencies that are possible in the cavity. We use the familiar relations

$$\nu = \frac{c}{\lambda} = \frac{kc}{2\pi} \tag{1.3.5}$$

and

$$k = |\mathbf{k}| = (k_x^2 + k_y^2 + k_z^2)^{1/2} \tag{1.3.6}$$

to obtain the following set of allowed frequencies for a cubical cavity:

$$\nu = \nu_{lmn} = \frac{c}{2L}(l^2 + m^2 + n^2)^{1/2} \tag{1.3.7}$$

An important distinction between masers and lasers can be explained on the basis of this formula for cavity frequencies. In the case of a maser (which typically operates at macroscopic wavelengths, $\lambda \approx 1$ cm) the cavity itself can be constructed on the scale of the wavelengths of interest. Consider, for example, a cubical cavity with

$$L = 1 \text{ cm}$$

From (1.3.7), then, the wavelengths of the cavity modes are

$$\lambda_{lmn} = \frac{c}{\nu_{lmn}} = \frac{2}{\sqrt{l^2 + m^2 + n^2}} \text{ cm} \tag{1.3.8}$$

Thus the wavelengths of the principal (lowest) modes of the cavity are

$$\lambda_{110} = \lambda_{101} = \lambda_{011} = \frac{2}{\sqrt{2}} \text{ cm} \approx 1.41 \text{ cm}$$

$$\lambda_{111} = \frac{2}{\sqrt{3}} \text{ cm} \approx 1.15 \text{ cm}$$

$$\lambda_{210} = \lambda_{201} = \lambda_{012} = \lambda_{021} = \lambda_{102} = \lambda_{120} = \frac{2}{\sqrt{5}} \text{ cm} \approx 0.89 \text{ cm}$$

$$\lambda_{211} = \lambda_{121} = \lambda_{112} = \frac{2}{\sqrt{6}} \text{ cm} \approx 0.82 \text{ cm}$$

and so on. (We do not list wavelengths of "modes" for which any two of the integers l, m, n are zero, because the field then vanishes. See Chapter 2.)

The point of this explicit catalog of modes is to show that the mode wavelengths are all on the order of the cavity size, $L = 1$ cm, and that they are also separated from each other by amounts not too much smaller than the cavity size. (See Figure 1.11.) If we want to have coherent—ideally single-mode—maser emission in the wavelength range around 1 cm, such a cavity is just right. Even if the maser gain medium is fairly broadband, able to amplify radiation over an unusually wide band of frequencies, say $\Delta\lambda/\lambda = 0.01$ around $\lambda = 1.41$ cm, the cavity will nevertheless give single-mode output at $\lambda = 1.41$ cm, because there are no other modes within 1% of the 1.41-cm mode. The distance between the modes is too large for any other mode to experience significant gain. In the case of microwave wavelengths, therefore, single-mode oscillation is easily accomplished by a propitious choice of the cavity dimensions.

In the case of optical or near-infrared *laser* wavelengths, however, this obviously cannot be done very easily. Optical cavity dimensions are enormously larger than an optical wavelength ($\lambda \approx 5 \times 10^{-5}$ cm). In this case, Eq. (1.3.7)

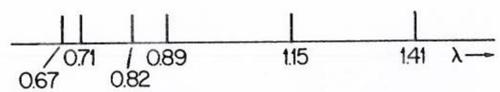

Figure 1.11 Wavelengths corresponding to the lowest modes of a cavity with $L_x = L_y = L_z = 1$ cm.

implies there can be a very large number of modes within the frequency range for which the active medium amplifies the field. The output laser radiation will then contain contributions from many different modes, differing in both their frequencies and their spatial distributions. Furthermore the relative proportions of the various modes might vary in time in a practically unpredictable way as a result of various perturbations such as mirror vibrations. Such unpredictable small variations would make the laser much ''noisier'' than we would like for many important applications.

It is not difficult to calculate the number of modes within a given frequency range when the cavity dimensions are much larger than the wavelength. The calculation is given in the next chapter. It is found that the number ΔN_ν of possible field modes (of a cavity of volume V) in the frequency interval from ν to $\nu + \Delta\nu$ is

$$\Delta N_\nu = \frac{8\pi\nu^2}{c^3} V \Delta\nu \tag{1.3.9}$$

Conversely, the interval $\Delta\nu$ between adjacent modes (i.e., for $\Delta N_\nu = 1$) is given by

$$\Delta\nu = \frac{c^3}{8\pi\nu^2 V} = \left(\frac{\lambda^3}{V}\right)\frac{\nu}{8\pi} \tag{1.3.10}$$

The modes are obviously extremely closely spaced if λ is optical and V is macroscopic. In a typical case ($\lambda \approx 600$ nm and $V \approx 1\ \text{cm}^3$) we have

$$\Delta\nu \approx \frac{100}{8\pi}\ \text{Hz} \approx 4\ \text{Hz} \tag{1.3.11}$$

Consider an example. In a 6328-Å He–Ne laser ($\nu = 4.7 \times 10^{14}$ Hz) the width of the gain curve (the frequency region over which stimulated emission is feasible) is about 1500 MHz. In this frequency interval a 1-cm^3 cavity has about 400 million modes available. A laser using this cavity would not produce anything resembling single-mode radiation.

It would appear from (1.3.10) and (1.3.11) that an extremely large number of modes must be present at optical frequencies. This is not the case. One way to reduce the mode density is to reduce the *dimensionality* of the cavity. A way to do this was suggested independently by Schawlow and Townes, R.H. Dicke, and Prokhorov in 1958: use an *open* resonator consisting of two parallel mirrors, as in Figure 1.1*d*.

The idea behind this suggestion is intuitively clear. A wave traveling at an angle to the axis joining the two mirrors will escape from the resonator after a few reflections. It does not represent a cavity mode. Also, it will not complete many

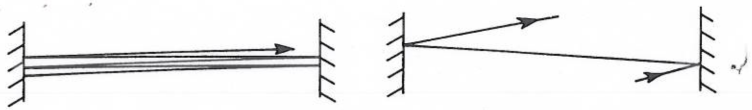

Figure 1.12 Sketch illustrating the advantage of a one-dimensional cavity. Stable modes are associated only with beams that are retroreflected many times.

round-trip passes through the gain medium, and therefore will not have a chance to be amplified very much. A wave traveling nearly exactly along the axis, however, will continue to bounce back and forth between the mirrors. It can set up a standing-wave pattern and does represent a cavity mode. Also, such a wave mode continues to extract energy from the gain medium as a result of stimulated emission. If the gain is large enough, exceeding a certain threshold value, laser radiation will build up in the form of such waves traveling back and forth along the axis. This is illustrated in Figure 1.12.

The shift to a one-dimensional cavity has a dramatic effect on mode spacing. Roughly speaking, in (1.3.10) the factor λ^3/V can be interpreted as $(\lambda/L)^3$, where L is a typical cavity side length. In a one-dimensional cavity we can anticipate that λ/L appears in the appropriate mode-spacing formula instead of $(\lambda/L)^3$. Thus we expect a mode spacing on the order of $\Delta\nu \approx (\lambda/L)(\nu/8\pi)$. A careful analysis (see Chapter 2) confirms this estimate with a slight change in the numerical factor 8π. It is found that in a one-dimensional cavity the successive modes are separated in frequency by

$$\Delta\nu = \frac{c}{2L} = \frac{\lambda}{L}\frac{\nu}{2} \tag{1.3.12}$$

which gives, in our He–Ne example, if we take $L = 50$ cm as is typical,

$$\Delta\nu = \frac{3 \times 10^{10}\ \text{cm/sec}}{(2)(50\ \text{cm})} = 300\ \text{MHz} \tag{1.3.13}$$

for the separation in frequency of adjacent resonator modes. As indicated in Figure 1.13, the number of possible frequencies that can lase is therefore at most

$$\frac{1500\ \text{MHz}}{300\ \text{MHz}} = 5$$

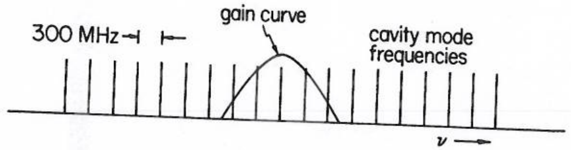

Figure 1.13 Mode frequencies separated by 300 MHz corresponding to a 50-cm one-dimensional cavity. A 1500-MHz gain curve overlaps only five modes.

The maximum number of possible modes, including polarization, is therefore 10, considerably smaller than the estimate of 4×10^8 modes obtained for three-dimensional cavities.

These results do not include the effects of diffraction of radiation at the mirror edges. Diffraction determines the x, y dependence of the field, which we have ignored completely. Accurate calculations of resonator modes, including diffraction, are often done with computers. Such calculations were first made in 1961 for the plane-parallel resonator of Figure 1.12 with either rectangular or circular mirrors. Actually lasers are seldom designed with flat mirrors. Laser resonator mirrors are usually spherical surfaces, for reasons to be discussed in Chapter 14. A great deal about laser cavities can nevertheless be understood without worrying about diffraction or mirror shape. In particular, *for most practical purposes, the mode-frequency spacing is given accurately enough by* (1.3.12).

1.4 LIGHT EMISSION AND ABSORPTION IN QUANTUM THEORY

The modern interpretation of light emission and absorption was first proposed by Einstein in 1905 in his theory of the photoelectric effect. Einstein assumed the difference in energy of the electron before and after its photoejection to be equal to the energy $h\nu$ of the photon absorbed in the process.

This picture of light absorption was extended in two ways by Bohr: to apply to atomic electrons that are not ejected during photon absorption but instead take on a higher energy within their atom, and to apply to the reverse process of photon emission, in which case the energy of the electron should decrease. These extensions of Einstein's idea fitted perfectly into Bohr's quantum-mechanical model of an atom in 1913. This model was the first to suggest that electrons are restricted to a certain fixed set of orbits around the atomic nucleus. This set of orbits was shown to correspond to a fixed set of allowed electron energies. The idea of a "quantum jump" was introduced to describe an electron's transition between two allowed orbits.

The amount of energy involved in a quantum jump depends on the quantum system. Atoms have quantum jumps whose energies are typically in the range 1–6 eV, as long as an outer-shell electron is doing the jumping. This is the ordinary case, so atoms usually absorb and emit photons in or near the optical region of the spectrum. Jumps by inner-shell atomic electrons require much more energy and are associated with X-ray photons. On the other hand, quantum jumps among the so-called Rydberg energy levels, those outer-electron levels lying far from the ground level and near to the ionization limit, involve only a small amount of energy, corresponding to far-infrared or even microwave photons.

Molecules have vibrational and rotational degrees of freedom whose quantum jumps are smaller (perhaps much smaller) than the quantum jumps in free atoms, and the same is often true of jumps between conduction and valence bands in semiconductors. Many crystals are transparent in the optical region, which is a sign that they do not absorb or emit optical photons, because they do not have quantum energy levels that permit jumps in the optical range. However, colored crystals such as ruby have impurities that do absorb and emit optical photons. These impurities are frequently atomic ions, and they have both discrete energy levels and broad bands of levels that allow optical quantum jumps (ruby is a good absorber of green photons and so appears red).

1.5 EINSTEIN THEORY OF LIGHT-MATTER INTERACTIONS

The atoms of a laser undergo repeated quantum jumps and so act as microscopic transducers. That is, each atom accepts energy and jumps to a higher orbit as a result of some input or "pumping" process, and converts it into other forms of energy—e.g., into light energy (photons)—when it jumps to a lower orbit. At the same time, each atom must deal with the photons that have been emitted earlier and reflected back by the mirrors. These prior photons, already channeled along the cavity axis, are the origin of the stimulated component to the atom's emission of subsequent photons.

In Figure 1.14 we indicate some ways in which energy conversion can occur. For simplicity we focus our attention on quantum jumps between two energy levels, 1 and 2, of an atom. The five distinct energy-conversion diagrams of Figure 1.14 are interpreted as follows:

a. Absorption of an increment $\Delta E = E_2 - E_1$ of energy from the pump; the atom is raised from level 1 to level 2. In other words, an electron in the atom jumps from an inner orbit to an outer orbit.

b. Spontaneous emission of a photon of energy $h\nu = E_2 - E_1$; the atom jumps down from level 2 to the lower level 1. The process occurs "spontaneously" without any external influence.

c. Stimulated emission; the atom jumps down from energy level 2 to the lower level 1, and the emitted photon of energy $h\nu = E_2 - E_1$ is an exact replica of a photon already present. The process is induced, or stimulated, by the incident photon.

d. Absorption of a photon of energy $h\nu = E_2 - E_1$; the atom jumps up from