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LASERS IN THE VACUUM ULTRAVIOLET AND IN THE X-RAY REGIONS OF THE SPECTRUM

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THE development of generators and amplifiers in the short-wave region of the spectrum, beyond the region of the visible light, is at present one of the promising trends in the development of quantum electronics, which is finding extensive applications in the control of chemical reactions, in holographic research on microscopic objects, and in investigations of the interaction of high-energy photons with matter. The recent first reports[1-5] of the development of lasers in the vacuum ultraviolet (UV) region of the spectrum (wavelength $\lambda < 1850$ Å) have confirmed the theoretical predictions[6,7] of the effectiveness of using, in this wavelength range, resonant transitions of the molecular type in condensed and gaseous phases. Amplification and generation of electromagnetic radiation were effected in [1,2,5] at ~1760 Å by exciting condensed xenon with a beam of fast electrons, while in [3,4] they registered amplification of radiation on electron-vibrational transitions of the hydrogen molecule, in the band from 1567 to 1613 Å, in an electric discharge in gas. These lasers had the shortest wavelength attained so far. Both in the case of condensed xenon and in the case of the hydrogen molecule, the transitions employed were those in which the minimum of the potential curve of the upper excited state did not coincide with the minimum of the potential curve of the lower ground state. When the Franck-Condon principle is taken into account, this leads to a relatively easy attainment of inverted population, and also to amplification of the radiation if the excitation rate of the upper level is high enough. In spite of this similarity, the processes that lead to laser action are essentially different from each other in these two cases and will be considered separately. In addition, we shall consider certain other still unrealized possibilities of obtaining generation in the vacuum ultraviolet and in the x-ray regions of the spectrum.

1. GENERATION OF VACUUM ULTRAVIOLET RADIATION IN CONDENSED INERT GASES

The first experimental investigations[8,9] of the luminescence spectra of the inert gases Ne, Ar, Kr, and Xe in the crystalline state and those in the liquid state, which are quite similar, revealed, in the vacuum ultraviolet region, emission bands located at much lower energies than the peaks of absorption by free excitons. This shift is due to the strong deformation of the crystal lattice, as a result of which the free exciton is captured by the lattice, forming an exciton complex R_2^*, which is similar in many respects to the diatomic molecule of the inert gas in the excited state. Capture of the exciton by the lattice is not connected with impurities, and is the result of a configuration instability of the degenerate band states of the perfect crystal relative to asymmetrical deformation of the crystal lattice.[10] As shown in [11], the development of such an instability and the transition of the free exciton into a self-captured state R_2^* occurs within a time $\sim 5 \times 10^{-12}$ sec, which is much shorter than the radiative lifetime $\tau_S \sim 10^{-9}$ sec of the free exciton, so that in the absence of impurities the luminescence in the vacuum UV region of the spectrum is connected mainly with radiative decay of the self-captured excitons. A similar property of self-capture in crystals of inert gases is possessed by free holes, which form, as a result of lattice deformation, charged centers R_2^+ analogous to diatomic molecular ions. Therefore the self-captured excitons can be produced not only from free excitons, but also by direct capture of an electron by an R_2^+ center. The lower levels of the self-captured excitons come from the ns series (n = 1, 2, 3, . . .) of the dipole-allowed free excitons Γ_{15} and form a hydrogen-like nσ series, the potential curves of which duplicate in form the potential curve of the self-captured hole. Below this series is located the potential curve of the self-captured exciton, corresponding to the molecular state $^3\Sigma_u^+$, which stems from dipole-forbidden free-exciton bands. Γ_{25} and Γ_{12}'. In spite of the fact that these exciton bands are not observed in the absorption spectra, the corresponding self-captured excitons make an appreciable contribution to the luminescence, since the transition forbidden in absorption becomes allowed in emission, owing to the strong asymmetrical deformation of the lattice. The overall picture of the dependence of the exciton energy on the lattice deformation parameter r, which is determined by the distance between the two nuclei in the self-captured hole, is shown for the case of crystalline Xe in Fig. 1 (for details see [12]).

We present here the potential curves of excited molecular states[13] and the scheme for populating the levels of the self-captured excitons from the free-exciton bands of the conduction band. The principal luminescence bands in inert-gas crystals in the vicinity of the vacuum UV are connected with the radiative transitions from the states 1σ, 2σ, and $^3\Sigma_u^+$ to the ground repulsion state $^1\Sigma_g^+$. Raising the temperature and the transition from the crystal to the liquid has little effect on the general shape of the potential curves, leading only to a broadening and to a shift of the exciton levels with $n \geq 2$, but a noticeable redistribution of the intensity among the bands occurs in this case in the emission spectrum,[1] owing to the increase of the rate of energy exchange between the upper and lower states of the self-captured excitons and to the temperature dependence of the quantum yield. At high temperatures one usually observes only one long-wave band connected with the $^3\Sigma_u^+$

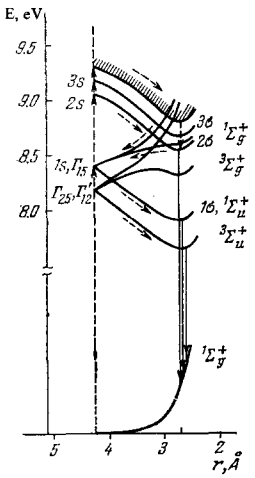

FIG. 1. Dependence of the exciton energy on the lattice deformation parameter and luminescence scheme in crystalline xenon.

$\rightarrow {}^1\Sigma_g^+$ transition (see Fig. 1). The generation condition for this transition can be calculated in accordance with a four-level scheme, since the time of relaxation from the level as the result of the repulsion of the atoms in the ground state is of the order of $\sim 10^{-12}$ sec, which is much shorter than the spontaneous lifetime of the upper level. This circumstance makes it possible to realize in principle a quasicontinuous generation regime in condensed inert gases, unlike in the molecular-hydrogen laser considered below, where the lower working levels are metastable.

For a resonator with flat mirrors, the threshold condition for the start of generation is in the general case

$$R_1 R_2 \exp[2(\alpha-\beta)L] = 1, \quad (1)$$

where R_1 and R_2 are the mirror reflection coefficients, L is the length of the resonator filled with the active matter, and α and β are the gain and the absorption coefficient. In condensed inert gases, the emission bands have a Gaussian shape, and the gain is therefore equal to

$$\alpha = \sigma_0 N_{ex} = \frac{1}{4\pi n^2} \sqrt{\frac{\ln 2}{\pi}} \frac{\lambda^2}{\tau_s \Delta\nu} N_{ex}, \quad (2)$$

where λ is the wavelength corresponding to the center of the band, τ_S is the spontaneous lifetime, N_{ex} is the concentration of the self-captured excitons, $\Delta\nu$ is the width of the emission band, and n is the refractive index. One of the features of the considered case of threshold generation conditions is that the losses increase with increasing excitation power. This increase is connected with photoionization of the self-captured excitons and absorption by the self-captured holes. This causes the absorption coefficient to depend linearly on the concentration of the self-captured excitons:

$$\beta = \beta_0 + \sigma_f N_{ex},$$

and we obtain for the threshold concentration of the excitons, with allowance for (1) and (2), the following expression:

$$N_{ex} = \frac{\beta_0 - (\ln R_1 R_2/2L)}{\sigma_0 - \sigma_f}. \quad (3)$$

Although the relation $\sigma_f \ll \sigma_0$ is usually satisfied, the coefficient σ_f, which takes into account the absorption by the self-captured excitons and the nuclei, can turn out in some cases to be equal to or larger than the cross section σ_0 of the working transition, and then generation can be obtained only at very high excitation levels on the interband transitions.[6]

When a condensed inert gas is excited by a beam of monochromatic electrons of energy E_0, the concentration of the self-captured excitons participating in the generation is connected with the current density j in the beam by the relation

$$N_{ex} = \frac{\tau_s \eta}{eI_i} \frac{E_0}{x_0} j, \quad (4)$$

where η is the quantum yield of the emission band under consideration, I_i is the energy necessary to produce one electron and hole pair, e is the electron charge, and x_0 is the effective depth of excitation connected with the range of the electron R and the density of matter ρ by the approximate relation $x_0 \approx R/5\rho$. For energies $E_0 \lesssim 1$ MeV the electron range can be represented by the following expression, which goes over into the Widdington formula at low energies:

$$R = \frac{1374}{24\pi a_0^2 N_A} \frac{A}{Z} \frac{E_0^2}{mc^2 (E_0 + mc^2)} = 0.274 \frac{A}{Z} \frac{E_0^2}{mc^2 (E_0 + mc^2)}, \quad (5)$$

where A is the atomic weight and Z the atomic number of the substance, N_A is Avogadro's number, a_0 is the Bohr radius, and mc^2 is the rest mass of the electron.

Equating the right-hand sides of (3) and (4) we obtain, with allowance for (5), the current density needed to initiate the generation. Under the experimental conditions,[5] when λ = 1760 Å, $\Delta\nu = 1.4 \times 10^{14}$ sec^{-1}, L = 1 cm, $R_{1,2} = 0.5$, $\beta_0 = 0.15$ cm^{-1}, and E_0 = 600 keV, using the numerical values of the parameters that enter in relations (1)–(5), namely n = 1.5, I_i = 22 eV, $\rho = 3.5$ g/cm^3, $\tau_S \approx 5 \times 10^{-9}$ sec, and $\eta \approx 0.1$ for liquid xenon, we obtain $x_0 = 240\ \mu$, $\sigma_0 = 5 \times 10^{-18}$ cm^2, $N_{ex} = 1.2 \times 10^{17}$ cm^{-3}, and the threshold current density is j $\approx$ 30 A-cm^{-2}.

The threshold current densities in the beam for other condensed inert gases turn out to be even higher.

The changeover from the amplification regime[1,2] to the generation regime was effected in [5] by using a powerful pulsed electron gun which yielded an electron current of density up to 300 A/cm^2 at an electron energy up to 1 MeV, with a current pulse duration $\sim 10^{-8}$ sec. The walls of the flat resonator in this study were aluminum mirrors sputtered on a lithium-fluoride substrate and covered with a protective coating of magnesium fluoride. The generation threshold, attained at current densities j $\sim$ 30–60 A/cm^2, was revealed by the narrowing of the 1760 Å line from a width $\sim$150 Å at low excitation powers to a value close to the spectrometer resolution limit $\sim$17 Å, and also by the narrowing of the angular directivity pattern of the radiation to a value $\sim 7°$. The shape and duration of the radiation pulse corresponded in this case to the shape and duration of the current pulse, thus demonstrating that the generation regime was quasicontinuous. In addition to liquid xenon, the described method can be used to produce lasing in other inert gases and their mixtures, both in the condensed state and in the high-pressure gas. In this case, by going over to Ne and He, one can hope in principle to extend the lasing wavelength to $\sim$600 Å.

2. AMPLIFICATION OF RADIATION ON RESONANT TRANSITIONS IN MOLECULAR HYDROGEN

A general scheme of obtaining population inversion between excited electronic states and the ground state

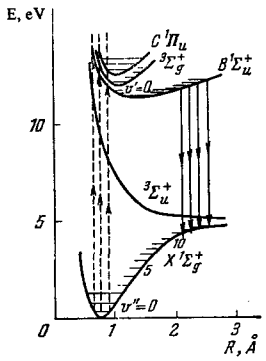

FIG. 2. Lower potential curves and transitions on which amplification of the radiation is observed in the H_2 molecule.

in diatomic molecules, for transitions lying the vacuum UV region of the spectrum, was proposed in [7]. In the simplest case of hydrogen molecules, such transitions include, first, $C^1\Pi_u \rightarrow X^1\Sigma_g^+$, the possible use of which was considered in [7,14], and $B^1\Sigma_u^+ \rightarrow X^1\Sigma_g^+$ (Fig. 2), on which inverted population was actually obtained and amplification of the radiation observed.[3,4] Inversion is produced for these transitions because the collisions between electrons and hydrogen molecules at the zeroth vibrational level $v'' = 0$ of the ground state cause predominant excitation of the vibrational levels $v' \geq 1$ of the dipole-allowed states $C^1\Pi_u$ and $B^1\Sigma_u^+$, and transitions to the upper vibrational levels of the ground state have a low probability, since they occur without a change of the electronic state. Owing to the shift of the potential curves in accordance with the Franck-Condon principle, the vibrational levels that are effectively populated are $v' = 1-4$ of the state $C^1\Pi_u$ and $v' = 3-7$ of the state $B^1\Sigma_u^+$. From these levels, besides the inverse transitions to the ground state on the lower vibrational level, a large probability is possessed by transitions to the upper vibrational levels $v'' = 1-8$ and $v'' = 10-13$, relative to which the states $C^1\Pi_u$ and $B^1\Sigma_u^+$ can turn out to be inversely-populated at the initial instants of the electric discharge. Such an inverted level population and amplification of radiation were attained in [3,4] on the transition $B^1\Sigma_u^+$ ($v = 1-4$) $\rightarrow X^1\Sigma_g^+$ ($v'' = 10-13$) (see Fig. 2) corresponding to the Lyman band in the wavelength interval from 1567 to 1613 Å.

To obtain a powerful pulse discharge in hydrogen, a high-voltage transmission line was used in these investigations, in the form of plane-parallel metallic plates between which the discharge gap was placed. Such a device, analogous to that employed by Shipman to excite generation of light in nitrogen,[15] makes it possible to attain within ~3 nsec currents of hundreds of kiloamperes in the discharge gap at a voltage ~100 kV. In [3] the discharge gap was a channel measuring $120 \times 1.2 \times 0.04$ cm, in which the transverse discharge was initiated simultaneously over the entire length, so that the radiation could be amplified in both directions along the channel. In this case the main proof of the amplification of the radiation in the ~1600 Å region was the anomalously large intensity of the P-branch in comparison with the R-branch for the observed vibrational-rotational transitions, and also the large peak radiation power, which at an optimal pressure of 60 Torr reached a value ~1.5 kW at a pulse duration ~2 nsec. The duration of the inverted population and the gain were so high that stimulated emission could be observed in this case without the use of mirrors.

In another experiment,[4] discharge in hydrogen was initiated from one end of a channel measuring $120 \times 1.2 \times 0.3$ cm. The discharge front propagated with the speed of light in the gas toward the other end of the channel, producing excitation and inverted-population waves traveling with the same velocity and in the same direction. Such a method of exciting gas produced, in addition to the effects observed in [3], also an uneven radiation power from both ends of the channel. This effect is connected with the preferred amplification of the photons that fall on the front of the traveling wave, where maximum inverted population is maintained over the entire duration of the electric discharge. In this case, at a pulse duration ~10^{-9} sec, the peak power of the radiation in the direction of the propagation of the excitation wave reached 100 kW, which is larger by one order of magnitude than the radiation power in the opposite direction. It should be noted in conclusion that the described method of generating radiation on resonant transitions in molecular hydrogen at a sufficiently short current pulse or when short laser pulses are used for the breakdown[16] should have a high efficiency, since the greater part of the energy transferred by the electron to the upper working level is usefully employed in the form of a quantum of stimulated radiation.

3. CERTAIN POSSIBILITIES OF PRODUCED VACUUM UV AND X-RAY LASERS

In addition to the cases considered above, searches are being carried out at present of other possibilities of attaining generation in the vacuum UV, and also in the x-ray regions of the spectrum, among which we can single out the following main trends:

1) excitation of the internal shells of the atoms,
2) use of transitions in multiply-charged ions,
3) the conversion of laser light in induced Compton scattering,
4) investigation of stimulated transitions in nuclei.

The possibility of attaining generation on characteristic emission lines of atoms was considered in [17-19]. In this case, following selective removal of one electron from the internal shell, say from the K shell of the Mo atom, the produced excited state of the atom with hole in the K shell is inversely populated relative to the transition of the hole to the L-shell levels (Fig. 3). Such transitions are accompanied by emission of the characteristic lines $K_{\alpha 1}$ and $K_{\alpha 2}$, on which it is possible to attain the amplification threshold in the case of a sufficiently intense and steep excitation pulse.

Selective removal of an electron can be realized in photoionization of the K shell by filtered bremsstrahlung or by the characteristic radiation of another element, in this case Rh, the characteristic lines of which fall in the K absorption edge of Mo (see Fig. 3). The lifetime of the hole in the L shell is usually larger than the lifetime in the K shell, and therefore here, just as in the case of molecular hydrogen considered above, where the lower working levels turn out to be metastable, amplification is possible only in the traveling-wave regime. The latter limitation is lifted, however, for elements with $Z \leq 36$ (for the $K_{\alpha 1}$ lines) and with

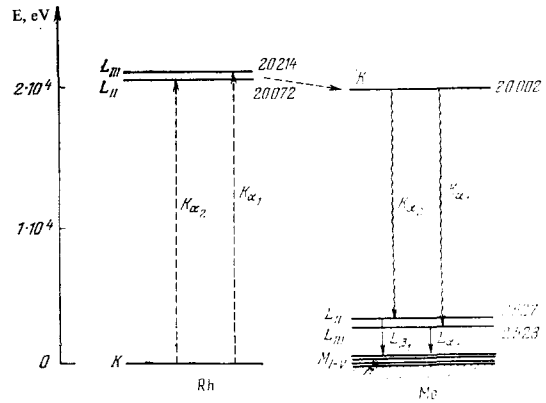

FIG. 3. X-ray levels, and the transitions between them, for the K and L series of the characteristic radiation in the Mo atom. On the left are shown, in inverse order, the levels of Rh, whose characteristic K_{α_1} and $K_{\alpha 2}$ lines fall in the K absorption edge of Mo.

$Z \leq 47$ (for $K_{\alpha 2}$ lines) at solid-state densities of matter, owing to the large probability of nonradiative departure of the hole from the L shell as a result of the Auger effect.[19] In this case, in a range of temperatures not exceeding 30–100 eV, at which the probability of nonradiative recombination on the valence levels is not less than that of the transition of a hole between the employed levels, it is possible to obtain a quasicontinuous amplification regime with multiple utilization of the same atom.

The power necessary to obtain amplification of the characteristic radiation in Na and Cu in the traveling-wave regime was estimated in [18]. In the case of Na, the hole transition $L_{III} \rightarrow M_I$ (λ = 372 Å) was considered, in which the spontaneous lifetime of the upper level was $\tau_S = 4 \times 10^{-10}$ sec. In this case the maximum amplification was obtained on the linearly-growing part of the front of the exciting pulse, at a time $0.55\tau_S$ after the pulse was turned on, and amounted to 40 dB/m at a pump power rate 4×10^{18} W/cm^2 sec. The active region was chosen in the form of a volume of dimensions $1 \times 1 \times 500$ cm filled with sodium vapor at a pressure 0.02 Torr and a temperature 310°C. At such parameters, which exclude the population of the lower working level by impact ionization, the efficiency of the utilization of the external source, with a maximum radiation at 50 eV, is only 0.3%. Such a system should yield under saturation conditions an output power of ~ 3 kW/cm^3 at a laser pulse duration 10^{-10} sec, and a maximum possible power of 1.5 MW when the amplification length is increased to 5 m.

An analogous calculation in the case of Cu atoms for the transition $K \rightarrow L_{III}$ (λ = 1.537 Å, $\tau_S = 4.5 \times 10^{-16}$ sec) leads to much more stringent requirements for the radiation-exciting pulse.

At an average photon energy in the pulse ~ 12 keV, the power in it should increase within a time of 10^{-15} sec to a value 2.5×10^{10} W/cm^2. In this case, for a solid strip of copper measuring $1\,\mu \times 1\,\mu \times 5$ mm, in which 10% of the pump power is absorbed, the gain at the wavelength λ = 1.537 Å should reach 400 dB/mm, which exceeds by 200 dB/mm the losses connected with absorption at this wavelength. These examples show that one of the main difficulties standing in the way of reaching the generation threshold in the vacuum UV and in the x-ray regions of the spectrum lies in the need of greatly increasing the power and decreasing the pulse duration of the pump resources when the generation wavelength is decreased.

Recent progress in the development of high-current sources of fast electrons and x-ray flash generators[20] has made it possible to develop powers of $\sim 10^{12}$ W in the considered energy region within a time $\sim 10^{-8}$ sec. Such sources of hard radiation, together with powerful γ-ray sources from nuclear fission, can be used to excite large values of active laser media in a wide region of the spectrum.[21] Even steeper growth fronts of the pump power can be obtained in a laser spark,[22,23] where in the case of ultrashort pulses of light the developed power $\sim 10^{12}$ W is reached within a time $\sim 10^{-12}$ sec and can be used in principle to pump lasers of the short-wave band. The pump sources in this band can also be proton beams, which can excite by charge exchange on alkali-element atoms, depending on the energy, generation on resonant transitions of hydrogen or in alkali ions.[24]

Another difficulty, connected with the absence of good reflecting surfaces for mirrors in the short-wave band, was overcome to a considerable degree in [25-27], where resonator schemes based on Bragg reflection were proposed for x rays. In the usual scheme of a flat ring resonator, the limitations connected with the Bragg condition and with the condition that the path along which the rays propagate be closed, have practically excluded any possibility of tuning the resonator to a given wavelength. In the indicated investigations, these limitations were eliminated in one case by using non-planar geometry of the ray path, and in the other by the possibility of self-intersection of the rays. One of the schemes of a flat tunable resonator with self-intersecting beams is shown in Fig. 4, where the Bragg angle θ need not be an integer fraction of 2π and depends on the rotation of the pairs of parallel planes relative to one another.[27] The requirement that the individual pairs of the crystal planes be parallel is easy to satisfy if the mirrors are parts of one and the same single crystal. Estimates show that the losses in such resonators do not exceed the losses in ordinary resonators for the visible band with plane-parallel mirrors.

In addition to the foregoing transitions between the excited levels of the internal shells of the atoms, the emission lines connected with transitions between excited levels of multiply-charged ions also fall in the short-wave region of the spectrum.

For ions with low ionization multiplicity, there are presently known several hundred transitions on which

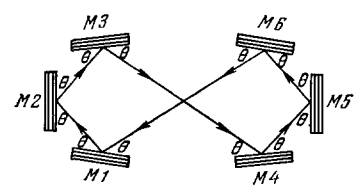

FIG. 4. Diagram of flat resonator for x rays. The beam is reflected at the Bragg angle θ from the pairwise-parallel crystal planes M1 and M6, M2 and M5, and M3 and M4.

generation in gas ion lasers was obtained in the visible and UV regions.[28] Among these transitions, the shortest wavelength, $\lambda = 2358$ Å, was obtained with the transition $3p^4D^0_{7/2} \rightarrow 3s^4P_{5/2}$ of triply-ionized neon.[29] Further progress into the short-wave region in ionic gas lasers entails difficulties of producing a plasma with sufficiently high ionization multiplicity in an ordinary electric discharge. The use of beams of fast electrons and of laser radiation to heat the plasma makes it possible to obtain a plasma in which the ion multiplicity reaches several times ten.[30-32] Transitions between excited levels of such ions lie in the vacuum UV and in the soft x-ray regions of the spectrum, and have so far not been sufficiently fully classified. Therefore, to predict the transitions and wavelengths on which lasing can be obtained it is natural to consider primarily multiply-charged ions connected by an isoelectronic sequence with ions of low multiplicity, on which generation was already observed. In ionic lasers, generation is usually observed for np → ns transitions, which occur without a change of the principal quantum numbers,[28] and therefore for these transitions the energy of the emission quanta in the isoelectronic sequence is a linear function of the ionization multiplicity.[33] By fixing two ions with known transitions in this sequence, we can therefore obtain the approximate values of the wavelengths for the entire isoelectronic series. Thus, for example, generation on the transition $(^3P^0)3p^4D_{5/2} \rightarrow (^3P^0)3s^4P^0_{3/2}$ was observed in N III at a wavelength 4510.8 Å and in O IV at a wavelength 3381.3 Å.[28] The following wavelengths values (A) are therefore obtained in the B I isoelectronic sequence, which includes the ions N III and O IV: 6763 C II, 4510.8 N III, 3381.3 O IV, 2705 F V, 2254 Ne VI, ..., 1503 Al IX, ..., 966 Ar XIV, ..., 615 Fe XXII, ..., 422 Kr XXXII, ... For the transition $(^3P)3p^4D^0_{7/2} \rightarrow (^3P)3s^4P_{5/2}$, the generation lines in O II and Ne IV give the following isoelectronic sequence of wavelengths: 4649 O II, 3129 F III, 2358 Ne IV, ..., 1356 Al VII, ..., 794 Ar XII, ..., 477 Fe XX, ..., 318 Kr XXX, ...

For the transition $(^2S)3p^3P^0_2 \rightarrow (^2S)3s^3S_1$ in the isoelectronic sequence of the atom Be I, where the generation was observed only in triply-ionized nitrogen ($\lambda = 3478.67$ Å), the remaining series is reconstructed from the known wavelength 2316.1 Å of this transition in five-fold ionized fluorine:[34] 6989 Be II, 4644 C III, 3478 N IV, 2781 O V, 2316 F VI, 1984 Ne VII, 1388 Al X, ..., 925 Ar XV, ..., 603 Fe XXIII, ..., 420 Kr XXXIII, ... Such wavelength sequences, the accuracies of which lie within several Angstrom units, can be constructed for a large number of isoelectronic series and make it possible to predict beforehand the most probable laser transitions in multiply-charged ions. It should be noted that whereas in ordinary ion lasers operating with ions of low multiplicity the efficiency reaches at best 0.1%, the efficiency for the transitions considered here should be even smaller, since the ratio of the quantum energy of the useful radiation to the total energy needed to produce the multiply-charged ion decreases with increasing charge. A larger efficiency can be obtained for the transitions (n + 1)s → np, which proceed with a change of the principal quantum number. Generation on such transitions can be realized in plasma lasers,[35-37] in which the upper working level is filled by recombining electrons at high densities N_e and relatively low temperatures T_e. In this case, for example, for the 3s → 2p transition in the ion Be II ($\lambda = 1776$ Å)[36] at values $N_e = 10^{17}$ cm^{-3}, $N_{Be\,II} = 10^{18}$ cm^{-3}, and $T_e = 8 \times 10^3$°K the gain can reach very large values, $\sim 5 \times 10^2$ cm^{-1}.

In addition to the transitions considered above in the discrete spectrum, generation in the short-wave region can be obtained with transitions in the continuous spectrum, by using stimulated Compton scattering of laser light by a beam of relativistic electrons.[38-40]

In this case the increase of the frequency of the incident light is due to the energy of the electrons in the beam. In the simplest case of a collinear collision, the incident laser beam is directed opposite to the motion of the electrons, the frequency of the scattered photons ν_2 is connected with the frequency of the incident photons ν_1 by the relativistic Doppler formula

$$\nu_2 \approx 4\nu_1 \left(\frac{E}{mc^2}\right)^2, \tag{6}$$

where E is the energy of the electrons in the beam.

The effective gain for the scattered radiation is connected with the density ρ_f, of the photons in the incident beam and the density ρ_e of the electrons in the beam by the following relation:[40]

$$\alpha \approx 0.7 r_0^2 \frac{E}{\Delta} \frac{h\nu_2}{\Delta} \lambda_1 \lambda_2^2 \rho_e \rho_f, \tag{7}$$

where Δ is the width of the energy scatter of the electrons in the beam, r_0 is the classical radius of the electron, and λ_1 and λ_2 are the wavelengths of the incident and scattered radiation. Substituting in (6) and (7) typical numerical values of the parameters of a neodymium-laser beam, namely $h\nu_1 = 1.17$ eV and $\rho_f = 1.8 \times 10^{22}$ cm^{-3}, and of the parameters of the electron beam, namely E = 2 MeV, $E/\Delta = 10^5$, and $\rho_e = 2 \times 10^{13}$ cm^{-3}, which corresponds to a current density $j \sim 10^5$ A/cm^2, we obtain for the scattered radiation $\lambda_2 \approx 166$ Å and a gain $\alpha = 2.2$ cm^{-1}. In addition to the relatively large gain in the short-wave band, lasers based on stimulated Compton scattering have also the advantage of being tunable, for in accordance with (6) their frequency is directly connected with the voltage in the electron accelerator.

The possibility of amplifying x-radiation on Mossbauer transitions in nuclei was investigated in [41-43]. Such a possibility is based on the fact that the cross section of stimulated emission at resonance $\sigma_0 \sim \lambda^2 \sim 10^{-18}$ cm^2 ($\lambda \sim 0.2$ Å, $h\nu \sim 50$ keV) greatly exceeds the cross section $\sigma_e \sim 10^{-22}$ cm^2 for photoabsorption on electrons in this region of the spectrum.

To start the amplification it is therefore sufficient to have the relative concentration of the active nuclei larger than a certain critical value $\sim \sigma_e/\sigma_0 \sim 10^{-4}$. However, the realization of a sufficiently narrow Mossbauer line, which is needed to prepare an isomer within ~1 hr, entails great difficulties connected with the elimination of the vanishingly small perturbations. In particular, as shown by estimates,[41] the excited nuclei should be located at the same level relative to the earth's gravitational field with accuracy to 10^{-6} cm, and the deviations of the temperature in the amplifying sample should not exceed 10^{-9}°K. In this connection, interest attaches to transitions with recoil, in which the nucleus acquires a kinetic energy after emission. In

this case the condition of inverted population is not obligatory,[44] for at a sufficiently low temperature the radiated photons cannot be absorbed on the same transition. It is necessary here that the Doppler width $\Delta\nu_D$ not exceed the frequency difference between the emitted and absorbed quanta:

$$\Delta\nu_D < \frac{h\nu^2}{2Mc^2},$$

where ν is the frequency corresponding to the difference between the energy levels and M is the mass of the nucleus. In spite of the fact that the emission and absorption lines do not coincide, the emitted quanta produce stimulated emission of the excited nuclei and can be amplified. In the case of positronium atoms[44] at a temperature 4°K, the generation in a resonator with Q = 10^5 sets in at an atom concentration $\sim 10^{16}$ cm^{-3}. Comparable concentrations of the excited nuclei can be obtained in (n, γ) reactions.

The examples considered above show that in the nearest time, in addition to the development of lasers for the vacuum UV band, it is realistic to expect generation to be attained in the soft x-ray region of the spectrum.

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