Effect of collisions on the intensities of the dielectron satellites of resonance lines of hydrogenlike ions

A. V. Vinogradov, I. Yu. Skobelev, and E. A. Yukov

P. N. Lebedev Physics Institute, USSR Academy of Sciences
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An explanation is offered for the anomalously high intensity, observed in a laser plasma, of the x–ray lines corresponding to the $2p^2\,{}^3P$-$1s2p\,{}^3P$ transition in multiply charged He-like ions. It is shown that in a plasma with $N_e > 10^{20}$ cm^{-3} the population of the $2p^2\,{}^3P$ levels is due to excitation by electron impact from the autoionization level $2s2p\,{}^3P$. The dependence of the intensity of these lines on N_e can be profitably used to determine plasma density, since there is no radiation capture in the case of dielectric satellites.

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1. INTRODUCTION

Dielectron satellites in x-ray spectra of multiply charged ions are the results of transitions from doubly excited states of ions having the preceding ionization multiplicity. The satellites in the spectra of H-like and He-like ions with charges $Z \gtrsim 10$ are almost as easy to register as the resonance lines themselves. Recently, systematic investigations and identification of the satellite lines of H-like and He-like ions have been carried out in the spectra of laser plasma and of the solar corona.[1–4] The interpretation of the intensities of the satellites[5–7] was carried out within the framework of the corona model for the populations[5,8,9] on the basis of relativistic calculations[8,10] of the probabilities of the radiative (A) and nonradiative (Γ) decays of the levels located beyond the boundary of the continuous spectrum. The satellite intensity I is in this case larger the large probability of the autoionization decay Γ, or, in other words, the faster the rate of dielectron recombination, which is proportional to Γ:

$$I=\frac{g}{g_1}4\pi^{3/2}a_0{}^3N_eN_{\mathrm{H}}(1s)\left(\frac{\mathrm{Ry}}{kT}\right)^{3/2}e^{-E/kT}\frac{A\Gamma}{\Gamma+\sum A}. \qquad (1)$$

Here $N_{\mathrm{H}}(1s)$ is the density of the H-like ions in the ground state, $g_1=2$ is the statistical weight of the ground state, $\alpha_0=0.53\cdot 10^{-8}$ cm is the Bohr radius, g is the statistical weight of the autoionization level, E is its energy reckoned from the ionization limit of the He-like ion, and N_e is the electron density. The probabilities of the radiative (A) and autoionization (Γ) decays of the $2l2l'\,{}^3P$ levels of the ion Mg XI for the group of satellites $2s2p^3P-1s2s^3S$ and $2p^2\,{}^3P-1s2p^3P$ are listed in Table I.

The values of Γ for different satellites can differ by a factor $\sim 10^2$.[11] It follows then from the corona model (1) that the analogous factor should be preserved also in the intensities of the satellites (primarily those for which $A \gtrsim \Gamma$). The last statement, however, contradicts the experimental data. Spectroscopic investigations of a plasma produced by neodymium and iodine lasers, carried out under different irradiation conditions and with different parameters of the heating radiation,[12–14] have shown that satellites with low values of Γ turn out to be much more intense than expected from the corona model.[15]

It is shown in the present paper that the cause of this discrepancy is the deviation from the corona model in a dense plasma. Namely, in the case of the Mg Xi ion at $N_e > 10^{20}$ cm^{-3} the population of the doubly excited states is caused not only by dielectron recombination, but also by collision-dominated transitions from the close-lying autoionization levels. As a result, at $N_e > 10^{20}$ cm^{-3} the ratio of the intensities of a number of satellites begins to depend on the electron density N_e. This suggests a new spectroscopic method for determining plasma density. An advantage of this method is that the radiation capture effects, which in general distort the intensity and waveform of the resonance lines themselves, do not play any role in the case of their dielectron satellites.

In the second section we consider the main excitation processes and obtain the populations of the autoionization levels of He-like ions in a tenuous and a dense plasma. The results are applied in the third section to the problem of laser-plasma diagnostics with the aid of the spectra of magnesium ions.

2. RELATIVE POPULATIONS OF DOUBLY EXCITED TRIPLET STATES OF He-LIKE IONS

The populations of doubly excited levels of He-like ions are determined by the balance equations (see the level scheme in Fig. 1)

$$N_i\Big(\Gamma_i+A_i+\sum_{k\neq i}W_{ik}\Big)=Q_iN_{\mathrm{H}}(1s)+\sum_{k\neq i}N_kW_{ki}, \quad i,k=1,\ldots,10, \qquad (2)$$

where N_i is the population of the i-th level, Γ_i, A_i, and Q_i are respectively the probabilities of its autoionization and radiative decay and the probability of dielectron capture on this level, $W_{ik}=N_eC_{ik}$, and C_{ik} is the rate of the $i-k$ transition in a collision with an electron. In the corona model (1) no account is taken of the terms

TABLE I.

Transition	$A\cdot 10^{-13}$, sec^{-1}	$\Gamma\cdot 10^{-13}$, sec^{-1}
$1s2s^3S_1-2s2p^3P_J$	1.22	1.36
$1s2p^3P_1-2p^2\,{}^3P_2$	0.598	0.22
$1s2p^3P_2-2p^2\,{}^3P_2$	1.82	0.22
$1s2p^3P_0-2p^2\,{}^3P_1$	0.813	0
$1s2p^3P_1-2p^2\,{}^3P_1$	0.609	0

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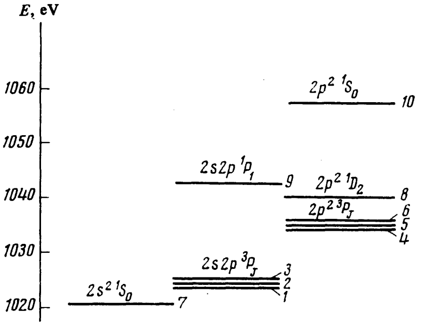

FIG. 1. Scheme of doubly excited levels $2l2l'$ of the Mg XI ion. The level energies are reckoned from the energy of the $1s$ level of the Mg XIII ion.

$$N_i\sum_{k\neq i} W_{ik}, \quad \sum_{k\neq i} N_k W_{ki},$$

which describe transitions between doubly excited states as a result of electron-ion collisions.

The satellites of the lines of H-like ions are intense only if their concentration is high enough, i.e., at $kT \gtrsim \frac{1}{4}Z^2\,\text{Ry}$. At these temperatures, the rates of the collision transitions within the singlet system and the triplet system is much higher than the rate C_{TS} of transitions between triplets and singlets.[16] Therefore at $N_e \ll N_0 = (\Gamma + A)/C_{TS}$ the balance equations for the triplet and singlet systems of levels can be written down independently. The satellites of interest to us start from the triplet levels $2s2p\,{}^3P_J$ and $2p^2\,{}^3P_J$. The balance equations for these levels are (the level designation is the same as in Fig. 1):

$$\begin{gathered} n_1(K_1+W)=q_1+n_5W, \quad n_2(K_2+W)=q_2+n_4\frac{W}{3}+n_5\frac{W}{4}+n_6\frac{5W}{12}, \\ n_3(K_3+W)=q_3+n_5\frac{W}{4}+n_6\frac{3W}{4}, \quad n_4(K_4+W)=q_4+n_2W, \\ n_5(K_5+W)=q_5+n_1\frac{W}{3}+n_2\frac{W}{4}+n_3\frac{5W}{12}, \quad n_6(K_6+W)=q_6+n_2\frac{W}{4}+n_3\frac{3W}{4} \end{gathered} \tag{3}$$

where $K_i = \Gamma_i + A_i$, $n_i = N_i/g_i$, $q_i = Q_i N_H(1s)/g_i$, $W = W_{15}$, and g_i is the statistical weight of the level i. In the system (3) we have neglected the quadrupole transitions between the sublevels of the fine structure, since the probability of such transitions is much less than the probability of the allowed dipole transitions, and in addition, we took into account a number of relations [see (4)] between the probabilities W_{ik}, relations that hold at $kT \gg |\Delta E_{ik}|$, where $\Delta E_{ik} = E_i - E_k$ is the energy difference between the levels i and k:

$$\begin{gathered} g_1W_{15}=g_5W_{51}=g_2W_{24}=g_4W_{42}, \quad g_2W_{25}=g_5W_{52}={}^3/_4W_{15}, \\ g_3W_{35}=g_5W_{53}=g_2W_{26}=g_6W_{62}={}^5/_4W_{15}, \quad g_3W_{36}=g_6W_{63}={}^{15}/_4W_{15}. \end{gathered} \tag{4}$$

Determining the values of n_4, n_5, and n_6 from the last three equations of the system (3) and substituting them in the first three, and also recognizing that for ions with charge $Z \sim 10-15$ the following relations are valid[17]:

$$K_1=K_2=K_3, \quad K_4=K_5\approx K_6, \quad q_1=q_2=q_3, \quad q_5=0, \quad q_4\ll q_1, \quad q_6\ll q_1, \tag{5}$$

we get

$$n_1=n_2=n_3=n=\frac{q_1}{K_1+K_6W/(K_6+W)}. \tag{6}$$

After this we obtain immediately

$$\begin{gathered} \frac{n_4}{n}=\alpha_4+\frac{x^2(1-\alpha_4)+x(1-\alpha_4+\beta_4)}{(1+x)^2}, \quad \frac{n_5}{n}=\frac{x}{1+x}, \\ \frac{n_6}{n}=\alpha_6+\frac{x^2(1-\alpha_6)+x(1-\alpha_6+\beta_6)}{(1+x)^2}, \end{gathered} \tag{7}$$

where

$$\alpha_i=\frac{\Gamma'_iK_1}{\Gamma_1K_i}, \quad \beta_i=\frac{\Gamma_i}{\Gamma_1}, \quad x=\frac{W}{K_6}=\frac{N_eC_{15}}{K_6}.$$

Formulas (7) describe both coronal and Boltzmann distributions of the populations within the triplet levels. In the corona limit $x \ll 1$ dielectron capture causes predominant population of the levels $2s2p\,{}^3P_J$ (levels 1, 2, 3) while the populations of the levels $2p^2\,{}^3P_J$ (levels 4, 5, 6) are small. With increasing density at $x \gtrsim 1$ the levels $2p^2\,{}^3P_J$ are now populated mainly be excitation from the levels $2s2p\,{}^3P_J$. At $x \gg 1$ the collisions establish the populations of all triplet levels in proportion to the statistical weights. The condition $x \sim 1$ means $N_e \sim (A+\Gamma)/C$, where $C \sim 10^2 C_{TS}$ is the rate of the $2s-2p$ transitions without change of spin. Thus, a Boltzmann distribution within the triplet levels is established already at $N_e \ll N_0 = (A+\Gamma)/C_{TS}$. The distribution among the triplet and singlet level still remains close to coronal in this case.

3. DETERMINATION OF PLASMA DENSITY FROM THE RELATIVE INTENSITIES OF SPECTRAL LINES OF DOUBLY EXCITED He-LIKE IONS

Radiative decay of the levels $2s2p\,{}^3P_J$ and $2p^2\,{}^3P_J$ leads to emission of seven most intense spectral lines with close wavelength. A convenient (and sometimes the only possible) procedure in experiment is to compare the total intensities of two groups of such lines:

1) $2s2p\,{}^3P_0-1s2s\,{}^3S_1$, $2s2p\,{}^3P_1-1s2s\,{}^3S_1$, $2s2p\,{}^3P_2-1s2s\,{}^3S_1$, 2) $2p^2\,{}^3P_2-1s2p\,{}^3P_2$, $2p^2\,{}^3P_2-1s2p\,{}^3P_1$, $2p^2\,{}^3P_1-1s2p\,{}^3P_1$, $2p^2\,{}^3P_1-1s2p\,{}^3P_0$.

Denoting the intensities of these groups by I_1 and I_2 and using formulas (7), we obtain for the ratio $\varkappa = I_2/I_1$ the following formula:

$$\varkappa=a\alpha_6+\frac{x^2[a(1-\alpha_6)+b]+x[a(1-\alpha_6+\beta_6)+b]}{(1+x)^2}, \tag{8}$$

where

$$\begin{gathered} a=\frac{5}{9}\,\frac{A(2p^2\,{}^3P_2-1s2p\,{}^3P_2)+A(2p^2\,{}^3P_2-1s2p\,{}^3P_1)}{A(2s2p\,{}^3P_1-1s2s\,{}^3S_1)} \\ b=\frac{1}{3}\,\frac{A(2p^2\,{}^3P_1-1s2p\,{}^3P_1)+A(2p^2\,{}^3P_1-1s2p\,{}^3P_0)}{A(2s2p\,{}^3P_1-1s2s\,{}^3S_1)}. \end{gathered}$$

Figure 2 shows plots of $\varkappa(N_e)$ calculated from formula (8) for the Mg XI ion. The probabilities of the radiative and autoionization transitions used in (8) were taken from[11]. The rates of collision transitions between the autoionization levels were calculated from the formula[17]

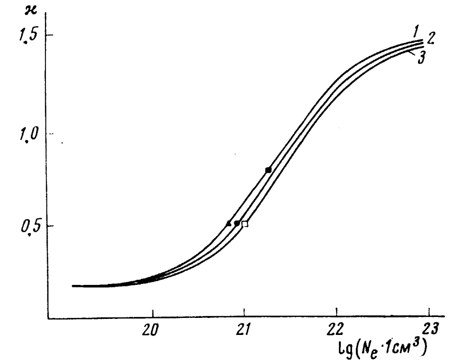

FIG. 2. Plots of $\varkappa(N_0)$ for the ion Mg XI at various temperatures: 1) $kT = 160$, 2) $kT = 250$, 3) $kT = 330$ eV. The symbols □, ●, ■, ▲ mark the values of $\varkappa$ obtained in the experiments of[7,12–14], respectively.

$$C_{15}=3.04\cdot 10^{-10}\sqrt{\frac{Z^2\mathrm{Ry}}{kT}}\left(4.6+\ln 11.6\sqrt{\frac{kT}{Z^2\mathrm{Ry}}}\right)\quad [\mathrm{cm}^3\cdot \mathrm{sec}^{-1}]\,. \quad (9)$$

When the electron density changed from 10^{20} to 10^{22} cm^{-3} the intensity ratio $\varkappa$ increased by almost one order of magnitude. This makes it possible to determine the plasma density by measuring the ratio $\varkappa$ of the satellite intensities. The temperature dependence of the ratio $\varkappa(N_e)$ turns out to be quite weak (see Fig. 2). We note that for ions with a larger nuclear charge the $\varkappa(N_e)$ curves shift towards larger values of the electron density, making diagnostics of a denser plasma possible.

The considered diagnostic method was used to determine the density of a plasma obtained by focusing laser radiation on a magnesium target under different irradiation conditions[7,12–14] (see Fig. 2). The measured densities agree with those determined by measuring the Stark broadening of the lines of H-like ions.[15,18]

4. CONCLUSION

We have elucidated in this paper the features of the excitation of the satellites of resonance lines of H-like ions and proposed a new method for the diagnostics of a plasma with $N_e \sim 10^{20}$–10^{23} cm^{-3}. Up to now, three methods of x-ray diagnostics of a plasma having so high a density have been used: a) determination of the ratio of the intensities of the resonance and intercombination line of He-like ions,[19,16] b) from the Stark broadening of the lines of H-like ions,[20,18] and c) from the intensity ratio of the fine-structure components of the resonance line of H-like ions.[21] A common shortcoming of these methods is that at least one of the lines used for the diagnostics is a resonant one, i.e., it is emitted in an optically allowed transition whose final state becomes the ground state of the ion. As a result, under real experimental conditions, the optical thickness of the plasma in this line is practically always larger than unity and effects of radiation capture in the resonance lines become important and can lead to noticeable changes in both the relative intensities of the spectral lines and in their waveforms.[22,23] At the same time, a theoretical calculation of the intensities and waveforms of the lines, with allowance for the capture of the radiation, is a very complicated problem, which can be solved at present only in a number of very simple cases.[24] It is therefore necessary to decrease in the experiments the concentration of the given ion, so as to ensure optical transparency of the plasma in the resonance line. This leads usually to a decrease of the intensities of the observed spectral lines and affects strongly the measurement accuracy.

The method proposed in this paper is free of this shortcoming. Indeed, the spectral lines used in this method are emitted in transitions whose final states are excited states of an He-like ion. Since the concentration of the ions in the excited state is quite small, the condition of optical transparency of the plasma in these lines is certainly satisfied, and no capture of the radiation takes place.

We note in conclusion that the obtained intensity ratios contain the local values of the plasma parameters. Particular interest attaches therefore to a measurement of the spatial distribution of the spectral-line intensity ratio, from which it is possible to reconstruct, with the aid of the procedure proposed in this paper, the distribution of the electron density. In particular, in the case of a laser plasma, a crystal joined with slit-type camera obscura can be used for this purpose in the 3–15 Å range, as well as various methods of obtaining monochromatic images with the aid of flat and bent crystals and gratings.[25–29]

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[1]E. V. Algitskiĭ, V. A. Boĭko, S. M. Zakharov, S. A. Pikuz, and A. Ya. Faenov, Kvantovaya Elektron. (Moscow) 1, 908 (1974) [Sov. J. Quantum Electron. 4, 500 (1974)].

[2]V. A. Boĭko, S. A. Pikuz, and A. Ya. Faenov, Preprint FIAN, No. 19, 1976.

[3]U. Feldman, G. A. Doshek, D. J. Nagel, W. E. Behring, and R. D. Cowan, Astrophys. J. 187, 417 (1974).

[4]U. Feldman, G. A. Dochek, D. J. Nagel, R. D. Cowan, and R. R. Whitlock, Astrophys. J. 192, 213 (1974).

[5]A. H. Gabriel and T. M. Paget, J. Phys. B 5, 673 (1972).

[6]L. A. Vaĭnshteĭn, I. A. Zhitnik, V. V. Korneev, and S. L. Mandel'shtam, Kratk. Soobshch. Fiz. No. 3, 35 (1972).

[7]V. A. Boĭko, O. N. Krokhin, S. A. Pikuz, and A. Ya. Faenov, Kvantovaya Elektron. (Moscow) 1, 2178 (1974) [Sov. J. Quantum Electron. 4, 1212 (1975)].

[8]A. H. Gabriel, Mon. Not. Roy. Astr. Soc. 160, 99 (1972).

[9]L. P. Presnyakov, Usp. Fiz. Nauk 119, 49 (1976) [Sov. Phys. Usp. 19, 387 (1976)].

[10]U. I. Safronova, A. I. Ivanova, and V. I. Kharitonova, Teor. Eksp. Khim. 5, 325 (1969).

[11]E. V. Aglitskiĭ, V. A. Boĭko, L. A. Vaĭnshteĭn, S. M. Zakharov, O. N. Krokhin, U. I. Safronova, and G. V. Sklizkov, Preprint FIAN, No. 113 (1973).

[12]G. V. Peregulov, E. N. Ragozin, and V. A. Chirkov, Kvantovaya Elektron. (Moscow) 2, 1844 (1975) [Sov. J. Quantum Electron. 5, 1012 (1975)].

[13]M. P. Vanyukov, V. I. Kryzhanovskiĭ, V. A. Serebryakov, V. N. Sizov, and A. D. Starikov, Opt. Mekh. Promst. No. 12, 82 (1973).

[14]V. A. Katulin, V. Yu. Nosach, and A. M. Petrov. Kvantovaya Elektron. (Moscow) 3, 1829 (1976) [Sov. J. Quantum Electron. 6, 998 (1976)].

[15]V. I. Vayanov, V. A. Boĭko, A. V. Vinogradov, *et al.*, Pis'ma Zh. Eksp. Teor. Fiz. 24, 352 (1976) [JETP Lett. 24, 319 (1976)].
[16]A. V. Vinogradov, I. Yu. Skobelev, and E. A. Yukov, Kvantovaya Elektron. (Moscow) 2, 1165 (1975) [Sov. J. Quantum Electron. 5, 630 (1975)].
[17]V. P. Shevelko, A. M. Urnov, and A. V. Vinogradov, J. Phys. 9B, 2859 (1976).
[18]V. I. Bayanov, S. S. Gulidov, A. A. Mak, G. V. Peregudov, I. I. Sobel'man, A. D. Starikov, and V. A. Chirkov, Pis'ma Zh. Eksp. Teor. Fiz. 23, 206 (1976) [JETP Lett. 23, 183 (1976)].
[19]E. V. Aglitskiĭ, V. A. Boĭko, A. V. Vinogradov, and E. A. Yukov, Kvantovaya Elektron. (Moscow) 1, 579 (1974) [Sov. J. Quantum Electron. 4, 322 (1974)].
[20]A. M. Malvezzi, E. Jamritt, and G. Tondello, Opt. Commun. 13, 307 (1975).
[21]A. V. Vinogradov, I. Yu. Skobelev, and E. A. Yukov, Fiz. Plazmy 2, 391 (1977) [Sov. J. Plasma Phys.].
[22]V. V. Ivanov, Perenos izlucheniya i spektry nebesnykh tel (Radiative Transfer and Spectra of Celestial Bodies), Nauka, 1969.
[23]I. L. Beĭgman, V. A. Boĭko, S. A. Pikuz, and A. Ya. Faenov, Zh. Eksp. Teor. Fiz. 71, 975 (1976) [Sov. Phys. JETP 44, 511 (1977)].
[24]V. V. Ivanov, in: Teoriya zvezdnykh spektrov (Theory of Stellar Spectra), ed. V. V. Sobolev, Nauka, 1966, p. 127.
[25]V. A. Boĭko, O. N. Krokhin, S. A. Pikuz, and A. Ya. Faenov, Kvantovaya Elektron. (Moscow) 1, 2178 (1974) [Sov. J. Quantum Elektron. 4, 1212 (1975)].
[26]G. V. Peregudov, E. N. Ragozin, and V. A. Chirkov, Kvantovaya Elektron. (Moscow) 2, 1844 (1975) [Sov. J. Quantum Electron. 5, 1012 (1975)].
[27]V. A. Boĭko, S. A. Pikuz, and A. Ya. Faenov, Priv. Tekh. Eksp. No. 5, 222 (1975).
[28]V. Yaakobi and L. M. Goldman, Phys. Rev. Lett. 37, 899 (1976).
[29]A. A. Ilyukhin, G. V. Peregudov, and E. N. Ragozin, Pis'ma Zh. Tekh. Fiz. 2, 681 (1976) [Sov. Tech. Phys. Lett. 2, 266 (1976)].

Instability of coherent propagation of light pulses in resonantly absorbing media

L. A. Bol'shov, V. V. Likhanskiĭ, and A. P. Napartovich

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We show that the standard form of 2π pulses (short and relatively powerful light pulses propagating in resonantly absorbing media without loss) is unstable with respect to transverse perturbations. A transverse structure develops when the pulse traverses a distance of the order of its length L in the medium. The characteristic scale length of the transverse structure arising is $\sim(\lambda L)^{1/2}$, where λ is the wavelength of the light.

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1. McCall and Hahn[1] showed experimentally and theoretically that a short light pulse with a sufficiently large energy can propagate in a resonantly absorbing medium without loss and retaining its shape (2π pulse). Subsequently this effect has been widely studied both theoretically and experimentally (see, for instance, the surveys by Poluéktov *et al.*[2]). Theoretically one usually considers the propagation of one-dimensional pulses which extend to infinity in the transverse direction. The evolution of the transverse structure of a 2π pulse connected with its transverse dimension being finite has been studied numerically (see, e.g.,[2-4]). In particular it was noted in[3,4] that a 2π pulse has a tendency for self-focusing. The problem of the evolution of a three-dimensional coherent light pulse is complicated for a complete analytical study. However, one can obtain a number of important conclusions by studying the stability of a one-dimensional 2π pulse with respect to transverse perturbations. The present paper is devoted to the solution of that problem.

2. To simplify the exposition we restrict ourselves to the two-level model of a medium without degeneracy, neglecting inhomogeneous broadening. In that case the reduced equations for the field E, the polarization P of the particles in the medium, and the difference n in the populations of the lower and upper levels have the form

$$\frac{\partial E_1}{\partial x}+\frac{\eta}{c}\frac{\partial E_1}{\partial t}+\frac{1}{2k}\nabla_\perp^2 E_2=-\frac{2\pi N\omega}{c\eta}P_1,$$
$$\frac{\partial E_2}{\partial x}+\frac{\eta}{c}\frac{\partial E_2}{\partial t}-\frac{1}{2k}\nabla_\perp^2 E_1=\frac{2\pi N\omega}{c\eta}P_2, \tag{1a}$$
$$\frac{\partial P_1}{\partial t}=\frac{\mu^2}{\hbar}E_1 n,\quad \frac{\partial P_2}{\partial t}=-\frac{\mu^2}{\hbar}E_2 n,\quad \frac{\partial n}{\partial t}=-\frac{1}{\hbar}(E_1P_1-E_2P_2). \tag{1b}$$

Here

$$E=E_1\cos(\omega t-kx)+E_2\sin(\omega t-kx),$$
$$P=P_1\sin(\omega t-kx)+P_2\cos(\omega t-kx),$$

η is the non-resonance refractive index, N the number of resonant particles, and μ the transition dipole moment. We assume that the light frequency is the same as the resonance frequency of the medium. Equations (1) have a well known stationary solution[1] in the form of a one-dimensional soliton (2π pulse):

