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ON THE PROBLEM OF QUASARS EMISSION LINES FORMATION BY CHERENKOV MECHANISM

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The possibility of explaining the broad emission lines formation in quasars spectra by Cherenkov radiation is considered. This idea was suggested and worked out in a series of papers [2–4] in which, however, a number of effects, being essential for the line profile, were ignored.

We solved a classical NLTE line–formation problem with allowance for the Cherenkov radiation. The results for the L_α–line and $T_e \sim 1.5 \cdot 10^4$ show that the effect of this mechanism becomes measureable merely for unreasonably high densities of relativistic electrons ($\tilde{n}_e \geq 10^{11} cm^{-3}$).

1. ***Introduction.*** **The unusual features of broad emission lines of quasars have, over the years, attracted wide attention. These lines are conventionally assumed to be formed through recombination–cascade and collisional exitation mechanisms, as with that in gaseous nebulae (see e.g. [1]). Quite different interpretation associated with the Cherenkov mechanism of radiation has recently been suggested by the group of astrophysists in Hefei (You, Cheng et al. [2–4]). They drew attention to the fact that, in the immediate neighbourhood of a spectral line, the refractive index can be greater than unity. It follows that over this limited range of frequencies Cherenkov emission may be observed providing that there exist definitely abundant relativistic electrons and sufficiently dense medium. The emission has the form of a redshifted, assymmetrical and broad line. Its profile was plotted by solving the equation of radiative transfer in which Cherenkov emission, and absorption in both continuum and line frequencies were taken into account. However, the usual emission processes in the line (due to spontaneous transitions) and continuous spectrum were ignored that is why the theory does not involve any parameter, estimating the efficacy of the**

Cherenkov mechanism. The only condition for the Cherenkov line appearance is that to be observable on the noisy background of continuum.

A model problem with geometry of emitting medium applicable to quasars was considered in [4]. The authors concluded that the number density of relativistic electrons of orders $10^4 - 10^6$ см$^{-3}$ are sufficient to make Cherenkov emission dominating. This is the only paper which intends to compare Cherenkov line intensity with that of line formed through traditional mechanisms. However, the authors confined themselves by comparing the emissivities of proper lines and leave out of account the multiple scattering in the line frequencies, while the gas is supposed to be optically thick. Apparently, they are under erroneous impression that the conventional line can not be formed in an opaque medium. Note, that the diffusion procceses must be taken into account even in the case when the only source of radiation in the medium is the Cherenkov mechanism. The statement that "... the Cherenkov line photon can avoid the resonance absorption because of the Cherenkov redshift and can escape easily from the deep inner part of the gas, i.e. gas appears more transparent for the Cherenkov line than for a conventional line" (see [4]) is not quite correct because this redshift is negligble, whereas the Lorenz profile decreases rather slowly in the wings and covers, as a matter of fact, the whole range of the Cherenkov line. The evident conclusion is that, in the dense medium, both these two lines are originated and that they are formed as a result of multiple scattering of radiation in the line frequencies. It will be also noted that the radiation diffusion is accompained with redistribution over frequencies which is very important in forming the line profile.

Bearing in mind the importance of the problem, we reconsider it with due allowance for all the effects discussed above. In principle, its correct formulation requires to treat the classical NLTE line–formation problem [5] with additional source of the Cherenkov radiation within a medium. In the simplest case of one–dimensional transfer of radiation (two–stream approximation), the problem admits an explicit analytical solution.

In the present paper we shall consider two typical situations:

i. the spectral line is formed in an isothermal medium as a result of collisional mechanism, and

ii. the line formation is due to illumination of a medium by strong ultraviolet radiation released by the external source of energy.

In both cases the necessary conditions for the Cherenkov emission are assumed to be satisfied. We shall see that extremely high densities of relativistic electrons are required to make the intensities of the Cherenkov and conventional lines comparable. In Section II, we revise some formulas obtained in the above–cited papers and formulate the problem in terms of quantities generally used in the radiative transfer theory. The solutions of the problem for coherent and completely non-coherent scattering are given in Sections III and IV, respectively. The numerical results and their physical consequences are presented and discussed in Section V.

2. *The formulation of the problem.* Consider a schematic atomic model consisting of two levels between which radiative and collisional transitions can occur. If we take into account the emission and absorption processes in continuum and assume that there is an additional emission source due to Cherenkov radiation, then the equation of transfer becomes

$$\mu \frac{dI_\nu}{dr} = -\alpha_\nu I_\nu - n_1 B_{12} \frac{h\nu}{4\pi} \Phi(x) I_\nu + n_2 \frac{h\nu}{4\pi} \Phi(x) (A_{21} + B_{21} I_\nu) + + B_\nu(T) + S_\nu^* \tag{1}$$

where I_ν is the radiation intensity per unit interval of ν, α_ν is the absorption coefficient in continuum, n_1 and n_2 are the number densities of atoms in the ground and excited states, $x = (\nu - \nu_0)/\Delta\nu_0$ is the so called dimensionless frequency, which is the frequency displacement from the center of the line measured in $\Delta\nu_0 = \Gamma/4\pi$ units, Γ is the damping constant, A_{21}, B_{21}, B_{12}, are Einstein's coefficients of transitions probability and $B_\nu(T)$ is the Planck function. The normalized profile of the line absorption coefficient $\Phi(x)$ assumed to be Lorenzian $\Phi(x) = (1/\pi)(1 + x^2)$ as it is the case in the You et al. treatment (see also further discussion). The last term in the right-hand side of Eq. (1) represents the emissivity of the Cherenkov radiation and can be given by the approximate formula (see e.g. [2])

$$S_\nu^* = \pi (e^2/c)(n^2 - 1 - \gamma^{-2}) \tilde{n}_e v \tag{2}$$

where n is the refractive index, $\gamma = 1/\sqrt{1-(v/c)^2}$ is the relativistic factor, v and $\tilde{n}_e$ are the velocity and number density of relativistic electrons, e is the charge of electron. The approximation is due to asumption that $n \approx 1$. The Cherenkov emission requires that $(vn/c) > 1$, or in our approximation $n^2 - 1 - \gamma^{-2} > 0$.

Turning again to Eq. (1), remind that it concerns the case of complete redistribution over frequencies when the emission and absorption coefficients profiles are assumed to be identical. The populations of the two levels are governed by the statistical equilibrium equation which can be written in the form

$$n_1 \left(B_{12} \int_{-\infty}^{\infty} \Phi(x)\, J_x dx + b_{12}\right) = n_2 \left(A_{21} + a_{21} + B_{21} \int_{-\infty}^{\infty} \Phi(x)\, J_x dx\right) \quad (3)$$

where $J_x = \oint I(r, \mu, x) \frac{d\omega}{4\pi}$ is the specific intensity, a_{21} and b_{12} are the coefficients of collisional de-excitation and excitation, respectively.

Combining Eqs.(1) and (3) and introducing the optical depth according to the formula

$$d\tau = k_o [n_1 - (g_1/g_2)\, n_2]\, dr$$

with $k_o = h\nu B_{12}/4\pi\, \Delta\nu_o$, we obtain after standard but cumbersome transformations [5]

$$\mu \frac{dI}{d\tau} = -\, [\Phi(x) + \beta_\nu]\; I(\tau, \mu, x) + S(\tau, x) \quad (4)$$

where

$$S(\tau, x) = \lambda \Phi(x)\, J + [(1-\lambda)\, \Phi(x) + \beta_\nu] B_\nu(T) + S_\nu^*(\tau) \quad (5)$$

$J = \int_{-\infty}^{\infty} \Phi(x)\, J_x\, dx$, $\beta_\nu = \alpha_\nu / k_o\, \Delta\nu_o\, [n_1 - (g_1/g_2)\, n_2]$ is the ratio of absorption coefficient in continuum to that in the line integrated over all frequencies (the dependence of α_ν, β_ν, and $B_\nu(T)$ on x or ν can be neglected troughout the line so that the appropriate subscripts of these quantities will be hereafter omitted),

$$\lambda = A_{21} \Big/ \left[A_{21} + a_{21}\left(1 - \exp\left(-\frac{h\nu}{kT}\right)\right)\right], \quad (6)$$

$$S_\nu^*(\tau) = \pi\, \bar{n}_e\, e^2\, \Delta\nu_o\, (n^2 - 1 - \gamma^{-2}) / h\, B_{12}\, [n_1 - (g_1/g_2) n_2]\,. \quad (7)$$

Note that the radiation intensity $I(\tau, \mu, x)$ in Eq.(4) is evaluated per unit interval of x.

For the further insight to the problem we need the frequency dependence of S_ν^* or $n^2 - 1 - \gamma^{-2}$. For this purpose we consider the well-known expression for the complex refractive index n which is a threshold of investigations performed in [2-4]:

$$(\hat{n}^2 - 1)/(\hat{n}^2 + 2) = n_d \varepsilon \quad (8)$$

where n_a is the number density of atoms and ε is the polarizability per atom. This expression, referred to as the Lorenz–Lorentz formula (see e.g. [6]), holds, in general, for isotropic cristalls, however it is often used for liquids and dense gases as well.

If we confine ourselves by considering the resonance line alone, we can write

$$n_a\varepsilon = \frac{e^2}{2\pi m}\,\frac{f_{12}}{2\pi(\nu_o^2-\nu^2)+i\,\Gamma}\,[n_1-(g_1/g_2)\,n_2] \tag{9}$$

where f_{12} is the oscillator strngth and m is the mass of electron. As it is known, $\hat{n}= n-i\,\kappa$ where κ is the extinction coefficient. Now to find out the frequency dependence of n and κ, we must separate the real and imagine parts of n. Making use of the relation (8), we get

$$n^2-\kappa^2=C\ ,\quad n\,\kappa=D \tag{10}$$

where

$$C=\frac{1+(x+2\bar{b})\,(x-\bar{b})}{1+(x-\bar{b})^2}\ ,\quad D=\frac{3}{2}\,\frac{\bar{b}}{1+(x-\bar{b})^2} \tag{11}$$

and $b= c^3 g_2\ [n_1-(g_1/g_2)\,n_2]\,/12\pi^2\nu_o^3\,g_1$ (in You et al. notations $x=z\,/g$ and $\bar{b}= b/g$).

The solution of (10) is

$$n^2=\left[\sqrt{C^2+4D^2}+C\right]/2\ ,\quad \kappa^2=\left[\sqrt{C^2+4D^2}-C\right]/2\,. \tag{12}$$

Now we note that for L_α we have $\lambda= c/\nu_o= 1.216\cdot10^{-5}$ cm and $\bar{b}<< 1$ if only n_1 is not unreasonably high ($\sim10^{17}$ cm^{-3}) (under normal conditions $n_2<< n_1$) so that expressions (11) and (12) after some simplifications yield

$$n^2-1=3\bar{b}x/(1+x^2)\ , \tag{13}$$

$$\kappa=(3/2)\,\bar{b}/(1+x^2)\ . \tag{14}$$

It is of interest to ask how accurate are these formulas. Comparing with calculations based on (12), one finds that the worst error in κ of about 2.75% for $\bar{b}= 0.1$ and 0.24% for $\bar{b}= 0.01$ occurs near $x\sim 1$. As for the formula (13), it fails in the close vicinity of the line center ($x\leq 0.3$ for $\bar{b}= 0.1$ and $x\leq 0.1$ for $\bar{b}= 0.01$). Outside of this interval the error at worst is only about 4% for $\overline{b}=0.1$ and 0.4% for $\overline{b}=0.01$ and reaches near $x\approx 1.5$. For greater frequencies the difference between formulas (13), (14) and exact solution is negligble which could be expected because of coincidence of their asymptotics for $x\to\infty$. These estimates show that expressions (13) and (14) can be safely used throughout the whole region of Cherenkov radiation without any effect on the final result, taking into account, in particular, the width of this frequency range for quasars ($\Delta x\sim 10^2\div 10^5$) and also the fact that the frequency band $0<x<\gamma^{-2}$ will be excluded in view of the condition $n^2-1-\gamma^{-2}>0$ (see further discussion). Hence, the requirement $x>1$ adopted in [2–4] (in You et al. notations $z>g$) is not essential at all and can be avoided to make it possible the rigorous and self–consistent treatment of the problem.

It is easily seen that the formula (14) leads to the Lorenzian profile of the line absorption coefficient. With use of relation between absorption and extinction coefficients, indeed, we find

$$k_\nu=(4\pi/c)\,\nu\,\kappa=k_o\Phi(x)\,[n_1-(g_1/g_2)\,n_2]\;. \tag{15}$$

This is the reason why under $\Phi(x)$ in Eqs.(1), (4) and (5) the Lorenzian profile was meant.

Now inserting the expression (13) into (7), we obtain

$$S_\nu^*=(n_e\,e^2/2)\,\psi(x) \tag{16}$$

where $\psi(x)=-\left[x\,\Phi(x)+(1/\overline{b}\,\gamma^2)\right]$.

The frequency limits of Cherenkov radiation can be found by setting the expression (16) to zero:

$$x_{\max,\min}=(1\pm\sqrt{1-\delta^2})/\delta \tag{17}$$

with $\delta=2/3\overline{b}\,\gamma^2$.

The values of γ for quasars estimated from observations are of the order of 10^2 or 10^3 [7] so that for $\overline{b}\sim 10^{-2}\div 10^{-1}$ we have $\delta\sim 10^{-2}\div 10^{-5}$ and a fairly wide range where the mechanism acts $\Delta x\sim 10^2\div 10^5$. We see that, to be stringent, we would multiply the expression (16) by a unit step–function $\Lambda(x)$ equaled to 1, when $x\in(x_{\min},x_{\max})$ and 0, otherwise. Therefore, hereafter we take

$$\psi(x) = -\left[x\,\Phi(x) + (1/\bar{b}\gamma^2)\right]\Lambda(x)\ . \qquad (18)$$

To simplify the further discussion, we note that, under normal conditions $n_2 << n_1$ so that negative absorption processes can be neglected. This assumption is not essential and has no influence on the final results. It is obvious that the transfer equation remains to be valid, however, now $\lambda = A_{21}/(A_{21} + a_{21})$ and only the first term in brackets will be keeped in expressions for $d\tau$, β and S_ν^*.

Let us now treat the situation when the medium is illuminated from outside by ultraviolet radiation (by Lyman continuum quanta, in particular), as in a model problem of spectral lines origination in quasars considered in [4]. The diffusion of L_c–radiation leads to generation of the L_α–quanta through recombination–cascade mechanism and serves as an additional source for them. The simplest way to take these processes into account is the application of an approach which treats the L_α– and L_c– radiation fields separately and permits one to avoid difficulties connected with the levels interlocking effects. According to this approach, we assume that the diffusion problem for L_c– quanta has been solved and the appropriate source function $S_c(\tau)$ is known (see e.g. [7]). Physically, this quantity characterizes the rate of L_α– quanta release at the expense of L_c – continuum.

It is easily seen that the equation of statistical equilibrium takes now the form

$$n_1\ [R(T_e) + b_{12} + B_{12}J] = n_2(A_{21} + a_{21}) \qquad (19)$$

with $R(T_e) = (n_e n^+/n_1)\sum_{i=2}^{\infty} c_i(T_e)$, where $c_i(T_e)$ are the recombination coefficients; T_e is the electron temperature and n^+ is the number density of ions. The quantities $R(T_e)$ and $S_c(\tau)$ are related by the formula

$$R\ (T_e) = 4\pi\,(1-p)\,k_{\nu_1} S_c(\tau)/p \qquad (20)$$

where k_{ν_1} is the absorption coefficient per hydrogen atom at the Lyman–limit frequency and p is the fraction of radiative recombinations on the ground state.

As one might expect, the new term in Eq. (19) leads to appearance of the new source term in the right–hand side of the transfer equation (4–5):

$$S_0(\tau,x) = \lambda R(T_e)\Phi(x)/B_{12} = 4\pi\lambda(1-p)k_{\nu_1}\Phi(x)\ S_c(\tau)/p\ B_{12} \qquad (21)$$

with

$$\lambda = A_{21}/\,[A_{21} + a_{21} - (g_1/g_2)R(T_e)]\ . \qquad (22)$$

The asymptotic behavior of $S_c(\tau)$, when $\tau \to \infty$, is well–known [7]:

$$S_c(\tau) \sim \frac{N_c}{2\pi} \frac{sp}{p+s^2-1} e^{-sq\tau} \tag{23}$$

where N_c is the number of L_c–quanta incident on the unit square of a surface of the medium per 1 sec, $q=k_{\nu_1}/k_o$ and s is the root of the characteristic equation $(p/2s)\ln[(1+s)/(1-s)]=0$. For sufficiently wide range of electron temperature being of interest for us, we have $p \approx 0.5$ and $s \approx 1$ so that

$$S_c(\tau) \sim (N_c/2\pi)\, e^{-q\tau} \tag{24}$$

and hence we can accept

$$S_o(\tau,x)=I_o\Phi(x)\, e^{-q\tau} \tag{25}$$

with $I_o=\lambda\, gh\, \nu_{12} N_c/2\pi\Delta\nu_o$. It is obvious that for $\tau_o \to \infty$ the effect of L_α–quanta generation is maximum.

3. *The solution of the problem.* Now we have at our disposal all the quantities which are necessary for mathematical description of the model problems considered in [2–4]. Each of them requires solving the radiative transfer equation with appropriate source term and boundary conditions.

Let us start with an isothermal and uniform plane–parallel medium of optical thickness τ_o which is not illuminated from outside. Under these assumptions S^* does not depend on optical depth and is a function of frequency alone. In this Section the scattering process is accepted to be coherent so that Eqs. (4–5) can be rewritten in the form

$$\mu\frac{dI}{d\tau} = -v(x)\, I(\tau,\mu,x)+\lambda\Phi(x)\, J+u(x)B(T)+S^*(x) \tag{26}$$

where, for breavity, the following notations are introduced: $v(x)=\Phi(x)+\beta$, $u(x)=(1-\lambda)\Phi(x)+\beta$.

In addition, we let the medium to be one–dimensional or, this is the same, we shall use two–stream approximation to the problem which is sufficient for our purpose. This is also approved by the fact that Eq. (26) admits in this case an analytical solution which can be easily obtained by standard procedure (see e.g. [8]). In particular, for the line profile, being of our interest, this technique gives

$$r(x,\tau_o) = [u(x)+\omega\psi(x)]\, F(x,\tau_o) \tag{27}$$

with

$$F(x,\tau_o) = \frac{2}{k(x)} \frac{a(x)e^{k(x)\tau_o} + \overline{a}(x)e^{-k(x)\tau_o} - 2}{a^2(x)e^{k(x)\tau_o} - \overline{a}^2(x)e^{-k(x)\tau_o}} \tag{28}$$

where $k(x)=\sqrt{u(x)v(x)}$; $a(x)=1+k(x)/v(x)$; $\overline{a}(x)=1-k(x)/v(x)$; $\omega = \tilde{n}_e e^2/2B\ (T_e)$.

In the particular case of semi–infinite medium this expression simplifies to

$$r(x)= 2w(x)/\,[u(x)+k(x)] \tag{29}$$

where $w(x) = u(x)+ \omega\,\psi\,(x)$.

We see that the operator accounting of multiple scattering processes is a multiplying operator so that the effectivity of the Cherenkov mechanism is completely determined by the value of parameter ω which, in turn, depends on the number density of relativistic electrons. This kind of parameter was not contained in the theory under consideration.

Assuming finally that there exists an additional external source of continuum radiation given by (25), in place of (29) we can write

$$r(x)= 2\left[w\,(x)+ \overline{\omega}\,\frac{\Phi\,(x)\,k(x)}{k(x)+q}\right]/\,[u\,(x)+ k(x)] \tag{30}$$

with $\overline{\omega} = I_o/B\ (T_e)$. The formula (27) also can be generalized over this case, however we do not dwell here upon this point.

4. *Completely non–coherent scattering.* The shortest way in obtaining the line profile with frequency redistribution is the probabilistic approach based on the Ambarzumian's invariance principle and proposed in [9,10]. Referring the reader to these papers for details of the method, we give here only the final expression for the residual intensity (again we are limited ourselves by considering the one–dimensional semi–infinite medium)

$$r(x)= \left[w\,(x)+ \int_{-\infty}^{\infty} w\,(x')\,\rho\,(x',x)\,dx' + \frac{\lambda}{2}\,r_1\varphi(x)\right]/\,v\,(x)+ \tag{31}$$
$$+\overline{\omega}\,(1+\frac{\lambda}{2}\,r_2)\,\varphi\,(x)/\,[v(x)+q]$$

where

$$r_1 = d/(1-\frac{\lambda}{2}\,\delta_1)\,;\ r_2 = \delta_2/(1-\frac{\lambda}{2}\,\delta_2)\,;\ \delta_1 = \int_{-\infty}^{\infty} \frac{\varphi\,(x)\,\Phi\,(x)}{v\,(x)}\,dx\ ;$$

$$\delta_2=\int_{-\infty}^{\infty}\frac{\varphi(x)\,\Phi(x)}{\nu(x)+q}\,dx\,;$$

$$d=\int_{-\infty}^{\infty}\left[w(x)+\int_{-\infty}^{\infty}w(x')\rho(x',x)\,dx'\right]\frac{dx}{\nu(x)}\,;$$

$\rho(x',x)$ is the so-called "reflection" coefficient of a medium and $\rho(x',x)=$ $=(\lambda/2)\,\varphi(x)\,\varphi(x')/[\nu(x)+\nu(x')]$; $\varphi(x)$ is the Ambarzumian's φ-function which satisfies the following functional equation

$$\varphi(x)=\Phi(x)+\frac{\lambda}{2}\int_{-\infty}^{\infty}\frac{\varphi(x)\,\varphi(x')}{\nu(x)+\nu(x')}\,\Phi(x')\,dx'\,. \tag{32}$$

The function $\varphi(x)$ is well-known in the radiative transfer theory (see, e.g., [7,8]) and can be readily computed and tabulated utilizing the iteration method. Knowing $\varphi(x)$, the line profile is evaluated immediately from (31). We see that now the operator of multiple scatterings is not a multiplying operator as it was in the case of coherent scattering (Eqs.(28)–(30)). The contribution of each of the acting mechanisms is not stipulated now by the values of ω and $\overline{\omega}$ alone, it is accounting of in much more complex manner. Physical consequences of this fact will be shown in the next Section.

5. Numerical results and discussion. Let us start with the profile (27) for the line originated in a medium of finite optical thickness with no redistribution over frequencies. We obviously have a superposition of the Cherenkov radiation profile with that of conventional line formed as a result of thermal processes. The latter may have either emission or absorption profile depending on τ_0. The relative strengths of these two lines are essentially determined by the value of ω listed in Table 1.

Table 1

PARAMETER ω FOR $\tilde{n}_e=10^{11}\,cm^{-3}$

T_e	ω	T_e	ω	T_e	ω	T_e	ω
2000	2.7 +18	5000	1.0 +3	8000	1.4 -1	15000	1.4 -4
3000	7.3 +9	6000	2.0 +1	9000	2.7 -2	20000	1.9 -5
4000	3.8 +5	7000	1.2 0	10000	7.3 -3	25000	5.9 -6

The estimation of T_e for an average quasar obtained from observations is $T_e \sim 1.5 \cdot 10^4$K [11]. For the quasar 3C 273 considered by way of example in [4], $T_e \sim 3.0 \cdot 10^4$K [12] for which $\omega = 2.7 \cdot 10^{-6}$. Taking into account that under normal physical conditions $\beta \sim 10^{-3} - 10^{-4}$ and $1-\lambda \sim 10^{-2} - 10^{-3}$ we get $\psi(x)/u(x) \leq 10 - 100$ throughout the whole frequency domain of the Cherenkov emission. We infer that even for the values of n as high as $10^{11} cm^{-3}$, the second item in expression of $w(x)$ is less than the first one so that the conventional line dominates. For lower $\tilde{n}_e$ the Cherenkov effect becomes negligble (since ω is proportional to $\tilde{n}_e$). We would have been assured of this mechanism for reasonable values of n only at extremly low temperature ($T_e \sim 3.10^3 - 4.10^3$K). Although the presence of relatively cold regions can not be completely excluded, the main contribution to the line radiation is stipulated by heated regions where conventional mechanisms are dominating. The calculated values of total energy emitted by these two ways lead to the same result. Typical profiles for various $\tilde{n}_e$ are shown in Fig.1 ($\tau_0 = 10^2$) and Fig.2 ($\tau_0 = 10^3$). Fig.3 reproduces the limiting case of semi–infinite medium ($\tau_0 \rightarrow \infty$). The calculations were performed for the L_α line. One can see that the role of the conventional line is significant if only the optical thickness of a medium is not much too small.

Turning now to the model problem proposed in [4], recall that according to this model the quasar is a sphericlike region around the central continuum powerhouse

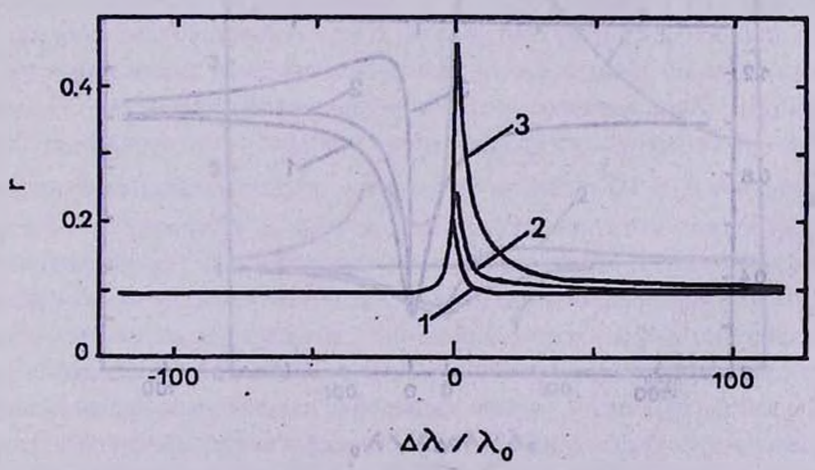

Fig.1. Line profiles for coherent scattering and a medium of optical thickness $\tau_0 = 10^{2}$ ($1-\lambda = 10^{-2}$, $\beta = 10^{-3}$, $T_e \sim 10^4$K): $\tilde{n}_e = 0$ (1), $\tilde{n}_e = 10^{11}$cm^{-3} (2), $\tilde{n}_e = 4 \cdot 10^{11}$ cm^{-3} (3).

with radius scale $R \sim 10^{19} - 10^{20}$ cm, and that there are a lot of spherical cloudlets with dense gas which spread over this region. The emission lines of hydrogen are assumed to be formed in neutral gas region of each cloudlet by the Cherenkov mechanism. Considering the model one must take into account the generation of L_α–quanta in the surface layers of cloudlets at the expense of L_c–radiation incident from the central source. This possibility was ignored in [4]. Recombination–cascade

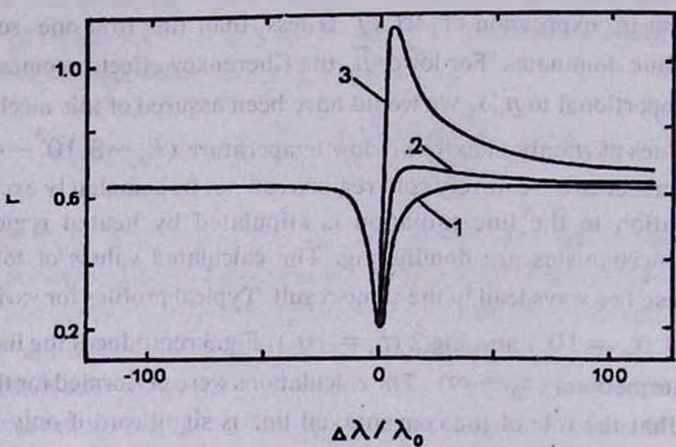

Fig.2. Same as Fig.1 for $\tau_0 = 10^3$.

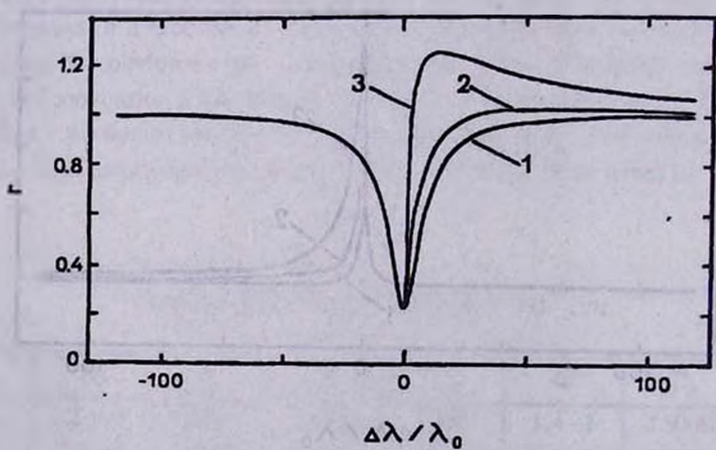

Fig.3. Same as Fig.1 for semi–infinite medium.

mechanism acted in this case is determined by the value of I_o or $\bar{\omega}$ (see Eqs. (30) and (31)). To compare the conventional line with the Cherenkov emission, we must estimate the ratio $\bar{\omega}/\omega$

$$\frac{\bar{\omega}}{\omega} = \frac{2I_o}{\tilde{n}_e e^2} = \frac{\lambda q h \nu_{12} L_c^*}{4\pi^2 \Delta\nu_o e^2 \tilde{n}_e R^2} \tag{33}$$

where L_c^* is the luminosity of the quasar in the Lyman continuum. In the order of magnitude L_c can be taken equal to total luminosity which is of order $10^{46} - 10^{47}$ erg/s [1]. Taking also $q = 10^{-4}$, $R = 10^{20}$ cm, $\Delta\nu_o/\nu_{12} = 10^{-2}$ we obtain $\bar{\omega}/\omega = 7.3 \cdot 10^{-5} \tilde{n}_e^{-1}$ so that for $\tilde{n}_e \sim 10^2 - 10^4$, $\bar{\omega}/\omega \sim 10^{-6} - 10^{-8}$ However for some quasars (e.g., 3C 345 and 3C 446) $R \leq 10^{15} - 10^{17}$ cm [1], and we have $\bar{\omega}/\omega \sim 10^{-2} - 10^4$ and the question of the Cherenkov radiation is a point at issue. It will be also noted that the inner parts of cloudlets, where this radiation is originated, are opaque for L_α– quanta and they undergo multiple scattering before escaping the medium. Owing to these processes, the electron temperature may become high enough to provide the thermal formation of Lyman lines of observable intensity.

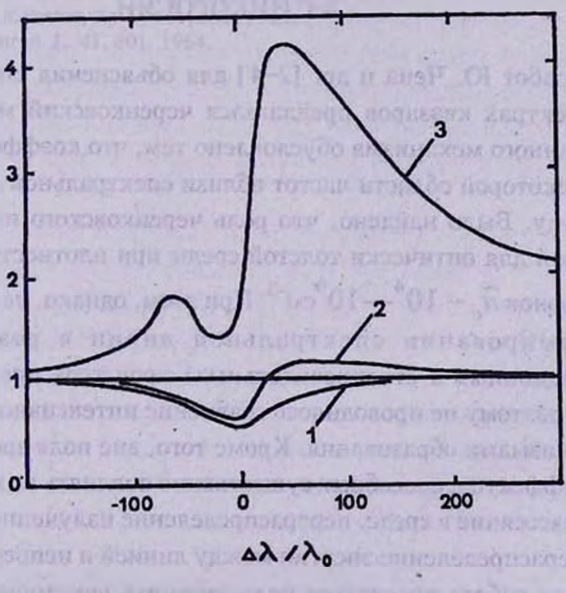

Fig.4. Line profiles for completely non–coherent scattering and semi–infinite medium; $\omega = 0$ (1), $\omega = 0.1$ (2), $\omega = 1$ (3).

Fig.4 shows the profiles of the L_α line for the complete redistribution of radiation over frequencies. The curves are obtained with use of (31) (for $r_1=0$, $\bar{\omega}=0$) and concern the semi-infinite medium. We see that for larger values of ω, the redistribution effects lead to appearing of emission in the shortwave wing of the line as well, though the conventional absorption line is still noticable. For optically thin medium when the conventional line is in emission, we have a broad assymmetrical but not redshifted (as in [2–4]) spectral line. Thus, we see that the difficulties encountered by the Cherenkov mechanism in interpretation of quasars emission lines are diverse and seen to be insuperable.

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К ЗАДАЧЕ ОБ ОБРАЗОВАНИИ ЭМИССИОННЫХ ЛИНИЙ КВАЗАРОВ ЧЕРЕНКОВСКИМ МЕХАНИЗМОМ

А.Г.НИКОГОСЯН

В серии работ Ю. Чена и др. [2–4] для объяснения широких эмиссионных линий в спектрах квазаров предлагался черенковский механизм излучения. Действие данного механизма обусловлено тем, что коэффициент преломления плазмы в некоторой области частот вблизи спектральной линии может превышать единицу. Было найдено, что роль черенковского излучения становится существенной для оптически толстой среды при плотности пучка релятивистских электронов $\tilde{n}_e \sim 10^4 - 10^6$ см$^{-3}$. При этом, однако, не рассматривался вопрос о формировании спектральной линии в результате обычных (рекомбинационных и столкновительных) процессов при тех же физических условиях и поэтому не проводилось сравнение интенсивностей линий с различными механизмами образования. Кроме того, вне поля зрения авторов остался целый ряд эффектов, способных существенно повлиять на профиль линии (многократное рассеяние в среде, перераспределение излучения по частотам внутри линии, перераспределение энергии между линией и непрерывным спектром).

Настоящая работа преследует цель уточнить некоторые из используемых в [2–4] формул и восполнить указанные пробелы. Ставится и дается точное решение классической задачи образования спектральной линии с учетом действия черенковского механизма излучения. Расчеты проводятся для линии L_α. Показывается, что при принятых в настоящее время значениях электронной температуры для среднего квазара ($T_e \sim 1.5 \cdot 10^4$K) влияние черенковского механизма становится заметным лишь при неправдоподобно высоких плотностях релятивистских электронов ($\geq 10^{11}$см$^{-3}$). С непреодолимыми трудностями, связанные с энергетикой квазара и размерами излучающих областей, сталкивается предложенная в [4] геометрическая модель излучения квазаров в линии.

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