Broad Red-Shifted Lines as a Signature of Outflow

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Abstract. We formulate and solve the diffusion problem of line photon propagation in a bulk outflow from a compact object (black hole or neutron star) using a generic assumption regarding the distribution of line photons within the outflow. Thomson scattering of the line photons within the expanding flow leads to a decrease of their energy which is of first order in v/c, where v is the outflow velocity and c the speed of light. We demonstrate that the emergent line profile is closely related to the time distribution of photons diffusing through the flow (the light curve) and consists of a broad redshifted feature. We analyzed the line profiles for the general case of outflow density distribution. We emphasize that the redshifted lines are intrinsic properties of the powerful outflow that are supposed to be in many compact objects.

INTRODUCTION

The problem of photon propagation in an optically thick fluid in bulk motion has been studied in detail in a number of papers. The idea that photons may change their energy in repeated scatterings with cold electrons in a moving fluid was suggested more than 20 years ago [1, 2, 3]. This process, often referred to as dynamical or bulk flow Comptonization, is similar in many ways to Comptonization by hot electrons once the thermal velocity is replaced by bulk velocity $\mathbf{v}$; there is however a qualitative difference, in that their energy gain is linear in the velocity v, rather than quadratic as is in the case of Compton scattering in a medium at rest. Photons diffusing and scattering in a medium with bulk flow can gain or lose energy to the flow depending on the divergence of its velocity field: they gain energy for $\nabla \cdot \mathbf{v} < 0$, while they lose energy in the opposite case. The case of a converging radial inflow, where $\nabla \cdot \mathbf{v} < 0$, has been treated in (among others) [2, 5, 4, 6, 7]. These investigations showed that, when monochromatic radiation with $\nu = \nu_*$ is injected at large Thomson depth in a spherical inflow, the emergent spectrum develops a broad, power-law tail at $\nu > \nu_*$. The power law index is related to a combination of the flow Thomson depth and its velocity gradient, and this has been invoked to explain the high energy spectra of accreting BH candidate sources in their "high" state [8, 9].

In this note we outline the treatment of the problem of an outflow and show that a similar broad spectrum is formed but now at energies $\nu < \nu_*$, with the power law index dependent on the velocity gradient (and the boundary conditions). This problem was first treated in [5] who found a particular solution of the equations of [2] for diverging flow using Fourier transformation to find that adiabatic expansion produces a drift of injected monochromatic photons towards lower frequencies and the formation of a power-law, low energy tail with spectral index 3.

This problem is relevant observationally because of the well known presence of red wings in the fluorescent Fe Kα line profiles observed in the AGN and galactic BH spectra [10]. These have been interpreted as due to the kinematic broadening of these transitions by motions appropriate to the vicinity of a black hole. As shown by [11] this arrangment could produce lines of the desired width and asymmetry provided that the disk extended to its innermost stable orbit ($\sim$ 3 Schwarzschild radii or less). Broadening by Thomson scattering on cold electrons was dismissed, as it would also introduce a (non-observed) break in the spectrum at energy $E \simeq m_e c^2/\tau_0^2 \simeq 20$ keV.

However plausible, this interpretation is not entirely without problems; for instance, the (occasionally) exceedingly broad red wing of the line observed requires (in such cases) that the disk illumination be concentrated very close to its inner edge ($F_x \propto r^{-8}$; [12]), more than most models would allow; in addition, as indicated by [13], the ionization of such a disk by the intense X-ray radiation might invalidate some of basic assumptions associated with this interpretation. Of additional interest is also the fact that there is a marked absence of a blue-shifted wing in the line, a feature expected for a random orientation of accretion disks and the observers' lines of sight.

Motivated by the above facts and the importance of the Fe K$\alpha-$line in probing the strong field limit gravity we believe it is important that all potential alternatives to line broadening be considered in detail. To this end we examine under what conditions these broad Fe$\alpha-$line profiles could be also attributed to the effects of an outflow rather than solely to accretion disk kinematics. The details of this treatment can be found in [14]. Herein we present only the final results and some conclusions.

RADIATIVE TRANSFER IN AN OUTFLOW: FORMULATION

Let $N(r) = N_0(r_0/r)^\beta$ be the radial number density profile of an outflow and let its radial outward speed be

$$v_b/c = (\dot{M}_{out}/4\pi c N_0 r_0^2)(r_0/r)^{2-\beta} = \dot{m}_{out}(r_0/r)^{2-\beta} \quad (1)$$

obtained from mass conservation in a spherical geometry (here $\dot{M}_{out} = 4\pi r^2 v_b N$). The Thomson optical depth of the flow from some radius r to infinity is given by

$$\tau = \int_r^\infty N_e(r)\sigma_\mathrm{T}\, dr = \sigma_\mathrm{T} N_0 r_0 (r_0/r)^{\beta-1}/(\beta-1), \quad (2)$$

where $N_e(r) = N(r)$ is the electron density, σ_T is the Thomson cross section, r_0 is a radius at the base of the outflow (this definition of optical depth makes the tacit assumption that $\beta > 1$) and $\tau_{T,0} = \tau(r_0) = \sigma_\mathrm{T} N_0 r_0/(\beta-1)$.

The transfer of radiation within the flow in space and energy is governed by the photon kinetic equation ([1], Eq. 18) for the photon occupation number $n(r,\nu)$, which in steady state reads

$$-\mathbf{v_b}\cdot\nabla n + \frac{1}{3}\nabla\cdot\left(\frac{c}{\kappa}\nabla n\right) + \frac{1}{3}(\nabla\cdot\mathbf{v_b})\,\nu\frac{\partial n}{\partial \nu} = -\tilde{j}(r,\nu), \quad (3)$$

where $\kappa = N_e(r)\sigma_\mathrm{T}$ is the inverse of the scattering mean free path, $\mathbf{v_b} = v_b\mathbf{e_r}$, is the flow velocity, $\mathbf{e_r}$ is the radial unit vector and $\tilde{j}(r,\nu)$ is the photon source term.

Because the operator of Eq. (3) is the sum of a space, L_r, and a frequency, L_ν, operators, for a RHS source $j(r,\nu)$, which is is factorizable, i.e.

$$L_r n + L_\nu n = -j(r,\nu) = -f(r)\varphi(\nu) \quad (4)$$

the solution of the above problem with boundary conditions independent of the energy ν, i.e.

$$L_r^{(1)} n = 0 \quad \text{as } \mathrm{r}\to\infty, \quad L_r^{(2)} n = 0 \quad \text{for } \mathrm{r} = \mathrm{r}_0\,, \quad (5)$$

is given by the convolution of the solutions of the time-dependent problem of each operator, namely

$$n(r,\nu) = \int_0^\infty P(r,u)X(\nu,u)du, \quad (6)$$

where $u = N_0\sigma_\mathrm{T} ct$ is the dimensionless time and $P(r,u)$ is the solution of the initial value problem of the spatial operator L_r

$$\frac{\partial P}{\partial u} = L_r P, \quad P(r,0) = f(r) \quad (7)$$

with boundary conditions

$$L_r^{(1)} P = 0 \quad \text{as } \mathrm{r}\to\infty, \quad L_r^{(2)} P = 0 \quad \text{at } \mathrm{r} = \mathrm{r}_0, \quad (8)$$

and $X(\nu,u)$ the solution of the initial value problem of the frequency operator L_ν

$$\frac{\partial X}{\partial u} = L_\nu X, \quad X(\nu,0) = \varphi(\nu) \quad (9)$$

with boundary conditions

$$\nu^3 X \to 0 \quad \text{when} \quad \nu \to 0,\, \infty. \quad (10)$$

The solution of the frequency problem is (ν_* is the injection frequency)

$$J_{\nu_*}(\nu,u) = \nu^3 X(\nu,u) = e^{-3u}\delta(\nu e^u - \nu_*). \quad (11)$$

Substitution of $J_{\nu_*}(\nu,u)$ from Eq.(11) into Eq.(6) gives us the Green's function

$$G_{\nu_*}(\nu) = \nu^3 n(r_{out},\nu) = \frac{1}{\nu_*}\left(\frac{\nu}{\nu_*}\right)^3 P[r_{out}, \ln(\nu_*/\nu)]. \quad (12)$$

The time distribution $P(r_{out},u)$ can be found from the solution of the time-dependent problem of the space operator and initial condition. The final result is

$$\mathscr{F}(\nu) \propto \frac{1}{\nu_*}\left(\frac{\nu}{\nu_*}\right)^{3(1+1/\beta)} \exp\left[-\frac{(\nu/\nu_*)^{3/\beta}\xi_*}{1-(\nu/\nu_*)^{3/\beta}}\right] \quad (13)$$

$$\times \left[1-\left(\frac{\nu}{\nu_*}\right)^{3/\beta}\right]^{-(\beta+2)}, \quad (14)$$

where $\xi_* = 3C_{\nu_b}(\beta-1)\tau_*^{1/(\beta-1)}$ and $C_{\nu_b} = (\dot{M}_{out}/4\pi c N_0 r_0^2)\tau_0^{(\beta-2)/(\beta-1)} = \dot{m}_{out}\tau_0^{(\beta-2)/(\beta-1)}$ and τ_* is the depth at which the photons are injected (we assumed the initial photon distribution to be $f(\tau) = \delta(\tau-\tau_*)$. One notes that asymptotically the flux is a power law $(\nu/\nu_*)^{3(1+1/\beta)}$. This is different from the index $\alpha = 3$ obtained in [4]. The difference of the two results can be traced to the different boundary conditions employed. The present treatment employs an absorptive boundary ($n(\tau_0) = 0$), while [4] employed a reflective one. Since the asymptotic slope is determined by the photons that spend the longest in the scattering flow, it is apparent that these photons have a greater chance of being lost (throug the boundary) in our case, leading to a steeper spectrum.

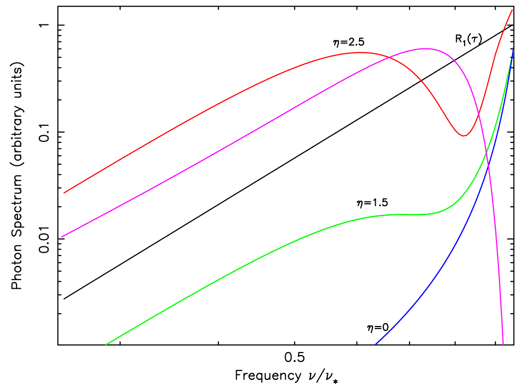

FIGURE 1. The emergent photon redshifted line spectra for four spatial monochromatic source distributions, $f(\tau) = \exp(\eta\xi)$, for $\eta = 0,\ 1.5,\ 2.5$ and $f(\tau) = R_1(\tau)$ the first eigenfunction of the spatial operator. These spectra are produced for $\beta = 2,\ \dot{m}_{out} = 0.8$ and $\tau_0 = 2$. The fifth photon spectrum is for the asymptotic case of $\tau_0 \gg 1$. The monochromatic $\delta-$function source is located at $\xi = \xi_* = 2$, and $\beta = 2$

These spectra are in general agreement with observations. Differentiation of these models with those of [11] can be obtained through the time lags of the line wings relative to the line core expected in the case examined here, compared to those of the lines are produced by reprocessing on a disk where these lags should have the opposite sign.

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