The role of Compton heating in radiation-regulated accretion on to black holes

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ABSTRACT

We investigate the role of Compton heating in radiation-regulated accretion on to black holes from a neutral dense medium using 1D radiation-hydrodynamic simulations. We focus on the relative effects of Compton-heating and photo-heating as a function of the spectral slope $\alpha$, assuming a power-law spectrum in the energy range of $13.6 \text{ eV–} 100 \text{ keV}$. While Compton heating is dominant only close to the black hole, it can reduce the accretion rate to $0.1\%$ ($l \propto \dot{m}^2$ model)–$0.01\%$ ($l \propto \dot{m}$ model) of the Bondi accretion rate when the BH radiation is hard ($\alpha \approx 1$), where $l$ and $\dot{m}$ are the luminosity and accretion rate normalised by Eddington rates, respectively. The oscillatory behaviour otherwise typically seen in simulations with $\alpha > 1$, become suppressed when $\alpha \approx 1$ only for the $l \propto \dot{m}$ model. The relative importance of the Compton heating over photo-heating decreases and the oscillatory behaviour becomes stronger as the spectrum softens. When the spectrum is soft ($\alpha > 1.5$), photo-heating prevails regardless of models making the effect of Compton heating negligible. On the scale of the ionization front, where the gas supply into the Strömgren sphere from large scale is regulated, photo-heating dominates. Our simulations show consistent results with the advection-dominated accretion flow ($l \propto \dot{m}^2$) where the accretion is inefficient and the spectrum is hard ($\alpha \approx 1$).

Key words: accretion, accretion discs – black hole physics – hydrodynamics – radiative transfer – methods: numerical.

1 INTRODUCTION

A realistic estimation of an accretion rate for a black hole (BH) is of great cosmological importance. It does not only determine its cosmological growth rate of the BH, but also it is linked directly to the BH luminosity, and thus radiative-feedback effects on the neighbouring gas of the host galaxies and intergalactic medium. Commonly, the Eddington-limited Bondi–Hoyle rate (Bondi & Hoyle 1944; Bondi 1952) is used as the prescription for the accretion rates of the BHs in cosmological simulations. However, this simple BH accretion recipe is not physically motivated since the gas accretion on to BHs is regulated by radiative and mechanical feedback from the BH. Due to the feedback, the thermal and dynamical structure of the accretion flow is far more complicated than the one described by the classical Bondi accretion calculation. The radiative feedback by BHs is important when the accretion activity of the BH peaks (i.e. quasar mode of active galactic nuclei, AGN), whereas the mechanical feedback is prominent in the low-accretion state (i.e., radio mode of AGN). Additionally, the BH spectrum in high-energy X-ray is known to display different characteristics switching between hard and soft states depending on the accretion luminosity.

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the BH when the BH is at rest (Park & Ricotti 2011, 2012) or in supersonic motion (Park & Ricotti 2013). Pressure equilibrium between the neutral and ionized gas across the I-front is the key factor for explaining reductions in the accretion rate (Park & Ricotti 2011, 2012, 2013). Thus, understanding the heating mechanisms or the instabilities (Williams 1999, 2002; Mizuta et al. 2005; Whalen & Norman 2008; Fender 2011; Park et al. 2014; Ricotti 2014) inside the H II region and near the I-front that might modify the thermal structure of the Strömgren sphere is extremely important in understanding the radiation-regulated growth of the BHs.

In the physical conditions close to an accreting BH, the strong radiation field and the ionized gas can exchange energy directly through Compton scattering. The direction of energy exchange is such as to bring the radiation-field and gas closer to thermal equilibrium. The subsequent Compton heating (or cooling) process increases (decreases) the gas temperature to the characteristic equilibrium. The subsequent Compton heating (or cooling) process is such as to bring the radiation-field and gas closer to thermal equilibrium. The Compton heating rate is determined by the number of Compton scattering events per unit time, and the energy exchange per interaction \(\Delta E(\nu, T)\), between the high-energy photons from a BH and the electrons in the neighbouring gas medium with a gas temperature \(T\). Under the assumption of spherical symmetry, the Compton heating rate \(\Gamma_C\) is given by:

\[
\Gamma_C = \frac{\sigma_T n_e}{m_e c^2} \int_{\nu_{\text{min}}}^{\nu_{\text{max}}} \nu F_\nu (\nu - 4k_B T) d\nu,
\]

where \(n_e\) is the electron number density, and \(\sigma_T\) is the Thomson cross-section. The radiation flux \(F_\nu\) at frequency \(\nu\) in Equation (1) depends on the distance from the BH and the optical depth \(\tau\), however the optical depth can be generally neglected in optically-thin situations \(F_\nu = L_\nu / 4\pi r^2\) such as in the vicinity of the BH. Assuming a power-law spectrum \(F_\nu = C \nu^{-\alpha}\), the radiation flux can be expressed \(F = C(\nu^{\max} - \nu^{\min})/(1 - \alpha)\) for \(\alpha \neq 1\) while \(F = C \ln(\nu^{\max}/\nu^{\min})\) for \(\alpha = 1\). Equation (1) can be integrated to give a simple expression:

\[
\Gamma_C = F \frac{\sigma_T n_e}{m_e c^2} \left( \frac{\nu^{\max} - 4k_B T}{\Lambda_C} \right),
\]

where \(\Lambda_C = \ln(\nu^{\max}/\nu^{\min})\) for the special case of \(\alpha = 1\). In general, \(\Lambda_C\) can be expressed for \(1 < \alpha < 2\) as

\[
\Lambda_C = \frac{2}{\alpha - 1} - \frac{\nu^{\max} - \nu^{\min}}{\nu^{\max} - \nu^{\min}}.
\]

From Equation (2), the characteristic Compton temperature \(T_C\) is defined as

\[
T_C = \frac{\hbar \nu^{\max}}{4k_B \Lambda_C},
\]

which is the temperature when the radiation field and the gas are in equilibrium state. The Compton temperature monotonically increases with the maximum photon energy \(E^{\max} = \hbar \nu^{\max}\) and decreases as a function of the spectral index \(\alpha\). In Fig. 1 circles and triangles show the Compton temperature as a function of the spectral index \(\alpha\) for \(E^{\max} = 100\) and 30 keV respectively with the fixed \(E^{\min} = 13.6\) eV. For example, with \(\alpha = 1\) and \(E^{\max} = 100\) keV, the Compton temperature reaches \(T_C \sim 3 \times 10^7\) K when \(T_C \sim 3 \times 10^6\) K when \(\alpha = 1.5\). Since \(T_C\) is approximately proportional to \(E^{\max}\), the \(T_C\) is lower when \(E^{\max} = 30\) keV. The maximum photon energy of \(E^{\max} = 100\) keV is adopted since Compton scattering becomes less efficient for photons with energies above \(E \sim 100\) keV, where the relativistic effects encapsulated by the Klein–Nishina correction reduce the scattering cross-section \(\sigma_T\) (Klein & Nishina 1929; Blumenthal 1974). In addition, the Nuclear Spectroscopic Telescope Array (NuSTAR) has shown that spectra of many (non-blazar) accreting BH systems appear to cut off at approximately 100 keV due to the physics of the accretion disk corona (Harrison et al. 2013).

2 METHODOLOGY

2.1 Compton heating as a function of spectral index

The Compton heating (or cooling) rate is determined by the number of scatterings \(N_{\text{scatt}}(\nu)\) per unit time, and the energy exchange per interaction \(\Delta E(\nu, T)\), between the high-energy photons from a BH and the electrons in the neighbouring gas medium with a gas temperature \(T\). Under the assumption of spherical symmetry, the Compton heating rate \(\Gamma_C\) is given by:

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2.2 Accretion feedback models

The bolometric luminosity of a BH is \(L = \eta \dot{m} \dot{m} \eta_{\text{edd}}\), where the luminosity and the accretion rate are normalised by the Eddington rates as \(\dot{L} \equiv \dot{L}_{\text{edd}} / \dot{m} \equiv \dot{M} / \dot{M}_{\text{edd}}\). The Eddington luminosity is \(L_{\text{edd}} = 4\pi G M \dot{m} \eta_{\text{edd}} \eta_{\text{edd}}\) and we define the Eddington accretion rate \(\dot{M}_{\text{edd}} \equiv L_{\text{edd}} / c^2\). For a thin disc (Shakura & Sunyaev 1973) model, the radiative-efficiency is constant \(\eta \sim 0.1\) and
Compton Temperature ($T_C$) as a function of spectral index ($\alpha$). Circles and triangles show $T_C$ for the maximum energy of the photon $E_{max} = 100$ and 30 keV, respectively. With a fixed $E_{max} = 100$ keV, the Compton temperature is $T_C \sim 3 \times 10^5$ K for $\alpha = 1$ and monotonically decreases with increasing $\alpha$, reaching $T_C \sim 3 \times 10^6$ K for $\alpha \sim 1.5$. With $E_{max} = 30$ keV, the Compton temperature is $T_C \sim 10^7$ K at $\alpha = 1$, which is lower than the $T_C$ for $E_{max} = 100$ keV. The difference between $T_C$ for different $E_{max}$ becomes reduced with increasing $\alpha$ due to the extra dependence on $E_{max}$ and $\alpha$ shown in Equation (3).

$0 < \dot{m} < \eta^{-1} \sim 10$ for Eddington-limited accretion. On the other hand, the efficiency increases as the accretion rate increases for spherical accretion or an advection-dominated accretion flow (ADAF), going as $\eta \propto \dot{m}$ (Shapiro 1973; Narayan & Yi 1995; Abramowicz et al. 1996; Park & Ostriker 2001). The ADAF model ($l \propto \dot{m}^\alpha$) is radiatively-inefficient and is applicable when the accretion rate is low. The soft synchrotron photons from the ADAF inverse Compton scatter off the hot ($\sim 100$ keV) electrons in the plasma and produce a hard spectrum ($\alpha = 1$) (Narayan et al. 1998).

Interestingly, it has been shown that the photo-heating/ionization alone does not produce a significant difference between $l \propto \dot{m}$ and $l \propto \dot{m}^2$ models when a same spectral index $\alpha$ is assumed for these two models (Park & Ricotti 2011). This is due to the fact that the thermal structure inside the H II region (described by the parameter $T_{in}$), the key for determining the accretion rate, is not affected by the difference between the two models but significantly altered by the spectral index as explained in the previous section. In this study, we examine how the inclusion of Compton heating affects the different accretion models ($l = 0.1\dot{m}$ and $l = 0.1\dot{m}^2$) focusing on the accretion rate and the oscillatory behaviour found in the previous works.

### 2.3 Radiation-hydrodynamic simulation setup

We run a suite of 1D radiation-hydrodynamic simulations for a fixed BH mass $M_{bh} = 10^6 M_\odot$ and a gas density $n_{H,\infty} = 10$ cm$^{-3}$. We use the non-relativistic hydrodynamic simulation code ZEUS-MP (Stone & Norman 1992; Hayes et al. 2006) with our 1D radiative transfer subroutine (Ricotti et al. 2001; Whalen & Norman 2006). The BH is assumed to be located at the origin of the computational domain which we cover with 256–512 logarithmically spaced radial grid zones. We assume spherical symmetry. We adopt an operator-split method that alternates hydrodynamic and radiative-transfer calculations, using the smaller of the hydrodynamical and ionization time steps at each time-step

\[ \Delta t = \min(\Delta t_{\text{hydro}}, \Delta t_{\text{ion}}). \]

The mass accretion rate is not prescribed but instead directly read at the minimum radius $r_{\text{min}}$ at every time-step and converted to the radiation luminosity depending on the accretion model. For the radiative-transfer calculation, the number of photons in 50 logarithmically spaced energy bins from 13.6 eV to 100 keV are calculated using the instantaneous BH luminosity and the spectral index $\alpha$. Our subroutine solves the 1D radiative transfer equation by calculating photo-heating/ionization, gas cooling, radiation pressures, and Compton heating as explained in Section 2.1.

The Compton heating rate $\Gamma_C$ is calculated directly by performing the integral shown in Equation (1). Our 1D radiative transfer subroutine updates the abundances of H I, H II, He I, He II, He III, and $e^-$ at each radius for every time-step. We run simulations with a different spectral index in the range of $1.0 \leq \alpha \leq 1.8$ with a fixed energy range from $E_{\text{min}} = 13.6$ eV to $E_{\text{max}} = 100$ keV. The spectral index $\alpha$ is kept constant in each simulation.

### 3 RESULTS

#### 3.1 Accretion rate as a function of spectral index

Fig. 2 shows the accretion rates ($\dot{\lambda}_{\text{rad}} = \dot{M}/\dot{M}_\text{Bondi}$) as a function of time for the simulations with Compton heating included. We normalize the accretion rate to the Bondi accretion rate $\dot{M}_B = \pi \rho c^2 \gamma M_\bullet r_{\text{BH}}$, where $\rho$ is the gas density at isothermal equation of state ($\gamma = 1$), $\rho_{\infty}$ is the density of the gas. For the assumed simulation parameters ($M_{\text{BH}}n_{H,\infty} = 10^7 M_\odot \text{cm}^{-3}$, $T_{\infty} = 10^5$ K, and $\eta = 0.1$), it can be shown that $\dot{\lambda}_{\text{rad}} \approx 2.5\dot{\lambda}_{\text{Bondi}}$. The left panels show the model $l \propto \dot{m}$ while the right panels show the model $l \propto \dot{m}^2$. From top to bottom, the spectral index decreases ($\alpha = 1.5, 1.2,$ and 1.0). Strong oscillations of the accretion rate are seen for $\alpha = 1.5$ for both models, however the oscillatory behaviour becomes suppressed with decreasing spectral index. The accretion amplitude between the maximum and minimum decreases and the mean accretion rate (shown as dashed lines) also decreases as the spectrum becomes harder for both $l \propto \dot{m}$ and $l \propto \dot{m}^2$ models. For $l \propto \dot{m}$ model at $\alpha = 1.0$, the oscillation becomes almost completely suppressed and the mean accretion rate becomes $\sim 10^{-3}$ of the Bondi rate. Mild oscillations are observed for the model $l \propto \dot{m}^2$ at $\alpha = 1.0$ and the accretion rate becomes $\sim 10^{-3}$ of the Bondi rate, higher than the rate for the $l \propto \dot{m}$ model.

Fig. 3 shows the time-averaged accretion rate $\langle \dot{\lambda}_{\text{rad}} \rangle$ as a function of the spectral index. Different symbols show different sets of simulations: $l \propto \dot{m}$ with photo-heating only (triangles), $l \propto \dot{m}^2$ model with Compton heating (circles), and $l \propto \dot{m}^2$ with Compton heating (squares). The average accretion rate for the photo-heating-only model shows mild dependence on the spectral index, while the both models with Compton heating show strong dependencies on the spectral index. $l \propto \dot{m}^2$ model shows that the accretion is suppressed as the spectrum becomes harder (0.1 per cent of the Bondi rate when $\alpha \sim 1$). The average accretion rate for the model $l \propto \dot{m}^2$ shows milder dependence on the spectral index than the model $l \propto \dot{m}$ does.

#### 3.2 Compton heating vs. photo-heating

Fig. 4 shows the distribution of the ratio between the Compton and photo-heating rates as a function of radius for $l \propto \dot{m}$ model. Since the ratio at each radius changes as a function of time (neutral fraction of H and He changes depending on the accretion luminosity)
due to the oscillatory behaviour of accretion, we show a 2D histogram of the combined outputs up to 135 Myr with regular time-step. Note that the relative distribution is relevant here since the number in each bin depends on the bin size and the frequency of the simulation outputs. The top panel shows the simulation with \( \alpha = 1.1 \) and the bottom shows the simulation with \( \alpha = 1.5 \). For both cases, Compton heating becomes dominant as the radial position inside the Strömgren sphere becomes nearer to the BH, while the ratio increases as a function of radius outside the Strömgren sphere (shown as narrow distribution in Fig. [3]). Note that both Compton and photo-heating rates are low outside the Strömgren sphere due to the high optical depth.

The ratio of the Compton-heating and photo-heating rates as a function of radius can be explained directly by the ratio between neutral and ionized fractions as mentioned in Sec. [1] Fig. [3] shows the ratio distribution between the abundances of \( \text{H II} \) and \( \text{H I} \) for the same simulations shown in Fig. [2]. The Compton heating rate is proportional to the electron number density as in Equation [1]

\[
\dot{\lambda} = \frac{M}{M_B} \dot{M}
\]

where the ionization fraction is close to unity inside the Strömgren sphere. In contrast, the photo-heating rate is proportional to the neutral fraction of the gas, which increases as a function of radius inside the Strömgren sphere. This fact explains well the radial behaviour of the relative strength of the two heating mechanisms.

Only a tiny fraction of the \( \text{H II} \) region, more than 2 orders of magnitude smaller in radius than the \( \text{I-front} \), is dominated by the Compton heating for \( \alpha = 1.1 \). Compton heating becomes more efficient at low \( \alpha \) due to the suppressed oscillatory behaviour as shown at the top panel of Fig. [3] whereas Compton heating is dominant over photo-heating only during the peak accretion phase for \( \alpha = 1.5 \) (bottom panel of Fig. [3]).

### 4 SUMMARY AND DISCUSSION

In this paper, we investigate the relative role of Compton-heating and photo-heating on radiation-regulated accretion on to BHs. We summarize our findings as follows.

- We find that the Compton heating is important when the spectrum of the BH radiation is hard (\( \alpha \approx 1 \)) or \( T_C = 3 \times 10^7 \) K. The oscillatory behaviour due to the radiative feedback loop becomes weak and the average accretion rate is well suppressed to 0.01 (\( l \propto \dot{m} \) model)–0.1 (\( l \propto \dot{m}^2 \) model) per cent of the Bondi rate. The accretion rate increases and the oscillatory behaviour becomes stronger with increasing spectral index. When the spectrum is soft (\( \alpha > 1.5 \)), photo-heating prevails regardless of models making the effect of Compton heating negligible.

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**Figure 2.** Evolution of the accretion rates for simulations with different spectral index (\( \alpha = 1.5, 1.2 \) and 1.0) for accretion models \( l \propto \dot{m} \) (left panels) and \( l \propto \dot{m}^2 \) (right panels). \( \lambda_{\text{rad}} \) shown here is the accretion rate normalised by the Bondi rate (\( \dot{M}_B \)). As the Compton temperature increases with decreasing spectral index \( \alpha \) from top to bottom panels, the accretion rate and the amplitude of the oscillations for both models decrease. For the case of \( \alpha = 1.0 \), oscillatory behaviour is almost suppressed for the \( l \propto \dot{m} \) model, while \( l \propto \dot{m}^2 \) model still displays oscillation of the accretion rate. The \( l \propto \dot{m} \) model shows significantly reduced time-averaged accretion rate (\( \lambda_{\text{rad}} \)), shown as a dashed line in each panel, while the \( l \propto \dot{m}^2 \) model displays the milder change.

**Figure 3.** Time-averaged accretion rates normalised by the Bondi rate (\( \lambda_{\text{rad}} \)) as a function of spectral index in the range \( 1.0 \leq \alpha \leq 1.8 \) for different feedback models: \( l \propto \dot{m} \) with photo-heating only (triangles), \( l \propto \dot{m}^2 \) with Compton heating (squares), and \( l \propto \dot{m} \) with Compton heating (circles). The average accretion rate for the photo-heating only model shows a mild dependence on the spectral index while the models with Compton heating show strong dependence on the spectral index. (\( \lambda_{\text{rad}} \)) is \( \sim 10^{-4} \) for the model \( l \propto \dot{m}^2 \) and \( \sim 10^{-3} \) for the model \( l \propto \dot{m} \) when \( \alpha = 1 \). The accretion rate for the model \( l \propto \dot{m} \) shows intermediate dependence on the spectral index between the models \( l \propto \dot{m} \) and photo-heating only.
Compton heating in accretion on to BHs

- Compton heating is important near the BH inside the Strömgren sphere, while the photo-heating still prevails in most part of the Strömgren sphere. In particular, we emphasize that the photo-heating dominates the Compton heating at the radial scale of the I-front where the gas supply to the inside the Strömgren sphere is regulated. The ratio of Compton-heating and photo-heating as a function of radius is well explained by the ratio between the ionized and neutral fraction.

- The accretion feedback model $l \propto \dot{m}^2$ shows physically consistent results with the radiatively inefficient ADAF model where the accretion rate is low and the spectrum is hard ($\alpha \sim 1$). $l \propto \dot{m}$ model shows intermediate results between the feedback model $l \propto \dot{m}$ and photo-heating-only model.

We limit the results of the current study to the case of accretion on to BHs from a neutral, dense, and warm medium with $T \sim 10^4$ K. When the gas is hot, ionized and optically thin as found in elliptical galaxies or at the central region of galaxy clusters, Compton heating is regarded as the critical heating mechanism for the BH feedback, especially when the spectrum is hard.

We fix the hardness of the spectrum during a simulation in the current work; however, the state transition of the BH spectrum should play an important role in determining the phenomena discussed here (Gan et al. 2014). Most of the gas accretion happens during the high luminosity phase where the spectrum is soft while the spectrum stays in a hard state during the low accretion phase (Yuan et al. 2009, Yuan & Narayan 2014). Dependence of the state transition on the BH luminosity should be considered in future work.

In this work, we assume an idealised power-law spectrum with a fixed energy range so that we can express the Compton temperature in an analytic manner. However, the high-energy observation of SMBHs reveals that the spectral energy distribution in the range of UV and X-ray shows a deviation from the single power-law. (Sazonov et al. 2004) found a double-peaked distribution from the weighted sum of obscured and unobscured quasar spectrum and calculated the effective Compton temperature accurately $T_c = (1.9 \pm 0.8) \times 10^9$ K. However they point out that the contribution of the blue bump to the Compton heating is less than 25 per cent, which is smaller than the error of the calculated Compton temperature. Additionally, the approximate estimation of the effective spectral slope of the averaged spectrum of the quasars in their work still falls within the range of $1 < \alpha < 1.5$. Considering the effective spectral slope and the Compton temperature, applying the double-peaked spectrum to the simulations is not likely to produce a significant qualitative difference from the current results.
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