BALQSO Spectra Explained by Shock Disruption of Galactic Clouds

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ABSTRACT

Blue-shifted Broad Absorption Lines (BALs) detected in quasar’s spectra are indicative of AGN outflows. We show, using 2D hydrodynamical simulations, that disruption of interstellar clouds by a fast AGN wind can lead to formation of cold, dense high speed blobs that give rise to broad absorption features in the transmission spectrum of the AGN continuum source. For a wind velocity of 0.1c and sufficiently high cloud density ($n_c \approx 10^4$ cm$^{-3}$, depending on size), absorption troughs with velocities up to about 3000 km s$^{-1}$ can be produced. For slower winds and/or lower cloud density the anticipated velocity of the absorbing clouds should be smaller.

Key words: galaxies: active, galaxies: kinematics and dynamics, galaxies: ISM, ISM: jets and outflows, ISM: clouds, quasars: absorption lines, quasars: general

1 INTRODUCTION

Quasars show ubiquitous outflows ($\approx 20$–50% of all AGN; e.g. Hewett & Foltz 2003; Dai et al. 2008), where blueshifted absorption lines are attributed to sub-relativistic ($\approx 10^3$ – $10^4$ km s$^{-1}$) mass ejection. These outflows are a prime candidate for producing various AGN feedback processes: curtailing the growth of the host galaxy, explaining the relationship between the masses of the central black hole and the galaxy’s bulge, and ICM and IGM chemical enrichment (e.g. Ostriker et al. 2010; Faucher-Giguère et al. 2012; Zubovas & Nayakshin 2014; Thompson et al. 2015; Yuan et al. 2018; Zeilig-Hess et al. 2019).

Most dynamical models of BAL winds assume that the outflow originates from the AGN accretion disk at $R \sim 0.01$ pc (about $\sim 10^{16}$ cm) from the central source (e.g., Murray et al. 1995; Proga et al. 2000; Contopoulos et al. 2017). These outflows typically reach 90% of their terminal velocity within 2–4 times the starting $R$.

Observationally, BAL troughs have widths of thousands of km s$^{-1}$, which cannot be explained by just the thermal width. Therefore, the above models must produce the troughs in the acceleration phase where $R \sim 0.1$ pc (e.g., Murray et al. 1995). However, over the past decade, analysis of more than 20 individual outflows measured $R \sim 10$ – 1000 pc (Moe et al. 2009; Dunn et al. 2010; Bautista et al. 2010; Aoki et al. 2011; Borguet et al. 2013; Chamberlain et al. 2015; Chamberlain & Arav 2015; Xu et al. 2018b; Miller et al. 2018). Moreover, two recent surveys demonstrated that half of the typical BAL outflows are situated at $R > 100$ pc (Arav et al. 2018; Xu et al. 2018a).

Since accretion disk wind models are inadequate for producing observed BAL troughs at $R \sim 10$ – 1000 pc, we turn our attention to models that produce the troughs at the distance where they are observed. The pioneering analytical work on such models was conducted by Faucher-Giguère et al. (2012) (henceforth FG12), where the BALs are “formed in situ in radiative shocks produced when a quasar blast wave impacts a moderately dense interstellar clump along the line of sight.” However, the complicated spatial and time-dependent behavior of such interactions limits the usefulness of the analytical approach in deriving the physical characteristics of such outflows.

In this paper we analyze the process of cloud disruption by a wind, and the subsequent evolution and properties of the cloud debris, using 2D hydrodynamical simulations. The output of the simulation is then used to produce synthetic spectra along different lines of sight that mimic observed spectra. In difference from FG12, we consider the interaction of a AGN wind with highly dense ISM clouds. FG12 envisaged that the cloud accelerates to its terminal velocity (roughly the wind velocity) before being broken into cold, dense fragments (cloudlets) that absorb the quasar radiation. However, our simulations indicate that under these conditions the cloudlets do not have time to cool substantially. Moreover, the mean number of absorbing cloudlets anticipated to be seen along a typical line of sight is too small to account for the broad absorption features detected in BAL spectra. As will be shown below, these problems are alleviated when the cloud is sufficiently dense, since (i) the cooling time of the shocked cloud is much shorter, and (ii) the mean number of cloudlets is vastly larger. While, as mentioned above, a moderately dense cloud is disrupted only after acquiring high velocity, a highly dense cloud is dispersed well be-
fore being accelerated significantly. We find that the cloud is nearly completely destroyed by the time the radiative shock has crossed it, and that its debris are accelerated by the wind to high velocities over a similar timescale. Our analysis indicates that for sufficiently dense clouds ($n_c > 10^4 \text{ cm}^{-3}$) this process can produce absorption troughs in the velocity range between 1000 and 3000 km s$^{-1}$, but not much higher.

2 BAL OUTFLOW MODEL

The model considered below posits that a substantial sub-class of BAL outflows are composed of debris of interstellar clouds that have been crushed by a fast AGN wind over a range of distances from the source (parsecs to kiloparsecs) that vary from object to object (see Fig. 1 for illustration). The exact nature of the wind’s origin is not specified and is not important for our analysis; it is merely assumed that the wind propagates from its injection point, near the central engine, to a relatively large distance where it reaches a region in which dense molecular clouds are abundant and existing in a pressure equilibrium with their surroundings. Once the wind (more precisely, the shocked wind bubble) encounters a cloud, a strong shock wave is generated inside the cloud. The shock sweeps the cloud over a crossing time $t_{cs}$, which will be calculated below, whereupon the cloud is broken into fragments (cloudlets). Each of these cloudlets is dragged and accelerated by the ram pressure of the wind to an asymptotic velocity comparable to that of its surrounding bulk flow, over an acceleration time $t_{drag}$, which will also be calculated below. During the dynamical evolution just depicted, the dense cloudlets cool radiatively, and at the same time are heated by repeated shocks and compression waves. This leads to a density-temperature relation down to a temperature at which dense enough cloudlets quickly cool to a temperatures between $10^3$ and $10^5 K$.

At these temperatures there is a significant population of ions that produce the observed BAL (e.g., C IV, see section 4).

A particular BAL system is formed when intervening cloudlets along the line of sight absorb the quasar radiation. The spread of cloudlets encompasses a relatively small volume around the original location of the progenitor cloud and, therefore, the dynamics of the cloudlets is affected only by a small section of the wind and the cloudlets system it is thus sufficient to specify the local wind quantities, rather than attempting to follow the global wind dynamics. This is the essence of the numerical model presented in section 3. Nonetheless, to relate the BAL outflow model to global observables we start with a brief description of the wind evolution in section 2.1.

The dynamics of cloud-wind interaction and cloudlets dragging is studied using 2D hydrodynamical simulations. Each simulation is run for long enough time to allow the cloudlets to accelerate to high velocities (a few thousands km s$^{-1}$). The sizes, densities, velocities and spatial distribution of cloudlets obtained from the simulation are then used to calculate observed spectra along different lines of sight. The details will be described in sections 3 and 4.

2.1 Wind dynamics

Consider a conical AGN wind with a total kinetic power $L_w = 10^{47}L_{\text{w}47}$ erg s$^{-1}$, propagation velocity $v_w = \beta_w c$, and opening angle $\theta_w$. The number density of the unshocked wind changes with distance $r = r_{pc}$ pc from the wind source as

$$n_w = \frac{L_w}{\pi (1 - \cos \theta_w) m_p v_w^2 r_{pc}^2} = 10^2 \frac{L_{\text{w}47}}{(1 - \cos \theta_w) c^3} r_{pc}^{-2} \text{ cm}^{-3},$$

in the region where the wind remains approximately conical (no collimation). The wind is supposed to propagate through an ambient medium with a power-law density profile of the form $n_a = n_{0a} r_{pc}^{-p}$, with the exponent ranging from $p = 0$ for a uniform background to $p = 2$ for an isothermal spherical distribution, and $n_{0a}$ is the density of the ambient medium at a radius of 1 pc, henceforth measured in c.g.s units. The interaction of the wind with the ambient gas leads to formation of a double-shock layer at the wind’s head, within which the shocked wind and the shocked ambient gas are separated by a contact discontinuity (Fig. 1). The shocked layer (more precisely the contact surface) propagates at a velocity that can be determined from local momentum balance to be

$$v_{w,\text{sh}} = v_w/ \left(1 + \sqrt{n_a/n_w}\right) \quad (\text{Zeilig-Hess et al. 2019}),$$

and it is seen that significant deceleration of the wind’s head commences at a radius $r_{dec}$, at which the ambient-to-wind density ratio exceeds unity. For the ambient density profile invoked above one finds:

$$r_{dec} \approx \sqrt[3]{\frac{10^2 L_{\text{w}47}}{(1 - \cos \theta_w) c^3} r_{pc}^{-2}} \text{ pc},$$

and it is seen that for a wind velocity $\beta_w \lesssim 0.1$ and opening angle $\theta_w \approx 30^\circ$ substantial deceleration is anticipated at radii $r \gtrsim 1$ kpc ($L_{\text{w}47}/n_{0a})^{1/2}$, adopting a uniform ambient density ($p = 0$) for illustration. The velocity of the wind’s head (i.e., the double-shock layer) can be expressed in terms of $r_{dec}$ as

$$v_{w,\text{sh}} = \frac{v_w}{1 + (r/r_{dec})^{-2}} \approx v_w (r/r_{dec})^{-2} \quad \text{at } r > r_{dec}. \quad (3)$$

It is worth noting that strong collimation of the wind is anticipated in the deceleration zone (e.g., Zeilig-Hess et al. 2019), that will alter the density profile of the unshocked wind and, consequently, the head velocity $v_{w,sh}$. Thus, at $r > r_{dec}$, $v_{w,sh}$ can be much larger than the value given by Eq. (3). Since observations of BAL outflows reveal cloudlets velocities up to about 0.1c, it seems more likely to assume that the wind encounters the cloud prior to decelerating, at $r \lesssim r_{dec}$, although strong collimation may alter this inference.
2.2 Characteristic timescales

A cloud encountered by the wind will first interact with the shocked ambient gas contained between the forward shock and the contact surface. For a highly supersonic wind (i.e., a strong forward shock), the width of this layer is approximately $\Delta r_f \approx r_s/4$ when the forward shock reaches a radius $r_s$, and its crossing time is

$$t_f \approx \Delta r_f/v_{ws} \approx 10^3 \text{yr} \left(\frac{r_s}{\text{1 kpc}}\right) \beta_{ws}^{-1}. \quad (4)$$

This should be compared with the cloud crashing time which, as confirmed in section 3 by numerical simulations, roughly equals the cross-over time of the shock generated inside the cloud by the cloud-wind collision.1 A strong shock is expected to form inside the cloud if the ram pressure of the shocked ambient flow, $m_p n_{ws} v_{ws}^2/2$, (assuming pure H composition for simplicity), largely exceeds the pressure of the unshocked cloud, $p_i = n_i k T_i$, where $n_i$ and $T_i$ are the initial density and temperature of the cloud. Expressed in terms of the sound speed of the unshocked cloud, $c_s = \sqrt{k T_i / m_p}$, here $\gamma = 5/3$ is the adiabatic index, and the ratio $\kappa \equiv n_i/n_d$ between the cloud density $n_i$ and the (unshocked) ambient medium density $n_d = n_{ws}/4$, the latter condition reads: $\kappa \ll 4 n_{ws}^2 / c_s^2$. For a spherical cloud of radius $R_c$, the shock crossing time of the cloud is given by

$$t_{sc} = \frac{2 R_c}{v_{cs}} \approx \left(\frac{R_c}{v_{ws}}\right) \sqrt{\kappa} \approx t_f \left(\frac{4 R_c}{r_s}\right) \sqrt{\kappa}, \quad (5)$$

noting that for a large density ratio ($\kappa >> 1$) the velocity of the shock inside the cloud is approximately $v_{cs} = v_{ws} \sqrt{n_{ws}/n_c} = 2v_{ws}/\sqrt{\kappa} > c_s$. Thus, as long as $\sqrt{\kappa} < r_s/4 R_c$ the cloud will be swept by the shock before it reaches the contact surface. Otherwise, the disruption will continue in the wind itself (either the shocked or unshocked wind), whereby in the expression for $\kappa$, $n_d$ must be replaced by $n_a$ when the cloud is inside the shocked wind layer or by $4 n_a$ when in the unshocked wind. Our numerical simulations show excellent agreement with Eq.(5).

Adopting for illustration $n_{ws}(R_d) = 1$, $\beta_{ws} = 0.1$ and $\theta_a = 30^\circ$ yields $r_{dec} = 1$ kpc and $t_f \approx 10^3$ yr at $r_{dec}$. A cloud of radius $R_c = 1$ pc will then be crashed at this distance within this time if its density satisfies $n_c \lesssim 6 \times 10^4 n_{ws}$.

Another important timescale is the acceleration time of a dragged cloudlet, defined as the time it takes the cloudlet to reach the velocity of the shocked ambient layer. To be precise, $t_{drag} = v_{ws}/a$, where $a = F/M_a$ is the acceleration due to the drag force exerted on a cloudlet of mass $M_a$ by the ram pressure of the shocked flow, $F = m_p n_{ws} v_{ws}^2 (\pi d^2)$. For a spherical cloudlet of radius $d$ and density $n_{ct}$ the mass is given by $M_a = m_p n_{ct} (4 \pi d^3 / 3)$, yielding

$$t_{drag} \approx \frac{4 n_{ct} d^2}{3 m_p n_{ws} v_{ws}} \approx t_f \left(\frac{16 n_{ct}}{3 \beta_{ws}^2}\right). \quad (6)$$

In practice the dragging time can be shorter by a factor of up to a few if the cloud is oblate. As an example, the acceleration time of a cloudlet having a density comparable to that of the progenitor cloud, $n_{ct} \sim n_c$, relative to the shock crossing time of the pre-crashed cloud is $t_{drag}/t_{sc} \sim (d/R_c)^2 / \sqrt{\kappa}$. In §3 we find that the density distribution of the cloudlets span a wide range, owing to large variations in the confining pressure in the region that contains the fragments of the crashed cloud.

2.3 Cooling

For a (forward shock) velocity of $\beta_{ws} \lesssim 0.1$ the temperature of the shocked ambient medium is $T_{ws} \approx 10^{10} K (\beta_{ws}/0.1)^2$. The dominant cooling process at such high temperatures is free-free emission.

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1 We find this result to hold even when the shock is very weak.
The corresponding cooling time is

\[ t_{\text{cool}} \approx 3 \times 10^8 \text{ yr} \left( \frac{T_{\text{as}}}{10^4 \text{ K}} \right)^{1/2} \frac{1}{n_{\text{as}}}, \quad (7) \]

much longer than the wind expansion time. However, the dense cloud material cools over a much shorter time. Assuming the shocked cloud material to be in pressure balance with its surroundings, its temperature is \( T_{\text{as}} = T_{\text{as}}(n_{\text{as}}/n_\text{e}) \approx T_{\text{as}}/\kappa \), and its free-free cooling time is shorter by a factor \( \kappa^{3/2} \) than the cooling time of the shocked ambient gas, adopting for illustration the jump conditions of a strong, non-radiative shock (\( n_\text{e} = 4n_\text{as} \)), for which

\[ n_{\text{as}}/n_\text{e} = 4n_\text{as}/4n_\text{e} = 1/\kappa. \]

The above estimate holds at temperatures \( T_{\text{as}} > 4 \times 10^4 \text{ K} \). At lower temperatures, \( 10^3 \lesssim T \lesssim 4 \times 10^4 \), the cooling rate is considerably enhanced. For an ionized, optically thin plasma the cooling rate at \( T = 10^6 \text{ K} \) is larger by nearly two orders of magnitude than the free-free rate, depending on metallicity (e.g., Sutherland & Dopita 1993; Gnat & Ferland 2012). To avoid unnecessary complications, we adopt in our simulations a cooling function of the form

\[ \Lambda(T) = 1.4 \times 10^{-27} \eta T^{1/2} \text{ erg s}^{-1} \text{ cm}^{-3}, \quad (8) \]

with \( \eta > 1 \) being a free parameter (\( \eta = 10 \) is adopted in the fiducial case discussed in §3). This function overestimates the cooling rate in the shocked ambient gas and underestimates the cooling rate of the cloud and cloudlets. However, even with this correction the cooling time of the shocked ambient gas is much longer than flow expansion time, so it doesn’t alter the dynamics. From Eqs. (5) and (8) with \( \eta = 10 \) we find that the cooling time of the shocked cloud will be shorter than the shock crossing time \( t_{\text{sc}} \), if the cloud-to-ambient density ratio satisfies \( \kappa > 500/(R_c/1 \text{ pc})^{-1/2}(\beta_{\text{as}}/0.1)^{1/2} \). This condition is fulfilled in all cases considered below. In those cases the shock inside the cloud is radiative and the shocked cloud is compressed to a density of \( n_{\text{as}} \approx 10^7 n_\text{as}(\beta_{\text{as}}/0.1)^2(T_{\text{as}}/10^4 \text{ K})^{-1} \).

2.4 Cloudlets statistics

A key question concerning the BAL model proposed here is how many cloudlets are expected to be detected along a given sightline on average. A simple estimate can be made upon assuming that most of the mass of the original cloud is divided between identical cloudlets, each having a radius \( d \) and density \( n_{\text{as}} \). The total number of cloudlets is then \( N_c = (n_c/n_\text{as}) (R_c/d)^3 \), where \( p = 2 \) for a 2D system (as in our simulations below) and \( p = 3 \) in 3D. Now, let us choose \( x \) to be the direction of the wind prior to collision with the cloud, viz., \( \beta_{\text{as}} = \beta_{\text{as}} x \), take the center of the cloud to be at the origin, and denote by \( D \) the distance from the \( x \)-axis (along the perpendicular direction) within which the majority of cloudlets are located long after accelerating to high velocities. Then, the mean number of cloudlets anticipated along any line of sight within \( D \) is:

\[ N_c(d/D)^{p-1} = (n_c/n_\text{as}) (R_c/d)^{p-1} (R_c/d). \]

Typically, \( D/R_c \sim a \) few \((R_c/D)\) ranges from about 2 to 5 in the simulations presented below, depending on \( \kappa \). Thus, for a size ratio of \( R_c/d = 10^2 - 10^3 \), between a few to a few tens cloudlets should be seen along any sightline that intersects the cloudlet cluster (that, is, within \( D \)) if \( n_c \approx n_\text{as} \). However, as will be shown in §3, the density ratio \( n_c/n_\text{as} \) depends on the distribution of the confining pressure behind the pre-collapsed cloud, which in turn depends on

\[ n_c/n_\text{as} \sim 10^{-3}, \quad (2) \]

In the model of Faucher-Giguère et al. (2012) \( n_c/n_\text{as} \sim 10^{-3} \), hence less than one cloudlet is expected to be seen along any line of sight on average.

Figure 3. Temperature (top) and pressure (bottom) distributions in the fiducial simulation at \( t = 65 \text{ kyr} \). The arrows indicate the velocity vectors of the flow.

3 NUMERICAL SIMULATIONS

The calculations of the cloud-wind collision process have been performed using 2D hydrodynamical simulations. A flow is injected from a planar boundary of a lateral extent large enough to contain the spread of cloud debris but small compared with the wind radius, and collides with a uniform heavy cloud initially embedded in a uniform, dilute ambient gas representing the galactic medium. We denote the velocity and density of that flow by \( \beta_{\text{as}} \) and \( n_\text{as} \), respectively, but note that this flow may represent either the shocked ambient flow, if cloud crash and cloudlets acceleration occur in that region, or the wind itself if the shock crossing time of the cloud is considerably longer than the wind expansion time \( t_f \) given in Eq. (4). The planar approximation of the initial flow is justified by the small ratio between the cloud radius to the putative cloud’s distance from the wind’s source, as explained in the previous section. Cooling of the gas is also included in the simulations in a manner specified below. Each simulation has been run for long enough time to allow complete destruction of the cloud and subsequent acceleration of the cloudlets to high velocities.

The output data of the simulation gives the spatial distribution of cloud fragments at any given time, their velocities, densities and temperatures. This data is used to produce column density histograms in velocity space along different lines of sight, that serve us later (section Eq. (9)) in creating synthetic spectra that can be compared with observed BAL spectra. It also enables us to get a grasp on the general dynamical and thermal evolution of the cloud fragments, and to understand which of the following processes happens faster in different stages of the system: further fragmentation accompanied by shock heating versus cooling.
We have made runs with a bremsstrahlung power law dependence of the free-free process (see section 2.4 for discussion). We performed several experiments differing in the size of the simulation box, the implementation of cooling and the ratio of densities, $\kappa$, between the cloud and the wind. In each of these experiments, the ranges of the Cartesian coordinates are of the form $0 < x < x_{\text{max}}$ and $|y| < y_{\text{max}}$, the unit length is equal to 1 pc and the resolution along each Cartesian direction is 0.01 pc. The initial condition of the simulation is an ideal gas with an adiabatic index $\gamma = 5/3$. We performed several experiments with different wind velocities, cloud densities, density ratio $\kappa$, and cooling parameter $\eta$. For our fiducial simulation we use $\beta_{\text{wr}} = 0.1$, $\eta = 10$, $\kappa = 4 \times 10^3$ and cloud density $n_c = 10^6$. Cooling is implemented by utilizing the Power_Law switch. The left panel in Fig. 4 shows the temperature-density phase diagram at $t = 65$ kyr for all fluid elements in the fiducial simulation box. The relation $n \propto T^{-1}$ revealed in the diagram indicates pressure confinement. The large scatter is due to the pressure variation inside the Mach cone, as seen in Fig. 3. The presence of points below the floor temperature is due to adiabatic expansion of some dense cloud fragments. Right: density-velocity diagram of the same matter. The red points designate fluid elements at temperature $T < 10^6$ K. As seen, the projected velocity of most fluid elements lies in the range $1000 \text{ km s}^{-1} < v_x < 3000 \text{ km s}^{-1}$.  

3.1 Setup

We make use of the PLUTO code (Mignone et al. 2007) version 4.0 in 2D Cartesian coordinates in the non-relativistic HD module on an ideal gas with an adiabatic index $\gamma = 5/3$. We performed several experiments differing in the size of the simulation box, the implementation of cooling and the ratio of densities, $\kappa$, between the cloud and the wind. In each of these experiments, the ranges of the Cartesian coordinates are of the form $0 < x < x_{\text{max}}$ and $|y| < y_{\text{max}}$, the unit length is equal to 1 pc and the resolution along each Cartesian direction is 0.01 pc. The initial condition of the simulation is an ideal gas with an adiabatic index $\gamma = 5/3$. We performed several experiments with different wind velocities, cloud densities, density ratio $\kappa$, and cooling parameter $\eta$. For our fiducial simulation we use $\beta_{\text{wr}} = 0.1$, $\eta = 10$, $\kappa = 4 \times 10^3$ and cloud density $n_c = 10^6$. Cooling is implemented by utilizing the Power_Law switch. The left panel in Fig. 4 shows the temperature-density phase diagram at $t = 65$ kyr for all fluid elements in the fiducial simulation box. The relation $n \propto T^{-1}$ revealed in the diagram indicates pressure confinement. The large scatter is due to the pressure variation inside the Mach cone, as seen in Fig. 3. The presence of points below the floor temperature is due to adiabatic expansion of some dense cloud fragments. Right: density-velocity diagram of the same matter. The red points designate fluid elements at temperature $T < 10^6$ K. As seen, the projected velocity of most fluid elements lies in the range $1000 \text{ km s}^{-1} < v_x < 3000 \text{ km s}^{-1}$.

3.2 Results

We performed several experiments with different wind velocities, cloud densities, density ratio $\kappa$, and cooling parameter $\eta$. For our fiducial simulation we use $\beta_{\text{wr}} = 0.1$, $\eta = 10$, $\kappa = 4 \times 10^3$ and cloud density $n_c = 10^6$. Cooling is implemented by utilizing the Power_Law switch. The left panel in Fig. 4 shows the $n - T$ relation of all cells in the simulation box. The slope of this relation, $n \propto T^{-1}$, indicates that cloudlets are in pressure balance with their local environment. However, there is a large spread in the density of the filaments, that may look suspicious at first glance. The reason for this is the...
large variation in the pressure behind the crashed cloud (see bottom panel in Fig. 3), which is caused by the deflection of wind streamlines that engulf the dense cloud material. From Eq. (6) we obtain \( t_{\text{drag}} = 3000 \, \text{yr} (\beta_{\text{drag}}/0.1)^{-1} (n_{c}/10^5 \, \text{cm}^{-3}) (d/0.01 \, \text{pc}) \) for a cloudlet of density \( n_{c} \) and size \( d \) at the floor temperature. Noting that the velocity of the background flow that drags the dense cloudlets is smaller than that of the injected wind by a factor of a few, we find that sufficiently small cloudlets of density \( n_{c} \lesssim 10^5 \) can accelerate to the terminal velocity within about 30,000 years, while the dragging time of considerably denser cloudlets is much longer, as indeed seen in Fig. 4. The \( n - T \) plot also reveal that all cloudlets with densities \( n_{c} \gtrsim 10^5 \, \text{cm}^{-3} \) quickly cool to the floor temperature. Cloudlets of lower densities exhibit a wide range of temperatures, presumably due to heating by repeated shocks. In practice, we expect this sharp drop in temperature seen at \( n \approx 10^5 \, \text{cm}^{-3} \) in Fig. 4 to occur at lower densities, since the cooling rate is likely to be substantially higher than that invoked in our simulations. The cooling of cloudlets to temperatures below the floor value, \( T_{f} = 3 \times 10^5 \, \text{K} \), might be caused by adiabatic expansion of dense cloudlets as they move into a low pressure zone.

The qualitative behaviour of the system seen in the fiducial simulation is quite typical. In all cases with \( \beta_{\text{drag}} = 0.1 \) we find that most cold (\( T_{c} \lesssim 10^6 \, \text{K} \)) dense cloudlets accelerate to velocities in the range \( 1000 - 3000 \, \text{km} \, \text{s}^{-1} \) right after shock crossing, at \( t \gtrsim 2t_{\text{cc}} \), with much slower acceleration at later times. A very small fraction of the cloudlets accelerated to even higher velocities, up to 5000 km s\(^{-1}\) in the fiducial case. For slower winds (injected velocity \( \beta_{\text{drag}} = 0.03 \)) the terminal velocity of cloudlets was found to be somewhat smaller. The total number of cloudlets depends also on the density of the pre-crashed cloud through the ratio \( n_{c}/n_{c_{f}} \), as explained in section 2.4 above.

As another example (case B) we show in Fig. 5 the density map and \( n - v_{c} \) diagram at time \( t = 6000 \, \text{yr} \) for a simulation with the same wind parameters and cooling (\( \beta_{\text{drag}} = 0.1, \eta = 10 \)), but a lower cloud density, \( n_{c} = 10^4 \, \text{cm}^{-3} \). The corresponding density ratio is \( \kappa = 4 \times 10^3 \). The cloud crashing time in this case is \( t_{\text{cc}} = 3000 \, \text{yr} \), as expected, and it is seen that as in the fiducial simulation, cloudlets reach their terminal velocities at \( 2t_{\text{cc}} \). We have run the simulation up to a longer time but didn’t observe a significant change in the velocity distribution of the cloudlets at times \( t > 6000 \, \text{yr} \). The vertical spread (along the y axis) is \( D \simeq 2R_{c} \) (Fig. 5), smaller than that found in the fiducial simulation. This is anticipated given the smaller density ratio \( \kappa \) in this run. The total number of dense cold cells is about 600, which translates to roughly 2 cloudlets along sight lines in the range \( -2 < y/R_{c} < 2 \), compared with hundreds in the fiducial case. However, as explained in section 2.4, this number is proportional to the initial cloud radius \( R_{c} \), hence, for a considerably larger pre-crashed cloud many more cloudlets should be seen in this case along sight lines that intersect the cloudlets cluster.

### 3.3 Resolution

The size of cloudlets in our simulations is limited by spatial resolution. With our grid spacing the smallest cloudlets have size of 0.01 pc. This means that further disruption of these cloudlets may be artificially suppressed. To check the effect of numerical resolution on our results we computed the size distribution of cloudlets in the simulation box in the fiducial case at time \( t = 65 \, \text{kyr} \). Fig 6 shows the mean size distribution of cloudlets at different densities, where the mean size \( d \) is defined as \( d = \sqrt{A_{c}} \) in terms of the cloudlet area \( A_{c} \), so that \( (d/0.01 \, \text{pc})^2 \) is essentially the number of cells contained in the cloudlet. We find that only 18% of the cloudlets contain one

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**Figure 5.** Density map (left) and \( n - v_{c} \) diagram (right) at time \( t = 6000 \, \text{yr} \) for case B simulation.

**Figure 6.** Mean size distributions of cloudlets at different densities. The numbers on the horizontal axis (i.e., the values of \( d/(0.01 \, \text{pc}) \)) give the square root of the number of cells in the cloudlet.
cell, and that over 50% contain at least 10 cells. This means that absorption occurs predominantly in resolved cloudlets\(^3\).

4 SYNTHETIC SPECTRA

This section outlines the method employed to calculate absorption spectra along given lines of sight, using the simulation data. Several representative examples of such spectra will be exhibited for illustration. A more comprehensive analysis of absorption spectra and comparison to observations will be presented in a follow-up paper. For simplicity, we suppose that the wind is conical (as opposed to strongly collimated), so that absorption of the quasar continuum radiation is seen along directions parallel to the wind streamlines. Within our planar approximation, this means sight lines along the \(x\)-direction (each marked by a value of the \(y\)-coordinate). Examples are shown in Fig 7. As explained in the preceding section, the planar approximation is justified by the fact that the observer is situated at a distance much larger than the size of both the AGN source, the original pre-shocked cloud and the system of spread-out cloudlets. We further assume that, in practice, fluid elements that have reached temperatures \(T < 10^6\) K in the simulation will cool quickly to lower temperatures, \(T \approx 10^4\) K, by virtue of the much faster cooling anticipated (see section 2.3 for discussion). When computing column densities that contribute to absorption we select only matter having \(T < 10^6\) K.

We shall focus here on the fiducial simulation, noting that the other cases lead to similar results, provided cloudlets are sufficiently abundant (see section 2.4 for a detailed discussion). The lines of sight indicated on the density map in Figure 7 will be used to illustrate the method and produce sample spectra.

The first step in the construction of synthetic BAL-spectra is the calculation of the column-density as a function of the projected velocity along the observed direction, here \(v_x\). Once the column-density is obtained, the optical depth for absorption at the corresponding wavelength will be calculated, out of which the normalized flux would be derived. For any given sight line we integrate the number density of cold matter (\(T < 10^6\) K) within a given velocity bin along the \(x\)-direction. In order to preserve the information of velocity dependence, we construct a 3D array of number-density values, distributed in 2D configuration-space (i.e., in the \(x-y\) plane) and in 1D projected velocity space, with a velocity bin size of \(\Delta v_x\). Only after obtaining the phase-space distribution of number-densities, we perform the horizontal integration to obtain the column density of matter moving in the velocity range \((v_x, v_x + \Delta v_x)\), along a particular \(y\)-line:

\[
N(y,v_x,\Delta v_x) = \int_0^{\Delta x} dx \int_{v_x}^{v_x + \Delta v_x} \frac{dn(x,y,v_x)}{dv_x} dv_x. \tag{9}
\]

In the integration we utilize a measure equal to the spatial resolution limit of the simulation, \(\Delta x = 0.01\) pc. The resulting column-density (henceforth \(N_H\)) distributions in horizontal-velocity space with a bin size of \(\Delta v_x = 20\) km s\(^{-1}\) (chosen to match the resolution of the COS spectrograph on HST) are shown in the left panels in Figure 8 for the four lines indicated in Fig. 7 \((y = -1.4, -4.0, -6.0, -6.6)\). It reveals substantial column densities up to \(v_x \lesssim 3000\) km s\(^{-1}\).

The sight-line analysis confirms that the typical values of \(N_H\) and speeds do not change much in the vertical direction, up until reaching large enough transverse coordinates \(|y| \gtrsim 6\), after which a significant drop in the number of blobs is seen. This means that at the late stage analyzed here, most of the cold-dense blobs are dispersed into a region of size about 6 times that of the original cloud, as indeed seen by eye in the density map, Fig. 7. The mean number of cloudlets identified along sight lines within the cloudlets cluster is consistent with the naive estimate in section 2.4.

The most ubiquitous BAL in quasar spectra arises from \(\text{C}IV\). Moreover, the defining criteria of a BAL depends on the \(\text{C}IV\) trough having velocity width larger than 2000 km s\(^{-1}\) at normalized residual intensity \(I/0.9\) (Weymann et al. 1991). We therefore produce a synthetic spectrum that shows the \(\text{C}IV\) absorption that is derived from our fiducial simulation. The observed \(\text{C}IV\) BAL arise from the doublet transitions at 1548.19\(\AA\) and 1550.77\(\AA\). To highlight the simulation’s results, we only use the stronger transition at 1548.19\(\AA\). In reality, when the velocity width of the outflow is larger than the equivalent 500 km s\(^{-1}\) velocity separation between the doublet components, the trough will be a blend of absorption from both transitions.

To produce a synthetic \(\text{C}IV\) BAL from our simulation we need to:

a) Determine \(N_{\text{CIV}}(y,v_x)\) from the simulation’s \(N_H(y,v_x)\) (see Eq. (9) and accompanied discussion).

b) Derive the optical depth \(\tau_{\text{CIV}}(y,v_x)\) associated with \(N_{\text{CIV}}(y,v_x)\).

c) Create a synthetic spectrum from the convolution of expected absorption associated with \(\tau_{\text{CIV}}(y,v_x)\) and the emission sources of the quasar.

To Determine \(N_{\text{CIV}}(y,v_x)\), we first need to take into account the abundance of carbon relative to hydrogen. For this purpose we use standard solar metallicity (Grevesse et al. 2010). Second, we need to determine the fraction of \(\text{C}IV\) to all carbon atoms. Ionization equilibrium in a quasar outflows is dominated by photoionization caused by the ionizing flux of the quasar (e.g., Arav et al. 2001). For illustrative purposes we use the photoionization solution of the outflow observed in quasar HE 0238–1904 (Arav et al. 2013). The solution was derived using the spectral synthesis code Cloudy [version c17.00, Ferland et al. (2017)]. As input we used the spectral energy distribution (SED) constructed from the observations of the quasar. The solution that best fits the data has \(\log(U_{\text{ion}}) = 0.3\) and \(\log(N_{\text{H}1}) = 20.8\) (log cm\(^{-3}\)) (Arav et al. 2013), the photoionization solution yields \(N_{\text{CIV}}/N_{\text{H}1} \approx 9.0 \times 10^{-8}\). Therefore, we use:

\[
N_{\text{CIV}}(y,v_x) = 9.0 \times 10^{-8} N_{\text{H}1}(y,v_x) \tag{10}
\]

---

\(^3\) Note that a large cloudlet can contribute to absorption along several lines of sight.
Figure 8. The left panels exhibit column densities of cold gas as a function of projected velocity $v_x$ along the lines of sight $y = -1.4, -4.0, -6.0, -6.6$ (from top to bottom). The width of each bar in the histogram corresponds to a velocity bin of $\Delta v_x = 20 \, \text{km s}^{-1}$. The right panels show the corresponding transmission spectrum.
The corresponding optical depth ($\tau$) of the CIV absorption can be expressed as (see equation (8) in Savage & Sembach 1991):

$$\tau_{\text{CIV}}(y, v_y) = 2.654 \times 10^{-15} f \lambda \times N_{\text{CIV}}(y, v_y)$$  \hspace{1cm} (11)

where for a CIV trough, $\lambda = 1548.19\\AA$ and $f = 0.19$ is the wavelength and the oscillator strength of the transition, respectively; and $N_{\text{CIV}}(y, v_y)$ is in ions cm$^{-2}$(km s$^{-1}$)$^{-1}$.

Quasar's ultraviolet emission is composed of a continuum source and a broad emission line (BEL) source, where the size of the latter is at least a hundred times larger than the former. Therefore, it is not surprising that in many cases the outflow covers the continuum source entirely, but only a negligible portion of the BEL emission (e.g., Arav et al. 1999). Our synthetic spectra assume this empirical result: the absorber fully covers the continuum, but does not cover the BEL.

The flux of the simulated CIV absorber is therefore given by the absorbed continuum (normalized to 1):

$$I_{\text{CIV}}(y, v_y) = e^{-\tau_{\text{CIV}}(y, v_y)}$$  \hspace{1cm} (12)

and an added unabsorbed BEL whose parameters are: velocity centroid at 0 km s$^{-1}$, $\sigma = 1000$ km s$^{-1}$, and the maximum intensity of 0.5.

The derived synthetic spectra for CIV troughs at specific line of sight $y$ are shown as the histograms in the right panels of Figure 8. The synthetic absorption spectra we obtained shows a diversity in the trough widths. Examples of broad absorption troughs can be seen in panels (a) and (d) of Figure 8, where the trough's width spans over more than 2000 km s$^{-1}$. However, many of the troughs we obtain are narrower, an example of which can be seen in panel (c) of Figure 8. We note that the velocity width of observed quasar outflows, with a width larger than 1000 km/s, has a strong peak around 1500 km/s (see section 4.3 and figure 6 in Trump et al. 2006). Finally, we also find marginally broad troughs of widths $\lesssim$2000 km s$^{-1}$, an example of which is seen in panel (c) of Figure 8. A more systematic analysis of the lines of sight in terms of BAL widths statistics as well as a direct comparison to observations will be presented in a follow-up paper (Xinfeng et al. in preparation).

5 CONCLUSIONS

In this work we constructed a numerical model for the interaction of AGN wind with galactic clouds in an attempt to explain the observed BAL QSO spectra. The consideration of this model is motivated by recent observations that seem to indicate that a substantial fraction of the BAL outflows are located at large distances from the central AGN - tens to hundreds parsecs.

The wind-cloud interaction was modelled using 2D hydrodynamical simulations that incorporate a simple cooling function. The evolution of the system was followed for sufficiently long time to allow complete destruction of the cloud by the engulfing wind and further acceleration of its debris to high velocities. The output data of the simulation was rearranged and integrated along lines of sight in the direction of the wind propagation, and then used to compute synthetic spectra by accounting for absorption of the AGN continuum radiation by the accelerated, cool cloud fragments.

We find that while a fraction of the cloud fragments reach velocities close to the injected wind velocity, the cool dense cloudlets that can absorb the quasar light reach significantly lower velocities (as seen in Figs. 4 and 5), owing to the complex flow pattern within the Mach cone. We also find that the mean number of absorbing cloudlets along lines of sight that intersect the cloud debris depends on the density and radius of the pre-crashed cloud. For a wind velocity $v_{\infty} \approx 0.1c$ and sufficiently dense cloud the synthetic spectra that we obtained mimic the features of observed BAL spectra in terms of acquired velocities (blue-shifted wavelengths) and in terms of column densities (depths of the absorption troughs). In some of the lines of sight tested we find broad trough widths up to $\sim 3000$ km s$^{-1}$, whereas along other lines we find narrower troughs. Further analysis and direct comparisons to observations will be presented in a follow-up paper (Xinfeng et al. in preparation).

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