New insight on Young Stellar Objects accretion shocks
– a claim for NLTE opacities –

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ABSTRACT

Context. Accreted material onto CTTSs is expected to form a hot quasi-periodic plasma structure that radiates in X-rays. Simulations of this phenomenon only partly match with observations. They all rely on the assumption that radiation and matter are decoupled, and use in addition a static model for the chromosphere.

Aims. We test the validity of these two assumptions in refining the physics included in existing 1D models, and we propose guides for further improvement.

Methods. We simulate accretion columns falling onto a stellar chromosphere using the 1D ALE code AstroLabE. This code solves the hydrodynamics equations along with the two first momenta equations for radiation transfer, with the help of a dedicated opacity table for the coupling between matter and radiation. We derive the total electron and ions densities from collisional-radiative NLTE ionization equilibrium.

Results. The chromospheric acoustic heating has an impact on the duration of the cycle and on the structure of the heated slab. In addition, the coupling between radiation and hydrodynamics leads to a dynamical heating of the accretion flow and the chromosphere, leading to a possible unburial of the whole column. These two last conclusions are in agreement with the computed monochromatic flux. Both effects (acoustic heating and radiation coupling) have an influence on the amplitude and temporal variations of the net X-ray luminosity.

Key words. Stars: pre-main sequence – Accretion, accretion disk – Methods: numerical – Hydrodynamics – Radiative transfer – Opacity

1. Introduction

1.1. Accretion onto CTTSs

Classical T Tauri Stars (CTTSs) are solar-type pre-main sequence stars surrounded by a thick disk composed of gas and dust (see e.g. Feigelson & Montmerle 1999). Disk material follows a near-Keplerian infall down to the truncation radius, at which thermal and magnetic pressures balance. Free-falling material flows then from the inner disk down to the stellar surface in magnetically confined accretion columns (Calvet & Gullbring 1998). Hot spots observations (Gullbring et al. 2000) suggest filling factors of up to 1% (Bouvier et al. 1995).

Accreted gas is stopped where the flow ram pressure and the thermal pressure of the stellar chromosphere balance: an accretion shock forms and the post-shock material accumulates at the basis of the column. The hot slab of post-shock material is separated from the accretion flow by a reverse shock. A typical simulated structure of an accretion shock can be found e.g. in Orlando et al. (2010) and is sketched in Figure 1.

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This shock disappears rapidly during the development of the reverse shock, and forms back during the collapse phase of QPOs (one may compare Figure 9 at $t = 97\ s$ and $t = 140\ s$ for instance). Most of the time, according to simulations, the accretion shock is replaced by a contact discontinuity. The remaining shock, i.e. the reverse shock, is sometimes confusingly called accretion shock in the literature.

![Fig. 1. Sketch of the basis of an accretion column and its three distinctive zones: the chromosphere (left, dark grey), the accretion flow (right, mid-grey) and the zone in between (middle, light grey) hereafter called hot slab or post-shock medium.](https://example.com/figure1.png)
1.2. X-ray observations and simulations

The main signature of the accretion process comes from the X rays emitted by the dense ($n_e > 10^{11} \text{ cm}^{-3}$) and hot ($T_e \sim 2–5 \text{ MK}$) post-shock plasma (see e.g. Kastner et al. 2002; Stelzer & Schmitt 2004 for TW Hya, Schmitt et al. 2005 for BP Tau, Günther et al. 2006 for V4046 Sgr, Argiroffi et al. 2007, 2009 for MP Muscae, Robrade & Schmitt 2007 for RU Lup and Huenemoerder et al. 2007 for Hen 3-600). Another signature is the UV - optical veiling, which is attributed to the post shock medium, the heated atmosphere and the pre-shock medium (Calvet & Gullbring 1998). In addition, Doppler profiles of several emission lines trace the high velocity in the funneled flow (Muzerolle et al. 1998).

1D hydrodynamical models (Sacco et al. 2008, 2010) predict Quasi-Periodic Oscillations (QPOs) of the post-shock slab with periods ranging from 0.01 to 1000 s, depending on the inflow density, metallicity, velocity and inclination of the stream with respect to the stellar surface. For a typical free-fall radial velocity of 400 km s$^{-1}$, Sacco et al. (2010) found for instance a period of 160 s at $10^{13} \text{ cm}^{-3}$. These oscillations are triggered by the cooling instability (for further details, see e.g. Chevalier & Imamura 1982; Walder & Folini 1996; Mignone 2005).

Although plasma characteristics derived from X-ray observations are consistent with the density and the temperature predicted by these numerical studies, there is no obvious observational evidence for such periodicity. Drake et al. (2009) studied thoroughly soft X-ray emission from TW Hydrae and found no periodicity in the range 0.0001–6.811 Hz. Günther et al. (2010) completed this study with optical and UV emission, and they come to the same conclusion in the range 0.02–50 Hz. However, a recent photometric study of TW Hya based on MOST satellite observations reports possible oscillations with a period of 650–1200 s which could be assigned to post-shock plasma oscillations (Siwak et al. 2018).

Observations thus raise the question of the existence of an oscillating hot accretion in the accretion context. Several numerical studies explored multi-dimensional magnetic effects, like leaks at the basis of the column (Orlando et al. 2010), the tapering of the magnetic field (Orlando et al. 2013), or perturbations in the flow (Matsakos et al. 2013). Although QPOs are still obtained in these numerical studies, the accretion funnel basis is either fragmented in out-of-phase fibrils, or buried under a cooler and denser gas layer that strongly absorbs X rays (Sacco et al. 2010; Bonito et al. 2014; Colombo et al. 2016; Costa et al. 2017). The global QPO behavior is therefore suppressed and the measured X-ray emission of the post-shock structure is then systematically low compared to the one derived from simulations (Curran et al. 2011).

1.3. Radiation effects

In these numerical works, the accretion is supposed to take place on a quiet medium (an isothermal atmosphere in the best cases). Moreover, the post-shock medium is assumed to be optically thin, and the coupling between radiation and matter is reduced to a gas cooling function (see e.g. Kirienko 1993, and Figure 3). Although this assumption can be justified to model the infalling gas and the post-shock plasma, it is inconsistent with any stellar atmosphere model. The energy balance between radiation and gas in the lower stellar atmosphere is then replaced by a non-physical tuning (heating function, off threshold, ...). Such an assumption may affect the burial of the post-shock structure as well as the accretion structure itself.

In this work, we focus and refine the physics encompassed in existing 1D models. We first explore the effect of chromospheric shocks perturbations on the accretion dynamics. In a second stage, we analyze the role of radiation absorption/emission on the column, and in the emerging spectra. In Section 2, we present the radiation hydrodynamics model and the numeric tool we used for the hydrodynamics and the spectra synthesis. We detail in Section 2.2.3 the two extreme radiative regimes encountered in this context, and a simple model for intermediate radiative regimes. Section 3 is dedicated to accretion simulations and to the corresponding discussions. The last section (Section 4) presents caveats and possible improvements to this work.

2. Physical and numerical models

2.1. Hydrodynamics model

2.1.1. Hydrodynamics equations

We consider a star of radius $R_\star$ and mass $M_\star$. The accreted and stellar atmospheric plasmas at position $r$ ($r = |\mathbf{r}|$), hereafter taken from the stellar surface, are characterized by a (volumetric mass) density $\rho$, a velocity $\mathbf{v}$, a thermal pressure $p$, and a volumetric internal energy density $e$. The plasma evolution is modeled by solving the hydrodynamics equations, written in the conservative form:

$$\partial_t \mathbf{q}(\mathbf{r},t) + \nabla(\mathbf{q}(\mathbf{r},t) \cdot \mathbf{v}(\mathbf{r},t)) = 0,$$

with $\mathbf{q}(\mathbf{r},t) = (\rho, \rho \mathbf{v}, \rho e, \mathbf{S}_M, \mathbf{S}_E, \mathbf{q}_E, \mathbf{q}_C, \mathbf{q}_K)$, where $\mathbf{S}_M$ and $\mathbf{S}_E$ are the total number density of free particles (neutrals, electrons, and ions, see Section 2.2.3). The closure relation for this system of equations – the equation of state – is adapted from the ideal gas law: $p = n g T$, $e = 3/2 p$, and the chemical composition of the plasma is given by $g = G M_\odot/(R_\star + r)^2$, $r/t$. The gas source terms $\mathbf{S}_M$ and $\mathbf{S}_E$ include the contributions of thermal conduction ($q_C$, Spitzer & Härm 1953; Vidal et al. 1995), artificial viscosity ($q_{vis}$ and $q_{vis}$, von Neumann & Richtmyer 1950) and the coupling with radiation ($\mathbf{S}_M$ and $\mathbf{S}_E$, see Section 2.2.3). The closure relation for this system of equations – the equation of state – is adapted from the ideal gas law: $p = n g T = \varepsilon = 3/2 p$, where $n g$ stands for the total number density of free particles (neutrals, electrons, and ions, see Section 2.1.2). $T$ represents their kinetic temperature $T$.

The effect of collisional ionization and 3-body recombination processes (see Section 2.1.2) on the gas energy density is included through their corresponding rates in the thermochemistry term $q_C$.

2.1.2. Collisional-radiative ionization

The accretion shock forms where the ram pressure is balanced by the local thermal pressure: the stellar chromosphere, that needs to be modeled. In contrary to the solar case, there is a very limited information about T Tauri chromospheres. Thus, as our goal is to propose a qualitative description of the dynamics of this chromosphere, and in absence of any reliable information, our chromospheric model (see Appendix B) is inspired from the solar case: therefore, we have chosen to use solar parameters in our simulations and the chemical composition $q_C$.

A sink is algebraically identified as negative source term. All particles are assumed here to have the same kinetic temperature, i.e. $T_{\text{neutral}} = T_{\text{ion}} = T_{\text{electron}}$. Accreted material is expected to be depleted in heavy elements (Fitzpatrick 1996). However, this phenomenon is not included in this study.
(solar abundances) is then taken from Grevesse & Sauval (1998).

To estimate the total free electron density \( n_e \), we only consider hydrogen (H i, H ii) and helium (He i, He ii, He iii). The chemical composition is completed by a "catch-all" metal "M"\(^5\). For the evaluation of \( n_e \), most simulations are performed using a modified Saha equilibrium model (Brown 1973); the last simulation presented in this paper (referred to as the Hybrid setup) uses a time-dependent collisional-radiative ionization model with:

- **collisional ionization** rates given by Voronov (1997);
- **radiative recombination** rates computed by Verner & Ferland (1996);
- **helium dielectronic recombination** rate proposed by Hui & Gnedin (1997);
- **photoionization** rates derived from Spitzer (1998) and Yan et al. (1998) cross-sections, and the local radiation energy density.

Ions and neutral number volumetric densities \( n_n \) are then computed by a conservative set of equations (see e.g. (1)). The electron number volumetric density is then derived from the neutrality conservation: \( n_e = n_{H^+} + n_{He^+} + 2 n_{He^{++}} \).

This calculation is performed independently from the opacity computation (see Appendix A), that uses a more refined version of the chemical composition (Grevesse & Sauval 1998).

### 2.2. Radiation model

#### 2.2.1. Radiation and hydrodynamics

The coupling between radiation and matter enters at different scales in astrophysical plasmas. At a microscopic scale, radiation affects the thermodynamical state of the matter through its contribution to the populations of the electronic energy levels of each plasma ion. The computation of these populations is based on a large set of kinetic equilibrium equations that take into account excitation and de-excitation processes due to collisions (interactions with massive particles, mostly electrons) as well as radiative processes (interactions with photons). This step allows to derive also the monochromatic absorption/emission coefficients \( k_i \) (also called monochromatic opacity) which, in turn, are used to compute the local radiation intensity by solving the equations of radiative transfer. Two limiting (and simplifying) cases are expected: at large electron densities, one recovers the Local Thermodynamic Equilibrium (LTE), whereas at low density and for an optically thin medium, the coronal limit is reached (Oxenius 1986).

The main issue in performing such calculations is an intricate coupling between the kinetic equilibrium equations (which would be easy to solve had the radiation field is known), and the radiative transfer equation (which would be easy if the atomic level populations, and hence the absorption and emission coefficients, are known). Since a mean free path of photons is typically much larger than the mean free path of massive particles, an explicit treatment of the radiation transport necessarily involves a significant non-locality of the problem. This issue was satisfactorily solved in the case of stationary stellar atmospheres (see, e.g. Hubeny & Mihalas 2014), using efficient iterative methods. However, this problem remains difficult in the case of a non-stationary plasma, where the equations of hydrodynamic need to be coupled, at each time, with the equations for the radiative transfer.

\(^5\) with a number abundance of 0.12%, and a mass (averaged over abundances) of 17 u.

Therefore, the previous kinetic equations have to be solved simultaneously with the monochromatic radiative transfer equations. This allows computing the frequency-averaged local radiation flux, energy and pressure, and helps including these quantities in the hydrodynamics equations (Eq. 1). In practice, this exact description would require extensive numerical resources: the difficulty is commonly reduced by averaging the radiation quantities by frequency bands. In the multi-groups approximation, the absorption and emission coefficients are averaged over several frequency band using adapted weighting functions. The larger the number of groups, the better the precision of the computation. The simplest and most commonly used approach is the monogroup approximation, which means that the radiation quantities are averaged over the whole frequency domain covered.

Besides these delicate issues, radiative transfer takes part in the computation of the spectrum emerging from this structure. This is usually done by a post processing of the hydrodynamic results by more detailed spectral synthesis tools, as detailed in Section 2.3.3.

#### 2.2.2. Momenta equations

The radiation field is described here by the momenta equations (see e.g. Mihalas & Mihalas 1984) for the frequency-integrated radiation energy \( (E_i) \) and momentum or flux \( (\mathbf{M}_i = F_i/c^2) \) volumetric densities, within the comoving frame (Lowrie et al. 2001):

\[
\begin{align*}
\partial_t F_i + \mathbf{v} \cdot \nabla F_i + \left( \mathbf{P} \cdot \mathbf{\nabla} \right) F_i + \left( \mathbf{b} \cdot \mathbf{\nabla} \right) E_i = \gamma E_i \\
\partial_t M_i + \mathbf{v} \cdot \nabla M_i + \left( \mathbf{P} \cdot \mathbf{\nabla} \right) P_i + \left( \mathbf{b} \cdot \mathbf{\nabla} \right) (\mathbf{M} \cdot \mathbf{v}) = \delta_{M_i}
\end{align*}
\]

(2)

The (monogroup) radiation quantities are integrated from 1 to \( 10^4 \) Å. The M1 closure relation allows then to derive the radiation pressure \( P_i \) from the radiation energy density: \( P_i = D E_i \). \( D \) and \( \chi \) are respectively the Eddington tensor and factor \((D \equiv \chi \text{ in ID})\) and are defined as follows:

\[
D = \frac{1 - \chi}{2} + \frac{3\chi - 1}{2} \mathbf{i} \otimes \mathbf{i}
\]

\[
\chi = \frac{3 + 4f^2}{5 + 2\sqrt{4 - 3f^2}}
\]

(3)

with the reduced flux \( F_i = F_i/(cE_i) \) (and \( f = \parallel f \parallel \)), the flux direction \( \mathbf{i} = f/f = F_i/F_i \) and \( L_2 \) the second-order identity tensor.

Depending on the expression of the radiation source terms, these equations can continuously model optically thin to thick propagation media (see e.g. Mihalas & Mihalas 1984).

#### 2.2.3. Radiation source terms - opacities & line cooling

This work aims at describing in a consistent way the system composed of three zones, which are coupled together through radiation but in different thermodynamical states (Figure 1): the dense and optically thick near-LTE chromosphere (Section 2.2.3.1) on the one hand, the optically thin coronal hot accretion slab and cold accretion flow (Section 2.2.3.2) on the other hand. We also expect, according e.g. to Calvet & Gullbring (1998), that the frequency distribution of the radiation flux varies strongly from the X rays to the infrared. We decided to work step by step, using a model which makes a continuous transition between the optically thick LTE approximation and the coronal limit, as described in Section 2.2.3.3.
2.2.3.1. Optically thick limit

The deep stellar atmosphere is optically thick and can be considered at LTE, i.e. each microphysics process is counter-balanced by the reverse process. In LTE and regimes close to LTE, the radiation energy and momentum source terms are defined by (see e.g. Mihalas & Mihalas 1984):

\[ \dot{s}_{E}^{\prime} = \kappa \rho c (a_{g} T^{4} - E_{s}) \quad \text{and} \quad \dot{s}_{M}^{\prime} = -\kappa_{R} \rho c M_{s} \quad (4) \]

Two radiation-matter coupling factors appear here (in cm\(^2\) g\(^{-1}\)). The Planck mean opacity \(\kappa\) is based on the frequency-integrated absorption and emission coefficients \(\kappa_{e}\) weighted by the Planck distribution function \(B_{\nu}\), while the Rosseland mean opacity \(\kappa_{R}\) is the harmonic mean of \(\kappa_{e}\) weighted by the temperature derivative of the Planck function \(\partial_{T}B_{\nu}\), as follows (Mihalas & Mihalas 1984):

\[ \kappa_{P} = \frac{\int \kappa_{e} B_{\nu} d\nu}{\int B_{\nu} d\nu} \quad \& \quad \kappa_{R}^{-1} = \frac{\int \kappa_{e}^{-1} \partial_{T}B_{\nu} d\nu}{\int \partial_{T}B_{\nu} d\nu} \quad (5) \]

In these frequency averages, the Planck mean is dominated by the contribution of strong (lines) absorption whereas the Rosseland mean is dominated by the regions in the spectrum of lowest monochromatic opacity. As a consequence, at large optical depths, \(\kappa_{P}\) correctly describes the energy exchange between particles and photons, while \(\kappa_{R}\) gives the correct total radiative flux (Hubeny & Mihalas 2014).

![Fig. 2. Planck (\(\kappa_{P}\), left) and Rosseland (\(\kappa_{R}\), right) opacities with respect to gas density and temperature, in log scale (cf. Appendix A). The black curve represents typical conditions met with accretion shock and flow.](image)

Several opacity tables are available for a variety of chemical compositions. However, they all fail to cover the full \((\rho, T)\) domain explored in our simulations (see solid black line in Figure 2). We constructed then with the SYNSPEC code (Section 2.3.3) our own LTE opacity table (see Appendix A for further details), presented in Figure 2. These opacities include atomic (high \(T\)) and molecular (low \(T\)) contributions.

2.2.3.2. Optically thin limit

Due to its very low density \((\rho \sim 10^{-13} \text{ g cm}^{-3}\)), the accreted plasma can be described by the limit regime where the gas density tends towards zero: the coronal regime. The coupling between radiation and matter boils down in this case to an optically thin radiative cooling function \(\Lambda(T)\) (in erg cm\(^{-3}\) s\(^{-1}\)). The radiation source/sink terms become then:

\[ \dot{s}_{E}^{\dagger} = n_{e} n_{H} \Lambda(T) \quad \text{and} \quad \dot{s}_{M}^{\dagger} = 0 \quad (6) \]

The term \(\dot{s}_{M}^{\dagger}\) is set to zero since there is no coupling between radiation and matter in this regime (see Appendix C for more details).

![Fig. 3. Optically thin radiative cooling for different metallicities \(Z\), adapted from Kirienko (1993).](image)

The present work is based on the cooling function provided by Kirienko (1993), reproduced in Figure 3, with \(Z/Z_{\odot} = 1\) (see Appendix B.1 for the explanation).

2.2.3.3. Intermediate regimes

The previous source terms describe two well-defined plasma situations. On the one hand, the basis of the stellar chromosphere is optically thick and can be described by the previous LTE radiation source terms. On the other hand, the low density and hot slab is mostly optically thin and can be described in the coronal regime.

It is physically expected and numerically compulsory to perform a smooth and continuous transition to encompass intermediate regimes. This could be done using adequate opacities and emissivities, as for instance obtained in a collisional-radiative model, unfortunately not available yet for the whole range of physical conditions of the present study. Thus we have preferred to follow the transition between LTE and coronal regimes with the probability for a photon (emitted from the column center) to escape sideways (see e.g. Lequeux 2005, equation 3.66):

\[ \zeta = \frac{1 - \exp(-3 \tau_{e})}{3 \tau_{e}} , \quad \tau_{e} = \kappa_{R} \rho L_{c} \quad (7) \]

\(\rho\) and \(\kappa_{R}\) values are taken at the photon emission position. The characteristic length \(L_{c}\) is here taken as the accretion column mean radius (i.e. 1000 km, see Section 3.1). Radiation source terms become then (see Appendix C for further details):

\[ \dot{s}_{E}^{\dagger} = (1 - \zeta) \dot{s}_{E}^{\prime} + \zeta \dot{s}_{E}^{\dagger} \quad \text{and} \quad \dot{s}_{M}^{\dagger} = \dot{s}_{M}^{\prime} \quad (8) \]

the star (*) and dagger (†) denoting respectively the LTE (Eq. (4)) and the coronal (Eq. (6)) expressions.

2.3. Numerical tools

2.3.1. One-dimensional approach

Observations indicate that, in general, the ambient magnetic field is of the order of 1 kG (Johns-Krull et al. 1999; Johns-Krull 2007). The resulting Larmor radius (1 mm) is very small, i.e. the plasma follows the magnetic field lines. Moreover, the Alfvén velocity reaches 3% of the speed of light and the magnetic waves behave thus like usual light waves. Therefore, focusing on the heart of an accretion column in strong magnetic field case, we
can model the accreted material along one field line, that will be
assumed to be radial relative to the stellar center. Since the ac-
retion process is expected to involve strong shocks, we chose a
numerical tool able to achieve very high spatial resolution.

2.3.2. AstroLabE – an ALE code

The present work is based on numerical studies performed with
the 1D Arbitrary-Lagrangian-Eulerian (ALE) code AstroLabE
(see e.g. de Sá et al. 2012; Chièze et al. 2013). It is based on
the Raphson-Newton solver (Numerical Recipes, Section 9) and
a fully implicit scheme (the CFL condition can then be ignored)
to compute primary variables at each time step.

This code solves, along with the adequate physics and chemistry
equations (see Sections 2.1 and 2.2), the equations describing the
behavior of the grid points. The space discretization can follow
an Eulerian or a Lagrangian description. Moreover, the grid can
freely adapt to hydrodynamics situations (the arbitrary descrip-
tion, Dorfi & Drury 1987): this helps us reach high resolution
around shocks with fixed cardinality ($\delta r/r_{\text{max}} \sim 10^{-7}$ with 150–
300 grid points).

Beside its application to stellar accretion (de Sá 2014), Astro-
LabE has been used in several astrophysical situations such as
the interstellar medium (Lesaffre 2002; Lesaffre et al. 2004),
experimental radiative shocks (Stehlé & Chièze 2002; Bouquet
et al. 2004) or type Ia supernovae (Charignon & Chièze 2013)
studies.

2.3.3. SYNOPSIS – a spectrum synthesizer

For the computation of the opacities and of the emerging spec-
tra, we used the public 1D spectrum synthesis code SYNOPSIS
(Hubeny & Lanz 2017). It is a multi-purpose code that can ei-
ther construct a detailed synthetic spectrum for a given model
atmosphere or disks, or generate LTE opacity tables. In this pa-
er, we used SYNOPSIS both for generating opacity tables (see
Section 2.2.3.1 and Appendix A), and for the snapshots spectra
presented in Section 3.4.4.

The resulting synthetic spectrum reflects the quality of the
input astrophysical model; using an LTE model results in an
LTE spectrum, while using a NLTE model results in a NLTE
spectrum. The snapshots of our hydrodynamic simulations pro-
vide temperature and density as a function of position; it is
therefore straightforward to compute LTE spectra for such struc-
tures. It would be in principle possible to construct approximate
NLTE spectra, keeping temperature and density fixed from the
hydrodynamic simulations (the so-called restricted NLTE prob-
lem). This could be done for instance by the computer program
TLUSTY (Hubeny & Lanz 1995, 2017), which would provide
NLTE level populations that can be communicated to SYNOPSIS
to produce detailed spectra. However, as previously mentioned,
such a study is computationally very demanding and is well be-
yond the scope of the present paper. Nevertheless, since NLTE
effects may be important, this will be done in a future paper.

This will allow to inspect the effect of the LTE approximation on our
results.

This synthetic spectrum, computed at different altitudes of the
accretion column will reveal the role played by the different
parts of the spectrum, from X ray to Visible ($1-10^4$ Å). However,
it is important to note that, as the accretion column is limited in
diameter, some effects, like the absorption by the coldest parts
are only pertinent for an observation along or near the direction
of the accretion column. A 3D radiative transfer post-processing
would then be more suitable to the geometry of the system (Ibgui
et al. 2013).

3. Accretion basis simulations

3.1. Strategy and common parameters

We simulated for this study several physical situations in order
to check the net effect on the QPOs of the chromospheric model
on one side and of the matter-radiation coupling on the other
side. We present first the reference case: a gas flow hits a fixed,
rigid and non-porous interface (W–A case, Section 3.2). We
check then the effect of a dynamically heated chromosphere on
the accretion process (Chr–A case, Section 3.3) and we finally
check the effect of the radiation feedback on matter (Hybrid
case, Section 3.4). The conditions and main results of each
simulation are resumed in Table 1.

The simulations presented in this paper share few param-
ters:

- the computational domain size is $r_{\text{out}} = 10^5$ km (the outer
boundary limit);
- the column radius is set to $L_{\text{c}} = 1000$ km;
- for the gravity magnitude, we use $R_\odot = R_{\odot}$ and $M_\odot = M_{\odot}$;
- the accreted gas enters the computational domain through the
outer boundary with $\rho_{\text{acc}} = 10^{-12}$ g cm$^{-3}$, $T_{\text{acc}} = 3000$ K\(^6\) and
$v_{\text{acc}} = 400$ km s$^{-1}$.

The velocity of the accreted gas is derived from the free-fall ve-
locity at $r = r_{\text{out}}$ above the stellar surface, considering a null
radial velocity at the truncation radius $R_{\text{tr}} = 2.2 R_{\odot}$ (taken here
from the center of the star).

When the M1 radiation transfer is used (either near-LTE
transfer or intermediate regime), one solar luminosity gets
through the domain from the inner to the outer boundaries.

3.2. Reference case (W–A)

3.2.1. Setup

In the reference case, we simulate the accretion stream using the
same physics and assumptions than in previous models (see e.g.
Sacco et al. 2008; Koldoba et al. 2008). The matter-radiation
coupling is then described by the coronal radiative cooling (Sec-
tion 2.2.3.2) and the plasma ionization is computed with the
modified Saha equation (Section 2.1.2). In order to simplify
the discussion, we focus on the post-shock structure and on the
global dynamics. The stellar atmosphere is modeled in the sim-
plest way, hereafter called the "window" model. It consists in a
fixed rigid non-porous transparent interface. The main param-
eters are resumed in Figure 4.

3.2.2. QPO cycle

Besides the fact that matter accumulates on the left (inner)
rigid boundary interface, the system is found to be perfectly periodic.

Figure 5 presents five snapshots of density, temperature and ve-
locity profiles during a QPO cycle far from the initial stages. The
accreted gas falls from right to left. A hot slab of shocked mate-
rial builds first ($t = 2750$ and $2884$ s) and cools down according

\(^6\) In the Hybrid case, the temperature of the accretion flow is radiatively
heated by the chromosphere up to 5730 K, before the accretion process
starts.
to the coronal regime. Below a threshold temperature\(^7\), the fast, quasi-isochoric, cooling of the slab basis causes the collapse of the post-shock structure \((t = 2994\) and \(3110\) s). Just after the full collapse of the slab, since the accretion process is still working, a new slab forms and grows \((t = 3156\) s).

This simulation is to be compared to the ones performed by Sacco et al. (2008); Table 2 resumes the main parameters and results for fast comparison. Despite few key differences (Sun vs. MP Muscæ parameters & "window" vs. chromospheric heating function), the results are in good agreement with each other.

\(^7\) i.e. the temperature at which the thermal instability is triggered \((- 8 \times 10^7\) K) as expected from the optically thin radiative cooling variations with respect to temperature, see Section 2.2.3.2 and references therein for further details.

### 3.2.3. X-ray luminosity

Although this case is relatively academic, it can be interesting to compute the instantaneous X-ray luminosity \(L_\Lambda\) and its time average \(\bar{F}_\Lambda\), to compare them with the values obtained in the two more realistic cases discussed in the next sections:

\[
F_\Lambda = \int n_a n_H \Lambda(T) \, dr \quad \text{&} \quad \bar{F}_\Lambda = \frac{1}{\tau_{\text{cycle}}} \int_0^{\tau_{\text{cycle}}} F_\Lambda(\tau) \, d\tau \quad (9)
\]

These quantities are commonly compared to the incoming kinetic energy flux. However, since the flow accelerates in its free-fall from the outer boundary down to the reverse shock, the plasma velocity and density may change between the outer boundary of the simulation box and the location of the reverse shock. To get round this issue, we compare \(F_\Lambda\) and \(\bar{F}_\Lambda\) to the incoming mechanical energy flux:

\[
F_M = \frac{1}{2} \rho \, u_{\text{inf}}^3 + \int_{r_a}^{r_m} \frac{G \, M_* \, \rho \, u}{r^2} \, dr = 4.2 \times 10^9 \text{ erg cm}^{-2} \text{ s}^{-1} \quad (10)
\]

where the gravitation energy potential origin is set at the mean accretion shock position \((r_m \approx 10^7\) km), and the velocity is counted positively.

Figure 6 shows the time variation of \(F_\Lambda\). As expected, this quantity increases during the propagation of the reverse shock and decreases during the collapse. The time averaged flux \(F_\Lambda\) is equal to \(1.5 \times 10^9\) erg cm\(^{-2}\) s\(^{-1}\), i.e. 36% of the incoming mechanical energy flux \(F_M\).

### 3.3. Effect of a dynamical chromosphere (Chr–A)

#### 3.3.1. Setup

In this second setup (see Figure 7), we aim at studying the effect of a dynamically heated chromosphere on the phenomenon described in the previous Section. To achieve this, we "divide" the computational domain into two zones separated by a transparent\(^8\) Lagrangian interface.

The outer zone is described as before, i.e. with modified Saha ionization and optically thin radiative cooling (coronal regime). However, the inner zone is now described by our chromospheric model (see Appendix B). Ionization is still described by the modified Saha equation, but we use the LTE radiation source terms as given in Equation (4). To get a dynamically heated chromosphere, we first compute a hydrostatic equilibrium, with the outer

\(^8\) Although the column plasma is expected to be at coronal regime, LTE radiation transfer is needed to build the chromosphere layer. It is therefore essential to allow radiation to escape from the first zone through the second (optically thin) one.

---

**Table 1.** Characteristics of 3 simulations used in this work and their main results. "W–A" corresponds to our reference case.

| Name   | Atmos.          | Chromos. heating | Radiation source terms | Ionization model | \(H_{\text{max}}\) (x10\(^3\) km) | \(\tau_{\text{cycle}}\) (s) | Section | Fig. |
|--------|-----------------|------------------|------------------------|-----------------|----------------------------------|---------------------------|---------|------|
| W–A    | "Window"*       | –                | \(\Lambda^*\)          | Modified Saha*  | 20                               | 400                       | 3.2     | 5    |
| Chr–A  | Equilibrium atmosphere | \(L_{\odot}^*\) & acoustic heating | LTE (chromos.) & \(\Lambda^*\) (acc. flow) | Modified Saha*  | 17                               | 350                       | 3.3     | 8    |
| Hybrid | Equilibrium atmosphere | \(L_{\odot}^*\) | Intermediate (transition: \(\zeta\)) | Time-dependent collisional radiative | 9                                | 160                       | 3.4     | 9    |

**Fig. 4.** "W–A" simulation setup and boundary conditions.

**Notes.** \(H_{\text{max}}\): maximum extension reached by the post shock medium; \(\tau_{\text{cycle}}\): cycle duration; "Window": fixed rigid non-porous transparent interface; \(\Lambda\): optically thin radiative cooling; \(L_{\odot}\): one solar luminosity enters the simulation box from the inner boundary; Modified Saha: Brown (1973).

**Table 2.** Comparison between our reference case ("W–A") and results obtained by Sacco et al. (2008).

| Parameters & quantities | Sacco et al. (2008) | "W–A" |
|-------------------------|---------------------|-------|
| Object                  | MP Muscæ            | Sun   |
| Atmosphere              | Heating function    | "Window" |
| Radiation               | \(\Lambda\)         | \(\Lambda\) |
| Ionization              | Modified Saha       | Modified Saha |
| \(\rho_{\text{acc}}\) (g/cm\(^3\)) | 10\(^{-13}\) | 10\(^{-13}\) |
| \(\nu_{\text{acc}}\) (km/s) | 450                | 400   |
| \(T_{\text{acc}}\) (K) | 10\(^7\)            | 3 \times 10\(^3\) |
| \(\tau_{\text{cycle}}\) (s) | 400                | 400   |
| \(H_{\text{max}}\) (Mm) | 18                 | 20    |
| \(n_a\) (cm\(^{-3}\)) | 10\(^{11}\)–10\(^{12}\) | 10\(^{11}\)–10\(^{11.5}\) |
| \(T_{\text{max}}\) (K) | 10\(^{6.5}\)       | 10\(^{6.5}\) |
zone inactivated, and with one solar luminosity crossing the entire domain (no effect on the outer zone). Acoustic energy is then injected in the form of monochromatic sinusoidal motion of the first interface (a "window") with a 60 s period to mimic solar granulation. Several snapshots of temperature profiles are presented in Figure B.1. Once the shock-heated chromosphere reaches its stationary regime, the accretion process is launched (in the outer zone).

3.3.2. Acoustic perturbations

Figure 8 shows seven snapshots of density and temperature profiles during the first QPO cycle (1–354 s) They are followed in the second line by 5 snapshots of the second QPO cycle (354–415 s). The second cycle differs from the first one only during the slab building (354–397 s). The sixth snapshot (415 s) is very close to the snapshot of the first cycle at t = 71 s. The (unchanged) end of the second cycle is then not reported.

During the installation phase (1–336 s) of the reverse shock, the post-shock structure follows more or less the same scenario than for the reference case W–A. After several periods of the acoustic waves, small differences occur. The transmission of these waves/shocks to the accretion column depends on the leap of the acoustic impedance between the upper chromosphere and the hot slab, which results in reflection/transmission of these waves/shocks at this interface. The smallest leap is reached at the end of the collapse, near 336 s, leading to a transmission increase, which however remains still low. Their effect lead to small perturbations in the post-shock density (e.g. at 168 s).

Table 3. Position of the old and new reverse shocks between t = 358 s and t = 397 s (see Figure 8).

| Time (s) | r_{old} (km) | r_{new} (km) |
|---------|--------------|--------------|
| 358     | 10^{3.40}    | 10^{3.35}    |
| 380     | 10^{3.25}    | 10^{3.70}    |
| 386     | 10^{3.10}    | 10^{3.80}    |
| 397     | 10^{2.95}    | 10^{3.90}    |

After this time, the transmitted waves start to feed with matter the hot collapsing layer behind the reverse shock. The thick-
Fig. 8. Snapshots of the density (green), temperature (red), and gas velocity (grey) profiles of the first QPO cycle with the "Chr-A" setup; the accreted gas falls from the right to the left. The first line (between 0 and 354 s) corresponds to the first cycle. The second and third lines correspond to the beginning of the second cycle. Snapshots at \( t = 71 \) s and 415 s are very close: from this time, the cycle behaves like the previous one. A typical sequence is: growth of a hot slab of shocked material \( (t = 21 \) s), quasi-isochoric cooling at the slab basis (thermal instability, \( t = 168 \) s), start of the collapse of the post-shock structure \( (t = 336 \) s), impact of the collapsing material on the chromosphere \( (t = 354 \) s), launch of a new shock before the end of the collapse \( (t = 358 \) s), passing of the two shocks \( (t = 380 \) s), end of the collapse of the "old" structure \( (t = 386 \) s) and growth of the new slab \( (t = 415 \) s).

Fig. 9. Snapshots of the mass density (green), gas temperature (red), velocity (grey), electron density (light green), escape probability \( \xi \) (dark blue) and absorption probability \( 1 - \xi \) (cyan, see Section 2.2.3.3) profiles of the first QPO cycle with the "Hybrid" setup; the accreted gas falls from the right to the left on an equilibrium atmosphere.
ness of this layer increases, as can be shown in Figure 8 at 351 s, compared for instance with our reference case (3110 s, Figure 5). This structure collapses and hits at 354 s the dense chromosphere, leading to a secondary ("new forward") reverse shock which propagates backwards inside the slab. This behavior is confirmed by the velocity variations shown in grey in Figure 8. The two reverse shock pass then each other: the positions of the new shock (or contact discontinuity) and the previous (old) one are resumed in Table 3. The end of one cycle is therefore overlapped by the beginning of a new one.

3.3.3. Observational consequences

This model implies two main observational consequences. First, compared to the reference case, the QPO cycle period is modified by the acoustic heating. The question of possible resonance is pointless regarding multi-mode acoustic heating by out-of-phase waves emitted in different locations. The period is slightly reduced (from 400 s for the W–A model to 350 s here) when using solar parameters. Since CTTSs’ atmospheres have a stronger activity than the Sun’s (that we use for the chromospherical model), the effect is expected to be enhanced in CTTSs.

The second effect deals with the X-ray luminosity variation during a cycle, as reported in green in Figure 6. The growth phase is comparable with the W–A setup. However, the acoustic perturbations from the chromosphere induce strong differences in the collapse phase. Moreover, the overlapping of the beginning and end of the cycles affect their X-ray luminosity and the overall amplitude of the variations (contrast) is reduced compared the the reference case. QPO observations may thus require both higher time resolution and improved sensitivity.

These results show that, compared to the reference case, the dynamical heating of the chromosphere impacts the duration of the QPO period and its observability. Of course, a more realistic description of the chromospheric heating would require at least a 2D MHD picture. For instance, we know that chromospheric perturbations may lead – inside the column – to the development of fibrils (see e.g. Matsakos et al. 2013, ChrFlx# models), which is one of the scenario explaining the absence of observation of QPO. In the acoustic description of the chromospheric heating, these fibrils, evolving out of phase, will also be strongly affected by the chromospheric perturbations.

3.4. Radiation effect on accretion (Hybrid)

3.4.1. Setup

In this Section, the plasma model includes collisional-radiative ionization (see Section 2.1.2). The radiation-matter coupling is described within the intermediate regime (see Section 2.2.3.3). The goal of this last setup (see Figure 10) is to inspect the net effect of the matter-radiation coupling. We have therefore chosen not to consider any chromospheric activity (i.e. the acoustic heating is kept switched off).

3.4.2. Ionization model

In addition, we test in this setup the effect of the time-dependent ionization through radiative ionization/recombination and collisional ionization with a time-dependent formulation (see Section 2.1.2 for more details).

The use of collision and radiative rates with a calculation of the electron density at equilibrium, i.e. the equilibrium value of \( n_e \) at each time step, brings differences in the transition between the (almost) neutral medium and the fully ionized plasma. This transition lays between \( 10^{3.6} \) and \( 10^{4.2} \) K. However, such temperatures are only reach by the accreted gas during the cooling instability. Its overall effect is hence negligible. The results presented in Figure 9 are based on this equilibrium calculation of \( n_e \).

The main difference brought by a time-dependent calculation of the electron density is a tiny ionization delay behind the reverse shock front, as shown in Figure 11.

At the shock front, the kinetic energy is converted into thermal energy, and then a part of this thermal energy is used to ionize the post-shock material with a time scale connected to the ionization rates; the affected gas layer is up to 0.2 km thick, and thus negligible compared to the whole structure (that is at least \( 10^4 \) km thick, see Table 1).

This justifies the use of the (time-independent) Saha-Brown model for ionization for the two other cases (W–A and Chr–A). Günther et al. (2007) and Sacco et al. (2008) obtain the same conclusion from different approaches.

3.4.3. Atmospheric heating and cycle reduction

The first cycle is presented in Figure 9; it shows the time variations over 160 s of the gas temperature and mass density, and of the photon escape (\( \zeta \)) and absorption (\( 1 - \zeta \)) probabilities (see Section 2.2.3.3) for the same snapshots. Although \( 1 - \zeta \) shows strong variations, its net value beyond the accretion shock remains small compared to unity, and the post-shock dynamics is roughly the same as in the test case (W–A, see Section 3.2). The next cycles only differ from this first one by the position of the interface between the slab and the chromosphere, as discussed in Section 3.4.3.2.

The global behavior follows the trends of the two previous models. However, several effects must be highlighted: a global
motion of the whole post-shock structure, the reduction of the oscillations period and of the post-shock extension. These effects are discussed below.

3.4.3.1. QPO reduction

Changing the radiation feedback on the plasma directly affects its cooling efficiency, and thus the cycle duration along with its maximum spatial extension. This simulation shows that a weak coupling between radiation and matter\(^9\) is able to significantly change the dynamics of QPOs: the cycle is reduced to 150 s and at most half the maximum extension of the previous simulations.

3.4.3.2. Chromospheric beating

Another effect is the heating of the upper chromosphere by the radiating post-shock plasma at the basis of the accretion column. For instance, between 9 and 70 s, its temperature varies from 7000 to 10 800 K at 800 km, and the pressure increases from 800 to 2600 dyn/cm\(^2\) at this location (Figure 12). As a consequence, the whole post-shock structure is pushed upwards from 875 km to 3150 km, thus out of the unperturbed chromosphere (by about 2000 km, see e.g. Vernazza et al. 1973).

3.4.3.3. Accretion flow pre-heating

We note for similar reasons, that the pre-shock is also heated up to ~ 8500K (at \(t = 93\) s). We believe that such preheating of the accretion flow, which was already pointed out in Calvet & Gullbring (1998); Costa et al. (2017), is attributed to the radiating hot slab.

These effects are quantified in Figure 13, which reports the time variations of the chromosphere/slab interface position and velocity, as also hot slab/accretion flow interface and velocity. The temperatures of the heated chromosphere and pre-shock are also reported.

3.4.4. Monochromatic radiative flux

As this simulation is performed using only one group of radiation frequencies, it is interesting to analyze more precisely the details of the previous radiative heating via its feedback on the monochromatic radiation flux.

To this purpose, the hydro structures has been post-processed with the SYNSPEC code (Section 2.3.3). To be consistent, we take the atomic data already provided for the calculation of the average opacities used previously (see Section 2.2.3 and Appendix A). We thus estimate the Eddington flux \(H_e\), i.e. the first moment of the specific intensity. Since the line profile behavior is not investigated here, velocity effects are neglected.

It is important to recall that a quantitative comparison of this synthetic flux with observations, especially in the X rays (see e.g. Güdel et al. 2007; Robrade & Schmitt 2007; Drake et al. 2009) would require NLTE and 3D radiative transfer post-processing. However, using 1D radiative transfer and the LTE approximation is here interesting as it corroborates or not the general accepted trends, e.g. a strong X-ray emission and an excess of luminosity in the UV-VIS range (Brickhouse et al. 2010; Ingleby et al. 2013; Calvet & Gullbring 1998).

A typical spectrum emerging from \(4.6 \times 10^8\) km (located within the accretion flow) is reported in Figure 15. It is computed from a snapshot (\(t = 70\) s) of the Hybrid model (see Fig.

---

\(9\) The absorption probability \(1 - \zeta = 1.3 \times 10^{-12}\) to \(5.5 \times 10^{-4}\) in the post-shock structure and \(5.5 \times 10^{-4}\) in the accretion flow.
ures 9 and 14). At this stage, the chromosphere extends up to $1.4 \times 10^3$ km, the hot plasma from $1.4 \times 10^3$ km to $8.3 \times 10^3$ km and the accretion flow from $8.3 \times 10^3$ km to $1 \times 10^5$ km. The flux that emerges from this layer presents three characteristic spectral bands:

- in the range 1–100 Å (X-rays), the bump is attributed to the hot plasma of the column, with intense lines up to $10^{12}$ erg cm$^{-2}$ s$^{-1}$ Å$^{-1}$;
- in the range 100–900 Å (EUV), radiation is efficiently absorbed by the inflow and will lead to its heating as observed in our hydro simulation;
- in the range 900–10 000 Å (UV + Vis + IR), the second bump is attributed to the stellar chromosphere and photosphere.

The strong absorption of the EUV radiation is due to the huge optical depth of the accretion flow.$^{10}$ This effect may then be attenuated in the case of a bent column or when the observation is performed side-on and not along the column. This absorption effect on the spectrum is illustrated in Figure 16, which presents the flux emerging right after the reverse shock front, at $r = 8.3 \times 10^3$ km. This figure shows that this absorption also affects, to a lesser degree, the visible spectrum originating from the chromosphere. This must be considered when interpreting the UV excess (see e.g. Calvet & Gullbring 1998; Hartmann et al. 2016).

If these synthetic spectra allow to explain the general trends of the radiation per frequency intervals, as already mentioned, they can’t be use for quantitative comparison with observations. Let us for instance consider the wavelength-integrated physical flux in the X-rays, $F_X$ in the interval 2–27 Å, observed by Chandra (Brickhouse et al. 2010), and its time average over the cycle $F_X$:

$$F_X = 4\pi \int \frac{27}{2} H_\lambda \, d\lambda$$

$$F_X = \frac{1}{\tau_{\text{cycle}}} \int_0^{\tau_{\text{cycle}}} F_X \, dt = 1.1 \times 10^{11} \text{ erg cm}^{-2} \text{ s}^{-1} \tag{11}$$

The time variation of $F_X$ (see Figure 17, blue curve) reveals four main phases: the building of the post-shock slab (0 – 5 s), the slow radiative cooling (5 – 70 s), the collapse (70 – 90 s) and the chaotic collapse (90 – 160 s). The peak at 90 s is due to the apparition of numerous strong lines between 10 Å and 15 Å. During the chaotic collapse, the accretion shock appears and disappears several times, as well as for the reverse shock, but not in phase: $F_X$ presents therefore strong time variations, as shown in Figure 17. The time averages of $F_X$ and $F_{\Lambda}$ (Eq. (9)) are expected to remain smaller than the incoming mechanical energy flux $F_{\text{mech}}$ (Eq. (10)). This condition is satisfied for $F_{\Lambda} \leq 30\%$ of $F_{\text{mech}}$.$^{11}$ However, $F_{\Lambda}$ is about 30 times higher than this limit, which is not physical. This discrepancy is imputable to the LTE approximation used in SYNSPEC: the blue curve in Figure 17 is up to 2.5 decades higher than the red curve.

From the mechanical energy flux, and assuming a stellar radius equal to 1 solar radius, we derive from the luminosity measured by Brickhouse et al. (2010) ($1.3 \times 10^{30}$ erg s$^{-1}$) a filling factor of 2.1%. To recover this luminosity from $F_X$, the filling factor would be even smaller. Using the optically thin approximation ($F_{\Lambda}$), we calculate a more realistic value of the filling factor (6.9%).

$^{10}$ at $\rho = 10^{-13}$ g cm$^{-3}$ and $T \sim 5 \times 10^5 – 8 \times 10^5$ K.

$^{11}$ $F_{\Lambda}$ should be slightly overestimated due to the larger wavelength interval used for its estimation.
4. Refining the models

4.1. A more realistic chromosphere

It should be pointed out that this study uses a solar model for the chromosphere with acoustic heating. Compared to the description of this heating, a more important improvement would be to consider a realistic T-Tauri chromospheric model, which is today not very well known. This may affect ionization (and then gas pressure with another chemical abundances) as well as slab characteristics (through gravity) and radiation effects (through opacities and incoming luminosity). Our results are then to be considered qualitatively and not quantitatively.

4.2. Improvements of the radiation model

We use in this work radiation momenta equations with the M1 closure relation. Although this is already a strong improvement compared to other approaches like the diffusion model, it could be improved by using half fluxes (i.e. the inward and outward components of the radiation flux). This should disentangle the radiation flux coming from the star and from the post-shock structure.

The M1 closure relation allows the radiation field to reach at most one direction of anisotropy; half-fluxes can extend it to two, i.e. the maximum number of anisotropy directions reachable in 1D. Half fluxes (along with M1) would then be equivalent to the momenta equations with the M2 closure relation (Feugeas 2004), without its prohibitive numerical cost. The M1 model and its limits have been thoroughly studied (see e.g. Levermore 1996; Dubroca & Feugeas 1999; Feugeas 2004). The behavior of this model along with half fluxes needs however to be examined.

More important is the approximation made with the monogroup approach used in this work. The whole spectrum is then approximated as a black body providing the adequate opacities characteristics (through gravity) and radiation emissivity (equivalent to $n_e n_H A$), radiation energy absorption ($k\rho$) and radiation flux sinking ($\alpha_k$). Moreover, all these quantities, computed with a radiative-collisional model, have to be averaged over adequate weighting functions. Due to recent progresses in this topic (Rodriguez et al. 2018), new results are expected in a near future. And a NLTE description should be used to compute the emerging spectra: this work is already in progress using TLUSTY code.

5. Conclusion

In this study, we used 1D simulations with detailed physics to check the validity of the two following common assumptions in accretion shock simulations: the stellar atmosphere can be either modeled by a hydrostatic or a steady hydrodynamic structure, and the dynamics of accretion shocks is governed by optically thin radiation transfer. We checked first that we are able to recover previous results (Sacco et al. 2008, W–A case, Section 3.2) and tested independently each of these assumptions (Chr–A case, Section 3.3, and Hybrid case, Section 3.4). Each of them proved to have a non-negligible impact on the typical characteristics of the accretion dynamics and on its observability.

This study could be completed with a simulation that would include both a dynamically-heated chromosphere and the hydro-brid setup. However, it appears at this stage more important to take into account a NLTE radiative description based on adapted opacities and radiative power losses. Another necessary improvement will be through a multi-group radiation transfer to catch at least the effect of EUV absorption and X-ray radiative losses on the structure of the column, and to analyse the possibility of a radiative precursor which could pre-heat the incoming flow. The study is also to be extended to multi-dimensional simulations in order to check the effects of both radiation and magnetic field closer to the real picture (Orlando et al. 2010, 2013; Matsakos et al. 2013, 2014).

Acknowledgements. We express our gratitude to Jason Ferguson for providing us with the molecular LTE opacity tables used in this work. We thank Rafael Rodriguez for sharing with us preliminary results about NLTE microscopic collisional-radiative data, Salvatore Orlando for reading the manuscript, Ziane Izri for its contribution at the beginning of this project and Christophe Santy & Rafael Rodriguez for helpful discussions. We extend thanks to the referee for his/her valuable feedback on this article and the rich discussions they initiated. I.H. thanks the Physics Department of Sorbonne Université for his visiting professorship.

This work was supported by the french ANR StarShock and LabEx Pla@Par projects (resp. ANR–08–BLAN–0263–07 and ANR–11–IDEX–0004–02), PICS 6838, Programme National de Physique Stellaire of CNRS/INSU and Observatoire de Paris.

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Appendix A: Opacity tables

The specificity of the accretion shock study led us to work on dedicated opacity tables. We expose in the appendix the reasons behind this choice and the creation process. The resulting opacity table is accessible upon request.

Appendix A.1: Motivation

![Fig. A.1. Planck (left) and Rosseland (right) opacities (in cm$^2$ g$^{-1}$) with respect to gas density and temperature – any in log scale, as provided by the Opacity Project (Opacity Project Team 1995). The black curve is a typical characteristic of an accretion column.](image)

Most available opacity tables are defined on a slanted ($\rho, T$) or ($n_e, T$) domain (see e.g. Figure A.1). However, one peculiarity of accretion shock structures is the presence of a low density hot post-shock plasma (black curve vertex in Figure A.1) that explores a domain uncovered by publicly available tables. More complete tables are thus mandatory for the present study.

Appendix A.2: Choice of primary tables

![Fig. A.2. SYNSEPEC (top) and Ferguson (bottom) Planck (left) and Rosseland (right) opacities (in cm$^2$ g$^{-1}$) with respect to gas density and temperature – any in log scale. The dotted lines show transition temperatures chosen for each table (see Appendix A.4.1.).](image)

To cover the density and temperature range corresponding to our conditions, we implement in the code SYNSEPEC (see Section 2.3.3), initially dedicated to stellar atmospheres, modules allowing to generate LTE monochromatic opacities at a given density and temperature. These monochromatic opacities were then averaged with the proper weighting functions to generate the adequate Rosseland and Planck mean opacities tables (hereafter called "SYNSPEC tables", see Figure A.2, top panels). These opacities are consistent with Opacity Project (see e.g. Opacity Project Team 1995) data, that we use as reference, for $T$ between 10$^{12}$ K and 10$^{12}$ K. The advantage of SYNSPEC comes from the high number of atomic species considered, since a very detailed chemical composition is necessary to model the radiation properties of a plasma at high temperatures.

However, below 10$^{1}5$ K, the molecular chemistry cannot be neglected, but is not included in this work on SYNSPEC. We completed thus the SYNSPEC tables with low temperature molecular opacities provided by Ferguson et al. (2005) between 10$^{3}$ K and 10$^{5}$ K ("Ferguson tables", see Figure A.2, bottom panels), that show excellent agreement with Opacity Project at upper temperatures. To facilitate the merging process, we obtained from the authors tables with compatible density and temperature grid (Ferguson, priv. comm.): mesh points from Ferguson and SYNSPEC tables are identical in the common domain (10$^{3.5}$–10$^{6}$ K and 10$^{14}$–10$^{6}$ g cm$^{-3}$).

Appendix A.3: Preliminary study

Appendix A.3.1: Analysis of primary tables

Considering opacity variations as well as temperature and density ranges, we decided to work with the logarithm of all these quantities. As first derivatives, we use then:

$$
\begin{align*}
\partial_{T}\kappa &= \frac{\partial \log_{10} \kappa}{\partial \log_{10} T} = \frac{T}{\kappa} \frac{\partial \kappa}{\partial T} \\
\partial_{\rho}\kappa &= \frac{\partial \log_{10} \kappa}{\partial \log_{10} \rho} = \frac{\rho}{\kappa} \frac{\partial \kappa}{\partial \rho}
\end{align*}
$$

where $\kappa$ stands for $\kappa_p$ or $\kappa_R$.

![Fig. A.3. $\partial_{T}\kappa_R$ from SYNSPEC table (top) and $\partial_{\rho}\kappa_p$ from Ferguson table (bottom) with respect to the logarithm of gas density and temperature. The dotted lines show transition temperatures chosen for each table. Some anomalies are revealed, especially around 10$^{12}$ K and high densities: this zone is cut during the merging process.](image)

Preliminary analysis of SYNSPEC and Ferguson tables revealed local aberrations, especially looking at the temperature or density derivatives (see Figure A.3). We may use the merging process to smooth most aberrations.
Appendix A.3.2: Physical and numerical constraints

In order to get a satisfying merging, several numerical and physical constraints must be respected:
- as far as possible, opacities must be of class $C^1$ (values and first derivatives must be continuous);
- the transition region should be as narrow as possible;
- the transition region must encompass anomalies encountered in both primary tables.

Such a table is composed of a limited number of discrete points: the first constraint can be reported to the interpolation method as far as opacity values in the transition present smooth variations.

To ensure a smooth transition between the molecular and the atomic (primary) tables, the transition must not take into consideration the values within the transition. The transition values loose then any physical meaning, and must be as few as possible.

Appendix A.4: Merging process

Appendix A.4.1: Method

In this Section, the index "A" refers to values taken at the lower transition temperature, as the index "B" for the upper ones. The transition temperatures chosen to merge SYNSPEC and Ferguson tables are:
- $T_A = 10^{3.71}$ K and $T_B = 10^{3.80}$ K for $\kappa_{p}$ (~1200 K wide);
- $T_A = 10^{3.65}$ K and $T_B = 10^{3.86}$ K for $\kappa_{R}$ (~2800 K wide).

The problem is decoupled in temperature and in density. First, we consider the merging at each mesh density as an isolated problem, and apply a correction – if needed – to improve smoothness along the density.

Appendix A.4.2: Merging along temperature

To satisfy the class $C^1$ constraint, we combined (see for instance Auer 2003 and Ibgui et al. 2013):
- piecewise cubic Hermite polynomials, which ensure continuity of values ($\kappa_A, \kappa_B$) and derivatives ($\partial_T \kappa_A, \partial_T \kappa_B$) at each transition limit;
- van Leer (1973) slopes to compute $\partial_T \kappa_A$ and $\partial_T \kappa_B$, so as to prevent the apparition of spurious extrema in forcing their location to the estimated closest mesh point.

For each grid density $\rho_i$, opacity at temperature $T_i \in [T_A, T_B]$ is estimated using the formula:

$$
\log_{10} \kappa(T_i; \rho_i) = u_i^3 (3 - 2u_i) \log_{10} \kappa_B + (u_i - 1) u_i^2 h \partial_T \kappa_B \\
+ (u_i - 1)^2 (2u_i + 1) \log_{10} \kappa_A + (u_i - 1)^2 u_i h \partial_T \kappa_A
$$

with $h = \log_{10} T_B - \log_{10} T_A$ and $u_i = (\log_{10} T_i - \log_{10} T_A)/h$. This expression can be rewritten as a 3rd degree polynomial in $u_i$.

12 We note that the Ferguson tables showed opacity discontinuities in their hottest and densest part, as SYNSPEC tables in their coolest and densest part, as SYNSPEC tables in their coolest and densest part (see Figure A.3). Since the values within the transition region are ignored, this is used to artificially remove anomalies: as far as possible, the transition region must be chosen so that it covers most of them.

13 Fritsch & Butland (1984) derivatives are used in these papers; they generalize van Leer slopes to non-regular grids.

Appendix A.4.3: Density correction

At this stage, we reached class $C^1$ along temperature, but there is no guarantee of continuity along density. However, in practice, it was $C^1$, except for few mesh temperatures $T'$. Since the dependency in density is held by the 3rd degree polynomial coefficients, we look at the behavior of each of them with respect to density. Every coefficient showed spurious variations nowhere but at densities $\rho_i^*$. We apply then piecewise cubic Hermite polynomials along with van Leer slopes (density derivatives) to estimate these coefficients for each $\rho_i^*$. These new coefficients are then used to reestimate opacity values along the temperature for the $\rho_i^*$.

Appendix A.4.4: Final tables – interpolation process

![Graph](image_url)

Fig. A.4. Merged table Planck (left) and Rosseland (right) opacity temperature (top) and density (bottom) first derivatives in the ($\rho, T$) plane, any in log scale; the grey shape represents the transition region.

We checked smoothness of the result by looking at the first derivatives. Figure A.4 shows no anomaly within the transition temperature range $[T_A, T_B]$ (grey shape). The remaining anomalies are not reached in our simulations.

The interpolation process is copied from the merging method, i.e. piecewise cubic Hermite polynomials along with van Leer slopes, since it satisfies criteria described in Section A.3.2. Interpolation is first performed along temperature at the $2 \times 2$ grid densities framing the requested density, so as to calculate van Leer slopes at the requested temperature and interpolate along density. Interpolating along temperature and then density showed to be slightly more accurate than interpolation along density first. This
is arguably due to stronger variations of opacities (especially Planck opacity) with respect to temperature.

**Appendix B: Chromospheric model**

One of our objectives is to describe the dynamics of the column and its impact on the chromosphere, as well as the feedback of the chromosphere on the column. This requires then to include an adequate description of the physical mechanism leading to the chromospheric heating. This appendix presents the simple but self-consistent model of a chromosphere used in this work.

**Appendix B.1: Motivations and limits**

The study of the solar chromosphere is a tough problem in itself. Its modeling is of interest for us since the base of the accretion column lies in the stellar chromosphere: the dynamics and observability of the column base may then depend on its structure and dynamics. Moreover, the chromosphere may be heated locally by the accretion process. The inner heating mechanism in the chromosphere is still subject of debates: it is mainly thought to originate either from acoustic waves dissipation (Biermann 1946; Schwarzschild 1948; or more recently Sobotka et al. 2016) or from MHD waves dissipation (Alfvén & Lindblad 1947; Jess et al. 2015).

Most accretion simulations model the stellar atmosphere – when it is modeled – as a hydrostatic plasma layer "tuned up" with ad-hoc sources to recover both temperature and pressure profiles (see e.g. the heating function empirically introduced by Peres et al. 1982). Although this must work for a static structure, it is delicate to predict the dynamic behavior of such a structure facing the continuous perturbation from an infalling plasma flow: such solution is not adapted to studies involving (in a self-consistent way) the dynamics of a perturbed atmosphere, like in the context of accretion.

We do not pretend to develop a "state of the art" model in this paper: we only aim at using a reasonable model that is both dynamic and self-consistent with our radiation hydrodynamics model. In our 1D model, we do not consider any magnetic effect but a very effective confinement of the accretion flow along the field lines. To allow fast qualitative comparison between our model and theoretical models & observations (see Figure B.1), we only used solar parameters (i.e. abundances, luminosity, mass and radius).

**Appendix B.2: Acoustic waves and shocks**

Acoustic waves are generated by photospheric granulation (see e.g. Judge 2006). These waves propagate upwards up to the height where their velocity overcome the local sound speed, and degenerate then into shocks. The nature of this mechanism is random: two different locations at the stellar surface will be crossed over by acoustic shocks that ought to be out of phase one with each other.

In our simulations, acoustic energy is supplied in the form of a monochromatic sinusoidal motion of the first Lagrangian interface \(T = 60\ s\) and \(f_{ac} = 10^8\ erg\ cm^{-2}\ s^{-1}\), see e.g. Rammacher & Ulmschneider 1992; Ulmschneider et al. 2005; Kalkofen 2007). Resulting acoustic waves propagate and degenerate into shocks. Figure B.1 shows several temperature snapshots of such simulation along with the chromospheric model from Vernazza et al. (1973). Below 300 km, acoustic waves are damped and hardly appear on snapshots. Above 500 km, waves are fully degenerated into shocks: their strength is then governed by the balance between steepening in the pressure gradient and dissipation. Since the corona and the upper chromosphere (above 10^3 km) are readily crushed by the accretion flow, the heating of these areas is not considered in our model.

**Appendix B.3: From solar to stellar chromosphere**

Observations of the solar chromosphere provide time and space averages of thermodynamics quantities (\(\rho, T, p, \ldots\)). Detailed observation of CTTS chromospheres would demand higher space and time resolution than the ones permitted by current observational technologies. Most works on this field rely then on scaling laws (see e.g. Ayres 1979; Calvet 1983) or ad hoc fittings to recover specific observational features (see e.g. Dumont et al. 1973; Cram 1979; Calvet et al. 1984; Batalha & Basri 1993).

**Appendix C: Radiation source terms in the Hybrid model**

This work encompasses several radiation regimes, from optically thick LTE radiation transfer (Section 2.2.3.1) to optically thin coronal NLTE regime (Section 2.2.3.2). The momenta equations (Section 2.2) can handle all of them, assuming the proper radiation source terms are provided.

In the LTE case, both radiation energy and momentum source terms are well defined (Eq. (4)). In coronal regime, this is not the case. Gas and radiation are decoupled in such a regime. Radiation only acts then as a gas energy sink: the radiation energy source term (the gas sink) boils down to a cooling function (Kirienko 1993, see e.g.). Computing the radiation flux is irrelevant in such regime and then no radiation momentum source term is provided. That is why we set \(s_{R}^\prime\) to 0.

In the "Hybrid" setup, we aim at modeling radiative conditions that are neither LTE nor coronal regimes but something in between. To determine if the situation is closer to one or the other, and how close, we choose to look at the probability for a photon to escape the accretion column (see Eq. (7)). We use it as a weighting factor to average the source terms, as shown in Section 2.2.3.3. The process is straightforward for the radiation energy source term, but not for the radiation momentum source term since \(s_{R}^\prime\) remains unknown. We assume then that the coronal Rosseland
mean opacity may not significantly differ from its LTE value. This intuition is reinforced by preliminary calculations concerning NLTE radiative collisional opacities (Pérez, priv. com.).