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of the hot halo gas

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Bastian Sander

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Der galaktische Materiezyklus umfasst die großräumigen Gasströmungen, die zum einen von Galaxien aus dem intergalaktischen Medium¹ in Form kalten, metallarmen Gases akkretiert und zum anderen durch Supernovae als heißes, hochionisiertes, prozessiertes Material aus der galaktischen Scheibe herausgeschleudert werden. Auch wird Gas aus Satellitengalaxien durch Gezeitenwechselwirkungen mit der Zentralgalaxie herausgelöst und ordnet sich als Gasschweif entlang des Orbits der Satellitengalaxie an. Alle diese Gasströmungen wechselwirken mit dem galaktischen Halo, der aus heißem, hochionisiertem und sehr verdünntem Gas besteht.

Ein substanzieller Anteil des einfallenden, metallarmen Gases besteht aus sogenannten „Hochgeschwindigkeitswolken“ (engl. high-velocity clouds, HVCs), deren Name sich von Radialgeschwindigkeiten ableitet, die viel zu hoch sind, um konsistent zu sein mit einer differentiellen Rotation der Galaxie. Selbst 60 Jahre nach ihrer Entdeckung durch Muller et al. (1963) entziehen sich HVCs in jeder Hinsicht einem umfassenden Verständnis. Weder ihr Ursprung noch der Grund für das völlige Fehlen von Sternentstehung sind geklärt. Ihre Entfernungen sind oftmals sehr ungenau bestimmt, wodurch sich Eigenschaften wie Masse, Dichte oder Größe ebenfalls unpräzise ableiten lassen. Diverse Einzelbeobachtungen im ultravioletten und optischen Bereich mithilfe der Absorptionsspektroskopie und von Röntgenemissionen, aber auch großangelegte Himmelsdurchmusterungen der 21-cm-Emissionslinie des neutralen Wasserstoffs wurden durchgeführt, um Säulendichten, Metallgehalte und Ionisierungsstufen, Geschwindigkeiten und Geschwindigkeitsdispersionen, den Staubgehalt, Morphologien und Gradienten, Magnetfelder und die Anzahl enthaltener Sterne zu bestimmen oder um die Existenz dunkler Materie nachzuweisen.

Obwohl viele hundert der dichtesten und schnellsten HVCs, die sog. kompakten und ultrakompakten HVCs², die Bedingungen für Sternentstehung erfüllen, wurden bis zum jetzigen Zeitpunkt keine Sterne entdeckt, die diesen Wolken zugeordnet werden können. Welches Wechselspiel von physikalischen Prozessen zu dieser Beobachtung führt, ist ungeklärt. In der vorliegenden Dissertation wird nach einer Antwort auf diese seit vielen Jahrzehnten wissenschaftlich diskutierten und bisher theoretisch unverstandenen Fragestellung gesucht. Dazu werden in dieser Arbeit multi-physikalische Simulationen auf Computerclustern durchgeführt. Es werden dazu die entsprechenden hydrodynamischen Gleichungen numerisch auf einem adaptiven Rechengitter gelöst. Zu

¹intergalactic medium (IGM)

²compact high-velocity clouds (CHVCs) und ultra-compact high-velocity clouds (UCHVCs)

den im Rahmen dieser Arbeit implementierten physikalischen Prozessen zählen saturierte Wärmeleitung, verschiedene Kühl- und Heizprozesse im Plasma, Massendiffusion, Dissoziation von Molekülen und Ionisierung von Atomen, eine variable Metallizität sowie ein temperaturabhängiges mittleres Molekulargewicht. Semi-implizite und semi-analytische numerische Integrationsmethoden werden verwendet.

Es zeigt sich, dass vereinfachende Annahmen in der numerischen Modellierung zu unrealistischen Entwicklungen solcher Wolken führen. Selbst bei Modellwolken mit Massen weit unterhalb ihrer Jeans-Masse darf die Eigengravitation nicht vernachlässigt werden, da sich in diesem Falle kein radialer Dichtegradient und somit keine Kern-Halo-Struktur ausbildet und die Rechenergebnisse zu falschen Schlussfolgerungen in Bezug auf Massenverlusten, Lebensdauern und Morphologien verleiten. Weiterhin wird gezeigt, dass selbst starke Magnetfelder, die eine von einem heißen Plasma umgebene Wolke gegen Wärmeleitung durch Elektronen abschirmen, den über die gesamte Wolkenoberfläche aufsummierten Wärmestrom nur um ca. $1/3$ abschwächen. Es bleibt die Erkenntnis aus dieser Dissertation, dass Eigengravitation und Wärmeleitung einen erheblichen Einfluss auf die Entwicklung selbst massearmer Wolken innerhalb eines heißen Gases haben.

Laut Theorie evaporieren kühle interstellare Wolken im heißen Plasma (Cowie & McKee, 1977). Die zeitabhängige Behandlung der Wärmeleitung sowie die Berücksichtigung von Plasmakühlung und Eigengravitation in den Simulationen führen jedoch dazu, dass stattdessen heißes Umgebungsgas auf die Wolken kondensiert. Im Rahmen dieser Dissertation werden zum ersten Mal die Kondensationsraten quantifiziert und eine Korrelation mit dem radialen Dichtegradienten gezeigt.

Weiterhin hängt die Dynamik des Mischprozesses aus heißem Umgebungsgas und Wolkengas vom radialen Dichtegradienten der Wolke ab. Je homogener die Wolke ist, desto schneller und gleichmäßiger verteilt sich kondensiertes Umgebungsgas und die darin enthaltenen Metalle in der Wolke. Nur bei steilem radialem Dichtegradienten in der Wolke sammeln sich diese Metalle in den Randbereichen der Wolke und werden nur sehr langsam ins Wolkeninnere transportiert. Weiterhin bildet sich eine Übergangszone zwischen Wolke und Umgebungsgas aus, die viel schmäler ist und auch näher an der Wolkenoberfläche liegt als nach Cowie & McKee (1977) erwartet wird. Da die Wärmeleitung den Heizeintrag in die Übergangszone dominiert, werden Temperaturen um ca. $10^{5,5}$ K erreicht, bei denen sich die in HVCs beobachteten hochionisierten Elemente bilden, vor allem O VI.

Die simulierten CHVCs bilden beobachtete Morphologien wie Kopf-Schweif-, Bugstoß- oder hufeisenförmige Strukturen. Wärmeleitung führt zu einem stetigen, faserigen

Massenabtrag von Wolkengas, der gering ist im Vergleich mit identischen Wolken ohne Wärmeleitung. Das modellierte heiße Umgebungsgas bildet einen extremen Zustand in Temperatur und Dichte des beobachteten galaktischen Halos. Es ist im Rahmen dieser Arbeit gelungen, einen physikalischen Grund für die Unterdrückung der Sternentstehung in HVCs nachzuweisen. Die saturierte Wärmeleitung wirkt permanent dem mit zunehmender Abbremsung der CHVCs verringertem dynamischem Druck entgegen. Die sich verdichtenden Gebiete innerhalb der Wolken werden durch den thermischen Gegendruck an einem Gravitationskollaps gehindert. Die Ergebnisse dieser Dissertation legen demnach nahe, dass keine Sterne in HVCs gebildet werden.

The Galactic matter cycle encompasses large-scale gas flows, which on the one hand are accreted from the intergalactic medium (IGM) in the form of cold, metal-poor gas by galaxies, and on the other hand are expelled from the Galactic disk as hot, highly ionized, processed material through supernovae. Gas from satellite dwarf galaxies is also stripped off through tidal interactions with the central galaxy and aligns itself as a gas tail along the orbit of the satellite galaxy. All these gas flows interact with the Galactic halo, which consists of hot, highly ionized, and very diluted gas.

A substantial portion of the infalling, metal-poor gas consists of so-called high-velocity clouds (HVCs), whose name derives from radial velocities that are much too high to be consistent with differential rotation of the galaxy. Even 60 years after their discovery by Muller et al. (1963), HVCs elude comprehensive understanding in every respect. Neither their origin nor the reason for the complete absence of star formation is clarified. Their distances are often very inaccurately determined, making it difficult to precisely derive properties such as mass, density, or size. Various individual observations by absorption spectroscopy in the ultraviolet and optical, by emission spectroscopy of X-rays, as well as large-scale sky surveys of the 21-cm emission line of neutral hydrogen have been conducted to determine column densities, metal contents and ionization stages, velocities and velocity dispersions, dust contents, morphologies and gradients, magnetic fields, and the number of contained stars, or to demonstrate the existence of dark matter.

Although many hundreds of the densest and fastest HVCs, the so-called compact high-velocity clouds (CHVCs) and ultra-compact high-velocity clouds (UCHVCs), fulfill the conditions required for star formation, no stars assigned to HVCs have been discovered to date. The combined interplay of physical processes leading to this observation is unclear. This thesis seeks an answer to this scientifically discussed and theoretically unexplained question, which is debated since many decades. Multi-physical simulations are conducted on computing clusters to address this question. Within this thesis, the corresponding hydrodynamic equations are solved numerically by using an adaptive mesh-refinement code. Physical processes implemented within this work include saturated thermal conduction, various cooling and heating processes in plasma, mass diffusion, dissociation of molecules and ionization of atoms, variable metallicity, and temperature-dependent mean molecular weight. Semi-implicit and semi-analytical numerical integration methods are used.

It is found that simplifying assumptions in numerical modelling lead to unrealistic evolutions of such clouds. Even in clouds with masses well below their Jeans mass, self-gravity must not be neglected, since in this case these low-mass clouds do not form a radial density gradient and thus no core-halo structure, leading to incorrect conclusions regarding mass-loss rates, lifetimes, and morphologies. Furthermore, it is shown that even strong magnetic fields, which shield a cloud surrounded by a hot plasma against thermal conduction by electrons, weaken the heat flux summed over the entire cloud surface by only $1/3$. The insight gained from this thesis is that self-gravity and thermal conduction have a significant impact on the evolution of even low-mass clouds within a hot gas.

According to analytical theory, interstellar clouds embedded in hot plasma evaporate (Cowie & McKee, 1977). However, treating thermal conduction time-dependently and considering plasma cooling and self-gravity in the simulations show that, instead, hot ambient gas condenses onto the clouds. Within this thesis, condensation rates are quantified for the first time and a significant correlation with the radial density gradient is shown.

Furthermore, the dynamics of the mixing process between hot ambient gas and cloud gas depends on the radial density gradient in the cloud. The more homogeneous the cloud, the faster and more evenly the condensed ambient gas and thus the contained metals distribute within the cloud. Only in the case of a steep radial density gradient in the cloud do these metals accumulate in the outskirts and are transported into the interior of the cloud only very slowly. Furthermore, a transition zone between cloud and ambient gas is formed, which is much narrower and closer to the cloud surface than expected by Cowie & McKee (1977). Thermal conduction dominates the heat input into the transition zone hence generating transient temperatures about $10^{5.5}$ K, at which the highly ionized elements observed in HVCs are formed, especially O VI.

The simulated CHVCs reproduce observed morphologies such as head-tail, bow-shock, or horseshoe-shaped structures. Thermal conduction leads to a continuous, filamentary mass removal of cloud gas, which is low compared to identical clouds without thermal conduction. The modelled hot ambient medium represents an extreme state of the Galactic hot halo in terms of density and temperature. Within this work a physical reason for the suppression of star formation is provided. Saturated thermal conduction permanently counteracts a ram pressure, which decreases with ongoing deceleration of the CHVCs. The compressed subregions within the clouds are thus prevented from gravitational collapse by thermal counter-pressure. The results of this thesis therefore suggest that no stars are formed in HVCs.

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1 INTRODUCTION

To our current understanding, the environments of galaxies are far from being equilibrated in every respect. The region outside of their respective Galactic disk but within their virial radius is commonly dubbed circumgalactic medium (CGM). The physics in the CGM is governed by a Galactic matter cycle (see reviews by Tumlinson et al., 2017; Faucher-Giguere & Oh, 2023), which is sketched in figure 1.1.

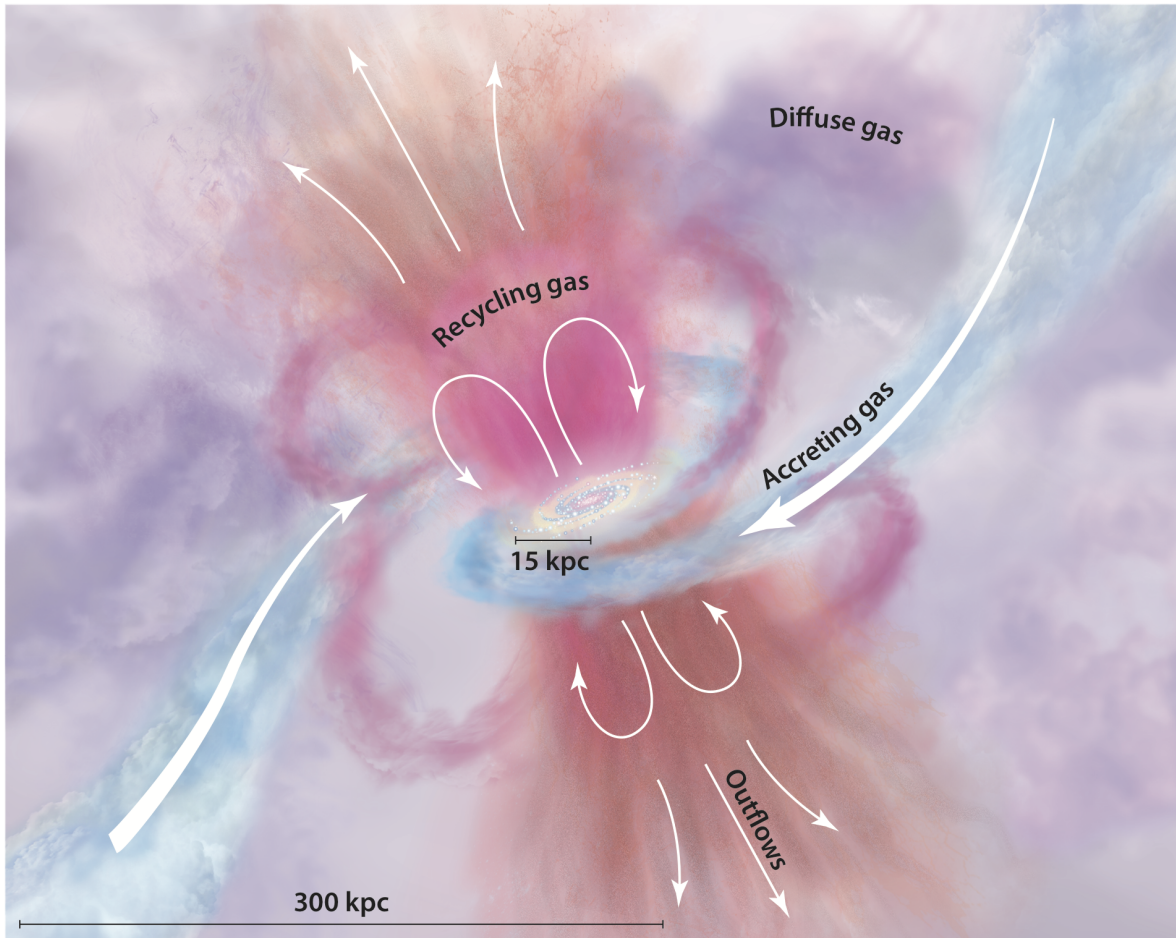


Figure 1.1: A sketch of the Galactic matter cycle with all its components feeding both the ISM and the CGM: outflows, fountains, inflows. [figure taken from Tumlinson et al. (2017)]

On the one hand, processed matter is continuously transported from the interstellar medium (ISM) to the CGM (Werk et al., 2019): Stars are formed in molecular clouds within the Galactic disk and they explode within supernovae at the end of their evolutionary stages. The hot gas is blown away and forms bubbles, which open in chimneys, leave the disk and feed the hot, tenuous gaseous halo around galaxies

(chapter 4), which provides the highly ionized part of the CGM. Due to the thermal instability (Field, 1965) subregions in the halo can precipitate and consequently form cool, dense, and gravitationally bound multi-phase clouds, which fall back to the disk while they undergo hydrodynamic interactions with the ambient halo gas. These clouds provide in turn the sites for star formation. This innergalactic part of the matter cycle is referred as the *Galactic fountain* (Shapiro & Field, 1976; Bregman, 1980; Fraternali, 2017). However, the rate of such infalling neutral gas at intermediate velocities ($\approx 30 \dots 100 \text{ km s}^{-1}$, see section 3.2.5) in the disk-halo interface ($z_{\text{gal}} < 2.5 \text{ kpc}$) amounts to only $\dot{M} = 0.01 \dots 0.05 \text{ M}_{\odot} \text{ yr}^{-1}$ (Richter, 2016), which is not sufficient to sustain the current star-formation rate $\dot{M}_{\text{SF}}$ deduced for the Milky Way. Several rates are inferred (e.g., Diehl et al., 2006; Robitaille & Whitney, 2010; Chomiuk & Povich, 2011; Licquia & Newman, 2015) with differences due to the initial mass function adopted. The most recent value is given by $\dot{M}_{\text{SF}} = 1.7 \pm 0.7 \text{ M}_{\odot} \text{ yr}^{-1}$ (Fox et al., 2019).

On the other hand, the CGM is filled with pristine, metal-poor gas that is being accreted from the intergalactic medium along cosmic filaments (Putman et al., 2012). Among infalling gas there is the manifold sample of high-velocity clouds (e.g., Gibson, 2002). The accretion rates for various compilations of high-velocity clouds (HVCs) are given in table 1.1. It is apparent that infalling high-velocity gas is able to compensate $\dot{M}_{\text{SF}}$ to a substantial amount.

Table 1.1: Accretion rates [$\text{M}_{\odot} \text{ yr}^{-1}$] of the Milky Way for various compilations of high-velocity gas mainly inside a galacto-centric distance of 20 kpc.

compilation	accretion rate	reference
neutral and ionized HVCs incl. Magellanic System ^(a)	> 6.1	Richter et al. (2017a)
Magellanic System	$\gtrsim 3.7^{(b)}$	Fox et al. (2014)
Magellanic Stream (MS)	$\gtrsim 1.9^{(b)}$	
ionized HVCs	$0.4 \dots 1.4$	Lehner & Howk (2011)
neutral (HVCs and MS)	$\sim 0.7^{(b)}$	Richter (2012)
HVC complexes	$0.08 \dots 0.4$	Putman et al. (2012)

^(a) The Magellanic System is the most prominent satellite orbiting the Milky Way. It consists of the Small and Large Magellanic Clouds (both visible to the naked eye), which are connected by the Magellanic Bridge, the Leading Arm, and the MS (see section 3.2.2).

^(b) At an assumed distance for the Magellanic System of 55 kpc.

Thus, extragalactic infalling HVCs provide a substantial reservoir to replenish gas consumed by star formation. Understanding the physics of their interactions with the CGM is crucial to understanding the Galactic cycle of matter. This thesis aims at investigating the still open question how HVCs could add to the current star-formation rate by providing potential gas reservoirs.

This thesis is structured as follows:

Chapter 2 motivates the research question addressed in this thesis while chapters 3 to 5 put the research of this thesis into the scientific context. In chapter 3 the current status of research regarding different categories of HVCs is constituted, the physical processes being key for HVCs are presented, and possible origins of HVCs are discussed. In chapter 4 the current knowledge of the Galactic hot halo is described with emphasis on observational signatures of its sub-virial, virial, and super-virial phases. Chapter 5 comprises a brief introduction to the theory of star formation, and summarizes the results of observational campaigns aimed at finding stars in HVCs.

Chapters 6 to 9 present novel scientific results based on research carried out within this thesis. Within chapter 6 the simulation code and the numerics of all implemented physical processes are introduced. The initial model setups, methods for analysis, and central results are subsumed in chapter 7. The three published papers about the research work within this thesis form the main part. Chapter 8 summarizes the most important results figured out within this thesis. Finally, in chapter 9 potential follow-up research is discussed.

Several appendices contain additional information on rest frames, methods for distance determination, spectral lines, and the von-Neumann stability analysis for numerical integration schemes.

2 AIM OF THIS WORK

Within the Galactic matter cycle fresh gas needs to replenish gas that is continuously consumed by star formation. Neutral gas falling from the intergalactic medium (IGM) into the virial radius of the Milky Way (and other similar galaxies at $z \sim 0$) represents an essential reservoir for this. Much of this gas exists in form of high-velocity clouds (HVCs) and the accretion rate (see table 1.1) is sufficient to compensate for the current Galactic star-formation rate $\dot{M}_{\text{SF}} = 1.7 \pm 0.7 \text{ M}_{\odot} \text{ yr}^{-1}$ (Fox et al., 2019). The most likely candidates for star-formation activity are compact high-velocity clouds (CHVCs) and ultra-compact high-velocity clouds (UCHVCs) (sections 3.2.3 and 3.2.4), as they best meet the required conditions (section 5.1). However, no stellar counterparts assigned to these HVCs have ever been detected (section 5.3). This contradiction consequently leads to the following general question regarding HVCs:

Why there is no star formation observed in HVCs?

The virial sphere around our galaxy is filled with a highly ionized medium (HIM) (chapter 4) so the environmental conditions (density, temperature, pressure) change for the HVCs when entering this sphere. While moving through this Galactic hot halo the HVCs are subject to hydrodynamic instabilities, ram pressure, drag, thermal conduction, dissociation and ionization, cooling and heating, mixing with ambient gas (section 3.3). Is it an unsolved issue yet whether CHVCs and UCHVCs are *sufficiently compressed locally* under such complex conditions, such that a gravitational collapse is initiated thus leading to a star formation site (section 5.1). Hence, the present thesis primarily aims at finding an answer to the central research question:

Is star formation possible in CHVCs and UCHVCs under extreme conditions of the hot halo gas?

To systematically investigate this central question numerical simulations are conducted by means of an adaptive mesh-refinement code that provides an appropriate numerical tool to investigate multi-physics at different timescales and ranging over several spatial scales. For this purpose, several model clouds of CHVCs and UCHVCs moving through a Galactic hot halo are set up assuming an extragalactic origin and at galacto-centric distances of $\sim 100 \text{ kpc}$. The internal evolution of the model clouds is traced over 100 Myr , which corresponds to many respective dynamical times. By varying parameters (e.g. metallicity) or considering relevant physical processes (e.g.

thermal conduction, self-gravity) the evolved model clouds are compared in terms of morphologies, mass loss rates, column densities and other features observed at CHVCs and UCHVCs. By this, processes acting in favour or against a local gravitational collapse are figured out in simulated CHVCs and UCHVCs and it is shown quantitatively why no stars are observed in HVCs.

Another main focus lies with saturated thermal conduction and its effect on condensation rates. According to the current state of knowledge there are no quantified condensation rates for interstellar clouds within an ionized medium provided in the literature. Knowing the condensation rates is pivotal to reveal mixing processes and the spatial distribution of accreted metals within the clouds. Moreover, the formation of a transition zone between cool cloud and hot ambient medium due to time-dependent thermal conduction is not investigated. Detailed comparisons with analytical theory are lacking hitherto. Understanding the physics within transition zones gives a clue on the formation of highly ionized elements, especially O VI. To explore the meaning of saturated thermal conduction a simulation suite of multi-phase clouds resting in a hot gas is carried out and analyzed in detail.

3 HIGH-VELOCITY CLOUDS

3.1 A HISTORICAL REVIEW

The first high-velocity clouds (HVCs) are discovered by Muller et al. (1963), who have used the 25 m radio telescope in Dwingeloo (the Netherlands) to detect the 21-cm emission line from neutral hydrogen (appendix D.1). They report on four clouds with radial velocities within the velocity range of $-116 \dots -175 \text{ km s}^{-1}$ relative to the Local Standard of Rest (LSR) at high Galactic latitudes ($|b| \gtrsim 40^\circ$, see appendix B). These HVCs are today known as parts of Complexes A and C (van Woerden et al., 2004; Wakker, 2004). The column densities of these four clouds are between $0.4 \dots 2 \times 10^{20} \text{ cm}^{-2}$ and their velocity dispersions are quite high at $\Delta v_{\text{FWHM}} \sim 25 \text{ km s}^{-1}$. Even though Muller et al. (1963) are unable to determine any distance, they conclude from statistical considerations that these clouds must be located in the hot halo of the Milky Way, a gaseous medium proposed shortly before by Spitzer (1956). Several others add to these findings of high-velocity 21-cm emissions (e.g. Smith, 1963; Dieter, 1964; Hulsbosch & Raimond, 1966) by successively extending the searches to lower latitudes, e.g. $|b| \gtrsim 30^\circ$ (Muller et al., 1966) and $|b| \gtrsim 20^\circ$ (Kerr & Sullivan, 1969). As one of the first, Kepner (1968) emphasizes that distance determination is vitally important for both deriving physical key properties and substantiating hypotheses on origins.

Within the 1960's, astronomers especially in the Netherlands conducted systematic searches for HVCs using the Dwingeloo telescope. Neutral hydrogen clouds at many different radial velocities are found. As van Woerden et al. (2004) point out the operational threshold of $v_{\text{LSR}} = -100 \text{ km s}^{-1}$ is decided based upon a division of labour: Leiden astronomers analyze the high-velocity gas at $v_{\text{LSR}} < -100 \text{ km s}^{-1}$, while in Groningen the intermediate-velocity clouds (IVCs) at $-100 < v_{\text{LSR}} / \text{km s}^{-1} < -50$ are studied. Equipped with a much more improved receiver at Dwingeloo telescope, Hulsbosch (1978) was able to scan a broader velocity range ($-500 \dots +500 \text{ km s}^{-1}$) and detected many new gas clouds at $v_{\text{LSR}} < -300 \text{ km s}^{-1}$, which are named VHVCs (Very HVCs).

Since then, HVCs are detected using certain other facilities. Observations by earth-bound telescopes are restricted to the detection of the 21-cm line emitted by H I and the $\text{H}\alpha$ line emitted by H II. With the rise of detectors aboard satellites in the early 1990's HVCs are observed in a broader wavelength range in both emission (e.g. soft X-rays emitted by multiple ionized elements as observed by ROSAT satellite, cf. Kerp et al.,

1996b) and absorption (e.g. ultraviolet (UV) radiation absorbed by O VI as observed by FUSE satellite, cf. Collins et al., 2006). As well, particular rotational transitions of ^{12}CO are observed in infrared (IR) emission using the IRAC instruments aboard Spitzer satellite (Kim et al., 2007). The most important observation facilities are listed in table 3.1.

Table 3.1: Some of the most important observation facilities that have been used to detect HVCs. The angular resolution (half-power beam-width, HPBW) depends on the particular frequency band used in an observation.

facility	operational period	spectral domain	angular resolution (HPBW)
earth-bound telescopes			
Dwingeloo single dish	1956 to 2000	radio	36'
Parkes single dish	since 1961	radio	11'.5
Arecibo single dish	1963 to 2020	radio	3'.5
IAR single dish ^(a)	since 1966	radio	34'
WSRT interferometer ^(b)	since 1970	radio	2'
Effelsberg single dish	since 1972	radio	9'
ATCA interferometer ^(c)	since 1988	radio	1'
WHAM ^(d)	since 1996	optical	1°
GBT single dish ^(e)	since 2000	radio	9'.1
detectors aboard satellites			
XRT ^(f)	1991 to 1999	X-rays	5''
FUSE ^(g)	1999 to 2007	ultraviolet	
XMM-Newton ^(h)	since 1999	X-rays	6''
CXO ⁽ⁱ⁾	since 1999	X-rays	0''.5
IRAC ^(j)	2003 to 2020	infrared	2''
COS ^(k)	since 2009	ultraviolet	0''.05

Full names of the facilities:

- ^(a) Instituto Argentino de Radioastronomía (Argentine Institute of Radio Astronomy)
- ^(b) Westerbork Synthesis Radio Telescope
- ^(c) Australia Telescope Compact Array
- ^(d) Wisconsin H α Mapper
- ^(e) Green Bank telescope
- ^(f) X-ray Telescope (aboard ROSAT satellite)
- ^(g) Far Ultraviolet Spectroscopic Explorer
- ^(h) X-ray Multi-Mirror
- ⁽ⁱ⁾ Chandra X-ray Observatory
- ^(j) Infrared Array Camera (aboard Spitzer satellite)
- ^(k) Cosmic Origins Spectrograph (aboard Hubble Space Telescope)

The above facilities are used to conduct a number of observation campaigns to reveal specific physical properties of HVCs. Some of the most extensive surveys are shown in table 3.2. The findings all contribute to our current knowledge on the nature of HVCs, which are proven to be common in the Milky Way (Westmeier, 2018).

Table 3.2: Major surveys for Galactic H I that are important to study HVCs (partially taken from Peek et al., 2011). The SECCO survey was specifically designed to search for stars in HVCs.

survey acronym	duration	main references
Dwingeloo All-Sky Survey	1976 to 1986	Bajaja et al. (1985), Hulsbosch & Wakker (1988)
LDS ^(a)	1988 to 1993	Burton & Hartmann (1994), Hartmann & Burton (1997)
LAB ^(b)	1989 to 2005	Kalberla et al. (2005)
HIPASS ^(c)	1997 to 2002	Putman & Gibson (1999a,b), Putman et al. (2002)
AGES ^(d)	2004 to 2013	Auld et al. (2006)
GASS ^(e)	2005 to 2006	McClure-Griffiths et al. (2009)
ALFALFA ^(f)	2005 to 2012	Haynes et al. (2011), Adams et al. (2013)
GALFA-H I ^(g)	2005, 2009 to 2013	Peek et al. (2011, 2018)
EBHIS ^(h)	2008 to 2013	Winkel et al. (2010), Kerp et al. (2011)
SECCO ⁽ⁱ⁾	2014	Bellazzini et al. (2015a,b), Beccari et al. (2016)
HI4PI ^(j)	2016	HI4PI Collaboration et al. (2016), Westmeier (2018)

Full names of the surveys:

- (a) **L**eiden/**D**wingeloo H I Survey
- (b) **L**eiden/**A**rgentine/**B**onn H I Survey
- (c) **H**I **P**arkes **A**ll-**S**ky Survey
- (d) **A**recibo **G**alaxy **E**nvironment **S**urvey
- (e) **G**alactic **A**ll **S**ky H I Survey
- (f) **A**recibo **L**egacy **F**ast **A**recibo **L**-band **F**eed **A**rray H I Line Survey
- (g) **G**alactic **A**recibo **L**-band **F**eed **A**rray H I Survey
- (h) **E**ffelsberg-**B**onn **H**I Survey
- (i) **S**tellar **C**ounterparts of **C**ompact high-velocity clouds
- (j) **H**I **4** π Survey

HVCs within the Galactic hot halo of the Milky Way are mainly infalling as indicated by negative LSR velocities (cf. figure 3.1). About 37 % to 40 % of the Galactic sky is covered by HVCs (Murphy et al., 1995), which mainly consist of H I and whose column densities are typically above $5 \times 10^{17} \text{ cm}^{-2}$ (Wakker, 1991). About 74 % of the sightlines analyzed by Lehner et al. (2012) pass through HVCs having $N(\text{H I}) < 3 \times 10^{18} \text{ cm}^{-2}$ whereas $N(\text{H I}) < 8 \times 10^{17} \text{ cm}^{-2}$ for 46 % of sightlines. They estimate the sky covering factor of UV-selected HVCs with $|v_{\text{LSR}}| \geq 90 \text{ km s}^{-1}$ to be about 68 % for the Galactic sky at $|b| \geq 20^\circ$. Richter et al. (2017b) conduct the so-far largest survey of metal absorption in Galactic HVCs using archival UV spectra of extragalactic background sources. The overall sky-covering fraction of high-velocity absorption is 77 ± 6 % for the

most sensitive ion in their survey, Si III, and for column densities $\log N(\text{Si III}) \geq 12.1$. This value is $\sim 4 - 5$ times higher than the covering fraction of 21-cm neutral hydrogen emission at $\log N(\text{H I}) \geq 18.7$ along the same lines of sight, demonstrating that the Milky Way's circumgalactic medium (CGM) is multi-phase and predominantly ionized. They estimate the total mass and accretion rate of neutral and ionized HVCs to be $M_{\text{HVC}} \geq 3.0 \times 10^9 M_{\odot}$ and $\dot{M}_{\text{HVC}} \geq 6.1 M_{\odot} \text{ yr}^{-1}$, respectively. The Magellanic Stream contributes with more than 90% to this mass and mass flow. Putman & Gibson (1999b) mention a striking asymmetry in the HVC velocity distribution between northern and southern hemisphere. The maximum $|v_{\text{LSR}}|$ seen in the north is $\sim 450 \text{ km s}^{-1}$, which is $\sim 100 \text{ km s}^{-1}$ greater than the maximum seen in the south.

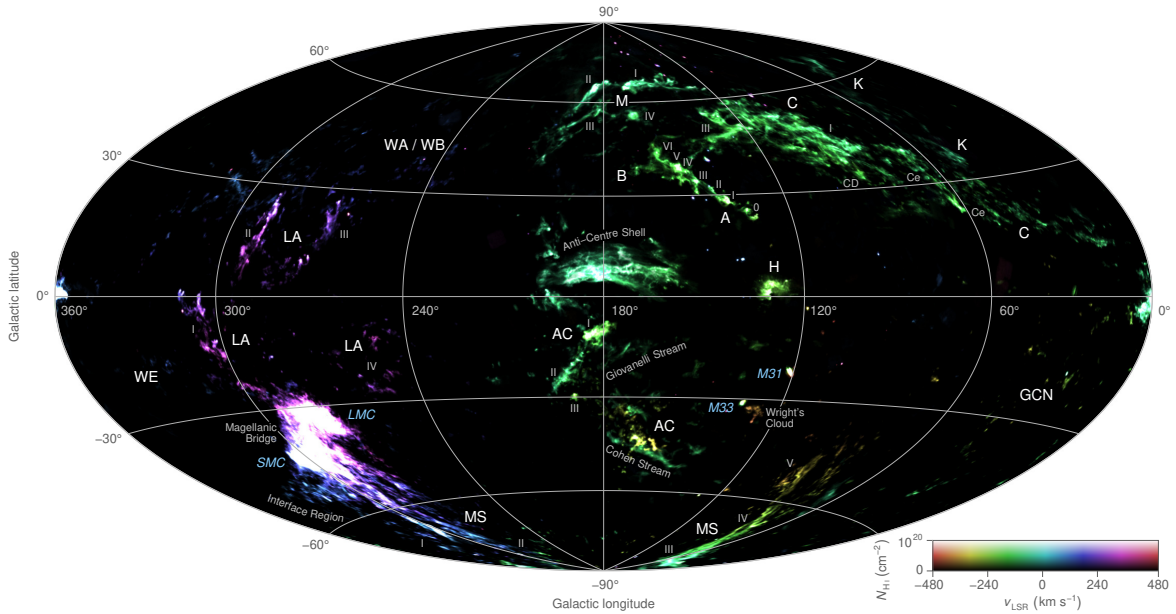


Figure 3.1: All-sky map of HVCs within the Galactic hot halo of the Milky Way based on the HI4PI. The velocities with respect to the LSR are shown, ranging from negative (in direction towards the sun) to positive values (away from the sun). Column densities range from the lower detection limit at $\sim 2.3 \times 10^{18} \text{ cm}^{-2}$ (HI4PI Collaboration et al., 2016) up to 10^{20} cm^{-2} . [figure taken from Westmeier (2018)]

Supported by numerical simulations (de Avillez, 2000; Sternberg et al., 2002; Heitsch & Putman, 2009) HVCs are observationally identified to be multi-phase clouds (Fox et al., 2005; Lehner et al., 2009b; For et al., 2016). That is, the observations in emission and absorption at different wavelengths reveal mainly neutral hydrogen and many elements at different ionization stages (such as Si II to Si IV, C II to C IV, O I to O VI, cf. appendix D), even molecules (CO, HCO⁺, for H₂ see Richter et al., 1999) and dust (Miville-Deschênes et al., 2005; Miville-Deschenes, 2006). Further, spatially distinct line broadenings (see appendix D.2) imply cold cores and warm envelopes (core-halo

structure, cf. Faridani et al., 2014). A variety of cloud morphologies (head-tail, bow-shock, horseshoe-shaped structures) trace hydrodynamic interactions with the hot halo (Brüns et al., 2000; Brüns & Mebold, 2004).

Stars are not yet found in HVCs at all. Consequently, precise distance measurements are not possible, which puts rather loose constraints on several key characteristics that are derived from the distance (see appendix C). It is yet not clear to which extent HVCs contribute to star formation. On the one hand, they provide a sufficient reservoir for replenishing gas that is converted into stars and some meet star-forming conditions. On the other hand, all hitherto surveys specifically searching for stellar counterparts assigned to HVCs failed (e.g., Ivezić & Christodoulou, 1997; Hopp et al., 2003a,b, see also section 5.3). Within the present thesis several environmental and cloud conditions are examined, and physical processes that support or suppress the formation of stars are studied by means of multi-physics numerical simulations (see results and discussion in section 7.6).

3.2 CATEGORIES

HVCs appear at the sky as single clouds or in groups of different sizes and properties. It turns out that two categories reflect different physical circumstances:

1. HVC complexes are several kiloparsecs in diameter and are groupings of smaller clouds (so-called *cloudlets*) showing both spatial and kinematic coherence. Single cloudlets can be separated by several degrees and are sometimes just loosely related (Wakker & van Woerden, 1991). Depending on the respective sightline, the observed ions, and the scanned LSR velocity range, a single complex usually reveals various metallicities, velocity dispersions, and column densities (section 3.2.1).

2. Single HVCs comprise the spatially extremely confined *compact high-velocity clouds (CHVCs)* and *ultra-compact high-velocity clouds (UCHVCs)*. They have sizes of a few ten to some hundred parsecs. They have higher peak column densities than HVC complexes, show typical morphologies indicating hydrodynamic interactions with their environments, and have multiple phases. Some of these clouds are far more distant than complexes and reside in the outer Galactic halo. The particular populations of CHVCs and UCHVCs is especially important within this thesis (see sections 3.2.3 and 3.2.4).

Contrary to HVCs there exist H I clouds that only slightly contradict to a motion of Galactic differential rotation. These intermediate-velocity clouds (IVCs) have common LSR velocities $|v_{\text{LSR}}| < 90 \text{ km s}^{-1}$. They are usually closer to the Galactic disk and show higher metallicities than HVCs implying a different formation scenario. A comparison between HVCs and IVCs is drawn in section 3.2.5.

The naming convention for single HVCs follows a certain scheme. Particular clouds are named HVC $\langle l \rangle \langle \pm b \rangle \langle \pm v_{\text{LSR}} \rangle$, with Galactic longitude l and Galactic latitude $\pm b$ of their peak H I column density (cf. appendix E), and $\pm v_{\text{LSR}}$ denotes their velocity measured with respect to the LSR (e.g. HVC 100-7+100, HVC 131+1-200).

Complexes are named with capital letters (e.g. Complex C, Complex L), because it is not clear whether one should use their brightest core, the average for the cores, a position-weighted or an intensity-weighted average to determine l , b , and v_{LSR} (Wakker et al., 2004).

3.2.1 HVC complexes

Several complexes have been identified over the past decades. The northern galactic sky is dominated by Complexes A and C, while in the southern hemisphere the Magellanic Stream, originating in the Magellanic Clouds, is major. The most massive and spatially most extended HVC complexes, which are located within the Galactic hot halo and contribute most to the mass infall towards the Galactic disk, are briefly characterized below. Their most relevant properties are listed in table 3.3.

Complex A is the first complex discovered (Muller et al., 1963), is one of the most prominent northern HVC complexes, and the most intensively studied HVC. It is also known as “Chain A” and extends over $(\Delta l, \Delta b) = (132 \dots 170, +32 \dots +45)$ deg (Welsh et al., 1996), as can be seen in figure 3.2 (left). Seven cloudlets have been identified in Complex A (Wakker & van Woerden, 1991) with particle densities and masses of $n(\text{H I}) = 1.3 \dots 3.2 \times 10^{-2} \text{ cm}^{-3}$ and $M(\text{H I}) = 0.5 \dots 4.5 \times 10^5 M_{\odot}$, respectively, which sum up to a total mass of $M(\text{H I}) = 1.6 \dots 2.3 \times 10^6 M_{\odot}$ (Barger et al., 2012).

The morphology (head-tail structure with multiple cores, Rayleigh-Taylor fingers on the low-latitude edge, two large wings generated by Kelvin-Helmholtz instabilities, see Welsh et al., 1996; Brüns et al., 2000; Barger et al., 2020) indicates ongoing hydrodynamic interactions with the ambient hot halo gas. Complex A is most likely accreted by the Milky Way from intergalactic medium (IGM) (van Woerden et al., 1998, 1999d; Barger et al., 2012).

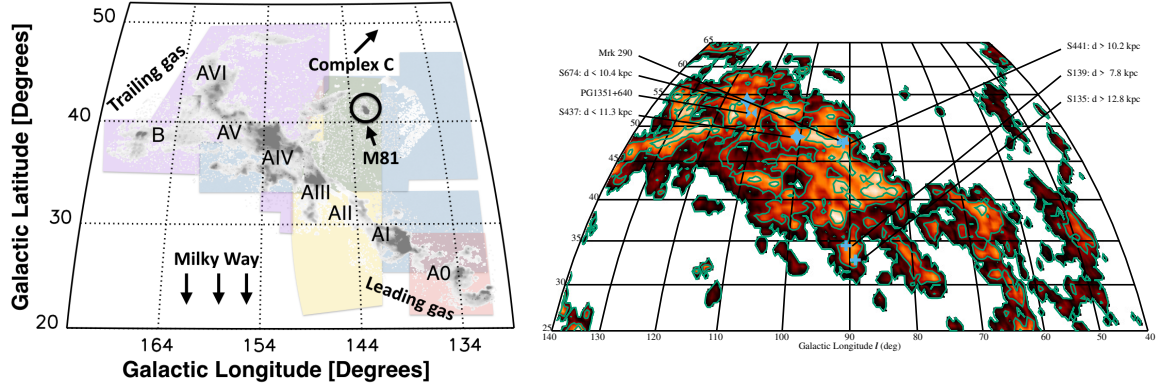


Figure 3.2: *Left plot:* A spectroscopically resolved H I 21-cm map of Complex A based on both new and archival Green Bank Telescope observations. The eight cloudlets with high $N(\text{H I})$ are labelled A 0 to A VI and B. [figure taken from Barger et al. (2020)] *Right plot:* Column densities in Complex C (*green isocontours*) in steps of $\log N(\text{H I}/\text{cm}^{-2}) = 18.2, 19.0, 19.7$. [figure taken from Thom et al. (2008)]

Complex C is first observed by Hulsbosch (1968) in H I. It is both the most massive HVC complex with $M(\text{H I}) \sim 5 \times 10^6 M_{\odot}$ (Thom et al., 2008) and the largest one covering $\sim 2000 \text{ deg}^2$ (Collins et al., 2007) in the northern Galactic sky (figure 3.2, right) that results in a mass inflow-rate of $0.1 \dots 0.25 M_{\odot} \text{ yr}^{-1}$ (Wakker et al., 2007). Abundances of observed ion species vary significantly between different sightlines (e.g., Gibson et al., 2001; Wakker et al., 2001; Sembach et al., 2004; Shull et al., 2011), thus Complex C is not well characterized by a single metallicity. The observed ratios of transition-temperature ions (such as C IV, N V, O VI) are well explained by either ionization by shocks or interaction with an ambient hot medium via turbulent mixing or conductive interfaces (Fox et al., 2004; Collins et al., 2007). Some sightlines reveal dust in Complex C (Miville-Deschênes et al., 2005; Fox et al., 2023).

Complex C consists of at least 90 cloudlets (Hsu et al., 2011) with masses ranging from $12.6 M_{\odot}$ to $6.3 \times 10^4 M_{\odot}$, sizes of $15 \dots 400 \text{ pc}$, and a median volume density of 0.065 cm^{-3} . It has multiple phases and shows signatures of ongoing phase transition from warm diffuse gas to cold dense gas (Marchal et al., 2021). Complex C likely interacts with the thick disk and lower halo of the Milky Way (Collins et al., 2007).

Complex H is one of the largest HVCs covering $\sim 480 \text{ deg}^2$ on the sky (Wakker & van Woerden, 1991; Simon et al., 2006). It is located close to the Galactic plane (figure 3.3, left), but in the very outskirts at a galacto-centric distance of $33 \pm 9 \text{ kpc}$ (Lockman, 2003).

Complex H separates into a central, bright core (Wakker et al., 1998) and a fainter, more diffuse tail (see figure 2 in Lockman, 2003). A narrow and a broad H I line

component reveal a two-phase structure in the core suggesting pressure equilibrium with the ambient low-density medium (Lockman, 2003).

The faint tail, which fills much of the gap in velocity between Complex H and the Galactic disk (Simon et al., 2006), provides strong evidence for ongoing interaction with the Milky Way. Regarding the integrated core mass of $M(\text{HI}) \approx 3 \times 10^6 M_\odot$ (Lockman, 2003) it is one of the most massive known HI clouds that does not contain any stars (star-formation rate $\lesssim 5 \times 10^{-4} M_\odot \text{yr}^{-1}$, Simon et al., 2006).

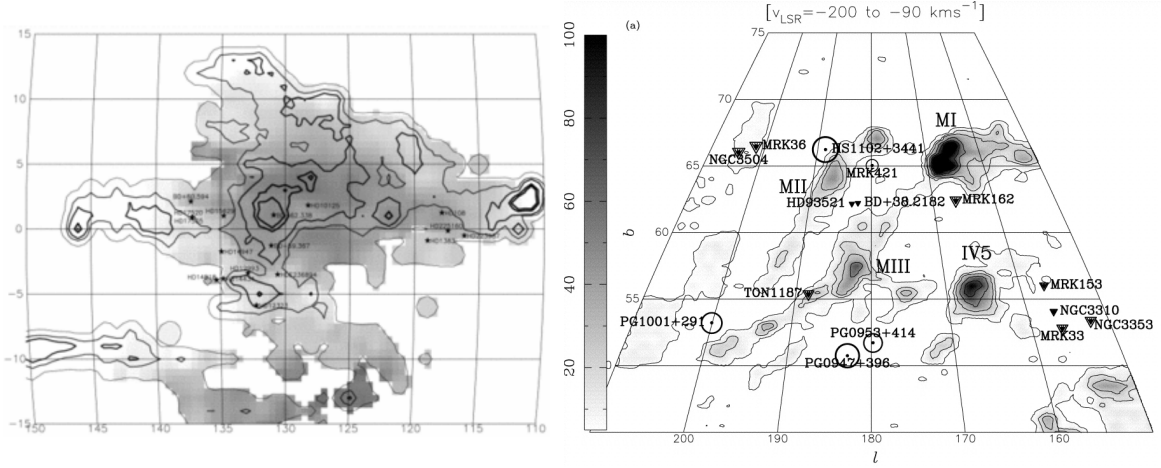


Figure 3.3: *Left plot:* Complex H as observed with International Ultraviolet Explorer (IUE) echelle spectrograph. The *greyscale map* are LSR velocities ranging from -200 km s^{-1} (*dark grey*) to -100 km s^{-1} (*light grey*) with *black isocontours* showing column densities in HI of $0.2 \dots 7.2 \times 10^{19} \text{ cm}^{-2}$. [figure taken from Wakker et al. (1998)] *Right plot:* Map of LSR velocities of Complex M (*greyscale map*) with overlaid $N(\text{HI})$ (*black isocontours* at $0.5, 1.5, 2.5, 5.0, 7.5 \times 10^{19} \text{ cm}^{-2}$) obtained from Leiden-Dwingeloo Survey. [figure taken from Savage et al. (2005)]

Complex M is located at very high Galactic latitudes spanning over $+45 \dots +70$ deg (Keenan et al., 1995, figure 3.3, right). It is composed of three distinct cloudlets with masses ranging from $M(\text{HI}) = 2.2 \dots 22 \times 10^3 M_\odot$ (Danly et al., 1993) to $M(\text{HI}) = 4 \times 10^3 D_{\text{sun}}^2 [\text{kpc}^2] M_\odot$ (Wakker, 2001).

Because of both its small distance to the Galactic disk and high metallicity, Complex M firmly resides in the Milky Way Halo and does not consist of primordial gas (van Woerden et al., 1998). It is proposed to be the high-velocity part of the IV Arch (for which near solar metallicity has been found, Kuntz & Danly, 1996; Wakker, 2001) and likely traces the condensed gas on return from a Galactic fountain (Yao et al., 2011). This interpretation is substantiated by an extended high positive velocity wing of O VI absorption extending to $v_{\text{LSR}} = +165 \text{ km s}^{-1}$ that may trace the outflow of hot gas into

the Galactic halo (Savage et al., 2005). Any far-IR emission from dust is not found Wakker & Boulanger (1986).

Complex Galactic Center Positive is also known as “Smith Cloud” or HVC 40-15+100. It is located in the region $(\Delta l, \Delta b) = (30 \dots 60, -10 \dots -40)$ deg (Wakker & van Woerden, 1991, figure 3.4, left) and has a size of 2×8 deg² corresponding to a radius of $0.5 \dots 1.8$ kpc (Wakker et al., 2008).

Hill et al. (2009) report on $M(\text{HI}) \gtrsim 6 \times 10^6 M_\odot$ (unequally distributed over head, tail, two clumps, and a diffuse component) and $M(\text{H II}) > 3 \times 10^6 M_\odot$. Wakker et al. (2008) deduce $n(\text{HI}) = 0.10 \dots 0.16 \text{ cm}^{-3}$. Dust is not found in this complex (Minter et al., 2024). A magnetic field $\gtrsim 8 \mu\text{G}$ is associated with Galactic-Centre-Positive (GCP) (Hill et al., 2013, based on measures of Faraday rotation). Stark et al. (2015) infer an average star formation rate surface density of $10^{-4.8 \pm 0.3} M_\odot \text{ yr}^{-1} \text{ kpc}^{-2}$ by assuming a constant star formation triggered ~ 70 Myr ago when the cloud first impacted the Galactic disk.

Leite et al. (2016) argue that GCP is a promising testbed for indirect detection of dark matter (DM) by synchrotron radiation associated with DM annihilations into electrons and positrons. The hypothesis of a DM confined GCP is also supported by the findings of Tepper-García & Bland-Hawthorn (2018b), whose simulations additionally explain the head-tail shape of GCP. Its origin is still speculative. On the one hand, it is difficult to accelerate such a massive cloud by a Galactic event thus ruling out a Galactic origin. On the other hand, the prograde and only mildly inclined orbit excludes an extragalactic origin (Lockman et al., 2008).

Complex Galactic Center Negative has a total angular extension of 134 deg² (Wakker & van Woerden, 1991, figure 3.4, right) and comprises 243 cloudlets, which belong to several subpopulations (distinct in the (l, b, v_{GSR}) space) that do not share a common origin (Winkel et al., 2011). Sembach et al. (1999) apply a photo-ionization model of ion ratios observed toward Markarian 509 and derive $n(\text{HI}) \sim 10^{-4} \text{ cm}^{-3}$ and gas temperatures $2 \dots 3 \sim 10^4 \text{ K}$. Kerp et al. (1999) correlate the ROSAT 1/4 keV all-sky survey with the Leiden/Dwingeloo HI survey and detect enhanced soft X-ray emission in the sightline towards Markarian 509, where Sembach et al. (1995, 1999) find highly-ionized gas associated with Complex Galactic Center Negative (GCN).

Ongoing interaction with an ambient halo gas is indicated by 33 cloudlets in GCN that show head-tail morphology (identified in both GASS and EBHIS, Winkel et al., 2011). In junction with reported highly-ionized gas, Complex GCN may serve as a

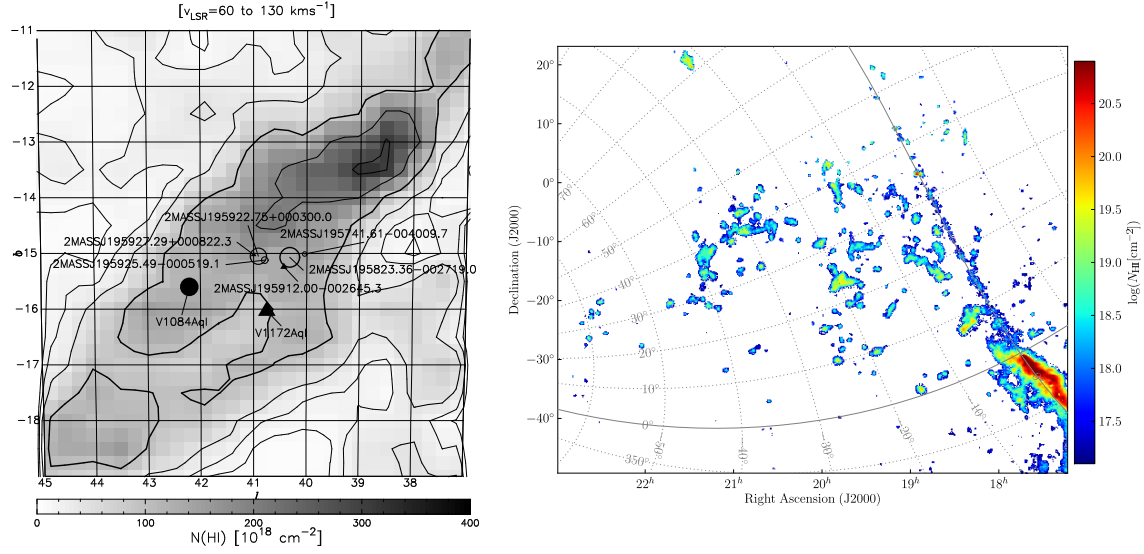


Figure 3.4: *Left plot:* Map of H I column density for Complex Galactic Center Positive (“Smith Cloud”) in the velocity range $+60 < v_{\text{LSR}}/\text{km s}^{-1} < +130$. The contours at $N(\text{H I}) = (1, 2, 4, 10, 20, 30) \times 10^{19} \text{ cm}^{-2}$ are based on LAB data. [figure taken from Wakker et al. (2008)] *Right plot:* H I column density of Complex Galactic Center Negative (*colored map*) based on EBHIS ($\delta > -5^\circ$) and GASS (remaining region) in the velocity range $-360 < v_{\text{LSR}}/\text{km s}^{-1} < -180$. [figure taken from Winkel et al. (2011)]

prime example for warm gas accretion onto the Milky Way (Winkel et al., 2011). The presence of C IV and Si IV provides strong evidence that GCN is not of primordial origin (Wakker, 2001), but instead represents an inflow of gas to the Galactic center (Mirabel, 1982; Mirabel & Morras, 1984). GCN was likely tidally stripped off a dwarf galaxy (Jin, 2010).

Table 3.3: Observed properties of most prominent HVC complexes in the Milky Way halo with given references: range of LSR velocities (*column 2*); HI column densities (*column 3*); intensities of H α emission (*column 4*) in terms of milli-Rayleigh, with 1 Rayleigh = $10^6/4\pi$ photons cm $^{-2}$ sr $^{-1}$ s $^{-1}$ $\approx 1.7 \times 10^{-6}$ erg cm $^{-2}$ sr $^{-1}$ s $^{-1}$ in H α ; metallicities (*column 5*); helio-centric distance (*column 6*); height above the Galactic plane (*column 7*).

Complex	v_{LSR} [km s $^{-1}$]	$N(\text{HI})$ [10^{20} cm $^{-2}$]	$I(\text{H}\alpha)$ [mR]	Z [$Z_{\odot}$]	D_{sun} [kpc]	z_{gal} [kpc]
A	$-210 \dots -140^{\text{[a]}}$	$0.2 \dots 0.4^{\text{[b]}}$	$29 \dots 90^{\text{[b,c]}}$	$0.1^{\text{[d]}}$	$7 \dots 11.5^{\text{[b]}}$	$2.6 \dots 6.8^{\text{[b]}}$
C	$-178 \dots -68^{\text{[e]}}$	$0.12 \dots 560^{\text{[e]}}$	$20 \dots 130^{\text{[c,f,g]}}$	$0.1 \dots 0.3^{\text{[g]}}$	$10 \pm 2.5^{\text{[h]}}$	$3 \dots 9^{\text{[i]}}$
H	$-230 \dots -120^{\text{[j,k,n]}}$	$0.2 \dots 2^{\text{[k]}}$			$27 \pm 9^{\text{[l]}}$	$\lesssim 0.4^{\text{[d]}}$
M	$-125 \dots -85^{\text{[m]}}$	$0.06 \dots 1.50^{\text{[c,n]}}$	$80 \dots 200^{\text{[c]}}$	$0.4 \dots 3.5^{\text{[l,q]}}$	$< 5^{\text{[o]}}$	$< 3.5^{\text{[p]}}$
GCP	$+99 \pm 1^{\text{[r]}}$	$1.1 \dots 1.7^{\text{[s]}}$	$430 \pm 4^{\text{[t]}}$	$0.15 \dots 0.53^{\text{[t,u]}}$	$12.4 \pm 1.3^{\text{[r]}}$	$2.9 \pm 0.3^{\text{[r]}}$
GCN	$-340 \dots -170^{\text{[v]}}$	$0.03 \dots 0.8^{\text{[w]}}$			$14 \dots 22^{\text{[w]}}$	

References to the observed values:

- | | | |
|--|--------------------------------------|--------------------------------------|
| ^[a] van Woerden et al. (1999c) | ^[b] Barger et al. (2012) | ^[c] Tufte et al. (1998) |
| ^[d] Wakker (2001) | ^[e] Hsu et al. (2011) | ^[f] Murphy et al. (2000) |
| ^[g] Tripp et al. (2003) | ^[h] Thom et al. (2008) | ^[i] Wakker et al. (2007) |
| ^[j] Wakker et al. (1998) | ^[k] Simon et al. (2006) | ^[l] Lockman (2003) |
| ^[m] Keenan et al. (1995) | ^[n] Schwarz et al. (1995) | ^[o] Wakker et al. (1996) |
| ^[p] Ryans et al. (1997b) | ^[q] Yao et al. (2011) | ^[r] Lockman et al. (2008) |
| ^[s] Wakker et al. (2008) | ^[t] Hill et al. (2009) | ^[u] Fox et al. (2016) |
| ^[v] Wakker & van Woerden (1991) | ^[w] Winkel et al. (2011) | |

3.2.2 The Magellanic System

The Magellanic System comprises particular high-velocity H I features around the Milky Way that are torn from the Large Magellanic Cloud (LMC) and Small Magellanic Cloud (SMC) by the gravitational tidal forces of the Milky Way (Mastropietro et al., 2005; Connors et al., 2006). The main such tidal components are the Magellanic Bridge (the part connecting the clouds), and two narrow filaments of H I called the (trailing) Magellanic Stream (MS) and the Leading Arm (LA). The Magellanic System is the only HVC complex to which stars are firmly assigned observationally (e.g., Reid & Parker, 2012; Price-Whelan et al., 2019; Sakowska et al., 2023). Hence, its distance is confirmed confidentially at $D_{\text{sun}} \approx 55 \text{ kpc}$ (D’Onghia & Fox, 2016). Column densities in H I range between $4.7 \times 10^{16} \dots 1.1 \times 10^{20} \text{ cm}^{-2}$ (Barger et al., 2017).

The MS is a very elongated high-velocity feature of H I emission in the southern sky and constitutes the largest and most massive Galactic HVC (Putman et al., 2003b) providing an inflow rate of $\dot{M}(\text{H I}) = 3.7 \dots 6.7 \text{ M}_{\odot} \text{ yr}^{-1}$ (Fox et al., 2014) that is larger than the inflow rate represented by all other HVCs (cf. table 1.1). Much of the MS gas is trailing the two Magellanic Clouds (see figure 3.5, left). The LSR velocity of H I ($\text{H}\alpha$) smoothly varies along the MS from $+200 \text{ km s}^{-1}$ ($+158 \text{ km s}^{-1}$) close behind the Magellanic Clouds to -361 km s^{-1} (-354 km s^{-1}) near the tail end (Barger et al., 2017). The MS is characterized by very wide emission line profiles $\Delta v_{\text{FWHM}} = 18 \dots 64 \text{ km s}^{-1}$ (indicative of a warm component, Barger et al., 2017) and mass portions $M(\text{H I}) = 2.7 \times 10^8 \text{ M}_{\odot}$ and $M(\text{H II}) \approx 10^9 \text{ M}_{\odot}$ (D’Onghia & Fox, 2016). Fox et al. (2013) find abundances $[\text{S}/\text{H I}] = -1.13$ and $[\text{O}/\text{H I}] = -1.24$ along the MS such that a uniform metallicity of $Z = 0.1 Z_{\odot}$ along the main body of MS is assumed. Barger et al. (2017) confirm these findings by observing respective $Z = 0.1 Z_{\odot}$ and $Z = 0.5 Z_{\odot}$ in each filament.

The LA spreads over 65 deg (corresponding to 52.5 kpc at assumed $D_{\text{sun}} = 50 \text{ kpc}$, Venzmer et al., 2012) and hosts many head-tail CHVCs, e.g., HVC 287.5+22.5+240 (Sembach et al., 2001, 2002), HVC 291+26+195 and HVC 297+09+253 (Ben Bekhti et al., 2006), HVC 306-2+230 (McClure-Griffiths et al., 2008). The LA is separated into four different complexes and many clumpy H I filaments (Ben Bekhti et al., 2006; Venzmer et al., 2012, see figure 3.5, right). The LSR velocities measured in the LA so far are invariably positive with $v_{\text{LSR}} = +75 \dots +300 \text{ km s}^{-1}$ (Fox et al., 2018). Within LA II Richter et al. (2018) detect molecular hydrogen with a high $N(\text{H}_2) = 1.66 \times 10^{18} \text{ cm}^{-2}$ and infer $n(\text{H}) = 16 \text{ cm}^{-3}$, $T_{\text{kin}} = 105 \text{ K}$, and $P/k = 1680 \text{ K cm}^{-3}$. They also conclude on dust due to strong depletion of Si, Fe, Al, and Ca. Observed abundance gradients provide strong evidence for a successive formation of the LA by gas being episodically

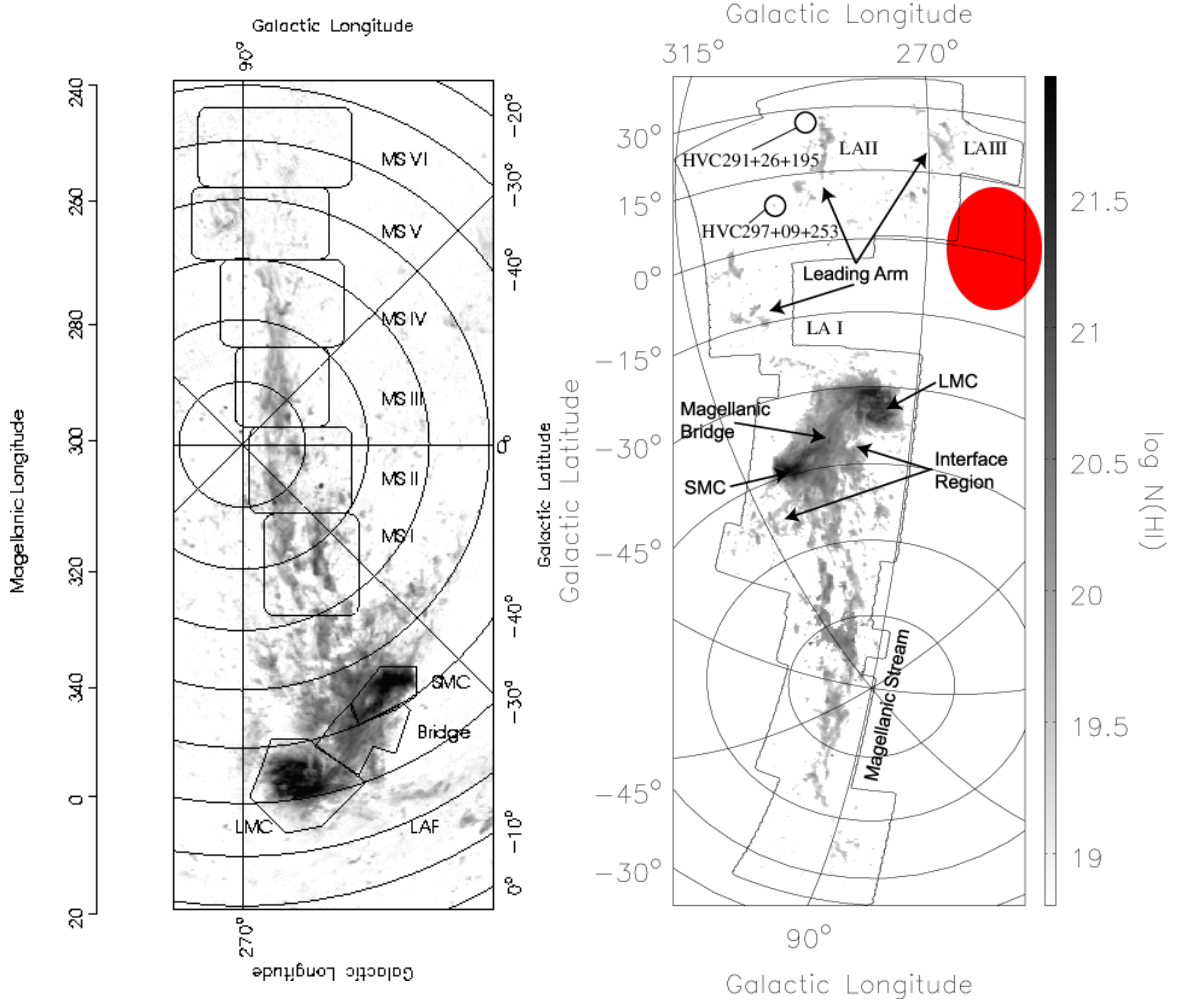


Figure 3.5: The Magellanic System with its main components as seen in direction towards the south Galactic pole ($b = 0$ deg). *Left plot:* LMC, SMC, Magellanic Bridge between both clouds, trailing MS, and the beginning of the LA feature (LAF) in the LSR velocity range $-450 \dots 400 \text{ km s}^{-1}$ as observed in HIPASS. The *greyscale map* visualizes $N(\text{HI})$ ranging from $2 \times 10^{18} \text{ cm}^{-2}$ to $> 6 \times 10^{20} \text{ cm}^{-2}$. The gas encircled by the *black boxes* is used for mass determination. [figure taken from Putman et al. (2003b)] *Right plot:* The most prominent HI substructures LA I, LA II, LA III, and two HVCs in the LA are shown in $N(\text{HI})$ between $6 \times 10^{18} \text{ cm}^{-2}$ and $9 \times 10^{21} \text{ cm}^{-2}$. [figure taken from Ben Bekhti et al. (2006)] The *red ellipse* (own representation) shows the region where Venzmer et al. (2012) localize the component LA IV.

removed from the SMC (Richter et al., 2018). For et al. (2013) suggest that ram pressure dominates tidal forces in forming the LA, based on 419 mainly head-tail HVCs observed in the region of the LA within HIPASS. However, Casetti-Dinescu et al. (2014) intercede for a tidal origin. They also find young stars in the LA, which give evidence for recent star formation triggered by an interaction between the Magellanic Clouds and the Milky Way.

3.2.3 Compact HVCs

Braun & Burton (1999) and Burton & Braun (1999) are the first to identify 66 HI clouds with angular sizes < 2 deg. Even though their LSR velocities are in the range of HVC complexes, their peak column densities are significantly higher and they are much more spatially confined. Hence, these clouds are named compact HVCs (CHVCs). Some of the reported CHVCs are surrounded by regions of absolutely no 21-cm emission down to the detection limit of $\sim 1.5 \times 10^{18} \text{ cm}^{-2}$ (chapter 15 in van Woerden et al., 2004, and see figure 3.6). They are thus called *isolated CHVCs* (Burton & Braun, 1999; Burton et al., 2005). Unlike HVC complexes, CHVCs distribute rather uniformly over the Galactic sky and neither tend to accumulate in the lower halo nor do they cluster along filamentary structures like the major HVC complexes (see fig. 1 in Burton & Braun, 1999). This lack of association constitutes CHVCs to be a subset of HVCs that may have originated in the same environmental circumstances, share a similar evolutionary history, and are in the same physical state (e.g., Williams et al., 2012).

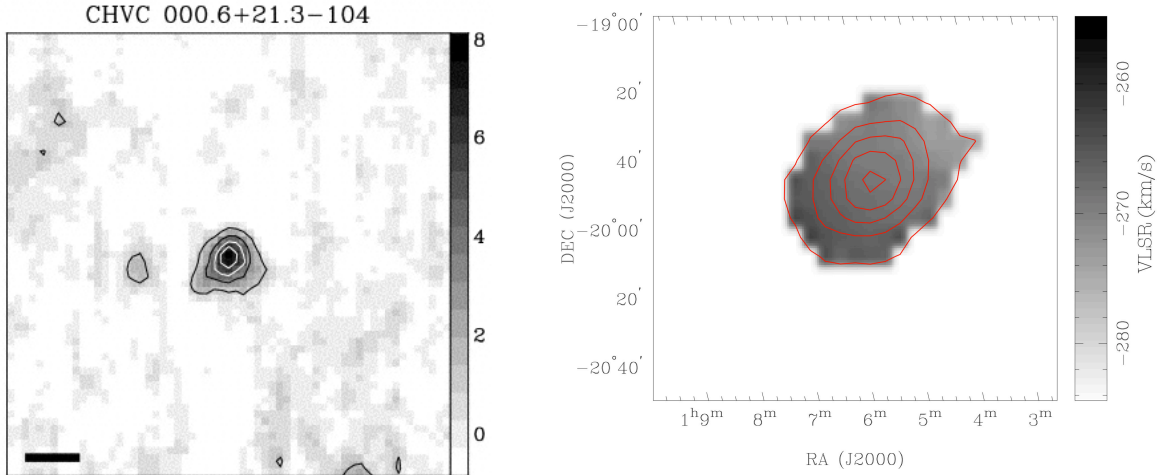


Figure 3.6: Examples of isolated CHVCs. *Left plot:* A CHVC taken from HIPASS. *Isocontours* show 30, 50, 70, and 90 % of peak $N(\text{HI})$ and the *black bar* in the lower left represents an angular extension of 2 deg. [figure taken from Putman et al. (2002)] *Right plot:* A CHVC extracted from GASS. The *greyscale map* shows LSR velocities while *red isocontours* mark column densities in HI. [figure taken from Moss et al. (2013)]

CHVCs show various morphologies like (multiple) core-halo structures (figure 3.7 and Burton et al., 2001), bow-shock and irregular shapes (figure 3.8), nearly spherical symmetry (CHVC 147.5-82.3-269, Westmeier et al., 2004, 2005b,c), and head-tail shapes.

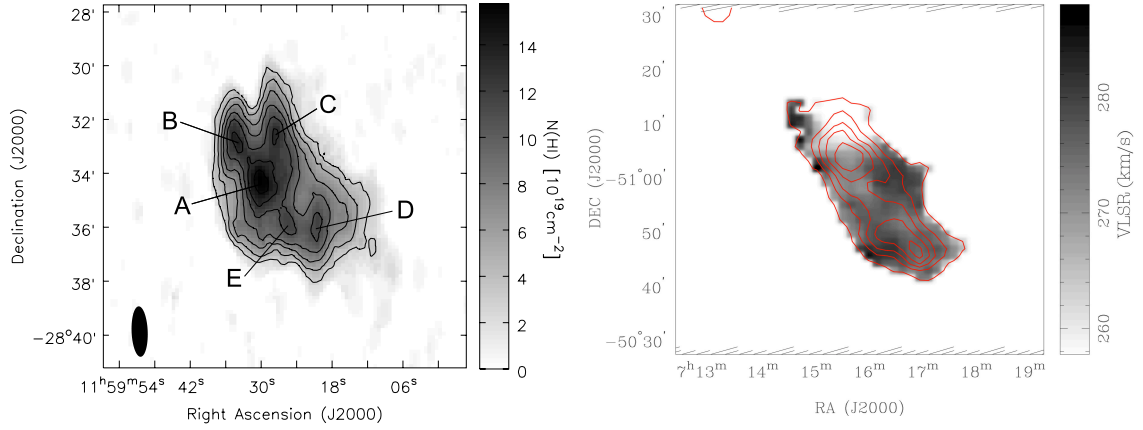


Figure 3.7: Examples of spatial multiple-component CHVCs. *Left plot:* Column density in H I (*black isocontours*) of CHVC 289+33+251 observed with ATCA telescope. The letters indicate 5 identified concentrations the cloud is resolved into. With a size of 4.4 arcmin this cloud is the most compact HVC that has ever found. [figure taken from Brüns & Westmeier (2004)] *Right plot:* LSR velocity (*greyscale map*) and column density in H I (*red isocontours*) with two distinct concentrations of a CHVC from GASS. [figure taken from Moss et al. (2013)]

A prime example for the latter is given by CHVC 125+41-207 showing $N(\text{H I}) = 5 \times 10^{20} \text{ cm}^{-2}$ (Brüns et al., 2001). Its characteristic head-tail morphology indicates an ongoing interaction with the surrounding gas (figure 3.9, left). Its distance of $\sim 130 \text{ kpc}$ thus confirms the existence of a very extended gaseous halo around the Milky Way (Brüns et al., 2001). At least 35 % of CHVCs exhibit head-tail morphologies (Putman et al., 2011) suggesting interaction with a diffuse hot Galactic halo medium (Williams et al., 2006, 2007; Bregman et al., 2009) and most of these are located in the proximity of HVC complexes. Head-tail CHVCs show distinct gradients in density (Putman et al., 2011), velocity (Brüns et al., 2000; Faridani et al., 2014), and in metallicity. The latter is due to a different mixing efficiency between tail gas and ambient hot gas than ongoing at the head(cf. figure 10 in Sander & Hensler, 2021). According to Brüns & Mebold (2004) a decreasing density, an increasing velocity dispersion and decreasing radial velocity along the tail observationally indicate that gas is being stripped from the HVC.

The discovery of the most compact known cloud CHVC 289+33+251 (Putman et al., 2002; Brüns & Westmeier, 2004) has led to the definition of ultra-compact HVCs (UCHVCs) with angular extensions below 20 arcmin, which is noticeably smaller than the typical sizes of CHVCs. CHVC 289+33+251 has an angular size of 4.4 arcmin (figure 3.7, left) and is resolved into five H I cores. It shows an overall velocity dispersion of 4.9 km s^{-1} , hence $T_{\text{kin}} < 532 \text{ K}$. Accordingly, $M(\text{H I}) = 5.66 \times 10^4 \left(\frac{D_{\text{sun}}[\text{kpc}]}{150 \text{ kpc}} \right)^2 M_{\odot}$ and the detected faint tail indicates an ongoing ram-pressure interaction with an ambient

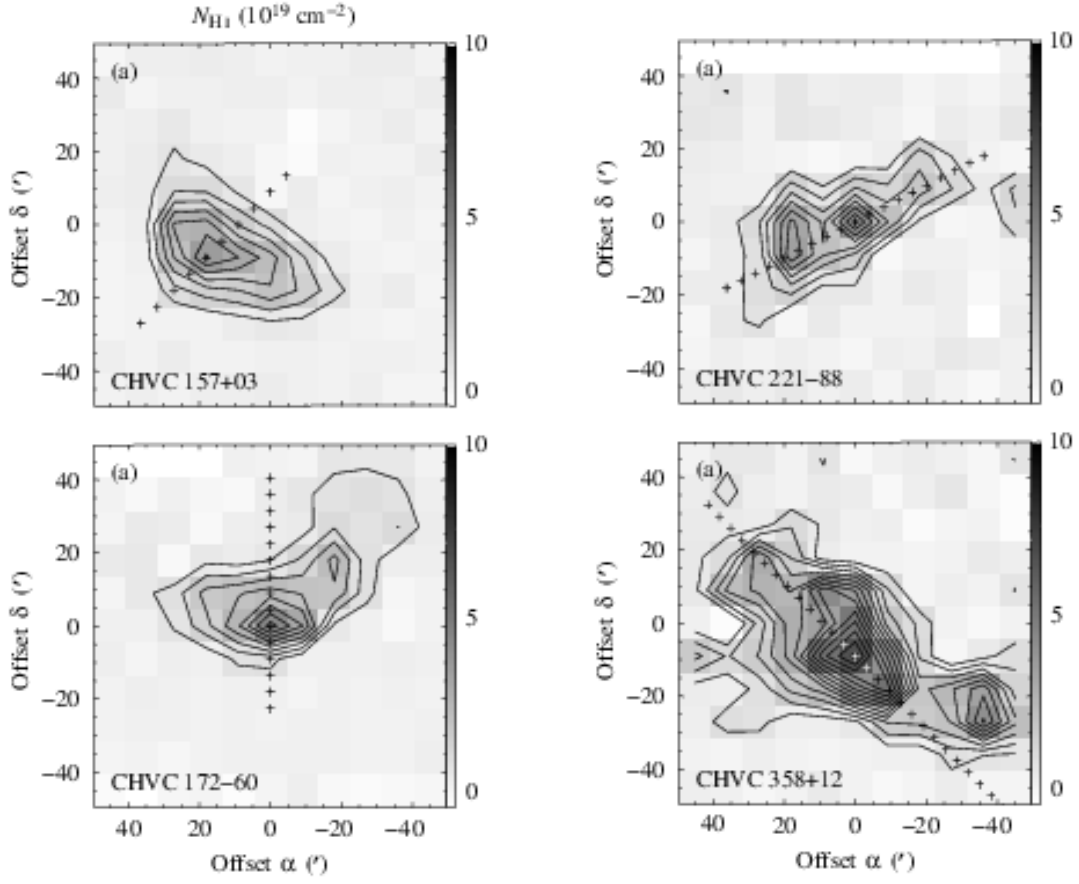


Figure 3.8: Column density in H I (*greyscale map*) of two bow-shock shaped (*left panel*) and two irregularly shaped CHVCs (*right panel*) as seen with Effelsberg telescope. The *black isocontours* start at $N(\text{H I}) = 5 \times 10^{18} \text{ cm}^{-2}$ and increase in steps of $N(\text{H I}) = 5 \times 10^{18} \text{ cm}^{-2}$. *Black crosses* mark positions of individually recorded spectra. [figures taken from Westmeier et al. (2005c)]

low-density medium. Williams et al. (2012) report $N(\text{H I}) > 10^{20} \text{ cm}^{-2}$.

The most complete list of CHVCs in the northern hemisphere is given by the survey GALFA-H I (Saul et al., 2012). They identify 692 CHVCs with inferred masses of $\sim 200 M_{\odot}$ and physical sizes of $\sim 10 \text{ pc}$ at $D_{\text{sun}} \sim 10 \text{ kpc}$ and masses of $\sim 2 \times 10^4 M_{\odot}$ and physical sizes of $\sim 100 \text{ pc}$ at $D_{\text{sun}} \sim 100 \text{ kpc}$. In the southern sky the surveys HIPASS (179 identified CHVCs, Putman et al., 2002), and GASS (1,417 CHVCs with angular extent $\leq 2 \text{ deg}$, Moss et al., 2013) provide the hitherto most complete sample of CHVCs. By inspecting figure 3.10 (upper left plot) one verifies the relatively high number of positive-velocity CHVCs as seen in GASS and HIPASS for the southern hemisphere. The fastest clouds are found in GASS with $|v_{\text{LSR}}| > 400 \text{ km s}^{-1}$. The broadest median linewidths are seen in HIPASS, but with lowest and topmost outliers in GALFA-H I. Vice versa, the most smallest median of linewidths are reported in GASS (figure 3.10, upper right) due to a much higher velocity resolution ($\Delta v = 1 \text{ km s}^{-1}$)

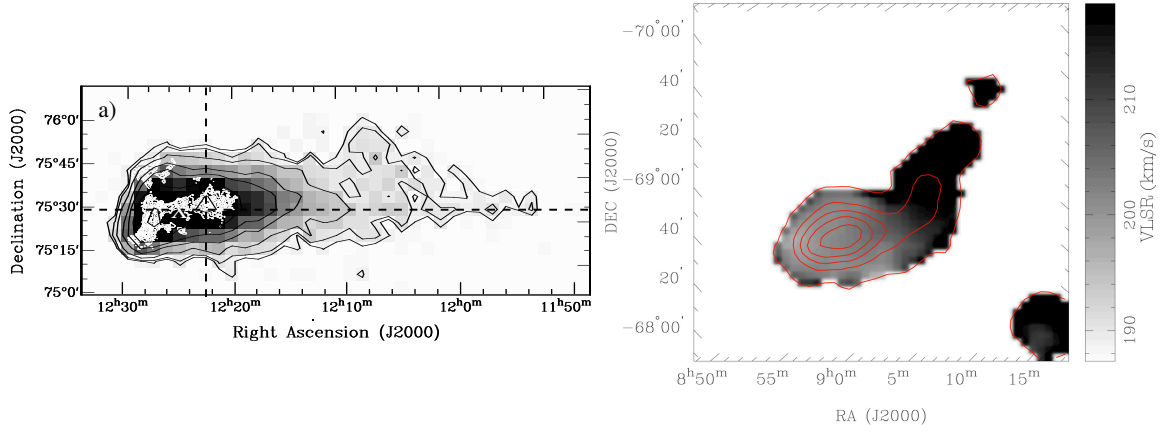


Figure 3.9: Examples of head-tail CHVCs. *Left plot:* Column density in H I (*black isocontours*) of CHVC 125+41-207 observed with Effelsberg telescope. The *white contours* denote 21-cm emission observed with Westerbork telescope. [figure taken from Brüms et al. (2001)] *Right plot:* LSR velocity (*greyscale map*) and column density in H I (*red isocontours*) of a CHVC from GASS. [figure taken from Moss et al. (2013)]

compared to HIPASS ($\Delta v = 26 \text{ km s}^{-1}$). The distributions of Δv_{FWHM} correspond to median Doppler temperatures of $\sim 3 \dots 7 \times 10^4 \text{ K}$ in the CHVCs. The medians of peak column densities are highest in HIPASS while most upper outliers are found in GASS. It is noticeable that the medians of peak column densities are similarly at $1 \dots 2 \times 10^{19} \text{ cm}^{-2}$ in all three surveys (figure 3.10, lower plot). To illustrate typical ranges of observed features some particular CHVCs are tabulated in table 3.4.

Explicit searches for stellar counterparts in CHVCs have not succeeded so far. Simon & Blitz (2002) are not able to provide evidence for stars in CHVCs down to a surface brightness of $1 \text{ mag arcsec}^{-2}$ fainter than normal Local Group dwarf galaxies by using digitized Palomar Observatory Sky Survey (POSS) data. Hopp et al. (2003a,b, 2007) search for an intrinsic stellar population in CHVCs using FORS at the VLT combined with near-IR photometry from the Two Micron All Sky Survey (2MASS). They do not detect any stellar signature thus concluding that CHVCs are implausibly the Local Group dwarf galaxies, which are plentiful in ΛCDM simulations (but missed so far in observations).

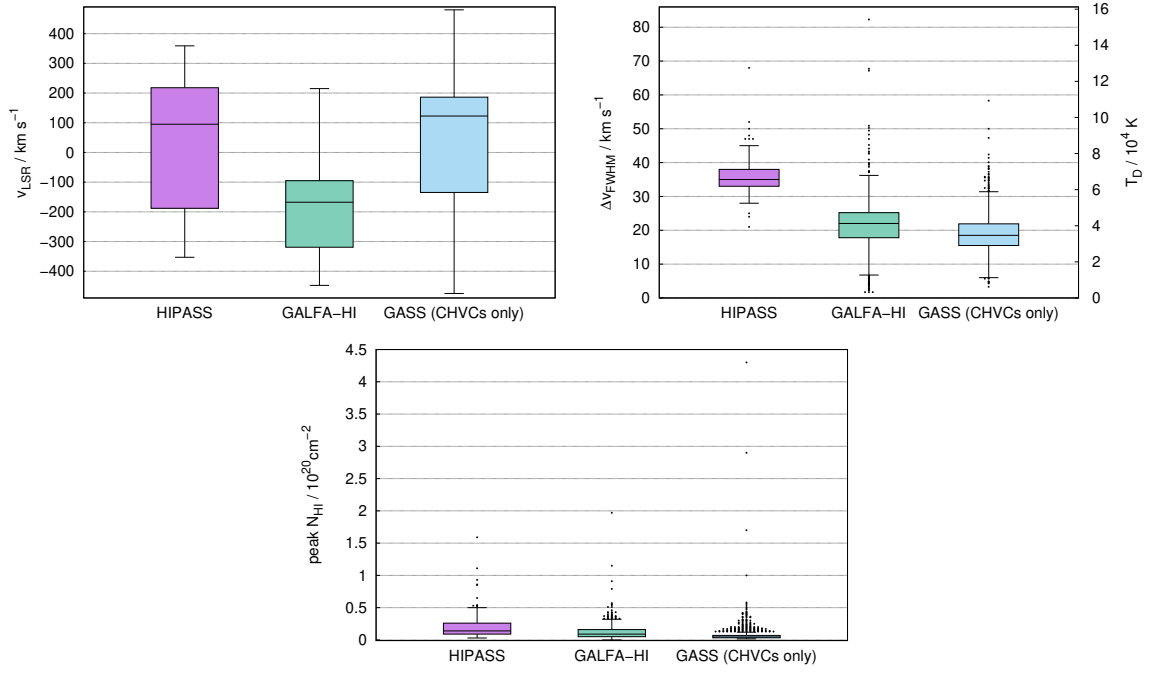


Figure 3.10: Boxplots for HIPASS (purple boxes, Putman et al., 2002), GALFA-H I Survey (green boxes, Saul et al., 2012), and GASS (blue boxes, Moss et al., 2013) showing the distributions of LSR velocities (*upper left*), velocity dispersions corresponding to certain Doppler temperatures (*upper right*), and peak column densities (*lower plot*). The middle of each boxplot indicates the median value. The upper and lower edges of each boxplot indicate the 75th and 25th percentiles, respectively. The ends of the vertical lines indicate upper and lower whiskers with the *black dots* showing outliers in the data sets.

Table 3.4: Some observed properties of Galactic CHVCs to illustrate typical ranges.

CHVC	v_{LSR} [km s ⁻¹]	$N(\text{H I})$ [cm ⁻²]	Δv_{FWHM} [km s ⁻¹]	T_{kin} [K]	D_{sun} [kpc]	$M(\text{H I})$ [M _⊙]
CHVC 006-03-474	-475 ^[a]					
CHVC 185-47+480	+480 ^[a]					
CHVC 125+41-207		10 ²¹ ^[b]	2...5 ^[c]	85...540 ^[d]	130 ^[d]	$2.4 \times 10^6 \left(\frac{D_{\text{sun}}[\text{kpc}]}{200 \text{ kpc}} \right)^2$ ^[d]
CHVC 262+23+349		> 10 ²⁰ ^[a]				
CHVC 239+30-103		≲ 10 ¹⁸ ^[e]				
CHVC 087+17-163			~ 80 ^[e]	≳ 10 ⁵ ^[e]		
CHVC 266-19+336					1.2...6.1 ^[f]	
CHVC 119-31-384					1.3...13.2 ^[f]	
CHVC 291+26+195					13...35 ^[g]	$5.6 \dots 5.9 \times 10^3 \left(\frac{D_{\text{sun}}[\text{kpc}]}{50 \text{ kpc}} \right)^2$ ^[g]
CHVC 297+09+253					19...37 ^[g]	$1.7 \dots 4.7 \times 10^4 \left(\frac{D_{\text{sun}}[\text{kpc}]}{50 \text{ kpc}} \right)^2$ ^[g]
CHVC 230+61+165					150 ^[h]	$5.9 \times 10^6 \left(\frac{D_{\text{sun}}[\text{Mpc}]}{1 \text{ Mpc}} \right)^2$ ^[h]

References to the observed values:

- ^[a]Moss et al. (2013) ^[b]Braun (2001) ^[c]Braun & Burton (2000) ^[d]Brüns et al. (2001)
^[e]Saul et al. (2012) ^[f]Putman et al. (2003a) ^[g]Ben Bekhti et al. (2006) ^[h]Burton et al. (2001)

3.2.4 Ultra-compact HVCs

The term “ultra-compact high-velocity cloud (UCHVC)” may be misleading as it is used ambiguously for two equitable scenarios in order to explain the stability of UCHVCs. Regarding 21-cm observations with WSRT, Adams et al. (2016) argue that these two categories of objects represent two completely different populations that are distinct in both H I morphology and kinematics.

1. Galactic halo clouds at distances of ~ 100 kpc have $M(\text{H I}) \sim 10^5 M_\odot$ and are stabilized by external pressure so only H I is needed, and this is supported by morphologies observed at UCHVCs (head-tail, bow-shock, irregular) that are most probably be explained by hydrodynamic interaction (see section 3.3.4) with an ambient gas at higher densities than in the Local Group (Westmeier et al., 2005a,c). This set of clouds have disordered velocity fields at all spatial scales and show up with core-halo morphologies (Adams et al., 2016). In this regard they resemble CHVCs, but with angular diameters $\Delta v_{\text{FWHM}} \lesssim 30$ arcmin (e.g. CHVC 289+33+251, Brüns & Westmeier, 2004). As is described in the previous section 3.2.3 they reveal observational signatures of mainly H I clouds without dark matter and being placed in the Galactic halo. They have typical masses $\lesssim 5 \times 10^5 M_\odot$ and distances of a few up to ~ 200 kpc that are inferred from surveys like HIPASS (Putman et al., 2002), GALFA-H I (Saul et al., 2012), and GASS (Moss et al., 2013). de Heij et al. (2002a) obtain $D_{\text{sun}} = 150 \dots 200$ kpc for both Milky Way and M 31 based upon combined HIPASS data (Putman et al., 2002) and a catalogue of 67 (U)CHVCs (de Heij et al., 2002b). By using the model of Bland-Hawthorn et al. (1998), Putman et al. (2003a) estimate $D_{\text{sun}} < 40$ kpc from deep H α spectroscopy (these results are supported by the findings of Tufte et al., 2002). Likewise, Westmeier et al. (2005b) infer $D_{\text{sun}} < 60$ kpc based on 11 observed (U)CHVCs implying sizes below 0.5 kpc and $M(\text{H I}) < 6 \times 10^4 M_\odot$. Regarding Maloney & Putman (2003), (U)CHVCs are more likely associated with the Galactic halo at $D_{\text{gal}} = 100 \dots 200$ kpc based on self-consistent models of gas within DM halos. Searches for clouds similar to the population of Galactic (U)CHVCs in nearby Local Group analogs indicate that (U)CHVCs are probably clustered within $D_{\text{gal}} = 90$ kpc and massive intragroup H I clouds are relatively rare (Braun & Burton, 2001; Zwaan, 2001b; Pisano et al., 2004, 2007; Chynoweth et al., 2009). Deep 21-cm observations of the M 31 halo reveal a deficiency of large H I clouds beyond galacto-centric distances of 50 kpc (Westmeier et al., 2007, 2008) and hence disprove the scenario by Blitz et al. (1999b) and Braun & Burton (1999) in which (U)CHVCs are regarded as gas-filled DM halos residing in the Local Group at $D_{\text{sun}} > 300$ kpc (Maloney & Putman, 2003;

Putman et al., 2003a; Richter, 2016). Plöckinger & Hensler (2012) obtain decelerated and multi-phase CHVCs by means of hydrodynamic simulations only if DM is not considered.

2. Potential galaxy candidates as being identified in the ALFALFA Survey (Giovanelli et al., 2005, 2010) revive the idea of DM halos placed in the Local Group or even Local Volume (Blitz et al., 1999a,b; Braun & Burton, 1999; Burton & Braun, 1999). There is observational evidence that these UCHVCs are the baryonic counterparts in shape of ultra-faint dwarf galaxies residing in isolated, low mass halos (so-called “mini-halos”) distributed across the Local Group. Giovanelli et al. (2010) derive characteristic values of $M(\text{H I}) = 5 \times 10^4 D_{\text{sun}}^2 [\text{Mpc}^2] \dots 10^6 D_{\text{sun}}^2 [\text{Mpc}^2] M_{\odot}$ and total masses $\sim 10^8 \dots 10^9 M_{\odot}$ from ALFALFA Survey. UCHVCs have well-ordered velocities down to the scale of resolution and higher column densities than the halo clouds (see figure 4 in Adams et al., 2016). Examples are Leo P (Giovanelli et al., 2013; Skillman et al., 2013), Leo T (tentatively, Irwin et al., 2007; Ryan-Weber et al., 2008; Faerman et al., 2013), and an isolated dwarf galaxy that has been recently discovered in the Arecibo Galaxy Environment Survey (AGES, Taylor et al., 2014). Brüns & Westmeier (2004) argue that at a distance of 150 kpc the virial mass of CHVC 289+33+251 is by a factor of 5.6 higher than the observed gas mass, which is consistent with the “missing” DM mini-halos that are predicted by cosmological Λ CDM simulations (e.g., Klypin et al., 1999; Moore et al., 1999; Knebe et al., 2008). At distances of ~ 1 Mpc they have $M(\text{H I}) \sim 10^8 M_{\odot}$ and $M_{\text{dyn}} \sim 10^9 M_{\odot}$ so dark matter is needed to account for their stability (Oort, 1966; Blitz et al., 1999b). A somewhat hybrid object is CHVC 274.68+74.70-123, which Bellazzini et al. (2015b) refer to as a nearly starless, faint dwarf galaxy at $D_{\text{gal}} > 3$ Mpc likely making it a missing link between normal dwarf galaxies and DM mini-halos. Adams et al. (2016) present follow-up H I observations with WSRT of 12 UCHVCs identified in the ALFALFA Survey. Three of them are excellent candidates for gas-bearing DM halos. They find $M(\text{H I}) = 0.6 \dots 48 \times 10^5 M_{\odot}$ and half-mass radius $R(\text{H I}) = 0.3 \dots 2.2$ kpc for a range $D_{\text{sun}} = 0.4 \dots 2$ Mpc.

In a subsequent analysis of the ALFALFA Survey Giovanelli et al. (2013) discover a previously unknown 21-cm H I line source identified as UCHVC and interpreted as a very faint, nearby, star forming dwarf galaxy called Leo P. By using the baryonic Tully-Fisher relation they infer $M(\text{H I}) = 3.1 \times 10^5 D_{\text{sun}}^2 [\text{Mpc}^2] M_{\odot}$, $R(\text{H I}) = 0.29 D_{\text{sun}} [\text{Mpc}] \text{ kpc}$, and an estimated $D_{\text{sun}} = 1.3_{-0.6}^{+0.9} \text{ Mpc}$. Skillman et al. (2013) figure out that Leo P is an extremely metal-deficient galaxy with $[\text{O}/\text{H I}] = -4.83$, $[\text{N}/\text{H I}] = -6.18$, $[\text{S}/\text{H I}] = -6.32$. Sand et al. (2015) report on CHVC 274.68+74.70-123 with $v_{\text{LSR}} = -123 \text{ km s}^{-1}$, a half-mass radius $R(\text{H I}) = 8.1 \text{ arcsec}$ (corresponding

to $117 \dots 707$ pc at $D_{\text{gal}} = 3 \dots 18$ Mpc), $M(\text{HI}) = 0.2 \dots 7 \times 10^7 M_{\odot}$. A faint, star-forming stellar system SECCO 1 is firmly assigned to CHVC 274.68+74.70-123 and a stellar mass of $\sim 1.6 \times 10^5 M_{\odot}$ is derived (Bellazzini et al., 2015b; Beccari et al., 2017). It represents the only stellar counterpart that could have been assigned to a UCHVC within the StEllar Counterparts of COmpact high-velocity clouds Survey (SECCO, Bellazzini et al., 2015a; Beccari et al., 2016).

Apart from these two examples the results of SECCO fully coincide with the various unsuccessful attempts performed in the last two decades in observing stellar counterparts of HVCs (e.g., Simon & Blitz, 2002; Willman et al., 2002a; Hopp et al., 2003b, 2007; Siegel et al., 2005).

3.2.5 Distinction to IVCs

The Magellanic Stream consists of gas tidally expelled from the SMC hence tracing the dark-matter potential of the Milky Way. While HVCs typically bear low metallicities of $\sim 0.1 \dots 0.3 Z_{\odot}$ (van Woerden et al., 2004) and appear to trace the sustained accretion of infalling intergalactic gas, IVCs with near-solar metallicity at $z_{\text{gal}} \sim 1$ kpc trace the circulation of material at disk-halo interface (Wakker et al., 2008). The latter circuit is known as “Galactic fountain” (Shapiro & Field, 1976).

IVCs deviate more mildly from differential Galactic rotation than HVCs do. Wakker (2001) uses a range in LSR velocities of $40 \dots 90 \text{ km s}^{-1}$ to kinematically demarcate IVCs from both HVCs and low-velocity clouds (LVCs, $|v_{\text{LSR}}| \lesssim 40 \text{ km s}^{-1}$) while others use a lower limit of $|v_{\text{LSR}}| = 20 \text{ km s}^{-1}$ (e.g., Albert & Danly, 2004) or an upper limit of $|v_{\text{LSR}}| = 100 \text{ km s}^{-1}$ (e.g., Richter, 2017). Together with their smaller distances to the Galactic plane ($z_{\text{gal}} = 0.5 \dots 1.5$ kpc, $D_{\text{sun}} < 5$ kpc, Wakker, 2001; Smoker et al., 2004; Lehner et al., 2022), nearly solar metallicities (featured by strong UV absorption in several metal lines, Wakker et al., 2001; Savage et al., 2007; Wakker et al., 2008), detectable amounts of dust (small, hot grains in compliance with shattering by shocks, Planck Collaboration et al., 2011) and, owing to dust grain destruction, reasonable amounts of H_2 ($N(\text{H}_2) = 10^{14} \dots 10^{17} \text{ cm}^{-2}$, Richter et al., 2003; Wakker, 2006), deuterium (Savage et al., 2007), and Ca (Keenan et al., 1995; Ryans et al., 1997b) indicate a close connection to Galactic fountains. An exception that proves the rule of Galactic origin is reported by Fukui et al. (2021b) for the Pegasus-Pisces Arch (a.k.a. IVC 86-36).

Sub-solar metallicities are exceptional, but not excluded for IVCs (Wakker, 2001; Hernandez et al., 2013). Hayakawa & Fukui (2022) find a range of $Z = 0.1 \dots 1.5 Z_{\odot}$

with 36 % of their analyzed IVCs showing $Z < 0.5 Z_{\odot}$. For a substantial fraction of these low-metallicity IVCs they propose accretion by the Milky Way instead of an origin in a Galactic fountain.

Putman et al. (2012) report on masses of $M(\text{H I}) = 0.5 \dots 8 \times 10^5 M_{\odot}$ for IVCs, which is on average less than masses of HVC complexes ($M(\text{H I}) > 6 \times 10^6 M_{\odot}$), but comparable to masses of CHVCs. High H I column densities of several 10^{20} cm^{-2} are found, which are sufficient to mitigate dissociating radiation such that CO (Röhser et al., 2016) or OH (Smith & Anish Rosh, 2023) is able to form. Prominent IVCs comprise IVC-South and Cloud g1 (Wakker, 2001; Barentine et al., 2008; Wakker et al., 2008), Pegasus-Pisces Arch a.k.a. IVC 86-36 (Wakker, 2001; Fukui et al., 2021b; Shelton et al., 2022), Intermediate-Velocity Arch (Kuntz & Danly, 1996; Ryans et al., 1997b; Richter et al., 2001c; Wakker, 2001), Low-Latitude Intermediate-Velocity Arch (Kuntz & Danly, 1996; Ryans et al., 1997a; Richter et al., 2001a; Wakker, 2001), and IVC Complex K (Kerr & Knapp, 1972; Haffner et al., 2001; Wakker, 2001).

3.3 COMPONENTS AND PHYSICAL PROCESSES

The following section comprises the observational signatures of the multi-phase nature of HVCs, the rare detections of magnetic fields, and illustrates the debate about enveloping DM halos. The most important physical processes, which are deduced from observations, are introduced: hydrodynamic interactions and their effect on the morphological evolution of HVCs, dissociation, ionization, heating and cooling of plasma, and saturated thermal conduction.

3.3.1 Multi-phase structure

High-velocity gas reveals a hierarchical core-halo structure with small, dense, cool, and bright condensations (cores) residing in larger regions of lower emissivity (Kim et al., 1989), which are warm and diffuse (halos). The cores of HVCs represent the cold neutral medium (CNM) while their envelopes constitute the warm neutral to slightly ionized medium (WNM) of a multi-phase Galactic halo. Their hot interfaces provide the transition to the highly ionized medium (HIM). Typical values for temperatures and densities in all the phases are given in section 4.1. This multi-phase structure of HVCs is observationally deduced from associated emission by H I at 21-cm and absorption by molecules and dust and successive infrared emission. Because molecules, dust, and neutral elements unconditionally need shielding to the dissociating and ionizing

extragalactic radiation they solely exist in dense cores of more than 10^{20} cm^{-2} (Burton et al., 2005) to 10^{21} cm^{-2} (Krumholz & McKee, 2008).

The multi-phase nature of HVCs is ubiquitous. Lehner et al. (2009b) report on neutral (O I, Fe I), weakly ionized (Fe II, N II), and highly ionized (O VI) components, some dust, but no molecules (similar findings are reported for Complex C, Murphy et al., 2000). Other observations indeed show molecular hydrogen (Wakker et al., 1997; Richter et al., 1999, 2001b; Bluhm et al., 2001; Wakker, 2006; Dessauges-Zavadsky et al., 2007; Crighton et al., 2013; Tchernyshyov, 2022), CO (Kim et al., 2005, 2007), or HCO^+ (Akeson & Blitz, 1999; Oka et al., 1999).

The amount of dust is commonly proposed from the abundance ratio $[\text{Ca II}/\text{H I}]$ (Robertson et al., 1991) or $[\text{Si II}/\text{H I}]$ (Sahu, 1999) acquired by optical absorption spectroscopy, because both elements easily deplete onto dust¹. Another method is to directly detect IR emitted by dust (Wakker & Boulanger, 1986; Miville-Deschênes et al., 2005; Miville-Deschenes, 2006). Bryukhanov & Trifalencov (1994) find a relative dust content in HVCs being lower by a factor of $2 \dots 4$ compared to the Galactic disk.

Neutral to weakly ionized elements are found in the shielding region around cores. These elements are observed by absorption spectroscopy in the optical (e.g., Robertson et al., 1991; Lehner et al., 1999; Ben Bekhti et al., 2008, 2009a,b, 2011) or UV (e.g., Bowen et al., 1995; Richter et al., 2001c, 2009) or by emission of the $\text{H}\alpha$ line. If the $\text{H}\alpha$ emission results from photoionization by the Galactic UV background, it is feasible to indirectly determine distances to HVCs (Weiner et al., 2001, 2002; Bland-Hawthorn & Maloney, 2002, and see appendix C). If, on the other hand, the ionization is due to a shock induced by the collision of a HVC with Galactic halo gas, the $\text{H}\alpha$ intensity yields information on the density and temperature of the halo gas (Tufte et al., 1998). Outstanding in studying the neutral components of HVCs is the 21-cm line emitted by H I. It is thus analyzed in numerous single observations (e.g., Giovanelli et al., 1973; Schwarz & Oort, 1981; Brüns et al., 2000; Blagrove et al., 2010) and in extensive surveys (see section 3.1).

One of the most prominent observational signatures of highly ionized regions in HVCs is given by the absorption line O VI, which is seen in the UV (Murphy et al., 2000; Fox et al., 2003). The O VI is produced in transition-temperature plasma (several 10^5 K) within the interface layers between the warm HVC gas-phase ($T < 5 \times 10^4 \text{ K}$) and the Galactic hot halo ($T > 10^6 \text{ K}$, Sutherland & Dopita, 1993; Narayanan et al., 2010). It is suspected, that O VI in HVCs is not produced in strong shocks, but instead forms in

¹Consider Winkel et al. (2011) who argument that Na I is a tracer for the cooler dense cores of clouds, while Ca II is usually found in the warmer envelopes of clumps.

interface layers between cool and hot gas, if the HVCs move with velocities comparable to the adiabatic sound speed of the hot, ambient medium (Sembach et al., 2000; Thom et al., 2011). In addition to O VI, the multiple ionization stages of silicon and carbon (Collins et al., 2004), sulphur (Savage et al., 2005), and several other elements (Lehner et al., 2012) are used to study the highly ionized interfaces of HVCs.

3.3.2 Magnetic fields

A systematic study is lacking to support a consistent theory on the evolution of magnetic fields in HVCs. Observations of magnetic fields in HVCs are very selective and rare.

Kazes et al. (1991) exploit the Zeeman splitting of H I and find magnetic fields in four HVCs with strengths between $-14.6 \dots +11.6 \mu\text{G}$, with only one detection being significant.

McClure-Griffiths et al. (2010) use the Faraday rotation² to measure the line-of-sight component of the coherent magnetic field in the HVC associated with the Leading Arm of the Magellanic System to be $\langle B_{\parallel} \rangle \gtrsim +6 \mu\text{G}$. The ratio of random to coherent magnetic field is found to be $B_r/B_{\parallel} < 0.8$ from a recent catalogue of extragalactic Faraday rotation derived from the NRAO VLA Sky Survey. By measuring Faraday rotation, a lower limit of $+5 \mu\text{G}$ (Betti et al., 2019) and $+8 \mu\text{G}$ (Hill et al., 2013) is found for the peak line-of-sight magnetic field strength in the Smith Cloud.

3.3.3 Dark matter

Both observations and simulations favour scenarios in which HVCs (in particular CHVCs and UCHVCs) are not associated with a notable amount of DM. Regarding Plöckinger & Hensler (2012) a head-tail structure is produced in clouds irrespective of a DM halo. Actually, in their simulations HVCs with DM approach the Milky Way disk as unrealistic “bullets” that lose almost no mass (less than 10 %). Westmeier et al. (2008) point out that HVCs at a considerable distance from the Milky Way are compatible with infall models without DM halos (Thilker et al., 2004; Chynoweth et al., 2011). The lack of detected DM halos with total gas masses $> 10^6 M_{\odot}$ within a galacto-centric radius of 300 kpc conflicts with cosmological simulations of structure formation (Kravtsov et al., 2004). Barger et al. (2012) argue that HVCs are unlikely to

²The Faraday rotation is the ability of a magnetic field to rotate the plane of polarization of a linear polarized electromagnetic wave if the field is oriented in parallel to the wave vector.

be embedded within DM halos as their corresponding H I virial distances would largely exceed 1,000 kpc (Oort, 1966; Freeman & Bland-Hawthorn, 2002).

Nichols et al. (2014) conclude that the present structure of the Smith Cloud (see section 3.2.1) can best be reproduced when assuming a DM halo. An alternative scenario, in which the Smith Cloud is free of DM, cannot be ruled out if it is on its first approach toward the Galactic disk (Galyardt & Shelton, 2016). Based on its high metallicity of $Z = 0.53 Z_{\odot}$, Fox et al. (2016) suggest that the Smith Cloud consists mainly of processed Galactic material, implying a negligible DM component.

3.3.4 Hydrodynamics

Galactic HVCs move through the Galactic hot halo gas thus experiencing hydrodynamic interactions, which result in particular cloud morphologies. Depending on their velocities and sizes, and on the conditions of the ambient gas, the morphologies are owing to drag (exerting ram-pressure), pressure conservation (expressed by Bernoulli's principle), and hydrodynamic instabilities.

Ram-pressure: HVCs that move through the Galactic hot halo are subject to a ram-pressure

$$P_{\text{rp}} = C_D \varrho_{\text{halo}} v_{\text{HVC}}^2, \quad (3.1)$$

where C_D is the drag coefficient, ϱ_{halo} is the halo density and v_{HVC} is the cloud's speed. The ram-pressure removes gas from the cloud if the criterion (Gunn & Gott, 1972)

$$P_{\text{rp}} > F_{\text{grav}}(r) = \left| \frac{\partial \Phi}{\partial r} \sum_{\text{gas}}(r) \right| \quad (3.2)$$

is fulfilled. By the radius-dependent gravitational restoring force per unit area $F_{\text{grav}}(r)$ the radius of the remaining cloud is inferred. At a given r , $\partial \Phi / \partial r$ describes the gradient of local gravitational potential Φ in radial direction and $\sum_{\text{gas}}$ is the local gas surface density. This mechanism of gas removal is called ram-pressure stripping. Extensive numerical studies of ram-pressure in astrophysical environments of different densities, temperatures, and chemical compositions are conducted by Roediger & Hensler (2005); Roediger & Brüggen (2006, 2007, 2008).

Brüns et al. (2000) and Putman et al. (2003b) distinguish prominent compact, cool condensations near the leading edge from trailing warm, diffuse gas in CHVCs (section 3.2.3) and the MS (section 3.2.2). Such a head-tail morphology is an asymmetric

density distribution with a sharp density gradient at the head and a shallow gradient in the tail and is a direct consequence of ram-pressure. At supersonic speeds, ram-pressure leads to bow-shock shaped HVCs or completely irregular structures (Putman et al., 2003b; Westmeier et al., 2005b,c).

Bernoulli’s principle: In streaming media, Bernoulli’s principle states that the static pressure decreases if the dynamic pressure increases and the sum of both (total pressure) is constant. The static pressure itself is a sum of different contributions, e.g., thermal pressure, magnetic pressure, photon pressure of radiation. If only thermal pressure in the hot halo is accounted for, Bernoulli’s principle reads

$$P_{\text{total}} = P_{\text{static}} + P_{\text{dyn}} = P_{\text{halo}} + P_{\text{dyn}} = P_{\text{therm}} + P_{\text{dyn}} = P_{\text{therm}} + C_D \rho_{\text{halo}} v^2. \quad (3.3)$$

Simulations of moving CHVCs by, e.g., Kwak et al. (2011, see their figure 2) reveal a much higher streaming velocity v at the cloud edges thus implying a higher P_{static} closely in front of and behind the cloud. Consequently, the cloud stretches perpendicularly to its direction of motion and leads to a horseshoe-shaped morphology (Vieser & Hensler, 2007b; Sander & Hensler, 2021). Such shape is observed in CHVCs (Brüns & Mebold, 2004, cf. their figure 10), HVC complexes (Winkel et al., 2011), and cloudlets related to the MS (Konz et al., 2002; Putman, 2004; Yagi et al., 2012).

Hydrodynamic instabilities: Because self-gravitating HVCs move relative to the Galactic hot halo, Kelvin-Helmholtz instability (KHI) and Rayleigh-Taylor instability (RTI) are likely to develop (see appendix A). Typical signatures for KHI are observed at CHVCs catalogued within the GALFA-H I survey (Begum et al., 2010). As a prime example, Complex A (section 3.2.1) reveals both KHI and RTI (see figure 6 in Barger et al., 2020). Fingers and globules (RTI signatures) point toward the Galactic gravitational potential at the lower side of Complex A. Cloud portions being ablated by KHI are observed on the opposite side. Both instabilities may also occur in Complex Galactic Center Negative (Winkel et al., 2011).

3.3.5 Chemistry

Metallicity: Observations mostly reveal sub-solar metallicities in HVCs and indicate a common metallicity of $0.1 \dots 0.3 Z_{\odot}$ (Wakker, 2001, and section 3.2), whereas the Galactic hot halo has solar metallicity in its super-virial phase (see chapter 4). On average, IVCs have higher metallicities than HVCs, because they originate most likely

from already enriched disk material (section 3.2.5).

Dissociation of hydrogen: Molecular hydrogen can absorb photons in the Lyman-Werner bands by means of the two-step Solomon process. Even for the lowest rotational-vibrational states there are 76 absorption lines ranging from 11.1 . . . 13.6 eV (Stecher & Williams, 1967; Dalgarno & Stephens, 1970). If by this process the electronic binding of H_2 gets broken, the molecule is dissociated. The existence of H_2 in cool cores of HVCs (see section 3.2.1) is only possible, because the column densities are sufficient to shield against dissociation (see section 3.3.1).

Ionization and recombination: Ionization is ubiquitously observed in HVCs that reside in the Galactic hot halo. The sources of ionization comprise soft X-rays diffusively emitted by the hot halo gas (Bregman & Harrington, 1986; Brüns et al., 2000; Mathur et al., 2021), Lyman continuum radiation from OB associations in the Galactic disk (photoionization, Weiner et al., 2002; Collins et al., 2007), and heat transfer into the clouds by ram pressure, turbulent mixing, or cosmic rays (collisional ionization, Brüns et al., 2000; Weiner et al., 2002). Photoionization convincingly describes the formation of neutral and low-ionized species, but fails in elucidating highly ionized elements. Instead, collisional ionization leads to the formation of high ions in the transition zones between HVCs and halo (e.g., O VI, Sembach et al., 2003; Fox et al., 2004). It is as well the most likely ionization mechanism for the MS (Weiner & Williams, 1996; Weiner et al., 2002). Observations with FUSE satellite and Hubble Space Telescope have unveiled highly ionized HVCs (Fox et al., 2006; Lehner & Howk, 2011).

Recombination of excited or ionized species is seen in optical and UV emission lines in HVCs (Tufte, 2004) with $\text{H}\alpha$ being the most prominent line (Tufte et al., 2002; Putman et al., 2003a).

3.3.6 Cooling and heating

The common multi-phase structure of HVCs (see section 3.3.1) is a result of different temperatures. The kinetic temperatures are derived from observed line widths (appendix D.2). In some HVCs even molecules and dust are observed (Richter et al., 2001b; Williams et al., 2012; Fox et al., 2023). In such molecular cores the cooling by radiative de-excitation of CO is important (Goldsmith & Langer, 1978; Falgarone & Puget, 1985; Whitworth & Jaffa, 2018) as well as by the depletion of molecules onto dust (Goldsmith, 2001). In the CNM, radiative line cooling by the fine-structure tran-

sition of C II at $\lambda = 158\mu\text{m}$ dominates, while in the WNM several lines contribute, e.g., C II at $\lambda = 158\mu\text{m}$, O I at $\lambda = 163\mu\text{m}$, and Ly α (Wolfire et al., 2003; Hsu et al., 2011). Highly ionized portions of HVCs are mainly observed by O VI absorption (Wakker et al., 2005; Benjamin et al., 2006; Fox et al., 2006), which is an excellent tracer of collisional processes at temperatures $\sim 1 \dots 5 \times 10^5 \text{ K}$ (Sutherland & Dopita, 1993). At temperatures above $\sim 10^5 \text{ K}$ the gas is fully ionized by means of collisions and only bremsstrahlung cools the gas significantly (Boehringer & Hensler, 1989). Radiative cooling at these temperatures is very likely to happen in interface layers between HVCs and CGM (section 3.3.1).

External UV and X-ray radiation heats gas by the photoelectric effect on dust grains (Weingartner et al., 2006). UV photons with $\lambda > 912 \text{ \AA}$ heat the gas by photoionizing its atoms (Bland-Hawthorn & Maloney, 1999a). Galactic HVCs are exposed to UV radiation emitted by OB associations or massive hot stars (Lehner et al., 2011), to the extragalactic UV background (Ganguly et al., 2005; Hill et al., 2013), and to the diffuse soft X-rays emitted by the Galactic hot halo (Snowden et al., 1998; Bregman, 2007; Mathur et al., 2021).

3.3.7 Thermal conduction

At the surfaces of many HVCs highly ionized species are observed (e.g., O VI, C IV, N V, Fox et al., 2005; Ganguly et al., 2005; Savage et al., 2010; Winkel et al., 2011) implying kinetic gas temperatures around $10^{5.5} \text{ K}$ that are unlikely to be produced by photoionization (Sembach et al., 2000; Nelson et al., 2018). Such high-temperature interfaces are proposed to result from thermal conduction caused by a hot ambient medium towards the cool HVCs (Fox et al., 2004; Savage et al., 2005; Wakker et al., 2005; Savage et al., 2010), and radiative turbulent mixing layers due to shear flows (Esquivel et al., 2006; Ji et al., 2019). Prominent examples for observed O VI absorption are highly ionized HVCs and Complex C (Fox et al., 2004), and the Magellanic Stream (Sembach et al., 2000). Savage et al. (2005) analyze sightlines crossing the thick disk and the hot halo of the Milky Way. They find a kinematic coupling between O VI and neutral to weakly ionized gas hence deducing a formation of O VI in the disk-halo interface.

Gnat et al. (2010) find that conductive interfaces greatly enhance highly ionized species (C III, N IV, O V, and above) compared to purely photoionized gas. Condensing conduction fronts agree better with observed narrow linewidths of O VI than evaporative interfaces (Borkowski et al., 1990). Fox et al. (2005) point out that conductive

interfaces reproduce column-density ratios of several ions like $N(\text{C IV})/N(\text{O VI})$ and $N(\text{N V})/N(\text{O VI})$, but can hardly explain observed $N(\text{Si IV})$.

Sander & Hensler (2023) extensively investigate the structure of the interface between multi-phase clouds and a hot gas by means of numerical simulations (cf. section 7.5). They show that hot gas condenses onto the cloud due to saturated thermal conduction. The gas is condensed faster than it distributes spatially within the cloud and therefore accumulates in the transition zone.

3.4 POSSIBLE ORIGINS

The origin of HVCs is still a matter of debate that began shortly after their discovery (Oort, 1966; Hulsbosch, 1975; Songaila et al., 1988). Not at least their high radial velocities with respect to the LSR, which are not consistent with a simple model of Galactic differential rotation, lead to assumptions on extragalactic origins. But also the findings of both coherent and individual velocity fields, which are mainly negative with respect to the LSR, has led to different categories of HVCs and provide a direct evidence for a net infall of low-metallicity gas towards the Galactic disk at diverse galacto-centric distances (de Boer, 1998; Wakker et al., 2007, 2008, see also section 3.2). Indirect evidence for a net infall of HVCs is provided by two arguments (Chiappini et al., 2001): (1) At the current rate of star formation, all of the interstellar medium (ISM) will be turned into stars within about a Gyr. (2) The narrowness of the distribution of metallicities of long-lived stars implies that the metallicity of the ISM remains more or less constant over a Hubble time, which is most likely realized by a continuous inflow of low-metallicity material with a present-day rate of about $1 \text{ M}_{\odot} \text{ yr}^{-1}$ (also known as *G-dwarf problem*, see van den Bergh, 1962; Pagel, 1994).

The large distance and metallicity ranges observed suggest that HVCs may have different origins (van Woerden et al., 1999a,b). Weiner et al. (2002) use the optical $\text{H}\alpha$ emission line from about 20 HVCs to separate them into two classes: $\text{H}\alpha$ -bright clouds with low velocity deviations from Galactic rotation, which are presumably affiliated with the Galactic disk, and $\text{H}\alpha$ -faint clouds with high velocity deviations, which are likely to be infalling gas. The most likely origins deduced from observations and simulations are briefly presented below (see also Richter, 2016).

Infalling gas: Galaxies accrete gas from the IGM as part of the Galactic matter cycle (cf. chapter 1) and thus are able to sustain their star formation. The infall of neutral gas onto galactic disks can be observed directly by 21-cm emission of H I from

HVCs that move through the hot halos of those galaxies. Based on their velocities, distances, and metallicities, the categories of CHVCs and UCHVCs (see sections 3.2.3 and 3.2.4) represent the prime candidates for neutral gas that is being accreted by the Milky Way (Richter, 2012) and by other galaxies like M 31 (Thilker et al., 2004), NGC 6946 (a spiral galaxy, Boomsma et al., 2008), or M 51 and M 83 (both in Miller & Bregman, 2005).

Outflowing gas: Thermal instability (Field, 1965) leads to condensation of hot gas expelled from the disk. Based on this, the concept of a Galactic fountain is introduced by Shapiro & Field (1976) and first applied to HVCs by Bregman (1980) in which gas is ejected from the disk by, e.g., supernovae and hot winds and subsequently condenses in the hot halo of the galaxy. Based on their velocities, distances, and metallicities (e.g., O/Fe ratio, Spitoni et al., 2008) the ubiquitously IVCs are the most likely candidates to result from processed galactic material (e.g., IV Arch, Richter et al., 2001c). Also some HVC complexes are reported to presumably originate in a Galactic fountain, e.g., Complex WB (Thom et al., 2006) and Complex M (Yao et al., 2011). As a consequence, a mixture between outflow and inflow is most realistic at the disk-halo interface of the Milky Way (Fraternali, 2017; Marasco et al., 2022).

Stripped gas: Satellite dwarf galaxies orbiting massive galaxies are stripped off gas due to tidal interactions with the host galaxy. So, this model is also called the “tidal debris model” (e.g. Simon & Blitz, 2002). The most popular example is the Magellanic System (LMC and SMC with the MS). At roughly 50 kpc, the Magellanic Clouds have experienced substantial gas stripping as evidenced by the MS which extends from them (e.g., Mathewson et al., 1974; Hsu et al., 2011). Another infalling satellite galaxy is the Sagittarius dwarf. Tepper-García & Bland-Hawthorn (2018a) conclude that the vast majority of the expelled gas was fully settled onto the Galaxy about 300 Myr ago. It is thus highly unlikely that any of the HVCs or IVCs originate from the Sgr dwarf.

Other origin scenarios: Formation of HVCs by cooling out of supernova shells (Yao et al., 2011), galactic winds (Cashman et al., 2023), or merger processes (Richter, 2006) are likewise eligible, because these formation scenarios cannot be rejected with certainty (as pointed out by, e.g., Wakker & van Woerden, 1997).

4 THE GALACTIC HOT HALO

4.1 MULTIPLE PHASES

While high-velocity clouds (HVCs) and intermediate-velocity clouds (IVCs) are representatives of the cold neutral medium (CNM) and warm neutral to slightly ionized medium (WNM) of the circumgalactic medium (CGM) the hot halo around galaxies constitutes its highly ionized part. A hot Galactic corona around galaxies is first suggested by Spitzer (1956), which he theoretically implies from two observational results:

1. Interstellar clouds at high Galactic latitudes would be too rarefied without any external pressure confinement to account for the observed absorption intensities of radiation emitted by background stars. Absorption spectroscopy by Münch (1957) of clouds above $b = 30$ deg is thus interpreted as a consequence of a surrounding gaseous medium with pressure comparable to the pressure in the Galactic plane confining the clouds. Since such absorbing clouds are ubiquitous, it is highly probable that they are already in pressure equilibrium with the ambient medium.
2. A remarkably constant background radio emission at $\lambda = 3.7$ m received from high Galactic latitudes indicates a rarefied gaseous medium surrounding the Milky Way with radial extension of $10 \dots 20$ kpc (Baldwin, 1955). Shortly before, Baldwin (1954) has observed radio emission in the Andromeda galaxy M 31 that stems from a galacto-centric region of size ~ 10 kpc and contributes major to the entire energy released at $\lambda = 3.7$ m.

Theoretical attempts to describe the feasibility of such a two-phase CGM (cool and dense clouds embedded in a warm and tenuous gas) followed a decade later. Based on stability considerations for a dilute gas in thermal equilibrium Field (1965) figures out combinations of temperature and density for which the state of the gas undergoes a *thermal instability* hence converging to either of two stable solutions (so-called *phases*):

1. a phase of density higher and of temperature lower than the original unstable medium,
2. a phase of density lower and of temperature higher than the original unstable medium.

Both phases coexist in pressure equilibrium. Field et al. (1969) apply these studies to the interstellar medium (ISM) to describe the observed CNM and WNM (for an application to the Galactic disk see Wolfire et al., 2003). This *two-phase model* is now well supported by theoretical, observational, and numerical investigations (Audit & Hennebelle, 2005). McKee & Ostriker (1977) extend the two-phase model by adding another hot, dilute and diffuse phase in mutual pressure balance fed by supernova explosions (see their figure 1). This *three-phase model* was motivated by direct detection of soft X-ray emission (Burstein et al., 1977) and absorption by O VI (York, 1974). McKee & Ostriker (1977) provide temperatures and densities for the CNM of $T_{\text{CNM}} \lesssim 100$ K and $n_{\text{CNM}} \sim 10 \dots 100 \text{ cm}^{-3}$, and for the WNM they give $T_{\text{WNM}} \lesssim 10^4$ K and $n_{\text{WNM}} \sim 0.1 \dots 1 \text{ cm}^{-3}$ (see also Cox, 2005).

The temperature of the hot phase corresponds to the virial temperature of a gaseous sphere (Binney & Tremaine, 1987)

$$T_{\text{vir}} \approx 6.8 \times 10^5 \mu \left(\frac{M_{\text{DM}}}{10^{11} M_{\odot}} \right) \left(\frac{10 \text{ kpc}}{R_{\text{Halo}}} \right) \text{ K}. \quad (4.1)$$

With a mean molecular weight of $\mu = 0.6$ (fully ionized gas) and assumed values for the mass of the dark matter (DM) halo ($M_{\text{DM}} = 10^{12} M_{\odot}$) and the halo radius ($R_{\text{Halo}} \sim 50 \text{ kpc}$) a Milky-Way like halo thus has a virial temperature of $\sim 10^6$ K according to equation (4.1).

4.2 OBSERVATIONAL EVIDENCES FOR A HOT PHASE

Tumlinson et al. (2017) review the multi-phase structure of the hot halo as demonstrated in the multiple ionization stages of several elements (cf. their figure 4). Direct evidence for a hot component ($T > 10^6$ K) in the multi-phase CGM comes from diffuse soft X-ray emission due to bremsstrahlung (Anderson & Bregman, 2010; Anderson et al., 2013), and in absorption along QSO sightlines for the Milky Way and external galaxies (Williams et al., 2005; Gupta et al., 2012). Hot and collisionally ionized gas emits soft X-rays that is being traced in metal emission-lines (mostly O VII, O VIII, and Fe lines in the CGM, Fresco et al., 2020). Kerp et al. (1999) find an emission excess in soft X-rays detected by *ROSAT* satellite close to several HVC complexes, which is mainly due to thermal plasma radiation from the Galactic halo and corresponds to a plasma temperature of $T \sim 1.5 \dots 2 \times 10^6$ K. Since the X-ray emissivity is proportional

to the square of gas density, Li (2020) suspects that the sensitivity of most telescopes is too low to detect faint X-ray emission in the outer Galactic halo beyond ~ 50 kpc. According to Mathur et al. (2021) future X-ray observation missions like XRISM¹ (contemplated by JAXA/NASA), Arcus² (recently accepted for Phase A study at NASA), and Athena³ (planned by ESA) will provide higher detection sensitivity to revolutionize the CGM science.

Regarding Tumlinson et al. (2017), an indirect evidence for a hot phase comes from highly ionized metals that correlate with the low-ionization HVCs (Sembach et al., 2003; Fox et al., 2006; Lehner et al., 2009a; Wakker et al., 2012), suggesting boundary layers between a hot medium and the cooler HVCs. As they point out, the term *multi-phase* relates to the observational fact, that many of the metal ions, which differentiate by an order of magnitude in their ionization-potential energies, are simultaneously present within the same absorber systems embedded in a Galactic halo. Damle et al. (2022) find in their HESTIA⁴ suite that ions tracing sub-virial temperatures (H I, Si III) are more clumpy than highly ionized elements (O VI, O VII, O VIII).

Bregman (2007) shows a positive correlation between the O VII equivalent width and the Galactic X-ray intensity at 3/4 keV. The correlation of the two quantities suggests that most of the O VII absorption has a Galactic origin, probably caused by a halo of size 15...100 kpc. The O VI absorption associated with the Magellanic Stream (MS) always corresponds to Stream H I thus implying a radius of at least 50 kpc for the Galactic halo.

Yao et al. (2009) fit a simple exponential radial profile to X-ray emission by O VII, O VIII and Ne IX yielding a gas temperature at the Galactic disk of $T_0 = 3.6 \times 10^6$ K (scale height 1.4 kpc, mean value of 2.4×10^6 K), which is consistent with the observed O VII/O VIII emission-line ratio. The corresponding midplane density is $n_0 = 1.4 \times 10^{-3} \text{ cm}^{-3}$ (scale height 2.8 kpc). Similar values are found by Yao & Wang (2007) and Hagihara et al. (2010). Recent studies favour a bimodal temperature distribution with a sub-virial component at $T \approx 3.2 \times 10^5$ K (seen in UV absorption with HST/COS, Savage et al., 2014) and a virial component at $T = 1.3 \dots 3.2 \times 10^6$ K (seen by broad O VI absorption lines using Chandra, Mathur et al., 2021). Das et al. (2021) add to these findings by revealing three distinct, coexisting, hot phases using deep Chandra observations of the Milky-Way halo toward the sightline Mrk 421. Different emission-

¹**X-Ray Imaging & Spectroscopy Mission**, see <https://www.nasa.gov/content/goddard/xrism-x-ray-imaging-and-spectroscopy-mission>

²See <http://arcusxray.org/>

³**A**dvanced **T**elescope for **H**igh **E**nergy **A**strophysics, see <https://sci.esa.int/web/athena>

⁴**H**igh-resolution **E**nvironmental **S**imulations of **T**he **I**mmEDIATE **A**rea, see <https://hestia.aip.de/>

line ratios imply a super-virial ($3.2_{-0.5}^{+1.5} \times 10^7$ K), a virial ($1.5 \pm 0.1 \times 10^6$ K), and a sub-virial phase ($3.0 \pm 0.4 \times 10^5$ K). Solar and super-solar metallicities (e.g. Mg, C, Si) are found in all three phases. Das et al. (2021) argue that the three hot phases might originate from active outflows driven by stellar winds and supernovae expanding into the halo.

The super-virial component is very hot and diffuse (Feldmann et al., 2013) thus it remains difficult to be detected in emission of X-rays. A large amount of this hot reservoir emits below the threshold of Galactic absorption at $0.2 \dots 0.3$ keV thus being substantially interfered by the Galactic X-ray background (Li et al., 2008; Rasmussen et al., 2009; Bogdán et al., 2015). It may serve as a reservoir for the missing baryons (Anderson & Bregman, 2010; Lara-DI et al., 2023), which are only partly covered by the sub-virial and virial components (Bregman, 2007; Gupta et al., 2012; Fang et al., 2013; Miller & Bregman, 2015). Bhattacharyya et al. (2022) present XMM-Newton observations around the sightline passing through Mrk 421. The emission spectrum of the hot halo shows that a two-phase temperature model is a better fit to the data compared to a single-phase model. In addition to the virial phase at $T \approx 2 \times 10^6$ K, a super-virial phase at $T = 6.5 \dots 9.1 \times 10^6$ K is required. Within sightlines passing the vicinity of Mrk 421 they even find $T \approx 1.2 \times 10^7$ K. Ponti et al. (2022) analyze the spectral properties of the integrated soft X-ray emission observed by eROSITA. The data are consistent with a super-virial component of the Galactic corona having $kT \sim 0.4 \dots 0.7$ keV (equivalent to $T = 4.6 \dots 8.1 \times 10^6$ K) and might be out of thermal equilibrium. Lara-DI et al. (2023) use the stacking technique through 55 sightlines with Chandra X-ray observations to study the super-virial component ($T \gtrsim 10^7$ K) of the Milky-Way halo. Regarding their findings this component is not restricted to reside inside Galactic bubbles.

Assuming hydrogen to be the most abundant element in the halo and to be fully ionized, the mean electron density as given by

$$n_{\text{elec}} \approx \frac{3M_{\text{Halo}}}{4\pi R_{\text{Halo}}^3 m_{\text{H}}} \text{ cm}^{-3} \quad (4.2)$$

yields $n_{\text{elec}} \sim 0.01 \text{ cm}^{-3}$ for a mass of the hot halo gas of $M_{\text{Halo}} \sim 10^{11} M_{\odot}$ ($M_{\text{gas}} = 2.4 \dots 8.7 \times 10^{10} M_{\odot}$, Martynenko, 2022) and $R_{\text{Halo}} \sim 50 \text{ kpc}$. Martynenko (2022) infers a radial profile $n_{\text{elec}} \propto r^{-0.9 \dots 1.1}$ ($r \gg 1 \text{ kpc}$) in regard to recent spectroscopy of O VII absorption and emission. Kerp et al. (2002) derive an electron density of $n_e = 5 \times 10^{-3} \text{ cm}^{-3}$, which is three times the halo density of $n = 1.5 \times 10^{-3} \text{ cm}^{-3}$. Gaensler et al. (2008) assume a hydrostatic stratification of ionized gas (see their equation 6) to provide a distance-dependent pulsar dispersion measure. The respec-

tive scale height has an updated value of 2.6 ± 0.4 kpc (derived for the distribution of warm gas in the CGM based on observations of Si IV and O VI by Qu & Bregman, 2019). Hsu et al. (2011) investigate cloudlets in the tail of Complex C (Stanimirović et al., 2006) and in the MS (Stanimirović et al., 2008), which are catalogued in the GALFA-HI Survey. By assuming pressure equilibrium of the Complex C cloudlets with the ambient hot halo at $z_{\text{gal}} \sim 3$ kpc, they derive halo densities $n_{\text{halo}} = 0.5 \dots 1 \times 10^{-3} \text{ cm}^{-3}$ (median $0.58 \times 10^{-3} \text{ cm}^{-3}$). When assuming the same for the MS at $z_{\text{gal}} \sim 40$ kpc, they find $n_{\text{halo}} = 0.8 \dots 2 \times 10^{-4} \text{ cm}^{-3}$ (median $0.8 \times 10^{-4} \text{ cm}^{-3}$). Feldmann et al. (2013) use hydrodynamic simulations to infer a halo density of the ionized gas of below 10^{-4} cm^{-3} for distances beyond 100 kpc (cf. their figure 1). Das et al. (2021) find $n \sim 10^{-4} \dots 10^{-3} \text{ cm}^{-3}$ for the virial phase and Ponti et al. (2022) find $n_e \approx 9 \times 10^{-4} \text{ cm}^{-3}$ (scale height is a few kpc) for the super-virial component. By analyzing cosmological simulations from the SMAUG⁵ project Fielding et al. (2020) find nearly constant radial temperature profiles, and broad pressure and radial velocity distributions in Milky-Way mass halos at low redshift. Their idealized simulations with strong Galactic feedback show similar profiles within half of the CGM's virial radius R_{vir} , whereas beyond R_{vir} the cold phase is much suppressed and pressure and velocity profiles are narrower.

Figure 4.1 displays all the phases of the CGM, which are covered by HVCs (chapter 3) and the Galactic hot halo (this chapter). The CNM comprises the cold and partly molecular inner region of multi-phase HVCs (cores and heads of compact high-velocity clouds (CHVCs) and ultra-compact high-velocity clouds (UCHVCs), cloudlets in complexes). The warm, slightly ionized envelopes (halos and tails of CHVCs and UCHVCs, complexes) consist of WNM and the highly to fully ionized transition zones between HVCs and ambient medium are part of the sub-virial phase of the highly ionized medium (HIM). The HIM itself separates into three fully ionized phases: sub-virial, virial, and super-virial. All these three hot phases are observed in the Galactic halo. The mutual pressure equilibrium (indicated by blue and black solid lines in figure 4.1) implies gradients in temperature and density between the phases. Hence, thermal conduction and mass diffusion set in. The yellow circle in figure 4.1 shows the Galactic hot halo that is used in the simulations performed in this thesis (section 7).

⁵Simulating Multiscale Astrophysics to Understand Galaxies, see <https://www.simonsfoundation.org/flatiron/center-for-computational-astrophysics/galaxyformation/smaug/>

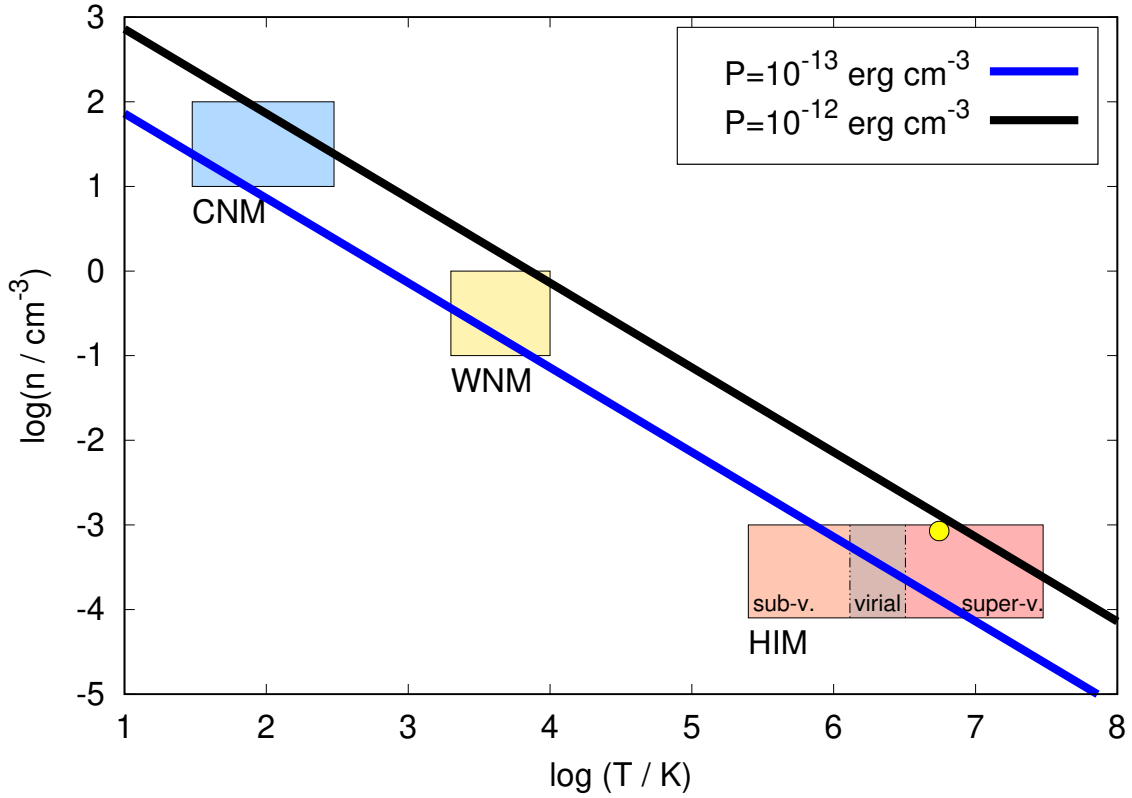


Figure 4.1: Phases of the CGM in the (T, n) -plane. Typical temperature-density regions for the phases are marked by *coloured boxes* and are labelled accordingly. The *diagonal solid lines* illustrate constant pressure values characteristic for the CGM. The *yellow circle* shows the hot halo used in the simulations carried out within this thesis (section 7.1).

4.3 PHASE TRANSITIONS

This somewhat ideal picture of coexisting static phases (CNM, WNM, sub-virial, virial, and super-virial hot phase) has been updated in the last decades based on improved observational facilities sensitive at many wavelengths and an increased computational power to perform numerical simulations. Consequently, we are now aware of a CGM of multiple temperatures, densities, metallicities, and ionization stages. The phases are transformed into one another by diverse physical processes at various timescales (e.g. turbulence, evaporation of clouds, supernovae, collisional and radiative ionization, thermal and hydrodynamic instabilities). Tumlinson et al. (2017) summarize several scenarios to produce an extended, multi-phase CGM:

1. massive inward cooling flows driven by local thermal instabilities (e.g., McCourt et al., 2012),

2. boundary layers between cool clouds and an ambient hot atmosphere they move through (e.g., Begelman & Fabian, 1990; Wakker et al., 2012),
3. continuous shocking and mixing of diffuse halo gas by galactic outflows (e.g., Thompson et al., 2016; Fielding et al., 2017).

Afruni et al. (2021) analyze the CGM in more than 40 low-redshift star-forming galaxies. They figure out that the cool ($T < 10^4$ K) part, which resides in the outer halos at distances $\sim R_{\text{vir}}$, is highly probably the result of cosmological inflow. An origin of cool gas in a supernova-driven outflow at these distances is ruled out according to unphysical coupling of the released energy to the gas. The same is observed in the Andromeda galaxy M 31 (using HST/COS absorption spectra within the AMIGA⁶ project, Afruni et al., 2022).

The ratio of the cooling time

$$\tau_{\text{cool}} \sim 0.63 \left(\frac{R_{\text{Halo}}}{10 \text{ kpc}} \right)^2 \text{ Myr} \quad (4.3)$$

and the free-fall time (both in Binney & Tremaine, 1987)

$$\tau_{\text{ff}} = \sqrt{\frac{3\pi}{32G\rho}} \sim 100 \left(\frac{10^{11} \text{ M}_{\odot}}{M_{\text{DM}}} \right)^{\frac{1}{2}} \left(\frac{R_{\text{Halo}}}{10 \text{ kpc}} \right)^{\frac{3}{2}} \text{ Myr} \quad (4.4)$$

is used to describe the dynamics of the hot halo gas. If $\tau_{\text{cool}} < \tau_{\text{ff}}$ then gas may condense from the halo, which is the key mechanism in the precipitation model to form HVCs or IVCs from a Galactic fountain (cf. section 3.4). McCourt et al. (2012) and Sharma et al. (2012) show numerically that the onset of thermal instability in hot, hydrostatically stratified atmospheres coincides with a ratio $\tau_{\text{cool}}/\tau_{\text{ff}} \approx 10$, below which isobaric density perturbations are thermally unstable and precipitate out of the hot phase thus forming a multi-phase medium (see also Esmerian et al., 2021). Condensation may be impeded only by additional heat sources (galactic feedback like star formation and supernovae) or halo magnetic fields (a value of $B \sim 0.1 \mu\text{G}$ is sufficient, Buie II et al., 2022). For a Milky-Way like halo ($M_{\text{DM}} = 10^{12} \text{ M}_{\odot}$, $R_{\text{Halo}} \sim 50 \text{ kpc}$) one finds $\tau_{\text{cool}} \sim 10 \text{ Myr}$ and $\tau_{\text{ff}} > 200 \text{ Myr}$, which is in line with the simulation results of Buie II et al. (2022) for the sub-virial to virial halo component. However, their simulated super-virial halo gas shows $1 < \tau_{\text{cool}}/\tau_{\text{ff}} < 100$ (see their figure 6), by this supporting a multi-phase nature of the CGM.

Esmerian et al. (2021) use zoom-in galaxy simulations conducted within the FIRE⁷

⁶Absorption Maps In the Gas of Andromeda, see Lehner et al. (2020)

⁷Feedback in Realistic Environments, see <http://fire.northwestern.edu>

project to investigate the origin of the multi-phase structure of the CGM. More than 70 % of their low-entropy gas ($T < 10^4$ K) is formed by thermal instability from the hot halo, which is triggered by density perturbations of strong amplitudes induced by extragalactic accretion, feedback-driven winds, and the debris of tidal interactions.

5.1 CONDITIONS FOR STAR FORMATION

Because stars form out of collapsing gas clouds (Larson, 1981, 1988, 1992, 1994, 1996, 1997, 2002) it is highly intuitive to assume a strong relation between the amount of available gas and the rate and number of stars to be formed. Kennicutt & Evans (2012) point out that the relation between the *volume density* of star formation and gas, respectively, is the physically meaningful quantity. However, most observations of external galaxies can only detect projected surface densities integrated along the sightline. The most commonly used empirical relation therefore maps the local *surface densities* of both star-formation rate (SFR), Σ_{SFR} , and gas, Σ_{gas} , at scales ranging from interstellar clouds to galaxies. The relation is formulated as a power law

$$\Sigma_{\text{SFR}} \propto \left(\Sigma_{\text{gas}} \right)^n \quad (5.1)$$

known as the *Kennicutt-Schmidt law* (Schmidt, 1959; Kennicutt, 1989). Based on values for the densities of helium and stars (young stars, white dwarfs) in the solar neighbourhood, Schmidt (1959) first derives an exponent of $n \approx 2$ (likely ranging between 1 and 3). By considering $\text{H}\alpha$ intensities from a set of 15 galaxies, Kennicutt (1989) proposes $n = 1.3 \pm 0.3$. He later extends the sample to 97 galaxies thus obtaining a value of $n = 1.4 \pm 0.15$ (Kennicutt, 1998). A scaling relation different from (5.1) may exist between star-formation rates and molecular gas masses (Lada et al., 2012).

A threshold value of $\Sigma_{\text{gas}} = 8 \text{ M}_{\odot} \text{ pc}^{-2}$ is derived as a sufficient precondition for star formation (Larson, 2010b), because above a column density of about $10 \text{ M}_{\odot} \text{ pc}^{-2}$ atomic hydrogen is shown to self-shield against dissociation (Krumholz et al., 2009) and thus forms molecules (and formation of molecules is necessary for star formation, Larson, 2010b). In this regard, dust is important for star formation in two ways: it catalyzes the formation of molecular hydrogen H_2 and additionally shields the H_2 from dissociation by the interstellar UV radiation field (Cortese et al., 2010). Shielding is pivotal since clouds that are too diffuse do not cool sufficiently to form stars (Glover & Clark, 2012). And especially H_2 is the cooling molecule needed to initiate cloud collapse thus regulating the star-formation (Grieco et al., 2023). Traficante et al. (2020) determine a threshold of $\Sigma_{\text{gas}} = 0.05 \dots 0.1 \text{ g cm}^{-2}$ for a region to form high-mass stars (stellar feedback would increase this threshold to $\sim 1 \text{ g cm}^{-2}$, Krumholz

& McKee, 2008). Clark & Glover (2014) show in numerical simulations that star formation is only possible in regions where the mean (area-averaged) column density exceeds $N(\text{H I}) \sim 10^{21} \text{ cm}^{-2}$ (see also de Blok & Walter, 2006).

If a cloud is only supported thermally (i.e., without radiative feedback, magnetic fields, turbulence), then linear perturbation theory leads to a gravitational collapse if the cloud's mass exceeds the temperature-dependent *Jeans mass* (Jeans, 1928; Binney & Tremaine, 1987)

$$M_J = \frac{\pi}{6} \left(\frac{\pi k \bar{T}}{G \mu} \right)^{3/2} \bar{\varrho}^{-1/2} = 18 M_\odot \bar{T}^{3/2} \bar{n}^{-1/2}. \quad (5.2)$$

If such a cloud is additionally embedded in an ambient medium of pressure P_{amb} , the threshold mass (5.2) changes accordingly to a new value

$$M_{\text{BE}} = 1.18 \left(\frac{k \bar{T}}{\mu} \right)^2 G^{-3/2} P_{\text{amb}}^{-1/2} \quad (5.3)$$

called *Bonnor-Ebert mass* (Ebert, 1955; Bonnor, 1956). The values $\bar{T}$, $\bar{\varrho}$, $\bar{n}$ and μ correspond to the mean values of kinetic temperature, mass density, particle density, and molecular weight in the cloud, respectively, and k_B and G are Boltzmann's constant and the constant of gravity. A gravitationally unstable cloud (or a region of it) collapses within a free-fall time (4.4). According to the first law of thermodynamics the cloud gas heats up at compression during collapse. It may thus happen that the respective threshold mass (5.2) or (5.3) re-exceeds the cloud's mass and gravitational collapse will be stopped. An efficient way to dissipate thermal energy from the collapsing region over decades of temperature and density is by radiative cooling (see section 6.2.6). The thermodynamic state of interstellar gas is thus a result of the interplay between heating and cooling processes, which in turn are determined by fundamental atomic and molecular physics and by chemical abundances (Klessen et al., 2008; Larson, 2012). Larson (2005) proposes that the thermal properties determine the almost universal low-mass stellar initial mass function (IMF).

So, two necessary conditions must be met in (a subregion of) a cloud for the gravitational collapse to initiate and evolve:

- (i) column density above 10^{21} cm^{-2} and
- (ii) efficient gas cooling over decades of temperature.

The above conditions for star formation are fulfilled in Giant Molecular Clouds (GMCs). These are molecular clouds with $M > 10^4 M_\odot$, but more than 80 per cent of observed

Galactic GMCs comprise $> 10^5 M_\odot$, and a substantial fraction even $> 10^6 M_\odot$ (Stark & Lee, 2006). Their filling factor in the circumgalactic medium (CGM) is below 1 per cent, but they contain the majority of gas in molecular form (Mihalas & Binney, 1981; Kulkarni & Heiles, 1988). The clouds consist of overdense regions (*clumps*) at several spatial scales with some of them form stellar clusters (*star-forming clumps*). Single stars or small stellar systems form out of gravitationally bound regions (*cores*) with masses around their Bonnor-Ebert mass (McKee & Ostriker, 2007). Molecular clouds are wrapped by atomic gas of typically $N(\text{H I}) \sim 2 \times 10^{20} \text{ cm}^{-2}$ in the solar neighbourhood (Bohlin et al., 1978), but a shielding column density of $N(\text{H I}) \sim 1.4 \times 10^{21} \text{ cm}^{-2}$ is needed for CO to form (van Dishoeck & Black, 1988). Beside simple diatomic molecules (H_2 , ^{13}CO) also more complex molecules have been found, e.g., NH_3 (Jijina et al., 1999), different types of PAHs (McGuire et al., 2021), and prebiotic organic molecules (Ohishi, 2019).

Molecular clouds cool mainly via radiative molecular de-excitation. Owing to the various molecules, that are present in GMCs, the gas temperature decreases well below 100 K with major contribution by the CO cooling rate. Additionally, dust cools the gas efficiently if collisions between dust grains and hydrogen molecules occur often enough to keep them thermally well-coupled, which is guaranteed for $n \sim 10^4 \text{ cm}^{-3}$ (Goldsmith, 2001). This is a somewhat larger value than typical densities of molecular gas in the Galactic plane ($3 \dots 12 \times 10^3 \text{ cm}^{-3}$, Liszt, 1993; Sanders et al., 1993). Hence, dust cooling becomes important if clumps and cores start to collapse.

The collapse passes through a fragmentation cascade (Larson, 2003; Field et al., 2006, 2008, 2009, 2010) with the characteristic mass scale for fragmentation and subsequent stellar birth being determined by the thermodynamic properties of the star-forming cloud material (Klessen et al., 2008). Simulations by Hocuk & Spaans (2010b) indicate that fragmentation increases with metallicity, but decreases through cosmic rays and gas-grain collisional heating. Strong X-ray feedback gets the star-forming cloud to fragment into larger clumps whereby fewer but more massive protostellar cores are formed (Hocuk & Spaans, 2010a). In either case, the so-called core-mass function (CMF) for core masses above $1 M_\odot$ seems to be as universal as the IMF and can be approximated by a power law (McKee & Ostriker, 2007). The initial angular momentum of the collapsing cloud region, which is conserved during the collapse phase, but redistributed in different ways, leads to either more or less massive protostellar disks and, by this, to successive fragmentation into binary or multiple systems of low-mass stars or clusters of massive stars (Goodman et al., 1993; Barranco & Goodman, 1998; Burkert & Bodenheimer, 2000; Larson, 2010a).

Evans (1999) points out that typical values of molecular clouds ($\bar{T} \sim 10 \text{ K}$, $\bar{n} \sim$

50 cm^{-3}) would lead to $M_J < 10^2 M_\odot$ and $\tau_{\text{ff}} < 7 \text{ Myr}$. Considering that most of the molecular gas in the Milky Way is contained in GMCs (Elmegreen, 1985), such Jeans masses and free-fall times imply a Galactic star-formation rate above $200 M_\odot \text{ yr}^{-1}$, much greater than the current value (chapter 1). This contradiction strongly implies support mechanisms other than thermal pressure. Most promising are magnetic fields (Larson, 2003; Price et al., 2009; Li et al., 2014a; Van Loo et al., 2015), (supersonic) turbulence (Larson, 2003; Mac Low & Klessen, 2004; Ballesteros-Paredes et al., 2006; Heitsch et al., 2008; Girichidis et al., 2020), and radiative feedback (Ostriker, 2009; Kim et al., 2013, 2014; Girichidis et al., 2020).

5.2 TRIGGERS OF STAR FORMATION

GMCs meet all conditions required for star formation (Shu et al., 1987). Therefore, understanding star formation always includes the consideration of formation and evolution of GMCs.

GMCs are formed by inelastic collision of cold H I clouds until the typical masses and sizes of self-gravitating GMCs are yield. This is called a „bottom-up“ process (McKee & Ostriker, 2007) and poses a problem in terms of very long formation times. The required time to accumulate masses of more than $10^5 M_\odot$ easily exceeds 10^8 yr (Kwan, 1979).

In a „top-down“ scenario the GMCs form out of the interstellar medium (ISM) by either a Parker instability, a Jeans instability, a magneto-Jeans instability, localized converging flows driven by stellar feedback or turbulence, or, in case of strong rotational shear (as is the case in interarm regions of spiral galaxies), by swing amplification (McKee & Ostriker, 2007; Dobbs et al., 2014). In fact, observations of the Milky Way and external spiral galaxies show that molecular gas is mostly located in the spiral arms (e.g., Solomon & Rivolo, 1989; Engargiola et al., 2003).

GMCs are also formed by tidal interactions (Poggianti et al., 2019; Price-Whelan et al., 2019; Pasha et al., 2021, and references therein). The faint tail of Complex H, which fills much of the gap in velocity between the complex and the Galactic disk (Simon et al., 2006), is tidally stripped off and very likely triggers the formation of Digel Cloud 1 (a molecular cloud at a galacto-centric distance $\sim 22 \text{ kpc}$) as well as star formation in that cloud (Izumi et al., 2014, who find two young embedded clusters at two CO peaks)

An externally induced event to form molecular clouds is ram pressure of the hot

intracluster medium (ICM). Gas is stripped from galaxies moving at high speed through galaxy clusters. The removed gas may form molecular clouds that in turn are the sites of star formation. Isolated extragalactic star clusters are the observed witnesses of such events (Hester et al., 2010; Smith et al., 2010; Yagi et al., 2010; Beccari et al., 2017; Bellazzini et al., 2022; Boselli et al., 2022; Jones et al., 2024).

The above presented formation sites for GMCs are generally very dynamic, such that very likely the formation of stars is triggered, too. Known events that locally trigger a Jeans instability (gravitational collapse) are pressure waves (Hennebelle et al., 2003), stellar winds or supernova-driven shock waves (Li et al., 2014b) all of which penetrate molecular clouds and generate density fluctuations. Both ejected material and turbulence created by stellar feedback (Ballesteros-Paredes et al., 2007; Luisi et al., 2021) as well as collisions with other clouds or with the Galactic disk (Fukui et al., 2021a) may lead to a local collapse. Thus, star formation is initiated by an externally induced compression of gas, which implies enhanced cooling. This process is called *triggered star formation* (Girichidis et al., 2020). The star-formation rate directly relates to the total hydrogen content (H I and H₂), which is quantified by equation (5.1).

5.3 STAR FORMATION IN HVCs

The conditions and triggers for star formation (sections 5.1 and 5.2) are theoretically well described, observationally confirmed, and numerically examined. Star-forming GMCs are primarily associated with dynamic regions of galaxies (interarm regions, disk-halo interface) or are formed in dynamic events (ram-pressure stripping in the ICM, tidal interactions with the Galactic gravitational potential). A pre-stellar collapse is triggered by various events as a result of environmental conditions.

Spectroscopic observations of high-velocity clouds (HVCs) reveal molecules like H₂ and CO and dust leading to efficient radiative cooling. High column densities, especially in compact high-velocity clouds (CHVCs), shield the molecules against dissociation. The high speeds of HVCs relative to the Galactic halo provide excellent environmental conditions, such that star formation could be triggered by means of pressure waves. When HVCs approach and collide with the Galactic disk, a triggered star formation may result likewise. Their multi-phase morphologies and high column densities combined with their ubiquitous presence in dynamic environments make CHVCs likely sites of star formation. This resulted in various attempts to observe stars in HVCs (see tables 3.1 and 3.2). Star-forming HVCs are of key interest within the Galactic matter cycle (see chapter 1), because their total infall rate is able to compensate for

the constant star-formation rate observed in the Milky Way (see table 1.1).

Star formation in HVC complexes has been investigated for many years. Ivezić & Christodoulou (1997) browse the IRAS¹ Point Source Catalog and HIRES maps for young stellar objects (YSOs) in the direction of five HI HVCs. Only one YSO is found in direction to core M1.1 in Complex M (section 3.2.1), but its physical relation remains highly tentative. Willman et al. (2002a,b) analyze the stellar content of 12 HVC complexes and CHVC 173+52-225 in an early data release of the Sloan Digital Sky Survey, but do not find any stellar overdensity. Their findings are substantial only for HVCs closer than 350 kpc. If they exist at all, stellar populations are either too faint or too small in angular extent to be discovered. Complex H has been investigated in both mid-IR and radio emission combined with a search in the database of the Two Micron All Sky Survey (2MASS, Lockman, 2003; Simon et al., 2006). No evidence of any assigned dwarf-galaxy-like stellar population (tracing current star formation) or a red giant population (tracing previous star formation) is found. The energy input needed to form the cavity structure observed within Complex H does not result from a stellar source (Morras et al., 1998). Stark et al. (2015) argue that star formation in HVC complexes is unlikely due to several reasons. Without dark matter (DM) support the complexes lose all their mass at timescales ~ 100 Myr, so the probability decreases continuously for exceeding their Jeans mass. The observed mean column densities of $N(\text{HI}) \sim 10^{19} \text{ cm}^{-2}$ (Putman et al., 2012) are hence far below the required 10^{21} cm^{-2} . A significant amount of molecular hydrogen is detected in only a few clouds (e.g., Richter et al., 2001b).

A collision between a HVC and the Galactic disk may lead to highly compressed condensations of low velocity dispersion (Mirabel, 1982), in which star formation may be initiated (e.g., the formation of the Orion Molecular Complex, Franco, 1986). Star formation, that is traced back to a collisional event, is observationally confirmed for the Magellanic System (McClure-Griffiths et al., 2008; Casetti-Dinescu et al., 2014, and section 3.2.2) and for the Smith Cloud (Stark et al., 2015, and section 3.2.1). For Complex C no such evidence is found for a collisionally triggered star formation (Kerp et al., 1996a).

The star-formation history in CHVCs is still very much a mystery. On the one hand, current peak HI column densities rarely exceed 10^{20} cm^{-2} thus being likely not sufficient to allow for extensive ongoing star formation (Braun & Burton, 2000). On the other hand, low metallicities have been measured in at least CHVC 125+41-207

¹InfraRed Astronomical Satellite, mission period for 10 months in 1983

(section 3.2.3) indicating some enrichment in former times. So, associated old stellar populations with a low surface density cannot be excluded. Several observations in the visible and the near-infrared failed in detecting any stars in CHVCs (Davies et al., 2002; Hopp et al., 2003a,b, 2007). By conducting a photometric search, Siegel et al. (2005) identify a number of faint late-type giant stellar candidates in CHVCs, but neither reveal a coherent red giant branch structure nor a clear excess of giant candidate counts. No evidence of any YSOs is found. Surprisingly, some Galactic halo CHVCs and UCHVCs have been observationally confirmed (sections 3.2.3 and 3.2.4) to meet the theoretical conditions to be the sites of star formation (section 5.1):

- (i) high peak column densities as being observed in cores and heads indicate sufficient gas density and
- (ii) a core-halo or head-tail structure indicates multiple phases of HVC gas thus accounting for efficient cooling.

To conclude, no evidences tracing either previous or ongoing star formation have been found in HVCs so far with statistical significance. The HVC cores do not appear to be active sites of star formation (Wakker et al., 1997). Indeed, low metallicities observed in CHVCs indicate enrichment in the past due to either old stellar populations or mixing of halo gas into the CHVCs.

6 NUMERICAL TECHNIQUES

Within this chapter the multi-physics grid code FLASH is briefly introduced, which is used to carry out the simulations in this thesis. The implementations of physical processes into FLASH are explained to such an extent that they complete the descriptions in Sander & Hensler (2021) (section 7.6) and Sander & Hensler (2023) (section 7.5).

6.1 THE FLASH CODE

Problems in fluid dynamics are solved numerically in either Eulerian or Lagrangian formalism. By the former approach the respective conservation laws are integrated to obtain local values for density, velocity, momentum, and pressure. This is performed practically on numerical grids on which the solutions are discretized. In contrast, the latter approach solves for the equations of motion for fluid elements in co-moving coordinates. The numerics is realized with smoothed particle hydrodynamics (SPH) codes (e.g., GADGET-2, Springel, 2005).

Within the present thesis the FLASH¹ code (Fryxell et al., 2000; Dubey et al., 2009) is utilized to simulate the evolution of compact high-velocity clouds (CHVCs) and ultra-compact high-velocity clouds (UCHVCs) in the Galactic hot halo in full three dimensions (i.e., without using spatial symmetries). This code is continuously developed by a research group at the University of Chicago and by a very lively community of users around the world. It is well documented, massively parallelized, and solves hydrodynamic problems over several orders of spatio-temporal magnitude by means of an adaptive grid (or, equivalently, a mesh). Within the version of FLASH used in this thesis implicit integration schemes are implemented to tackle heating and cooling of plasma and thermal conduction in a feasible amount of computing time. FLASH is extensively tested and compared to other grid-based codes (Tasker et al., 2008) and SPH codes (Mitchell et al., 2009; Federrath et al., 2010; Price & Federrath, 2010).

FLASH is written in the FORTRAN programming language. It is structured in many *units* to capture multi-physics problems. Each unit is responsible for either input and output of data, grid management, timestep calculation and solution forwarding, or calculating physical processes. The most important physics units used within the simulations handle hydrodynamics (flux conservation, updating solutions), source terms

¹Website: <https://astro.uchicago.edu/research/flash.php>

(cooling, heating, diffusion, conduction), material properties (viscosity, mass diffusivity, thermal conductivity), the equation of state, and self-gravity. Beside the numerical integration, FLASH uses external libraries for several tasks.

The Paramesh library (Olson et al., 1999; Olson, 2006) structures the computational domain into *blocks* and adaptively refines and derefines these blocks based upon the local density gradient. So, the domain is only refined if necessary which saves a lot of memory and processor power and speeds up the calculations. Each block consists of $8 \times 8 \times 8$ *interior cells* being wrapped by four layers of *guard cells* (cf. figure 6.1).

The interior cells contain the numerical solutions of all updated state variables. The guard cells contain boundary values needed to update the interior cells. The guard cells receive their values from adjacent blocks or externally specified boundary conditions via PARAMESH. The number of guard cells needed per dimension depends upon the interpolation schemes and the differencing stencils used by the various physics units (usually hydrodynamics).

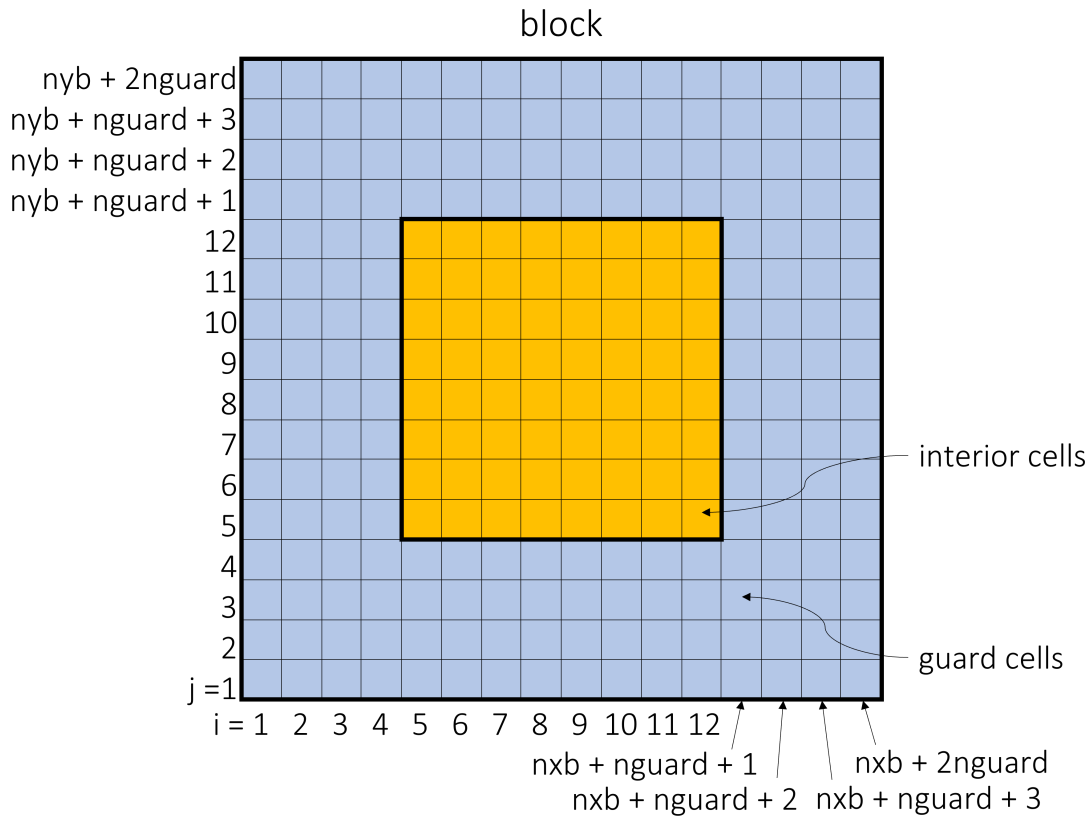


Figure 6.1: Indexing scheme for a two-dimensional grid block with four guard cells ($nguard = 4$) at each block boundary as needed by the numerical PPM scheme (see text). Interior cells are coloured yellow, guard cells are coloured blue.

If a *parent* block is refined by one level, it is divided into two new blocks per spatial dimension and every of the 8 new *leaf* blocks still consists of $8 \times 8 \times 8$ grid cells of $1/8$ volume each. Thus, blocks at different refinement levels have different grid spacings and thus cover different fractions of the computational domain. Any two adjacent blocks differ by not more than one level of refinement. This conservative criterion ensures flux conservation across the block surfaces. The grid is staggered, meaning that scalar quantities are defined in the middle of one grid cell while vector quantities are defined on the surfaces. In D spatial dimensions there are $2^{D(3+j-1)}$ grid cells at refinement level j .

Upon refinement (derefinement) a numerical solution on the coarser (finer) grid needs to be prolonged (restricted) by means of a V-cycle (Briggs et al., 2000; Zumbusch, 2003). Prolongation means interpolating coarse-gridded data on a parent block and mapping this solution to the new leaf blocks generated. Restriction means averaging fine-gridded data on leaf blocks and mapping this solution to their respective parent block. Global block numbers are based on Morton order approximating a space-filling behaviour. Storage order within each processor follows this ordering. After each (de)refinement the blocks are re-distributed among processors to achieve an even load balancing. The blocks are arranged within an oct-tree in three dimensions, a quad-tree in 2D (see figure 6.2), and a binary-tree in 1D. In 3D that simply means that a parent block has either 8 leaf blocks or none (equivalently, 4 and 2 leaf blocks in 2D and 1D, respectively).

The Message Passing Interface library (MPI)² ensures the efficient communication between processors on computational clusters as well as load balancing of computational tasks among the processors.

The Hierarchical Data Format library (HDF5)³ manages the output of simulation data. In FLASH there are two different kinds of output files: *plot files* and *checkpoint files*. A plot file contains a snapshot of the simulation at a certain simulation time. The file size is rather small and this is the preferred file type to generate scientific plots for analyzing the simulations (see chapter 7). In addition, the checkpoint files store *all* data of a simulation at a certain timestep. They are used to restart a simulation from any evolutionary stage.

The numerical scheme: The FLASH code uses a strang-split evolution (Strang, 1968) in order to integrate the Euler equations by using the piecewise parabolic method (PPM, Colella & Woodward, 1984; Woodward & Colella, 1984). The PPM scheme is

²Website: <https://www.intel.com/content/www/us/en/developer/tools/oneapi/mpi-library.html>

³Website: <https://www.hdfgroup.org/solutions/hdf5/>

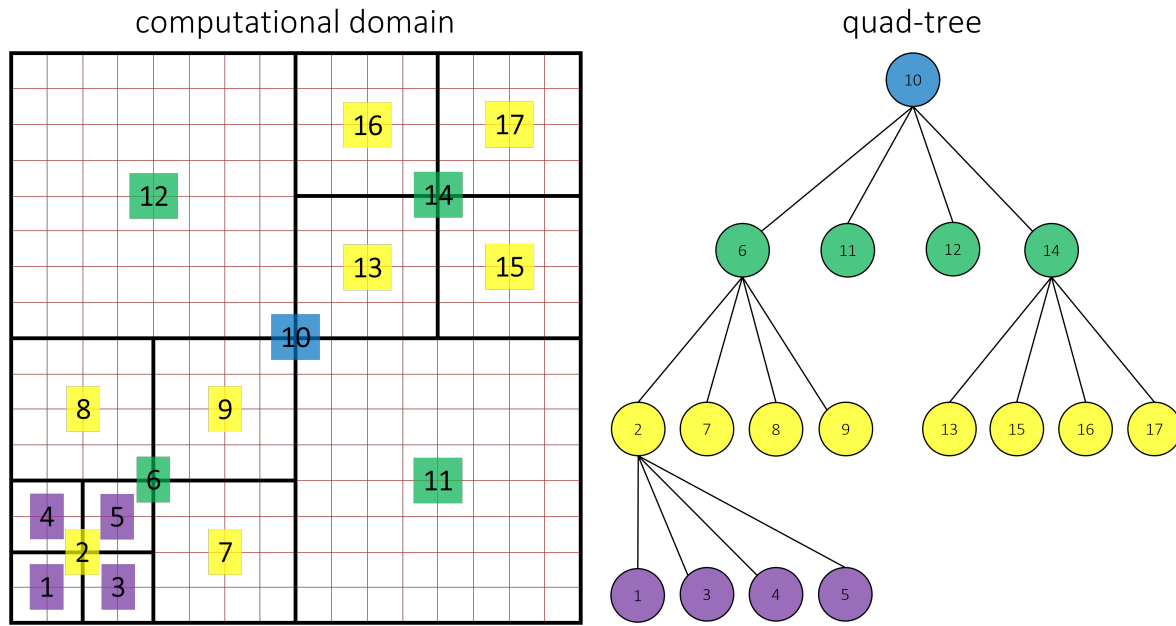


Figure 6.2: *Left plot:* Adaptive mesh refinement of the computational domain into blocks at different refinement levels as realized by PARAMESH. For the sake of clarity, a two-dimensional grid is shown. *Right plot:* The corresponding quad-tree is based on space filling by Morton order.

second-order accurate in both space and time, and greatly reduces numerical viscosity, so all fluid interfaces are sharply preserved and small-scale features can be resolved. The PPM algorithm requires four guard cells in each direction. Figure 6.3 illustrates the strang-split method as executed by FLASH: During an initial source step all physical quantities are written into the grid thus providing their initial values (see section 7.1 for more details). A subsequent evolution loop comprises the following sub-steps:

1. Advection 1: All physical quantities are updated in time for half an integration timestep $\Delta t_{\text{int}}/2$ in “sweep” order $x-y-z$, i.e., physical quantities are transported along the spatial dimensions in the given order by means of fluxes across the grid-cell boundaries.
2. EOS: The equation of state is solved to ensure that all physical quantities relate to each other, i.e., they are thermodynamically consistent in each grid cell.
3. Galilean transformation: This sub-step ensures that the simulated CHVCs are not „pushed out“ of the computational domain by drag of the ambient medium. The relative velocity between CHVCs and the computational domain is set to zero.
4. Source terms: The absolute values of the fluxes, which alter the physical quan-

tities, are updated based on the current plasma properties of viscosity, thermal conductivity, mass diffusivity, and cooling and heating functions.

5. Gravity: Poisson’s equation is solved using current density values.
6. Advection 2: Subsequently, the same advection is performed in “sweep” order $z - y - x$ for a second half-step $\Delta t_{\text{int}}/2$.
7. All the sub-steps EOS, Galilean transformation, Source terms, and Gravity are identically performed as above.
8. Source step: At the very end of the evolution loop the updated state variables (both scalars and vectors) are mapped to the grid, the refinement criteria are checked, and, where necessary, the grid is refined or de-refined, and the timestep is recalculated.

The evolution loop is executed until either the user-specified maximum simulation time or the wall-clock time specific to the computational cluster, the simulations are conducted on, is reached. The calculated quantities in figure 6.3 are explained within the subsequent sections.

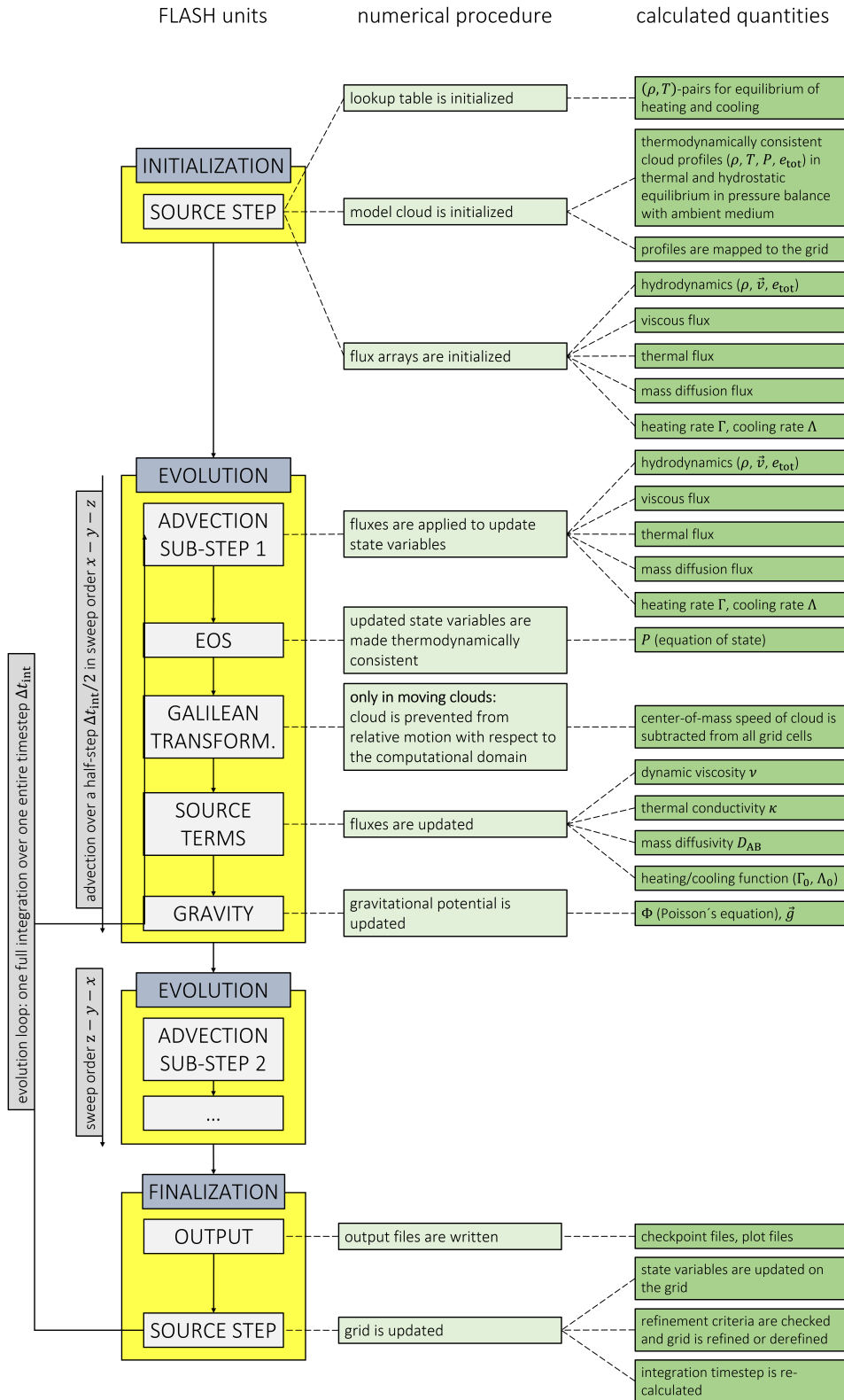


Figure 6.3: Strang-split method used by FLASH. The variable quantities are updated on the grid only after one entire integration timestep Δt_{int} . Please refer to the text for explanation of calculated quantities.

6.2 IMPLEMENTATION OF PHYSICAL PROCESSES

6.2.1 Chemistry

The chemistry network in FLASH is implemented within this thesis to account for calculating the time-dependent particle densities of H_2 , H I , H II , He , metals, and electrons. They are based on dissociation of H_2 , on several ionization and recombination processes, and on the local metallicity. The mean molecular weight is calculated to switch between particle and matter density.

Metallicity: High-velocity clouds have average metallicities of $0.1 \dots 0.3 Z_\odot$ (Wakker, 2001), with intermediate-velocity clouds (IVCs) showing a nearly solar metallicity on average. Metallicity plays a crucial role in controlling both cooling and heating rates of a multi-phase gas (section 3.3.6).

The solar mass fractions of hydrogen (X), helium (Y), and metals (Z) are given by (table 4 in Asplund et al., 2009)

$$(X_\odot, Y_\odot, Z_\odot) = (0.7314, 0.2485, 0.0201). \quad (6.1)$$

Assuming an arbitrary metallicity ratio $\Psi_{\text{metal}} \equiv Z/Z_\odot$, the respective mass fractions within this thesis are calculated according to

$$X(\Psi_{\text{metal}}) = X_\odot + (1 - \Psi_{\text{metal}})Z_\odot \left(1 - \frac{1}{X_\odot/Y_\odot + 1}\right) \quad (6.2)$$

$$Y(\Psi_{\text{metal}}) = Y_\odot + (1 - \Psi_{\text{metal}})Z_\odot \frac{1}{X_\odot/Y_\odot + 1} \quad (6.3)$$

$$Z(\Psi_{\text{metal}}) = \Psi_{\text{metal}}Z_\odot, \quad (6.4)$$

which correctly converge to (6.1) if $\Psi_{\text{metal}} = 1$. The equations (6.2) to (6.4) are derived by assuming a constant ratio X/Y . Let $X_\odot$, $Y_\odot$, $Z_\odot$ be the solar values for the mass fractions. Let $Z = Z_\odot + \Delta Z$, where Z is the actual mass fraction of metals and ΔZ either increases Z above the solar value ($\Delta Z > 0$) or decreases it accordingly ($\Delta Z < 0$). Respectively, define $X = X_\odot + \Delta X$ and $Y = Y_\odot + \Delta Y$. From the constraint $X_\odot + Y_\odot + Z_\odot = 1$ it follows

$$\begin{aligned} Z - \Delta Z &= 1 - X_\odot - Y_\odot = 1 - (X - \Delta X) - (Y - \Delta Y) \\ \iff -\Delta Z &= \Delta X + \Delta Y. \end{aligned} \quad (6.5)$$

Hence, a decrease (increase) in metals results in an increase (decrease) in hydrogen and helium. In terms of Ψ_{metal} it yields

$$\Delta Z = Z - Z_{\odot} = \Psi_{\text{metal}} Z_{\odot} - Z_{\odot} = (\Psi_{\text{metal}} - 1) Z_{\odot}. \quad (6.6)$$

Substituting (6.6) into (6.5) finally yields

$$(1 - \Psi_{\text{metal}}) Z_{\odot} = \Delta X + \Delta Y. \quad (6.7)$$

The constraint $X/Y = \text{const.}$ must hold for all Ψ_{metal} . So,

$$\frac{X}{Y} = \frac{X_{\odot}}{Y_{\odot}} = \frac{\Delta X}{\Delta Y}. \quad (6.8)$$

Hence, from (6.5) the change in Y is obtained

$$\Delta Y = -\frac{\Delta Z}{\Delta X/\Delta Y + 1} = (1 - \Psi_{\text{metal}}) Z_{\odot} \frac{1}{X_{\odot}/Y_{\odot} + 1}, \quad (6.9)$$

where (6.8) is used in the last identity. Analogously,

$$\Delta X = -\Delta Z - \Delta Y = (1 - \Psi_{\text{metal}}) Z_{\odot} \left(1 - \frac{1}{X_{\odot}/Y_{\odot} + 1} \right). \quad (6.10)$$

One easily verifies

$$X + Y + Z = X_{\odot} + \Delta X + Y_{\odot} + \Delta Y + Z_{\odot} + \Delta Z = 1 \quad (6.11)$$

for the values obtained in equations (6.6), (6.9) and (6.10).

In FLASH the metallicity ratio Ψ_{metal} is defined as a mass scalar (cf. section 7.3) ranging between 0 and 1. The local values of Ψ_{metal} are calculated during runtime in line with equation (7.14). The local mass fractions (6.2) to (6.4) are based on Ψ_{metal} . In this thesis, the local chemical composition thus depends on current hydrodynamics.

Dissociation of hydrogen: Regarding Black & Bodenheimer (1975), the fraction of dissociation y is given by

$$y = \frac{\varrho_{\text{H}}}{\varrho_{\text{H}} + \varrho_{\text{H}_2}}, \quad (6.12)$$

where ϱ denotes the respective mass density. Now, y is given by the quadratic equation

$$\frac{y^2}{1 - y} = \frac{2.11}{\varrho X} \exp \left\{ \frac{-52,490 \text{ K}}{T} \right\}, \quad (6.13)$$

based on the dissociation constant for H_2 , and X is given by equation (6.2). Solving (6.13) yields

$$y_1 = -\frac{F}{2} + \sqrt{\frac{F^2}{4} + F} \quad (6.14)$$

as a physically meaningful solution (see figure 6.4, left) with

$$F := \frac{2.11}{\rho X} \exp \left\{ \frac{-52,490 \text{ K}}{T} \right\}. \quad (6.15)$$

The dissociation fraction (6.14) is a very steep, but continuous function of plasma temperature and equals 100 % for $T \gtrsim 800 \dots 1,200 \text{ K}$ while the densities range from $10^{-6} \dots 10^4 \text{ cm}^{-3}$.

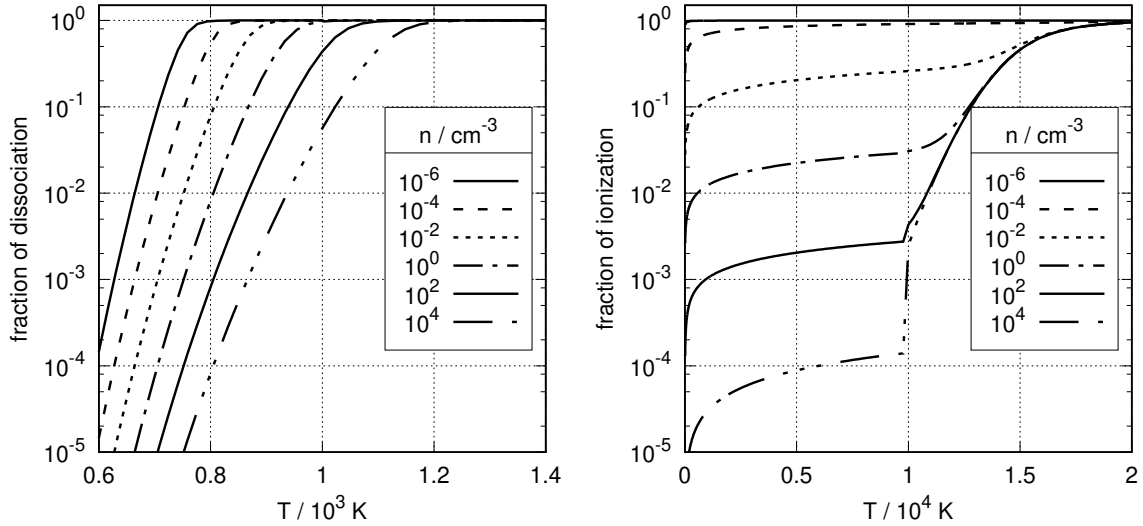


Figure 6.4: Dissociation of molecules (*left plot*) and ionization of atoms (*right plot*) for wide ranges of density and temperature. Above $800 \dots 1,000 \text{ K}$ a relevant fraction of molecules is dissociated. Both mechanisms are the more sensitive to the gas density the lower the temperatures are.

Ionization and recombination: The ionization level of hydrogen is 13.595 eV. The necessary energy for ionization can be provided by radiation (photons) or collisions (electrons). Vice versa, the ionized hydrogen can recombine again when it loses energy via radiation or collisions. Let n_i and n_{i+1} be the particle densities of i - and $i + 1$ -times ionized atoms, $R_{i \rightarrow i+1}$ and $C_{i \rightarrow i+1}$ be the rates of ionization by photons and collisions, and $R_{i+1 \rightarrow i}$ and $C_{i+1 \rightarrow i}$ be the rates of recombination by photons and collisions, respectively. Assuming thermodynamic equilibrium, recombination cancels

ionization and

$$n_i(R_{i \rightarrow i+1} + C_{i \rightarrow i+1}) = n_{i+1}(R_{i+1 \rightarrow i} + C_{i+1 \rightarrow i}) \quad (6.16)$$

holds. In the multi-phase circumgalactic medium (CGM) as modelled within the simulations in this thesis the following ionization processes are used:

- The ionization by soft X-rays emitted by the Galactic hot halo (cf. chapter 4) is a radiative process. The ionization rate is $R_{0 \rightarrow 1}^x = 7 \times 10^{-17} \text{ s}^{-1}$ for photon energies between $0.1 \dots 1 \text{ keV}$ (Dalgarno & McCray, 1972; Scheffler & Elsässer, 1992; Spitzer, 1998). Because the extinction coefficient scales $\propto \nu^{-3}$ only photons with $h\nu < 1 \text{ keV}$ are essential for photoionization.
- The ionization by cosmic rays is a collisional process. The ionization rates are not clearly determined. van Dishoeck & Black (1986) deduce $C_{0 \rightarrow 1}^{cr} = 5 \times 10^{-17} \text{ s}^{-1}$ from observations in the ζ Per cloud. Wolfire et al. (1995) use a primary ionization rate of $1.8 \times 10^{-17} \text{ s}^{-1}$ to obtain a total ionization rate of $C_{0 \rightarrow 1}^{cr} = 3 \times 10^{-17} \text{ s}^{-1}$. van der Tak & van Dishoeck (2000) obtain a rate of $2.6 \times 10^{-17} \text{ s}^{-1}$ based on H^{13}CO^+ fluxes of massive, young stars that are still embedded in molecular clouds. Semi-empirical estimates agree on values of $1 \dots 2 \times 10^{-17} \text{ s}^{-1}$, which are much too small to correspond to an efficient heating of the CGM at temperatures of a few 100 K (Lequeux, 2005). Within this thesis an ionization rate of $C_{0 \rightarrow 1}^{cr} = 7 \times 10^{-17} \text{ s}^{-1}$ is used (Vieser, 2001).
- Ionization by electrons with an ionization rate $C_{0 \rightarrow 1}^{el} = n_e Q_{0 \rightarrow 1}^{el}$ is a collisional process that scales with electron density. The rate coefficient for collisional ionization $Q_{0 \rightarrow 1}^{el}$ by thermal electrons is given by (Mihalas, 1967)

$$Q_{0 \rightarrow 1}^{el}(T) = 5.465 \times 10^{-11} \sqrt{T} \exp \left\{ -\frac{13.543 \text{ eV}}{kT} \right\} (-0.435 + 0.3 \log T) \text{ cm}^3 \text{ s}^{-1}. \quad (6.17)$$

The only meaningful recombination process is radiative recombination with a recombination rate $R_{1 \rightarrow 0}^{\text{rad}} = n_e \sum_{i=2}^{\infty} \alpha_i$. Captures to the $n = 1$ level are excluded, since the released photons do not contribute to the net recombination (“on the spot approximation”). Hence, only coefficient α_2 , which considers all recombinations to level $n = 2$, is taken into account. It is given by (Scheffler & Elsässer, 1992; Spitzer, 1998)

$$\alpha_2(T) = \frac{2.06 \times 10^{-11} f^2}{\sqrt{T}} \phi_2 \left(\frac{h\nu}{kT} \right) \text{ cm}^3 \text{ s}^{-1}, \quad (6.18)$$

with ionization fraction

$$f := \frac{n_{\text{H II}}}{n_{\text{H I}} + n_{\text{H II}}} = \frac{n_{\text{H II}}}{Xn}, \quad (6.19)$$

and X is defined in equation (6.2). The function ϕ_2 in (6.18) is tabulated in Scheffler & Elsässer (1992); Spitzer (1998), but for only a few temperature values. Because ϕ_2 slowly varies with temperature a simple linear approximation

$$\phi_2(T) = \left(\frac{\phi_2^{(a)}}{T^{(a)}} - \frac{T^{(a)}\phi_2^{(b)} - \phi_2^{(a)}T^{(b)}}{T^{(a)}(T^{(a)} - T^{(b)})} \right) T + \frac{T^{(a)}\phi_2^{(b)} - \phi_2^{(a)}T^{(b)}}{T^{(a)} - T^{(b)}} \quad (6.20)$$

is applied for interpolation between any two tabulated values $(T^{(a)}, \phi_2^{(a)})$ and $(T^{(b)}, \phi_2^{(b)})$. Coefficient α_2 is thus be readily available for all temperatures arising within the simulations conducted in this thesis.

Taking $i = 0$ in equation (6.16) yields

$$n_0(R_{0 \rightarrow 1} + C_{0 \rightarrow 1}) = n_1(R_{1 \rightarrow 0} + C_{1 \rightarrow 0}), \quad (6.21)$$

which is to be solved for f considering all ionization and recombination processes introduced above. So, equation (6.21) yields

$$n_{\text{H I}} (R_{0 \rightarrow 1}^x + C_{0 \rightarrow 1}^{\text{cr}} + n_e Q_{0 \rightarrow 1}^{\text{el}}) = n_{\text{H II}} n_e \alpha_2. \quad (6.22)$$

The electron number density n_e is approximated by $n_e \approx n_{\text{H II}} + \chi_0 (n_{\text{H I}} + n_{\text{H II}})$. Elements with ionization potentials lower than hydrogen add a portion $\chi_0 (n_{\text{H I}} + n_{\text{H II}})$ to the total electron density, with $\chi_0 = 5 \times 10^{-4}$ (Dalgarno & McCray, 1972). Using equation (6.19) in equation (6.22) yields

$$0 = f^2 + Af - B, \quad (6.23)$$

with

$$\begin{aligned} A &:= \left[\chi_0 \alpha_2 + (\chi_0 - 1) Q_{0 \rightarrow 1}^{\text{el}} + \frac{R_{0 \rightarrow 1}^x + C_{0 \rightarrow 1}^{\text{cr}}}{Xn} \right] \left[\frac{1}{\alpha_2 + Q_{0 \rightarrow 1}^{\text{el}}} \right] \\ B &:= \left[\frac{R_{0 \rightarrow 1}^x + C_{0 \rightarrow 1}^{\text{cr}}}{Xn} + \chi_0 Q_{0 \rightarrow 1}^{\text{el}} \right] \left[\frac{1}{\alpha_2 + Q_{0 \rightarrow 1}^{\text{el}}} \right] \end{aligned} \quad (6.24)$$

subsuming the rates of all ionization and recombination processes. It turns out, that

$$f_1 = -\frac{A}{2} + \sqrt{\frac{A^2}{4} + B} \quad (6.25)$$

provides a solution of (6.23) with physical meaning. By inspecting the right plot in figure 6.4 one verifies that f moderately and smoothly increases until hydrogen is fully ionized at $T \gtrsim 10^4$ K and increases dramatically for higher temperatures if particle densities are above 1 cm^{-3} . This increase is the more pronounced the higher the gas density is.

Mean molecular weight: The mean molecular weight μ is a function of temperature, density, and metallicity of the respective gas and thus implicitly depends on the fractions of both dissociation y and ionization f . Accounting for a gas that consists of hydrogen, helium, and metals, its particle density and mass density are related via

$$\frac{1}{\mu} = \frac{n}{\varrho} = \frac{n_X + n_Y + n_Z}{\varrho} = \frac{X}{\mu_X} + \frac{Y}{\mu_Y} + \frac{Z}{\mu_Z}. \quad (6.26)$$

If y and f are additionally considered, the complete expression for μ is given by

$$\frac{1}{\mu} = \frac{(1-y)(1-f)X}{\mu_{X,\text{mol}}} + \frac{y(1-f)X}{\mu_{X,\text{neu}}} + \frac{yfX}{\mu_{X,\text{ion}}} + \frac{(1-f)Y}{\mu_{Y,\text{neu}}} + \frac{fY}{\mu_{Y,\text{ion}}} + \frac{(1-f)Z}{\mu_{Z,\text{neu}}} + \frac{fZ}{\mu_{Z,\text{ion}}}. \quad (6.27)$$

The values of μ for molecular, neutral, and ionized gas components are listed in table 6.1.

Table 6.1: Mean molecular weights for the gas components in molecular, neutral, and ionized phase. Helium and metals are not considered to be present in molecular form. The atomic mass unit reads $u = 1.67 \times 10^{-24}$ g.

	hydrogen X	helium Y	metals Z
μ_{mol}	$2u$	$/$	$/$
μ_{neu}	$1u$	$4u$	$17u$
μ_{ion}	$\frac{1}{2}u$	$\frac{4}{3}u$	$2u$

The inverse of equation (6.27) is plotted in figure 6.5 for a gas composed of solar metallicity as given in equation (6.1). It is apparent that μ converges to $2.4u$ in the molecular phase of the cold neutral medium (CNM) ($T < 10^3$ K) for increasing density. The limit in the phase of warm neutral to slightly ionized medium (WNM) (fully dissociated, partly ionized, $T \lesssim 10^5$ K) reads $1.2u$. For $T > 10^5$ K it holds $\mu = 0.6$ irrespective of density.

By considering the mass fractions (6.2) to (6.4), the fractions of dissociation (6.14) and ionization (6.25), and the mean molecular weight (6.27) the particle densities of

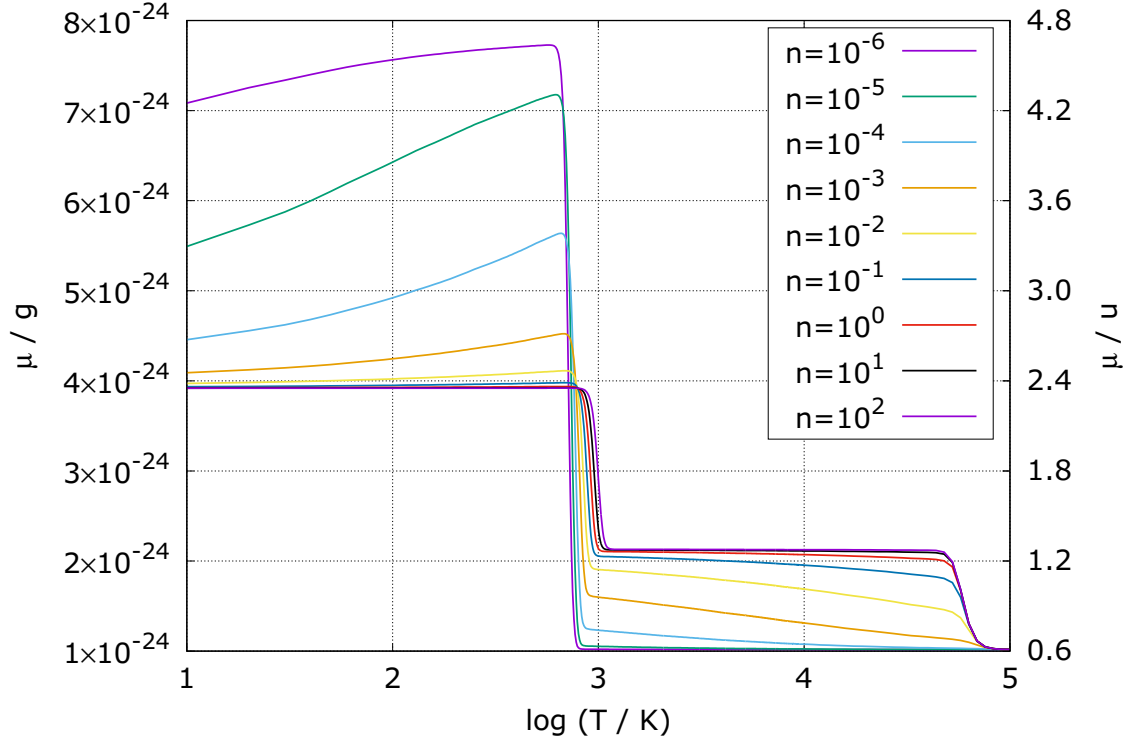


Figure 6.5: Mean molecular weight as a function of temperature for different densities.

H_2 , H I , H II , He , metals, and electrons are given by

$$n_{\text{H}_2} = (1 - y) \frac{n_X}{2} = (1 - y) \frac{Xn}{2} = (1 - y) X \frac{\rho}{2\mu} \quad (6.28)$$

$$n_{\text{H I}} = y(1 - f)n_X = y(1 - f)Xn = y(1 - f)X \frac{\rho}{\mu} \quad (6.29)$$

$$n_{\text{H II}} = yfn_X = yfXn = yfX \frac{\rho}{\mu} \quad (6.30)$$

$$n_{\text{He}} = (1 - f)n_Y = (1 - f)Yn = (1 - f)Y \frac{\rho}{\mu} \quad (6.31)$$

$$n_{\text{metals}} = (1 - f)n_Z = (1 - f)Zn = (1 - f)Z \frac{\rho}{\mu} \quad (6.32)$$

$$n_e = n_{\text{H II}} + \chi_0 (n_{\text{H I}} + n_{\text{H II}}). \quad (6.33)$$

The above densities are calculated in every grid cell during runtime in the simulations conducted with FLASH.

6.2.2 Hydrodynamics

The beginnings of numerical treatment of hydrodynamic shocks date back to von Neumann & Richtmyer (1950). Since FLASH is an Eulerian grid code it solves for

the successive states of a fluid at each position $\vec{x}$ and time t by calculating the *local* mass density ρ , total energy density ρe_{tot} , pressure P , temperature T , and velocity $\vec{v}$. The governing equations are derived from corresponding conservations laws for mass, momentum, and energy. These *Euler equations* are handled by FLASH in the conservation form

$$\frac{\partial \rho}{\partial t} + \nabla (\rho \vec{v}) = 0 \quad (6.34)$$

$$\frac{\partial \rho \vec{v}}{\partial t} + \nabla (\rho \vec{v} \vec{v}) + \nabla P = \rho \vec{g} \quad (6.35)$$

$$\frac{\partial \rho e_{\text{tot}}}{\partial t} + \nabla [(\rho e_{\text{tot}} + P) \vec{v}] = \sum_{i=1}^n S_i, \quad (6.36)$$

with S_i representing all source terms modifying the total energy e_{tot} of the fluid. From the fluid's e_{tot} and $\vec{v}$, as provided by equations (6.35) and (6.36), the internal energy

$$e_{\text{in}} = e_{\text{tot}} - \frac{1}{2} |\vec{v}|^2 \quad (6.37)$$

is gained. Thus, the fluid's pressure is being calculated from the caloric equation of state

$$P = (\gamma - 1) \rho e_{\text{in}}. \quad (6.38)$$

At every integration timestep Δt_{int} the Euler equations together with equations (6.37) and (6.38) are solved hence updating the physical state of the fluid at each position $\vec{x}$ in the computational domain for all times t . The evolution of the fluid is thus fully described (Landau & Lifshitz, 2013). Within all simulations an ideal gas is assumed, so $\gamma = 5/3$.

The numerical integration according to hydrodynamics is unconditionally stable if the integration timestep satisfies the Courant-Friedrichs-Lewy (CFL) condition (Courant et al., 1928)

$$\Delta t_{\text{int}} \leq \Delta t_{\text{cfl}} = C_{\text{cfl}} \frac{\min\{\Delta x, \Delta y, \Delta z\}}{\max\{|\vec{v}| + c_s\}}. \quad (6.39)$$

The CFL safety factor $0 < C_{\text{cfl}} < 1$ ensures that during one integration timestep a fluid parcel is being advected by not more than one grid cell. The volume of one grid cell at finest refinement level is $\Delta x \times \Delta y \times \Delta z$. The amplitudes of the fluid's velocity $|\vec{v}|$ and sound speed c_s in the denominator of equation (6.39) are global maximum values in the entire computational domain. Within this thesis, $C_{\text{cfl}} = 0.8$ in simulations of clouds at rest with respect to the ambient medium (section 7.5) and $C_{\text{cfl}} = 0.2$ in simulations of CHVCs (section 7.6).

If more physical processes enter the problem, FLASH conservatively selects Δt_{int} to be the minimum of all timesteps, which is provided by the fastest of all included physical processes acting as source terms in equation (6.36). Even though this ensures a numerically stable solution (see appendix F), the integration timestep Δt_{int} used for solving the Euler equations can be very small compared to Δt_{eff} . As described in Sander & Hensler (2021) the timescales of cooling (see section 6.2.6) and thermal conduction (cf. section 6.2.4) can drastically reduce Δt_{int} . Hence, (semi-)implicit integration schemes are implemented in FLASH within this thesis to avoid very long computation times.

6.2.3 Artificial viscosity

In FLASH, artificial viscosity is taken into account to smooth discontinuities at shock fronts across multiple grid cells. Thus, numerical instabilities are prevented at steep gradients (von Neumann & Richtmyer, 1950). A viscous flux

$$\vec{j}_{\text{visc}} = -\eta \nabla |\vec{v}|, \quad [\text{g cm}^{-1} \text{ s}^{-2}] \quad (6.40)$$

depends on the velocity gradient between shear flows and on the dynamic viscosity η . The latter relates to the kinematic viscosity via

$$\nu = \frac{\eta}{\rho}. \quad (6.41)$$

Viscosity implies friction thus releases heat that modifies the energy of the gas. The viscous heating is considered as a source term

$$-\nabla |\vec{v}| \vec{j}_{\text{visc}} = \eta (\nabla |\vec{v}|) (\nabla |\vec{v}|), \quad [\text{erg cm}^{-3} \text{ s}^{-1}], \quad (6.42)$$

on the right-hand side of equation (6.36). Consequently, the energy flux (6.40) modifies the total energy in the advection sub-steps within FLASH (cf. figure 6.3).

6.2.4 Thermal conduction

Particles of different velocities collide and, by this, exchange energy and momentum. This process is called thermal conduction and substantial work in the context astrophysical plasmas was carried out by Lyman Spitzer. Within the present work thermal energy is transported by unbound electrons and transferred to other unbound

electrons and ions. Thermal conduction by neutral particles and ions is negligible due to their large masses. According to Spitzer (1962) a temperature gradient ∇T in a plasma leads to both a heat flux $\vec{q}$ and an electric current $\vec{j}$:

$$\vec{j} = \frac{1}{\eta} \vec{E} + \alpha \nabla T \quad (6.43)$$

$$\vec{q} = -\beta \vec{E} - \kappa \nabla T. \quad (6.44)$$

The coefficient of thermal conduction for a Lorentz gas⁴ reads

$$\kappa_L = 20k_B \left(\frac{2}{\pi} \right)^{3/2} \frac{(k_B T)^{5/2}}{e^4 Z \ln(\Lambda_C) m_e^{1/2}}, \quad (6.45)$$

where the Coulomb logarithm $\ln(\Lambda_C) = \ln(\lambda_D/r_c) = 29.7 + \ln[10^{-6} T n_e^{-1/2}]$ defines the ratio between the Debye length λ_D and the classical mean distance of closest approach r_c in a Lorentz gas. For an arbitrary gas it holds

$$\kappa = \delta_T \kappa_L, \quad (6.46)$$

with a correction factor δ_T depending on the atomic number of ions. Spitzer & Härm (1953) determine δ_T numerically (cf. table 6.2).

Table 6.2: Conductivity correction δ_T for an arbitrary gas (characterized by atomic number Z) with corresponding reduction factor ϵ .

Z	1	2	4	16	$\rightarrow \infty$
δ_T	0.225	0.356	0.513	0.791	1.000
ϵ	0.419	0.410	0.401	0.396	0.400

In the present work no electric currents should flow in the direction of ∇T , such that electrons only transport thermal energy along the gradient. To this regard, a secondary electric field must be accounted for which guarantees that no electric energy is transported along ∇T , i.e., $\vec{j} = \vec{0}$ in equation (6.43). But this secondary field as well reduces the thermal conductivity (6.46) such that an effective thermal conductivity

$$\kappa_{\text{class}} = \epsilon \kappa = \epsilon \delta_T \kappa_L \quad (6.47)$$

results (Cohen et al., 1950). Spitzer & Härm (1953) show that the constraint $\vec{j} = \vec{0}$ is

⁴In order to describe dilute electron gases in metals, Hendrik Lorentz proposed in 1905 a model with fixed heavy atoms, whereas the electrons are interacting with the atoms but not with each other. For simplicity, Lorentz approximated the atoms by elastic spheres.

satisfied for $\epsilon = 0.4$ (table 6.2) such that (6.47) can more conveniently be written as

$$\kappa_{\text{class}} = 1.84 \times 10^{-5} \frac{T^{5/2}}{\ln(\Lambda_C)}. \quad (6.48)$$

The effective thermal conductivity (6.48) yields the classical heat flux

$$\vec{q}_{\text{class}} = -\kappa_{\text{class}} \nabla T, \quad (6.49)$$

in astrophysical plasmas for *shallow* temperature gradients. However, in the presence of *steep* temperature gradients the classical description (6.49) greatly overestimates the thermal energy conducted (Gray & Kilkenny, 1980; Campbell, 1984), which has been verified in laser-fusion experiments (Malone et al., 1975; Gary et al., 1978; Rogers et al., 1989). Levermore & Pomraning (1981) introduced the method of flux-limited diffusion that considers a maximum amount of thermal energy being conducted by electrons at thermal velocity v_e . Hence, the maximum possible heat flux amounts to

$$\vec{q}_{\text{max}} = n_e \left(\frac{3k_B T}{m_e} \right)^{1/2} \left(\frac{3}{2} k_B T \right) \hat{v}_e. \quad (6.50)$$

By introducing a flux-limiting parameter $\gamma = 0.25$ (Shirazian & Steinhauer, 1981) and considering $\vec{j} = \vec{0}$, the maximum flux (6.50) turns to be saturated. Cowie & McKee (1977) find

$$\vec{q}_{\text{sat}} = -\kappa_{\text{sat}} \nabla T = \epsilon \gamma \vec{q}_{\text{max}} = 0.1 \vec{q}_{\text{max}} \approx 5 \Phi_{\text{sat}} \varrho c_s^3 \hat{v}_e, \quad [\text{g s}^{-3}], \quad (6.51)$$

with sound speed c_s and a parameter Φ_{sat} that comprises uncertainties in the flux-limited diffusion and the impact of magnetic fields. Within this thesis, $\Phi_{\text{sat}} = 1$, because magnetic fields are not accounted for. A continuous transition between regions of shallow temperature gradients (classic domain) and steep gradients (saturated domain) is given by (Slavin & Cox, 1992)

$$\kappa_{\text{eff}} = \kappa_{\text{sat}} \left(1 - \exp \left\{ 1 - \frac{\kappa_{\text{class}}}{\kappa_{\text{sat}}} \right\} \right). \quad (6.52)$$

The heat input by thermal conduction is considered as a source term on the right-hand side of equation (6.36) by the specific energy flux

$$-\nabla \vec{q}_{\text{eff}} = \nabla (\kappa_{\text{eff}} \nabla T), \quad [\text{erg cm}^{-3} \text{ s}^{-1}] \quad (6.53)$$

and is added to the total energy in the advection sub-steps within FLASH (cf. fig-

ure 6.3). Please refer to Sander & Hensler (2021, discussed in section 7.6) for an extended discussion of the saturated heat flux.

The timescale for thermal conductivity is given by

$$\Delta t_{\text{cond}} = \frac{\min\{\Delta x, \Delta y, \Delta z\}^2}{\max\{\delta\}}, \quad \delta := \kappa_{\text{eff}}/(\varrho c_V), \quad (6.54)$$

and is shortest where ∇T is largest. The numerator comprises the smallest spatial scale in the computational domain, i.e., grid cells at finest refinement level. The denominator considers local thermal conductivity as defined in equation (6.52) with $c_V = \partial e_{\text{in}}/\partial T$ being the isochoric heat capacity. Within the simulations analyzed in this thesis, the shortest times are obtained at the interfaces between clouds and ambient hot gas, as both spatial refinement and thermal conductivity are highest here. So, in practice, it always applies $\Delta t_{\text{cond}} \ll \Delta t_{\text{eff}}$ in the simulations and computation times would be unbearably long as FLASH always applies the shortest of all possible timescales for integration. To take this circumstance into account the semi-implicit Crank-Nicolson integration scheme (Crank et al., 1947) is implemented into FLASH within this thesis.

If pure thermal conduction is present and no other source terms alter the internal energy of the fluid, then only equation (6.53) adds a source term and equation (6.36) reduces to

$$\frac{\partial e}{\partial t} = -\nabla \vec{q} = \nabla (\kappa \nabla T). \quad (6.55)$$

By taking the time dependence of T into account, the latter equation can be rewritten and rearranged thus yielding

$$\frac{\partial e}{\partial t} = \frac{\partial e}{\partial T} \frac{\partial T}{\partial t} = \nabla (\kappa \nabla T) \iff \frac{\partial T}{\partial t} = \frac{\partial T}{\partial e} \nabla (\kappa \nabla T). \quad (6.56)$$

By substituting (u being the atomic mass unit)

$$P = \varrho k_B T \mu u = (\gamma - 1)e \iff e = \frac{k_B}{(\gamma - 1)u} \frac{\varrho T}{\mu} \implies \frac{\partial e}{\partial T} = \frac{k_B}{(\gamma - 1)u} \frac{\varrho}{\mu} = \frac{e}{T}, \quad (6.57)$$

it follows from (6.56)

$$\frac{\partial T}{\partial t} = \frac{T}{e} \nabla (\kappa \nabla T). \quad (6.58)$$

In order to numerically integrate equation (6.58) in three dimensions $\vec{r} = (x, y, z)$ it is split into three ordinary differential equations (ODEs). Each of these ODEs is being solved separately along one of the spatial dimensions x, y, z while the respective other

two spatial dimensions are kept constant. Hence,

$$\frac{dT}{dt}(x) = \frac{T(x)}{e(x)} \frac{d}{dx} \left[\kappa \frac{dT}{dx}(x) \right] \Big|_{y,z=\text{const.}} \quad (6.59)$$

$$\frac{dT}{dt}(y) = \frac{T(y)}{e(y)} \frac{d}{dy} \left[\kappa \frac{dT}{dy}(y) \right] \Big|_{x,z=\text{const.}} \quad (6.60)$$

$$\frac{dT}{dt}(z) = \frac{T(z)}{e(z)} \frac{d}{dz} \left[\kappa \frac{dT}{dz}(z) \right] \Big|_{x,y=\text{const.}} \quad (6.61)$$

The numerical solution of the above equations is shown below exemplarily for the x -sweep (y - and z -sweep are handled equivalently). Regarding the Crank-Nicolson scheme the discretization of equation (6.59) is given by

$$\begin{aligned} \frac{T_{ijk}^{(n+1)} - T_{ijk}^{(n)}}{\Delta t} = & (1 - \alpha) \left\{ \frac{T_{ijk}^{(n)}}{e_{ijk}^{(n)}} \frac{1}{\Delta x} \times \right. \\ & \left[\left(\kappa_{i+1,jk}^{(n)} \frac{T_{i+1,jk}^{(n)} - T_{ijk}^{(n)}}{\Delta x} \right) - \left(\kappa_{ijk}^{(n)} \frac{T_{ijk}^{(n)} - T_{i-1,jk}^{(n)}}{\Delta x} \right) \right] \Big\} \\ & + \alpha \left\{ \frac{T_{ijk}^{(n+1)}}{e_{ijk}^{(n+1)}} \frac{1}{\Delta x} \times \right. \\ & \left[\left(\kappa_{i+1,jk}^{(n+1)} \frac{T_{i+1,jk}^{(n+1)} - T_{ijk}^{(n+1)}}{\Delta x} \right) - \left(\kappa_{ijk}^{(n+1)} \frac{T_{ijk}^{(n+1)} - T_{i-1,jk}^{(n+1)}}{\Delta x} \right) \right] \Big\}, \end{aligned} \quad (6.62)$$

where i is the index that counts the grid cells in x -direction, j in y -direction, and k in z -direction. The factor α determines the method to be explicit ($\alpha = 0$) or fully implicit ($\alpha = 1$). The Crank-Nicolson method sets $\alpha = 0.5$ and thus weights the contributions of the current timestep n and the next timestep $n + 1$ equally. Within this thesis, the weighting factor respects the ratio of Δt_{cfl} (6.39) to Δt_{cond} (6.54). So,

$$\alpha = \frac{1 - \beta - e^{-\beta}}{\beta(e^{-\beta} - 1)} = \begin{cases} 0.5 & , \alpha < 0.5 \\ \alpha & , 0.5 \leq \alpha \leq 1 \\ 1 & , \alpha > 1 \end{cases} \quad , \quad \beta = \frac{\Delta t_{\text{cfl}}}{\Delta t_{\text{cond}}}, \quad (6.63)$$

guarantees a minimum weight of 50 % for the implicit contribution to (6.62). Rear-

ranging equation (6.62) leads to

$$\begin{aligned}
& \left[-\frac{\alpha\Delta t}{(\Delta x)^2} \frac{T_{ijk}^{(n+1)}}{e_{ijk}^{(n+1)}} \kappa_{ijk}^{(n+1)} \right] T_{i-1,jk}^{(n+1)} + \left[1 + \frac{\alpha\Delta t}{(\Delta x)^2} \frac{T_{ijk}^{(n+1)}}{e_{ijk}^{(n+1)}} \left(\kappa_{i+1,jk}^{(n+1)} + \kappa_{ijk}^{(n+1)} \right) \right] T_{ijk}^{(n+1)} \\
& + \left[-\frac{\alpha\Delta t}{(\Delta x)^2} \frac{T_{ijk}^{(n+1)}}{e_{ijk}^{(n+1)}} \kappa_{i+1,jk}^{(n+1)} \right] T_{i+1,jk}^{(n+1)} \\
& = \left[\frac{(1-\alpha)\Delta t}{(\Delta x)^2} \frac{T_{ijk}^{(n)}}{e_{ijk}^{(n)}} \kappa_{ijk}^{(n)} \right] T_{i-1,jk}^{(n)} + \left[1 - \frac{(1-\alpha)\Delta t}{(\Delta x)^2} \frac{T_{ijk}^{(n)}}{e_{ijk}^{(n)}} \left(\kappa_{i+1,jk}^{(n)} + \kappa_{ijk}^{(n)} \right) \right] T_{ijk}^{(n)} \\
& + \left[\frac{(1-\alpha)\Delta t}{(\Delta x)^2} \frac{T_{ijk}^{(n)}}{e_{ijk}^{(n)}} \kappa_{i+1,jk}^{(n)} \right] T_{i+1,jk}^{(n)}. \tag{6.64}
\end{aligned}$$

With equation (6.57) the latter equation is rewritten as

$$\begin{aligned}
& \left[-\frac{\alpha\Delta t}{(\Delta x)^2} \frac{(\gamma-1)u}{k_B} \frac{\mu_{ijk}^{(n+1)}}{\varrho_{ijk}^{(n+1)}} \kappa_{ijk}^{(n+1)} \right] T_{i-1,jk}^{(n+1)} \\
& + \left[1 + \frac{\alpha\Delta t}{(\Delta x)^2} \frac{(\gamma-1)u}{k_B} \frac{\mu_{ijk}^{(n+1)}}{\varrho_{ijk}^{(n+1)}} \left(\kappa_{i+1,jk}^{(n+1)} + \kappa_{ijk}^{(n+1)} \right) \right] T_{ijk}^{(n+1)} \\
& + \left[-\frac{\alpha\Delta t}{(\Delta x)^2} \frac{(\gamma-1)u}{k_B} \frac{\mu_{ijk}^{(n+1)}}{\varrho_{ijk}^{(n+1)}} \kappa_{i+1,jk}^{(n+1)} \right] T_{i+1,jk}^{(n+1)} \\
& = \left[\frac{(1-\alpha)\Delta t}{(\Delta x)^2} \frac{(\gamma-1)u}{k_B} \frac{\mu_{ijk}^{(n)}}{\varrho_{ijk}^{(n)}} \kappa_{ijk}^{(n)} \right] T_{i-1,jk}^{(n)} \\
& + \left[1 - \frac{(1-\alpha)\Delta t}{(\Delta x)^2} \frac{(\gamma-1)u}{k_B} \frac{\mu_{ijk}^{(n)}}{\varrho_{ijk}^{(n)}} \left(\kappa_{i+1,jk}^{(n)} + \kappa_{ijk}^{(n)} \right) \right] T_{ijk}^{(n)} \\
& + \left[\frac{(1-\alpha)\Delta t}{(\Delta x)^2} \frac{(\gamma-1)u}{k_B} \frac{\mu_{ijk}^{(n)}}{\varrho_{ijk}^{(n)}} \kappa_{i+1,jk}^{(n)} \right] T_{i+1,jk}^{(n)}. \tag{6.65}
\end{aligned}$$

By defining

$$A_{ijk} := -\frac{\alpha\Delta t}{(\Delta x)^2} \frac{(\gamma-1)u}{k_B} \frac{\mu_{ijk}^{(n+1)}}{\varrho_{ijk}^{(n+1)}} \kappa_{ijk}^{(n+1)} \quad (6.66)$$

$$B_{ijk} := 1 + \frac{\alpha\Delta t}{(\Delta x)^2} \frac{(\gamma-1)u}{k_B} \frac{\mu_{ijk}^{(n+1)}}{\varrho_{ijk}^{(n+1)}} \left(\kappa_{i+1,jk}^{(n+1)} + \kappa_{ijk}^{(n+1)} \right) \quad (6.67)$$

$$C_{ijk} := -\frac{\alpha\Delta t}{(\Delta x)^2} \frac{(\gamma-1)u}{k_B} \frac{\mu_{ijk}^{(n+1)}}{\varrho_{ijk}^{(n+1)}} \kappa_{i+1,jk}^{(n+1)} \quad (6.68)$$

$$\begin{aligned} D_{ijk} := & \left[\frac{(1-\alpha)\Delta t}{(\Delta x)^2} \frac{(\gamma-1)u}{k_B} \frac{\mu_{ijk}^{(n)}}{\varrho_{ijk}^{(n)}} \kappa_{ijk}^{(n)} \right] T_{i-1,jk}^{(n)} \\ & + \left[1 - \frac{(1-\alpha)\Delta t}{(\Delta x)^2} \frac{(\gamma-1)u}{k_B} \frac{\mu_{ijk}^{(n)}}{\varrho_{ijk}^{(n)}} \left(\kappa_{i+1,jk}^{(n)} + \kappa_{ijk}^{(n)} \right) \right] T_{ijk}^{(n)} \\ & + \left[\frac{(1-\alpha)\Delta t}{(\Delta x)^2} \frac{(\gamma-1)u}{k_B} \frac{\mu_{ijk}^{(n)}}{\varrho_{ijk}^{(n)}} \kappa_{i+1,jk}^{(n)} \right] T_{i+1,jk}^{(n)} \end{aligned} \quad (6.69)$$

the set of non-linear equations

$$A_{ijk}T_{i-1,jk}^{(n+1)} + B_{ijk}T_{ijk}^{(n+1)} + C_{ijk}T_{i+1,jk}^{(n+1)} = D_{ijk} \quad (6.70)$$

is obtained. Thus, the x -sweep is obtained in matrix representation (the indices j and k are dropped for the sake of clarity):

$$\begin{pmatrix} B_1 & C_1 & 0 & 0 & 0 & 0 & \dots & 0 & 0 & 0 \\ A_2 & B_2 & C_2 & 0 & 0 & 0 & \dots & 0 & 0 & 0 \\ 0 & A_3 & B_3 & C_3 & 0 & 0 & \dots & 0 & 0 & 0 \\ 0 & 0 & A_4 & B_4 & C_4 & 0 & \dots & 0 & 0 & 0 \\ \vdots & \vdots & \vdots & \vdots & \vdots & \vdots & \ddots & \vdots & \vdots & \vdots \\ 0 & 0 & 0 & 0 & 0 & 0 & \dots & A_{n_i-1} & B_{n_i-1} & C_{n_i-1} \\ 0 & 0 & 0 & 0 & 0 & 0 & \dots & 0 & B_{n_i} & C_{n_i} \end{pmatrix} \begin{pmatrix} T_1^{(n+1)} \\ T_2^{(n+1)} \\ T_3^{(n+1)} \\ T_4^{(n+1)} \\ \vdots \\ T_{n_i-1}^{(n+1)} \\ T_{n_i}^{(n+1)} \end{pmatrix} = \begin{pmatrix} D_1 \\ D_2 \\ D_3 \\ D_4 \\ \vdots \\ D_{n_i-1} \\ D_{n_i} \end{pmatrix}. \quad (6.71)$$

In order to uniquely solve the above set of non-linear equations (6.71), the values $A_1, B_1, C_1, D_1, A_{n_i}, B_{n_i}, C_{n_i}, D_{n_i}$ need to be fixed by particularly chosen boundary conditions. In the simulations within this thesis the boundary condition $\nabla T = \vec{0}$ must be satisfied, which implies for the left boundary ($i = 1$)

$$T_{1jk}^{(n+1)} = T_{2jk}^{(n+1)}. \quad (6.72)$$

The coefficients are determined from equation (6.70) by respecting the boundary condition (6.72):

$$A_{1jk} = 0 \quad ; \quad D_{1jk} = 0 \implies B_{1jk} = -C_{1jk}, \text{ where } B_{1jk} = 1 \quad ; \quad C_{1jk} = -1. \quad (6.73)$$

Equivalently, the right boundary ($i = n_i$) reads

$$T_{n_i-1,jk}^{(n+1)} = T_{n_i,jk}^{(n+1)}. \quad (6.74)$$

Therefore,

$$A_{n_i,jk} = 0 \quad ; \quad C_{n_i,jk} = 0 \quad ; \quad B_{n_i,jk} = 1 \quad ; \quad D_{n_i,jk} = T_{n_i-1,jk}^{(n+1)}. \quad (6.75)$$

The above set of non-linear equations (6.71) together with boundary conditions (6.73) and (6.75) is solved for the unknown vector $T_{ijk}^{(n+1)}$, $j, k = \text{const.}$, by means of the Thomas algorithm (Press et al., 1992). Figure 6.6 (left plot) shows the x -sweeps.

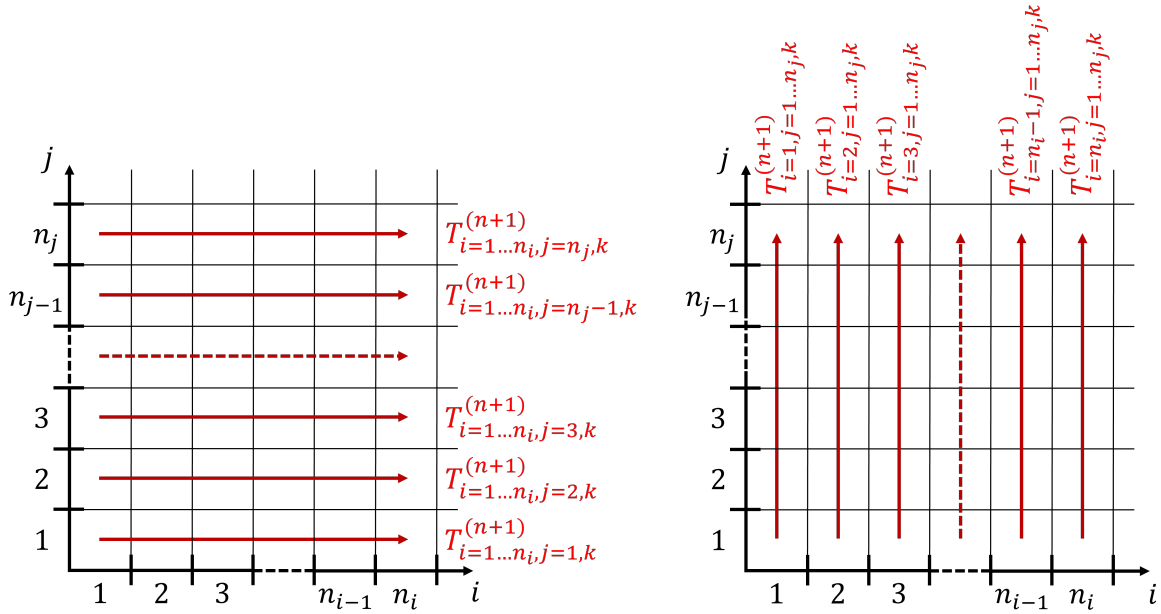


Figure 6.6: Sweeps in both x -direction (*left plot*) and y -direction (*right plot*) to solve for the new temperature distribution $T^{(n+1)}$ by means of the Crank-Nicolson discretization. The z -sweeps are executed analogously.

If for one x -sweep the new temperature distribution $T_{i=1 \dots n_i, jk}^{(n+1)}$ is found for fixed j and k by means of solving (6.71), then j is increased to $j+1$ and the x -sweep is executed again. Successively, all x -sweeps are executed for each $j = 1, \dots, n_j$. Thus, the entire (x, y) -plane for $k = 1$ is passed through. In what follows, all such (x, y) -planes for $k = 1, \dots, n_k$ are passed through, such that all x -sweeps are executed yielding

$T_{i=1\dots n_i, j=1\dots n_j, k=1\dots n_k}^{(n+1)}$. In an equivalent way, all the y -sweeps (figure 6.6, right plot) and then the z -sweeps are executed one after the other.

The temperature dependence of thermal conductivity (6.52) is too complicated to be explicitly taken into account in the discretization (6.65). Instead, an iterative procedure is used to calculate $\kappa^{(n+1)}$, which is shown in figure 6.7.

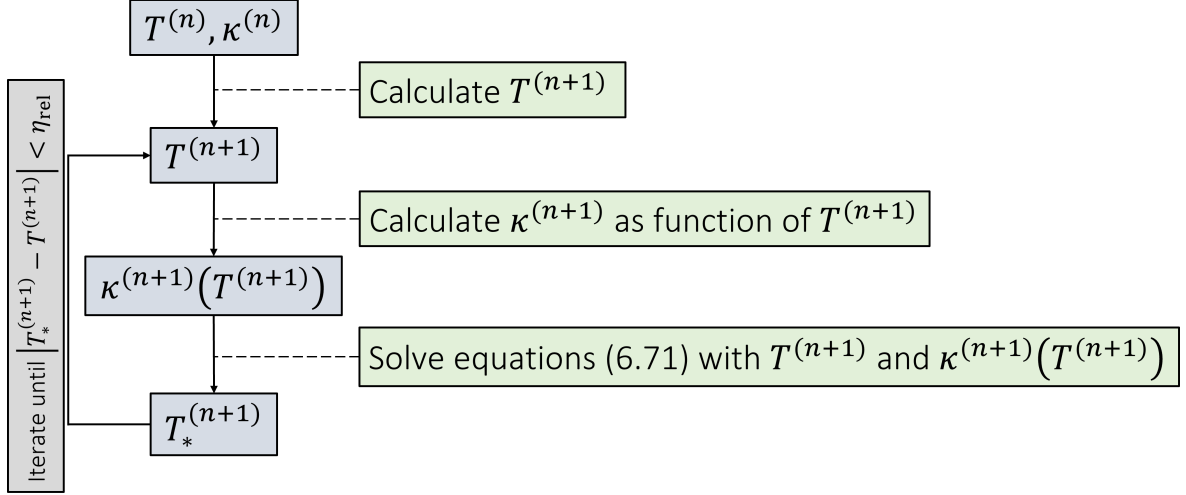


Figure 6.7: Numerical update of thermal conductivity is performed iteratively.

The iteration loops until $T_*^{(n+1)}$ converges to $T^{(n+1)}$, i.e., $\left| T_*^{(n+1)} - T^{(n+1)} \right| < \eta_{\text{rel}}$, with relative error $\eta_{\text{rel}} = 10^{-2}$. Upon convergence, the updated value $\kappa^{(n+1)}$ is determined, which in turn is used to solve equations (6.71).

6.2.5 Mass diffusion

The diffusive transport of material is described by *Fick's first law of diffusion*. The associated diffusion flux is of the same form as the heat flux (6.51). The *molar diffusion flux* reads

$$\vec{j}_{\text{mol}} = -c\mathcal{D}_{\text{AB}}\nabla n_{\text{A}}, \quad (6.76)$$

where $\mathcal{D}_{\text{AB}}$ [$\text{cm}^2 \text{s}^{-1}$] denotes the diffusivity of a gaseous species A through a stationary medium B , and ∇n_{A} [cm^{-1}] is the *number gradient* of particles belonging to species A . According to Bird et al. (1960) a Chapman-Enskog expansion yields very good results for $\mathcal{D}_{\text{AB}}$ for gases at sufficient low density. They provide

$$c\mathcal{D}_{\text{AB}} = 2.2646 \times 10^{-5} \frac{\sqrt{T \left(\frac{1}{\mathcal{M}_{\text{mol,A}}} + \frac{1}{\mathcal{M}_{\text{mol,B}}} \right)}}{\sigma_{\text{AB}}^2 \Omega_{\text{AB}}}. \quad (6.77)$$

If the molar density c [mol cm⁻³] is approximated by the equation of state for an ideal gas

$$c = \frac{P}{\mathcal{R}T}, \quad (6.78)$$

with $\mathcal{R}$ being the gas constant, then equation (6.77) turns into

$$\mathcal{D}_{AB} = 0.0018583 \frac{\sqrt{T^3 \left(\frac{1}{\mathcal{M}_{\text{mol},A}} + \frac{1}{\mathcal{M}_{\text{mol},B}} \right)}}{P \sigma_{AB}^2 \Omega_{AB}}. \quad (6.79)$$

with molar masses $\mathcal{M}_{\text{mol},A}$ and $\mathcal{M}_{\text{mol},B}$ for gaseous species A and B , pressure P , average collisional diameter $\sigma_{AB} = 0.5(\sigma_A + \sigma_B)$, and the temperature-dependent collisional integral Ω_{AB} . The values for σ_{AB} and Ω_{AB} are tabulated in Bird et al. (1960).

Equivalently, equation (6.76) can be expressed as a *mass diffusion flux*

$$\vec{j}_{\text{mass}} = -\varrho \mathcal{D}_{AB} \nabla n_A, \quad (6.80)$$

with matter density ϱ [g cm⁻³]. Both fluxes can be converted into one another by means of the molar mass $\mathcal{M}_{\text{mol}}$

$$\vec{j}_{\text{mass}} = \mathcal{M}_{\text{mol}} \vec{j}_{\text{mol}}. \quad (6.81)$$

Within this thesis, only one gaseous species is handled. Hence, equation (6.79) reduces accordingly (cf. Sander & Hensler, 2021). The mass diffusion flux (6.80) adds to equation (6.34) by

$$-\nabla \vec{j}_{\text{mass}} = \nabla (\varrho \mathcal{D}_{AB} \nabla n_A), \quad [\text{g cm}^{-3} \text{ s}^{-1}] \quad (6.82)$$

within the advection sub-steps within FLASH (cf. figure 6.3). This means that the contribution to the updated value of the local density provided by mass diffusion is given by equation (6.82).

Some authors argue that mass diffusion is negligible compared to numerical diffusion in simulations of the interstellar medium (ISM) (de Avillez & Breitschwerdt, 2004; Pan & Scannapieco, 2010).

6.2.6 Cooling and heating of plasma

Molecular, atomic, and ionized gases are subject to both cooling and heating processes, which means a loss or a gain of thermal energy through various processes (see

chapter 8 in Lequeux, 2005). The extent and efficiency of energy transfer through these mechanisms in astrophysical gas reservoirs depends on their density n , temperature T , dust content, the specific gas phase, and chemical composition (Gnedin & Hollon, 2012). Hence, in order to cool from highly ionized medium (HIM) over WNM to CNM several elements having different levels of excitation and ionization, and molecules with many rotational and vibrational transitions are needed (Dalgarno & McCray, 1972; Boehringer & Hensler, 1989; Gnat & Sternberg, 2007). Within this thesis, cooling rates are used in the form

$$\Lambda(T, n, Z) = n^2 \Lambda_0(T, Z) \quad (6.83)$$

and, equivalently, heating rates are given by

$$\Gamma(T, n, Z) = n \Gamma_0(T, Z), \quad (6.84)$$

with the cooling function $\Lambda_0(T, Z)$ and heating function $\Gamma_0(T, Z)$ depending on temperature and metallicity (Boehringer & Hensler, 1989). In what follows, the most relevant processes of plasma cooling and heating are presented that are implemented in FLASH within this thesis.

Cooling

Bremsstrahlung: At temperatures above $\sim 10^5$ K the gas is fully ionized by means of collisions (Sutherland & Dopita, 1993). Free-free transitions like bremsstrahlung are most prominent for cooling the gas (Boehringer & Hensler, 1989, and figure 6.8). Plasma cooling at temperatures $T > 10^4$ K is well reproduced by the cooling rate

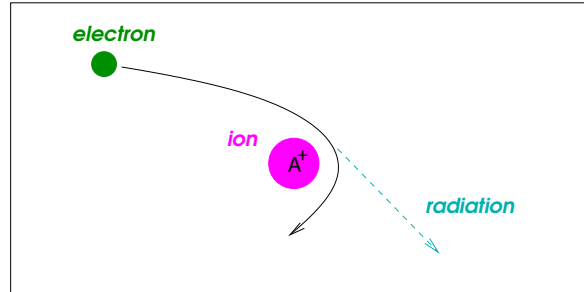


Figure 6.8: Radiative cooling by means of bremsstrahlung.

provided by Boehringer & Hensler (1989). They calculate the cooling rate for different metallicities by considering bremsstrahlung and atomic line radiation of excited atoms

and ions. Therefore,

$$\Lambda^{(\text{BH})} = n^2 \Lambda_0^{(\text{BH})} \text{ erg cm}^{-3} \text{ s}^{-1} \quad (6.85)$$

The values for the cooling function $\Lambda_0(T)$ are tabulated elsewhere. Because at $T > 10^4$ K the gas is both fully dissociated and ionized the rate (6.85) is weighted by the respective degrees y and f . The gas metallicity Z is applied to those parts of (6.85) that account for metals. Hence,

$$\Lambda^{(\text{BH})} = n^2 y f \left(\Lambda_0^{(\text{H})} + \Lambda_0^{(\text{He})} + Z \Lambda_0^{(\text{metals})} \right) \text{ erg cm}^{-3} \text{ s}^{-1}, \quad (6.86)$$

where

$$\Lambda_0^{(\text{metals})} = \Lambda_0^{(\text{C})} + \Lambda_0^{(\text{O})} + \Lambda_0^{(\text{Si})} + \Lambda_0^{(\text{S})} + \Lambda_0^{(\text{N})} + \Lambda_0^{(\text{Ne})} + \Lambda_0^{(\text{Fe})} + \Lambda_0^{(\text{Mg})}. \quad (6.87)$$

Because the above cooling rates are tabulated, a spline interpolation is used for intermediate values (Press et al., 1992).

Atomic de-excitation: Warm gas ($T \approx 10^3 \dots 10^4$ K) cools via atomic line radiation with major contribution by de-excitation of low-lying electronic states of tracer elements (mostly C II and O I, Dalgarno & McCray, 1972; Gaetz & Salpeter, 1983; Landi & Landini, 1999; Smith et al., 2008; Schure et al., 2009, and see figure 6.9, left). In such warm gas dissociation is total and ionization is partly (figure 6.4). Within this

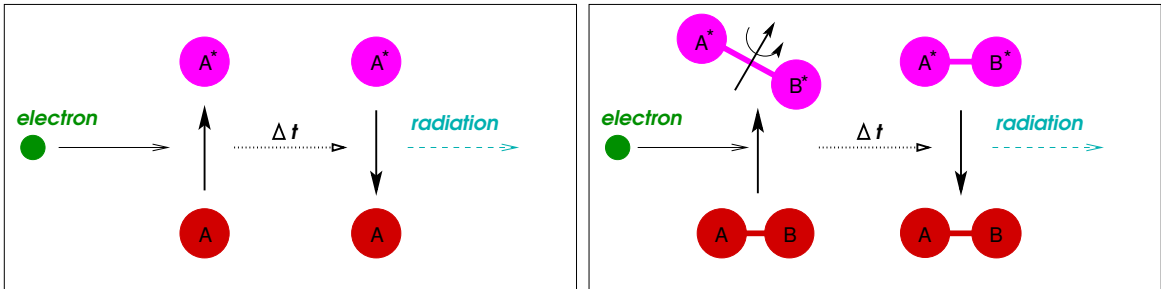


Figure 6.9: *Left plot:* The process of atomic line cooling by excitation and de-excitation of low-lying electronic states is a bound-bound transition. *Right plot:* At low temperatures molecules serve as cooling agents by radiative emission during de-excitation of higher rotational and vibrational states.

thesis, cooling in the range of warm temperatures $T \approx 900 \dots 10^4$ K is considered by the rate given by Dalgarno & McCray (1972). This rate only considers cooling from electron impact excitation and reads

$$\begin{aligned} \Lambda^{(\text{DM})} &= n^2 y (1 - f) \Lambda_0^{(\text{DM})} \text{ erg cm}^{-3} \text{ s}^{-1} \\ &= n^2 y (1 - f) \left(\Lambda_0^{(\text{H})} + Z \Lambda_0^{(\text{metals})} \right) \text{ erg cm}^{-3} \text{ s}^{-1}, \end{aligned} \quad (6.88)$$

with metallicity operating on the respective metal lines only.

Molecular de-excitation: Below temperatures of $\sim 10^3$ K molecules start to form. Radiative de-excitation of molecular rotational-vibrational states provides an efficient way of cooling (Goldsmith & Langer, 1978; Falgarone & Puget, 1985, cf. figure 6.9, right). In this coolest phase no gas is ionized and only molecules with lowest binding energies are dissociated. The function for molecular radiative cooling is provided by Falgarone & Puget (1985) and reads

$$\Lambda^{(\text{mol})} = n^2 \Lambda_0^{(\text{mol})} = 6.0 \times 10^{-29} T^{2.5} n^2 \left[n^{1.35} + 1.15 \times 10^3 \frac{T^{0.67}}{f_{\text{CO}}} \right]^{-1} \text{ erg cm}^{-3} \text{ s}^{-1}, \quad (6.89)$$

where $f_{\text{CO}} = n_{\text{CO}}/n_{\text{C}}$ represents the fraction of carbon nuclei in the gas phase in molecular form within dense molecular clouds and has an adopted value of $f_{\text{CO}} > 0.99$. The function for fine-structure line cooling is given by the analytical form

$$\Lambda^{(\text{C II})} = n_{\text{C II}} [0.5 n_{\text{H}_2} + n_{\text{H I}}] \times \begin{cases} 2.15 \times 10^{-37} T^{6.7} \text{ erg cm}^{-3} \text{ s}^{-1} & 10 \text{ K} < T \leq 20 \text{ K} \\ 7 \times 10^{-34} T^4 \text{ erg cm}^{-3} \text{ s}^{-1} & 20 \text{ K} < T \leq 30 \text{ K} \\ 8.9 \times 10^{-31} T^{1.9} \text{ erg cm}^{-3} \text{ s}^{-1} & 30 \text{ K} < T \leq 100 \text{ K} \end{cases}, \quad (6.90)$$

where $n_{\text{H}_2} = (1 - y)n_{\text{H I}}$ and $n_{\text{C II}} = f n_{\text{C I}}$. For the molecular cooling rates (6.89) and (6.90) the fractions of dissociation and ionization are considered. Thus,

$$\Lambda^{(\text{FP})} = \begin{cases} (1 - y)(1 - f) [\Lambda^{(\text{mol})} + \Lambda^{(\text{C II})}] & , T \leq 100 \text{ K} \\ (1 - y)(1 - f) \Lambda^{(\text{mol})} & , T > 100 \text{ K} \end{cases}. \quad (6.91)$$

The total cooling rate covering the entire temperature range in the simulations analyzed in this thesis thus sums up the contributions from (6.86), (6.88), and (6.91) and reads

$$\Lambda_{\text{total}} = \Lambda^{(\text{BH})} + \Lambda^{(\text{DM})} + \Lambda^{(\text{FP})}. \quad (6.92)$$

Heating

Photoelectric effect: At temperatures below 10^4 K an efficient way to heat up gas is given by the photoelectric effect on both dust grains and large polycyclic aromatic hydrocarbon (PAH) molecules (Lequeux, 2005). If the photons that are absorbed by PAHs or dust are energetic enough, an electron is released and adds to the kinetic energy of the surrounding gas by the amount of its excess energy (figure 6.10). For PAHs, the excess energy is the photon energy $h\nu$ reduced by the ionization potential.

For dust, the excess energy is given by $h\nu$ lowered by the work function and, if the grain is positively charged, Coulomb potential. Large molecules like PAHs are much more efficient in heating (Tielens, 2005). For a given gas density n the heating rate $\Gamma^{(\text{pe})}$ combined for PAHs and small dust grains reads (Wolfire et al., 1995; Tielens, 2005)

$$\Gamma^{(\text{pe})} = n\Gamma_0^{(\text{pe})} = 10^{-24} (n_{\text{H I}} + 2n_{\text{H}_2}) \epsilon G_0 \text{ erg cm}^{-3} \text{ s}^{-1}. \quad (6.93)$$

The incident far-ultraviolet (UV) field in terms of the local interstellar value (Habing, 1968) yields $G_0 = 1.7$ (Draine, 1978). The fraction of far-UV radiation absorbed by grains is approximated by the heating efficiency (Bakes & Tielens, 1994)

$$\epsilon = \frac{4.9 \times 10^{-2}}{1.0 + 4.0 \times 10^{-3} \left(\frac{G_0 \sqrt{T}}{n_e} \right)^{0.73}} + \frac{3.7 \times 10^{-2} (T/10^4)^{0.7}}{1.0 + \frac{2.0 \times 10^{-4} G_0 \sqrt{T}}{n_e}}, \quad (6.94)$$

with electron density n_e in ionization equilibrium.

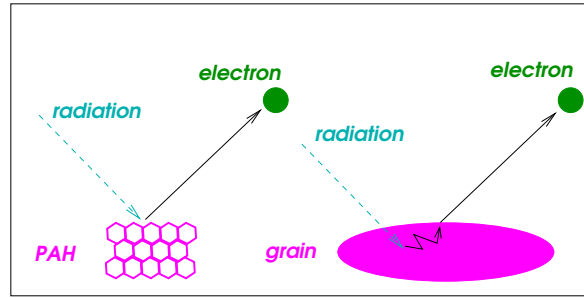


Figure 6.10: Gas heating by the photoelectric effect on PAH molecules and dust grains is most efficient at low temperatures.

Formation of H_2 : The most important chemical reaction that releases some kinetic energy to heat the gas is the formation of molecular hydrogen H_2 on the surfaces of dust grains. Two hydrogen atoms stick onto a dust grain, encounter one another and form an H_2 molecule. The excess energy from this reaction is successively transferred non-radiatively to phonons in the grain, to excitation of the H_2 molecule, and some fraction ϵ_{chem} is available for gas heating. Jura (1974) suggests $\epsilon_{\text{chem}} = 0.04$. Given the formation rate of H_2 on grain surfaces (Lequeux, 2005)

$$R^{(\text{H}_2)} = 3 \times 10^{-17} n_{\text{H I}} (n_{\text{H I}} + 2n_{\text{H}_2}) \text{ cm}^3 \text{ s}^{-1}, \quad (6.95)$$

the heating rate from the formation of H_2 is (Tielens, 2005)

$$\Gamma^{(\text{H}_2)} = n\Gamma_0^{(\text{H}_2)} = 3 \times 10^{-11} R^{(\text{H}_2)} \epsilon_{\text{chem}} = 3.6 \times 10^{-29} n_{\text{HI}} (n_{\text{HI}} + 2n_{\text{H}_2}) \text{ erg cm}^{-3} \text{ s}^{-1}. \quad (6.96)$$

Dust is needed as catalyst for this reaction, because the H_2 molecule is unable to radiate away its excitation energy so that it cannot form in the gas phase at interstellar densities (see chapter 9 in Lequeux, 2005). Dust is confirmed observationally in Complex C (Williams et al., 2012) and other high-velocity clouds (HVCs) (see section 3.3.1). Hence, at least in some HVCs the photoelectric effect (6.93) and H_2 formation (6.96) may contribute to gas heating essentially in their cores (because of sufficient gas density for the required shielding of H_2 and dust against UV radiation).

Thermalization: Incompressible fluids tend to transition from smooth laminar flow to turbulent flow when the ratio of inertial forces to viscous forces applies

$$Re = \frac{vl}{\nu} \gtrsim 3000. \quad (6.97)$$

The ratio is called the Reynolds number of the flow with mean velocity v , typical length scale l , and kinematic viscosity ν . Turbulent flows consist of eddies of descending sizes. The energy input to the flow at the largest scale cascades into successively smaller scale eddies all the way down to the molecular dissipation level where the energy is finally converted into heat. Tielens (2005) provides a heating rate of

$$\Gamma^{(\text{turb})} = n\Gamma_0^{(\text{turb})} = \frac{0.5uv^2}{l/v} n \quad (6.98)$$

by thermalization of turbulent motions (u is the atomic mass unit). Considering some typical values for l and v in the multiple phases of the CGM (Ferrière, 2020) the respective heating rate in cool cores, warm envelopes, and hot interface layers of HVCs can be estimated. So,

$$\Gamma^{(\text{turb})} = \begin{cases} 7.3 \times 10^{-28} n \text{ erg cm}^{-3} \text{ s}^{-1} & , \text{CNM } (v \sim 3 \text{ km s}^{-1}, l \sim 10 \text{ pc}) \\ 9.0 \times 10^{-27} n \text{ erg cm}^{-3} \text{ s}^{-1} & , \text{WNM } (v \sim 10 \text{ km s}^{-1}, l \sim 30 \text{ pc}) \\ 7.3 \times 10^{-26} n \text{ erg cm}^{-3} \text{ s}^{-1} & , \text{HIM } (v \sim 30 \text{ km s}^{-1}, l \sim 100 \text{ pc}) \end{cases} \quad (6.99)$$

Photoionization: UV photons with $\lambda > 912 \text{ \AA}$ may photoionize atoms and molecules. The heating contribution is provided by the released electrons, whose kinetic energy equals the difference of photon energy and ionization energy (figure 6.11, left). Photoionization dominates gas heating in H II regions. In the WNM the main source of

photoelectrons is carbon with an ionization potential of 11.2 eV. Because no photon with an energy larger than 13.6 eV exists, the mean photoelectron energy is $(13.6 - 11.2)/2 \text{ eV} \sim 1 \text{ eV}$ (Lequeux, 2005). This is much less than for hydrogen ionization. Lequeux (2005) obtains for the heating rate in the solar neighbourhood

$$\Gamma^{(\text{C I})} = 2 \times 10^{-22} n_{\text{C I}} \text{ erg cm}^{-3} \text{ s}^{-1}. \quad (6.100)$$

The number density of neutral carbon $n_{\text{C I}}$ follows from ionization equilibrium. Hence,

$$n_{\text{C I}} \chi = n_{\text{C II}} n_e \xi, \quad (6.101)$$

where n_e is the electron density, χ and ξ are the probabilities for ionization and recombination per second, respectively, with $\xi \sim T^{-0.7}$. For the CNM one may adopt $n_{\text{C I}}/n_{\text{C II}} = 0.0002$, $n_e = 0.0075 \text{ cm}^{-3}$, thus $n_{\text{C I}} = 2 \times 10^{-6} \text{ cm}^{-3}$ (Lequeux, 2005). So, from (6.100)

$$\Gamma_{\text{CNM}}^{(\text{C I})} = 4 \times 10^{-28} \text{ erg cm}^{-3} \text{ s}^{-1}. \quad (6.102)$$

Ionization: Cosmic rays transfer energy to the bound electrons of atoms and molecules of the gas by ionization and excitation, and as well to free electrons via Coulomb interactions and excitation of plasma waves (Lequeux, 2005). Regarding Black (1987) the heating rate due to cosmic rays in a molecular cloud is

$$\begin{aligned} \Gamma^{(\text{CR,mol})} &= 4 \times 10^{-28} \left(\frac{\zeta_p}{4 \times 10^{-17}} \right) \left(\frac{\epsilon_h}{6 \text{ eV}} \right) n_{\text{H}_2} \text{ erg cm}^{-3} \text{ s}^{-1} \\ &= 1 \times 10^{-11} \zeta_p n_{\text{H}_2} \text{ erg cm}^{-3} \text{ s}^{-1}, \end{aligned} \quad (6.103)$$

where $\zeta_p = 4 \times 10^{-17} \text{ s}^{-1}$ is the primary ionization rate and $\epsilon_h = 5.7 \dots 7.3$ is the mean heat input per primary ionization including secondary processes. So, $\epsilon_h/6 \text{ eV} \approx 1$. The corresponding heating rate in an atomic cloud of low ionization is

$$\begin{aligned} \Gamma^{(\text{CR,atm})} &= 3.5 \times 10^{-28} \left(\frac{\zeta_p}{4 \times 10^{-17}} \right) n_{\text{H I}} \text{ erg cm}^{-3} \text{ s}^{-1} \\ &= 8.75 \times 10^{-12} \zeta_p n_{\text{H I}} \text{ erg cm}^{-3} \text{ s}^{-1} \end{aligned} \quad (6.104)$$

and the heating rate in a highly ionized gas is

$$\begin{aligned} \Gamma^{(\text{CR,ion})} &= [2 \times 10^{-27}(1-f) + 2 \times 10^{-26}f] n \left(\frac{\zeta_p}{4 \times 10^{-17}} \right) \text{ erg cm}^{-3} \text{ s}^{-1} \\ &= [5 \times 10^{-11}(1-f) + 5 \times 10^{-10}f] n \zeta_p \text{ erg cm}^{-3} \text{ s}^{-1}, \end{aligned} \quad (6.105)$$

where $n = n_{\text{HI}} + 2n_{\text{H}_2} + n_{\text{HII}}$. Therefore, the total heating rate due to cosmic rays is

$$\Gamma^{(\text{CR})} = (1 - f)(1 - y)\Gamma^{(\text{CR},\text{mol})} + (1 - f)\Gamma^{(\text{CR},\text{atm})} + f\Gamma^{(\text{CR},\text{ion})} \quad (6.106)$$

The X-ray photons release a primary photoelectron, which is sufficiently energetic to collisionally ionize another atom and release a second photoelectron (figure 6.11, right). This process dominates heating for $T > 10^4$ K. Atomic clouds are nearly optically thick to soft X-rays as they are absorbed quickly in the outskirts (Wolfire et al., 1995). Scheffler & Elsässer (1992) determine the heating rate by X-rays to

$$\Gamma^{(\text{X})} = (1 - f)\Gamma^{(\text{CR},\text{atm})} + f \left[5 \times 10^{-11} (1 - f) \zeta_p n_{\text{HI}} \right] \text{ erg cm}^{-3} \text{ s}^{-1}. \quad (6.107)$$

The total heating rate within the simulations in question comprises contributions from (6.93), (6.96), (6.99), (6.100), (6.106), and (6.107) therefore yields

$$\Gamma_{\text{total}} = \Gamma^{(\text{pe})} + \Gamma^{(\text{H}_2)} + \Gamma^{(\text{turb})} + \Gamma^{(\text{CI})} + \Gamma^{(\text{CR})} + \Gamma^{(\text{X})}. \quad (6.108)$$

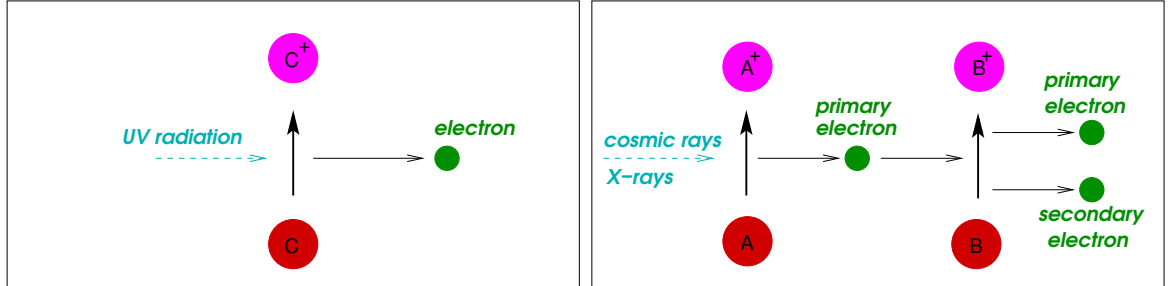


Figure 6.11: *Left plot:* Gas heating by UV radiation. Below 10^4 K mainly carbon is ionized. *Right plot:* Heating by cosmic rays and X-rays. Above 10^4 K a secondary ionization by the primary electron is very likely.

In figure 6.12 (upper two plots) the cooling and heating functions Λ_0 (6.92) and Γ_0 (6.108) are shown to evaluate their respective significance for different values of temperature. The effect of metallicity on cooling and heating efficiencies is exemplarily shown for three different values. The cooling function increases strikingly at $T = 10^4$ K and has local maxima at $\sim 2 \times 10^5$ K and 10^6 K. Here the cooling is enhanced by emission lines of H, C, and He, and ions of O and Fe.

The major contribution to the heating rate at $T < 10^4$ K is made by the photoelectric effect. If density is low enough to not shield from cosmic rays and X-rays their ionizing effects contribute as well remarkably to the heating rate. Above 10^4 K cosmic rays and X-rays are major in heating the gas.

The lower plot in figure 6.12 illustrates the combinations of temperature and density for which cooling and heating are in equilibrium. For $T > 10^6$ K the cooling time is several 10^8 years and therefore much longer than any other relevant timescales in the CGM.

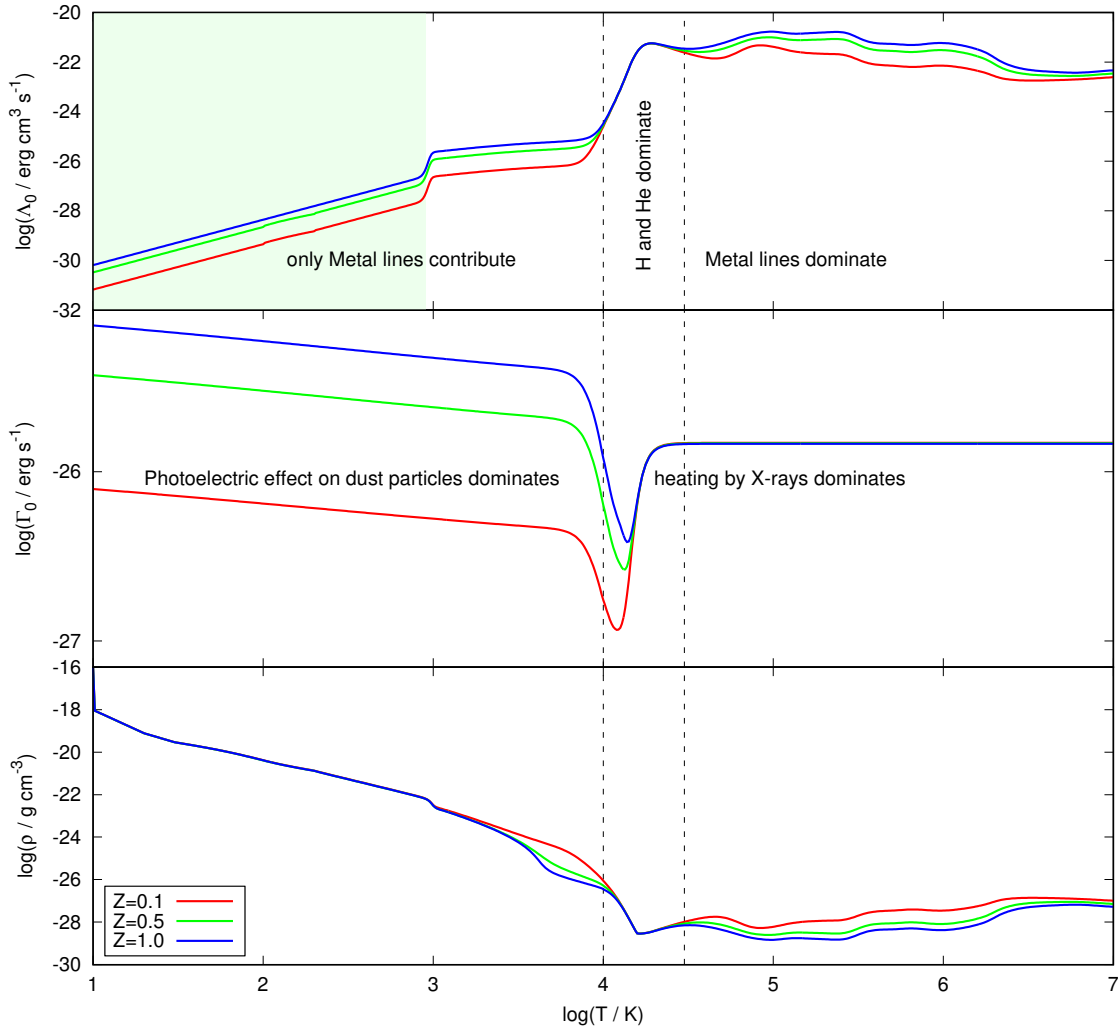


Figure 6.12: Cooling (*upper plot*) and heating (*middle plot*) in the CGM both depend on metallicity. The differences between the three exemplary metallicities ($Z = 0.1, 0.5, 1.0 Z_\odot$) is the larger the more dominant metal lines are in the respective temperature range. The differences vanish if only H and He contribute. The *green shaded area* shows the range of molecular cooling. The curve in the *lower plot* shows the pairs of temperature and density for mutually balanced energy input by heating and energy loss by cooling.

Within FLASH the various cooling and heating rates are added as a source term on the right hand side of equation (6.36), by this changing the internal energy of one grid

cell during the advection sub-steps (figure 6.3) by

$$\dot{e}_{\text{in}} \equiv \frac{\partial e_{\text{in}}}{\partial t} = \Gamma_{\text{total}} - \Lambda_{\text{total}}, \quad [\text{erg cm}^{-3} \text{ s}^{-1}]. \quad (6.109)$$

In the latter equation, Λ_{total} and Γ_{total} are given by (6.92) and (6.108), respectively. Regarding (6.109), the timescale for the net input of thermal energy from heating and cooling is given by

$$\Delta t_{\text{cool}} = \frac{\varrho e_{\text{in}}}{\dot{e}_{\text{in}}} = \frac{\varrho e_{\text{in}}}{|\Gamma_{\text{total}} - \Lambda_{\text{total}}| + \epsilon}, \quad (6.110)$$

where ϵ is a very tiny number (in **Fortran**: $\epsilon = \text{TINY}(1.0)$) to avoid division by zero for the case of thermal equilibrium. In regions with $T > 10^4 \text{ K}$ or of high density cooling can greatly excess heating thus $\Delta t_{\text{cool}} < \Delta t_{\text{eff}}$. As FLASH choses the global minimum of all timesteps for integration, only one such a region within the computational domain can reduce Δt_{int} . Thus, a bisection procedure is implemented, which converges the current temperature to the new value. By following the steps provided in Engeln-Müllges et al. (2011) the new temperature distribution $T^{(n+1)}$ is calculated and the internal energy e_{in} is updated accordingly. A loop over all grid cells in all blocks is performed and per grid cell (i, j, k) the following procedure is executed:

- (i) Take the current values of density ϱ_{ijk} , metallicity Z_{ijk} , $T_{ijk}^{(n)}$, e_{in} .
- (ii) Define $T_1 = 0.99T_{ijk}^{(n)}$, $T_2 = 1.01T_{ijk}^{(n)}$.
- (iii) Calculate $T_{\text{mid}} := T_2 + 0.5(T_1 - T_2)$ that divides the interval $[T_1, T_2]$ in half.
- (iv) Calculate the cooling rate $\Lambda_{ijk} = n_{ijk}^2 \Lambda_0(T_{ijk}^{(n)}, Z_{ijk})$ and the heating rate $\Gamma_{ijk} = n_{ijk} \Gamma_0(T_{ijk}^{(n)}, Z_{ijk})$.
- (v) Calculate the functions

$$\begin{aligned} f_1(e_{\text{in}}) &= e_{\text{in}}(T_1) - \left[e_{\text{in}}(T_{ijk}^{(n)}) + \frac{\Gamma_{ijk} - \Lambda_{ijk}}{\varrho_{ijk}} \Delta t_{\text{int}} \right], \\ f_{\text{mid}}(e_{\text{in}}) &= e_{\text{in}}(T_{\text{mid}}) - \left[e_{\text{in}}(T_{ijk}^{(n)}) + \frac{\Gamma_{ijk} - \Lambda_{ijk}}{\varrho_{ijk}} \Delta t_{\text{int}} \right], \\ f_2(e_{\text{in}}) &= e_{\text{in}}(T_2) - \left[e_{\text{in}}(T_{ijk}^{(n)}) + \frac{\Gamma_{ijk} - \Lambda_{ijk}}{\varrho_{ijk}} \Delta t_{\text{int}} \right]. \end{aligned}$$

- (vi) If $f_{\text{mid}} = 0$, then set $T_{ijk}^{(n+1)} = T_{\text{mid}}$ and the iteration is stopped.
- (vii) If $f_2 f_{\text{mid}} < 0$, then $T_{ijk}^{(n+1)}$ is in between T_{mid} and T_2 and then set $T_1 = T_2$, $T_2 = T_{\text{mid}}$, $f_1 = f_2$, $f_2 = f_{\text{mid}}$. If, oppositely, $f_2 f_{\text{mid}} > 0$, then $T_{ijk}^{(n+1)}$ is in between T_1 and T_{mid} and then set $T_2 = T_{\text{mid}}$, $f_2 = f_{\text{mid}}$. In both cases $T_1 < T_{ijk}^{(n+1)} < T_2$.

- (viii) If $|T_1 - T_2| > \eta_{\text{rel}} T_2$, then repeat the iteration starting at step (iii). If $|T_1 - T_2| \leq \eta_{\text{rel}} T_2$, then set $T_{ijk}^{(n+1)} = T_2$ for $|f_2| \leq |f_1|$, and set $T_{ijk}^{(n+1)} = T_1$ else.

The particle density n_{ijk} needed in step (iv) is calculated from matter density ρ_{ijk} by means of the mean molecular weight μ_{ijk} (see section 6.2.1). The relative error in step (viii) is chosen to be $\eta_{\text{rel}} = 10^{-2}$. The idea behind the procedure above is the following: in each grid cell (i, j, k) the current inner energy $e_{\text{in}}(T_{ijk}^{(n)})$ is updated by heating and cooling. If $\Lambda_{ijk} > \Gamma_{ijk}$ ($\Lambda_{ijk} < \Gamma_{ijk}$), then the updated inner energy is increased (decreased). The functions f_1 , f_{mid} , f_2 defined in step (v) are the differences between the energy at chosen temperatures T_1 , T_{mid} , T_2 and the updated inner energy as modified by heating and cooling. This difference must approach zero in order to converge the current to the updated energy. The steps (vi) to (viii) check whether either f_1 , f_{mid} , or f_2 approaches zero first. This identifies either $e_{\text{in}}(T_1)$, $e_{\text{in}}(T_{\text{mid}})$, or $e_{\text{in}}(T_2)$ to be the updated inner energy in cell (i, j, k) . Thus, the updated temperature $T_{ijk}^{(n+1)}$ is given by either T_1 , T_{mid} , or T_2 . The new value $T_{ijk}^{(n+1)}$ is thus used to update $e_{\text{in}}(T_{ijk}^{(n)})$ by heating and cooling via the equation of state at an integration timestep $\Delta t_{\text{int}} = \Delta t_{\text{eff}}$ (cf. equation (6.39)).

6.2.7 Self-gravity

By solving Poisson's equation

$$\nabla^2 \Phi(\vec{r}) = 4\pi G \rho(\vec{r}) \quad (6.111)$$

one yields the local gravitational potential Φ at point $\vec{r}$ for the respective mass distribution ρ at that point. The local gravitational acceleration can then be obtained by

$$\vec{g}(\vec{r}) = -\nabla \Phi(\vec{r}), \quad (6.112)$$

which in turn determines the gravitational force exerted on fluid parcels.

The numerical solution of equation (6.111) is obtained in FLASH by using a multi-grid approach optimized for oct-tree adaptive meshes (Ricker, 2008), which is briefly described here. The second-order derivative of the potential $\Phi_{b,l}^{(n+1)}(i, j, k)$ at timestep $n+1$ in grid cell (i, j, k) with dimensions $(\Delta x, \Delta y, \Delta z)$ in block b at refinement level l is directly calculated by means of a discretized Laplacian operator (equation 4 in Ricker,

2008)

$$\begin{aligned} \nabla^2 \Phi_{b,l}^{(n+1)}(i, j, k) = & \frac{1}{(\Delta x)^2} \left[\Phi_{b,l}^{(n)}(i+1, j, k) - 2\Phi_{b,l}^{(n)}(i, j, k) + \Phi_{b,l}^{(n)}(i-1, j, k) \right] \\ & + \frac{1}{(\Delta y)^2} \left[\Phi_{b,l}^{(n)}(i, j+1, k) - 2\Phi_{b,l}^{(n)}(i, j, k) + \Phi_{b,l}^{(n)}(i, j-1, k) \right] \\ & + \frac{1}{(\Delta z)^2} \left[\Phi_{b,l}^{(n)}(i, j, k+1) - 2\Phi_{b,l}^{(n)}(i, j, k) + \Phi_{b,l}^{(n)}(i, j, k-1) \right] \end{aligned} \quad (6.113)$$

yielding the numerical solution of the left-hand side of equation (6.111). The potentials on the right-hand side of equation (6.113) are given by the densities ϱ in the respective grid cells. The left-hand side of equation (6.113) is restricted to interior block cells (cf. figure 6.1), but guard cells are needed around the computational domain and around every block to solve for $\nabla^2 \Phi_{b,l}^{(n+1)}(i, j, k)$ also in the grid cells at the boundary of a block. The guard cells communicate with the grid thus keeping the densities of adjacent blocks. The multigrid solver implemented in FLASH applies to arbitrary mass distributions $\varrho(\vec{r})$.

By substituting equation (6.113) into equation (6.112) finally yields the gravitational acceleration $g_{b,l}^{(n+1)}(i, j, k)$ in each grid cell, which couples to hydrodynamics via the 6.35. Thus, the velocity in an arbitrary grid cell is modified by the self-gravitating forces of all other grid cells. The gravitational potential is updated at the end of each of the two evolution loops integrated over $\Delta t_{\text{int}}/2$ each (cf. figure 6.3).

6.2.8 Model simplifications

Here, the simplifications applied in the simulations within this thesis are briefly listed. They are detailed in Sander & Hensler (2021) and Sander & Hensler (2023, section 3.5). Thus,

- magnetic fields are neglected,
- stellar feedback is neglected,
- radiative transport is not computed in determining the fractions of ionization and dissociation,
- particle densities for the different species H_2 , H I , H II , He , electrons, and metals are calculated semi-analytically by means of local fractions of dissociation and ionization,
- the above different species are handled by a single-fluid approach,

- the Galactic hot halo has constant density and temperature,
- dust models (e.g. Draine, 2003; Zubko et al., 2004) are not handled explicitly, but are considered implicitly by the heating mechanisms of photoelectric effect on dust particles and condensation of molecular hydrogen on dust particles.

6.3 SLOPE LIMITER

Within the simulations conducted in this theses large and steep temperature gradients appear at the interfaces between clouds and ambient medium. Nearby such steep gradients spurious oscillations of the numerical solution compared to the analytical solution may arise (figure 6.13).

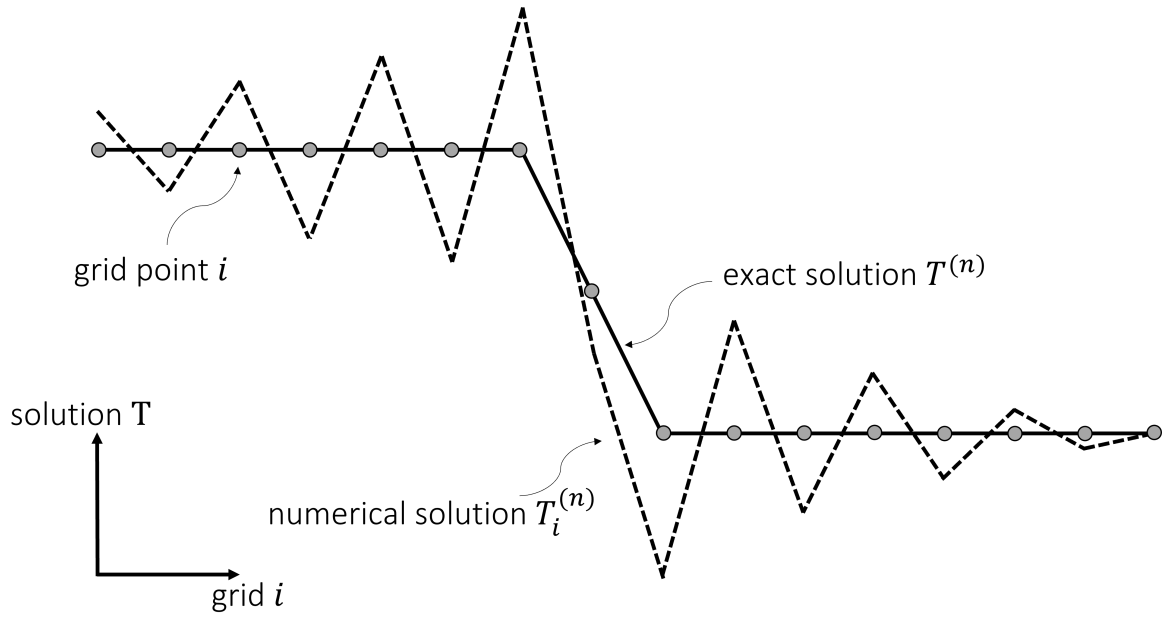


Figure 6.13: Spurious oscillations of the numerical solution T_i (*dashed line*) in the vicinity of a steep temperature gradient at a certain timestep n . Convergence to the exact solution $T^{(n)}$ (*solid line*) is not reached near the gradient. [adapted from Toro (2009)]

In such cases, numerical solutions do not converge to the analytical, exact solutions. Convergence of a function $u^{(n)}$ defined on a grid is given if its total variation (Toro, 2009)

$$TV(u^{(n)}) = \sum_{i=-\infty}^{\infty} |u_{i+1}^{(n)} - u_i^{(n)}| \quad (6.114)$$

is bounded. That is,

$$TV(u^{(n+1)}) \leq TV(u^{(n)}) \quad (6.115)$$

for all timesteps $n+1 > n$. In practice, such a desired property (6.115) of the grid-based numerical solution $u^{(n)}$ is obtained by *methods of total variation diminishing* (TVD). One way to construct TVD methods is provided by the flux-limiter (or: slope-limiter) approach (Sweby, 1984). Within this thesis the numerical solution in question is the temperature, hence $u_i = T_i$ in equation (6.115). To diminish the variations in T_i , the MINMOD slope limiter

$$\phi_{\text{MM}}(r_{i+\frac{1}{2}}) = \max \left\{ 0, \min \left\{ 1, r_{i+\frac{1}{2}} \right\} \right\} \quad (6.116)$$

is used (sketched in figure 6.14). The function

$$r_{i+\frac{1}{2}} = \frac{T_i - T_{i-1}}{T_{i+1} - T_i + \epsilon} \quad (6.117)$$

denotes the ratio of upwind change in temperature, $\Delta T_{i-\frac{1}{2}} = T_i - T_{i-1}$, to local change in temperature, $\Delta T_{i+\frac{1}{2}} = T_{i+1} - T_i$, for a positive wave-propagation speed. The safety factor ϵ has the same meaning as in equation (6.110). The function (6.117) is applied to the temperature gradient in the heat flux with conductivity (6.52) yielding

$$\vec{q}_{\text{eff}} \rightarrow \vec{q}_{\text{eff}}^{(\text{TVD})} = -\kappa_{\text{eff}} (\phi_{\text{MM}} \nabla T). \quad (6.118)$$

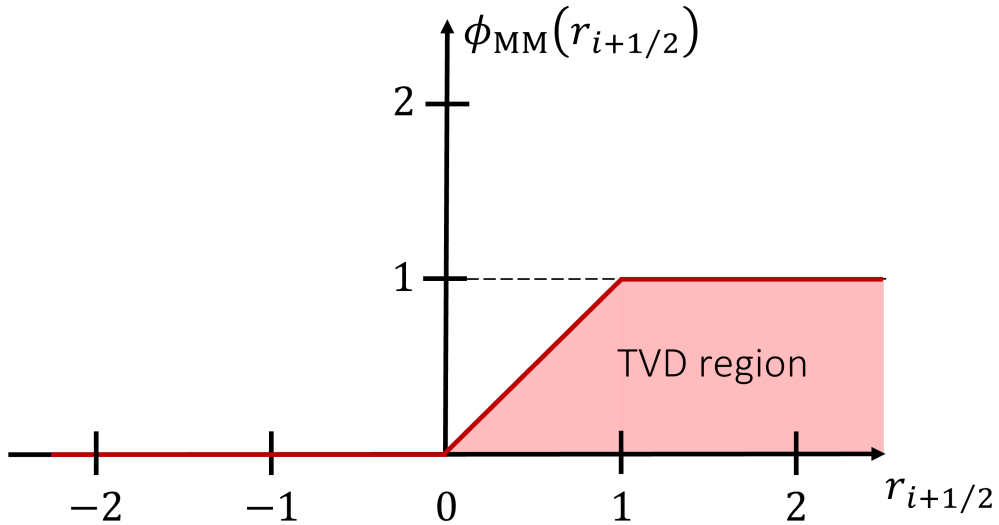


Figure 6.14: MINMOD slope-limiting function (6.116) as constructed from TVD region (Sweby, 1984) to ensure convergence of the numerical temperature solution to the exact one. [adapted from Toro (2009)]

7 THE SIMULATIONS

In the simulations conducted within this thesis the possibility of star formation in compact high-velocity clouds (CHVCs) is investigated. The simulated CHVCs resemble extreme candidates like HVC 289+33+251 (Brüns & Westmeier, 2004, see section 3.2.3), which are occasionally called ultra-compact high-velocity clouds (UCHVCs). When referring to the simulated model clouds only the term CHVCs is used to not confound with UCHVCs that are found in the ALFALFA Survey (Giovanelli et al., 2010) and which are different in many regards (cf. section 3.2.4).

The following chapter presents the initial conditions for the model clouds and the post-processing of the simulation data. The key results of the simulations are shown and how they relate to the aims of this thesis.

7.1 INITIAL MODELS

The initial setups of the different model clouds are part of the source step in the initialization phase of a simulation with FLASH as shown in figure 6.3.

The cloud profiles meet two initial conditions:

- (i) the cloud is gravitationally stable,
- (ii) the cloud is in equilibrium between heating and cooling.

Condition (i) above is met by hydrostatic equilibrium

$$\frac{\partial P}{\partial r}(r) = -g\rho(r) = -\frac{GM(r)}{r^2}\rho(r), \quad (7.1)$$

where $0 \leq r \leq R_{\text{cloud}}$. Hence,

$$R_{\text{cloud}} = \int_{r=0}^{R_{\text{cloud}}} dr \quad (7.2)$$

$$M_{\text{cloud}} = \int_{M=0}^{M_{\text{cloud}}} dM = \int_{r=0}^{R_{\text{cloud}}} \frac{dM(r)}{dr} dr = 4\pi \int_{r=0}^{R_{\text{cloud}}} r^2 \rho(r) dr. \quad (7.3)$$

Using the equation of state for an ideal gas

$$P = \frac{\varrho k_B T}{\mu u} \quad (7.4)$$

turns equation (7.1) into

$$\frac{GM(r)}{r^2} \varrho(r) = -\frac{\partial}{\partial r} \left[\frac{k_B T(r) \varrho(r)}{\mu(\varrho, T, Z) u} \right]. \quad (7.5)$$

Here, u is the atomic mass unit and μ is the mean molecular weight (see section 6.2.1). If the radial density profile $\varrho(r)$ is known, the radial temperature profile $T(r)$ can be uniquely determined.

Condition (ii) must be satisfied in every grid cell within the cloud. A balance between heating input Γ and cooling energy losses Λ is only fulfilled for particular (T, ϱ) -pairs that depend on metallicity (cf. figure 6.12, lower plot). Thus, the feasible values for ϱ at every cloud radius r are limited, because

$$\forall T \exists! \varrho : \Gamma = \Lambda. \quad (7.6)$$

Therefore,

$$\varrho(r) \longrightarrow \varrho(T, Z). \quad (7.7)$$

The above function $\varrho(T, Z)$ is determined by the condition $\Gamma = \Lambda$, such that the radial temperature stratification is calculated by

$$\begin{aligned} \frac{GM(r)}{r^2} \varrho(T, Z) &= -\frac{\partial}{\partial r} \left[\frac{k_B T(r) \varrho(T, Z)}{\mu(\varrho, T, Z) u} \right] \\ &\stackrel{(*)}{=} -\frac{\partial}{\partial T} \left[\frac{k_B T(r) \varrho(T, Z)}{\mu(\varrho, T, Z) u} \right] \frac{\partial T}{\partial r}(r) \\ \iff \frac{\partial T}{\partial r}(r) &= -\frac{GM(r)}{r^2} \varrho(T, Z) \left\{ \frac{\partial}{\partial T} \left[\frac{k_B T(r) \varrho(T, Z)}{\mu(\varrho, T, Z) u} \right] \right\}^{-1}. \end{aligned} \quad (7.8)$$

In the above identity $(*)$ the chain rule for derivatives is used. Equation (7.8) yields the radial temperature profile for the model clouds, such that simultaneously conditions (i) and (ii) above are satisfied. Within the initialization process a lookup table is generated, which contains the density values corresponding to temperature values $1 \dots 10^8$ K for heating-cooling balance. So, for every $T(r)$ the matching $\varrho(T, Z)$ is assigned with intermediate values being linearly approximated.

To solve equation (7.8) uniquely, the central cloud temperature $T_{\text{centre}} = T(r = 0)$ (inner boundary condition), the metallicity in the cloud, and the pressure of the

ambient highly ionized medium (HIM) P_{HIM} (outer boundary condition) need to be provided. For cloud initialization the outer boundary condition is replaced equivalently by two related boundary conditions: temperature T_{HIM} and particle density n_{HIM} . To calculate pressure $P(r)$ and internal energy $e_{\text{in}}(r)$ the equation of state (7.4) is solved for each (T, ϱ) -pair. So, the cloud is thermodynamically consistent at every radius. The boundary conditions used in the simulations are listed in table 7.2.

Equation (7.8) is integrated outwards until $P(r) \equiv P_{\text{HIM}}$. That is, pressure equilibrium between cloud and HIM defines the cloud radius R_{cloud} , which is quantified as the sum of all those radial integration steps dr in equation (7.2). Accordingly, the cloud mass M_{cloud} is calculated by summing up all the differential mass elements dM as is done by equation (7.3).

Metallicity: Metallicity plays a crucial role since both cooling and heating processes of the plasma depend on it (section 3.3.6) for all temperatures being relevant in the simulations analyzed within this thesis. According to Wakker (2001) the mean metallicity of high-velocity clouds (HVCs) is $\mathcal{M} = Z/Z_{\odot} = 0.3$, but the abundances of elements varies strongly from sub-solar to super-solar values. In order to bracket the observed values and to study distinct effects for the rates of heating and cooling, two different cloud metallicities of $\mathcal{M} = 0.1$ and $\mathcal{M} = 1.0$ are accounted for in the simulations (see table 7.2). Table 7.1 contains the corresponding mass fractions calculated with equations (6.2) to (6.4). Metallicity determines the cloud radius. As can be seen

Table 7.1: Mass fractions of hydrogen (X), helium (Y) and metals (Z) for the simulated metallicities of $\mathcal{M} = 0.1$ and $\mathcal{M} = 1.0$.

$\mathcal{M}$	$X(\mathcal{M})$	$Y(\mathcal{M})$	$Z(\mathcal{M})$
0.1	0.7449	0.2531	0.0020
1.0	0.7314	0.2485	0.0201

in figure 6.12 (lower plot) the equilibrium density is higher for $\mathcal{M} = 0.1$ in the temperature range $10^3 \dots 10^4$ K. These temperatures appear in the model clouds at larger radii. Hence, the sub-solar clouds show up with higher pressure profiles and pressure balance with the ambient hot gas is met further out. Vice versa, model clouds with solar metallicity are smaller. The metallicity is thus a boundary condition for integrating equation (7.8).

Resting clouds: Once the one-dimensional temperature profile is calculated using equation (7.8), the density profile is assigned via a lookup table, and both pressure and internal energy are calculated via the equation of state (7.4). Depending on initial

metallicity and inner and outer boundary values, four resting reference model clouds are obtained (figure 7.1): a low-mass and a massive model cloud with either $\mathcal{M} = 0.1$ or $\mathcal{M} = 1.0$. The radial profiles are mapped to the centers of the grid cells throughout the computational domain. Starting from a central cell $(i, j, k)_c$, the profiles are mapped isotropically thus yielding three-dimensional spheres constituting the model clouds. This starting grid cell $(i, j, k)_c$ determines the cloud centre and is taken to be the central cell in the computational domain. The side lengths L of the computational domain are $L = 160$ pc (low-mass models) and $L = 260$ pc (massive models). So, $(i, j, k)_c = (80, 80, 80)$ pc for low-mass model clouds and $(i, j, k)_c = (130, 130, 130)$ pc for massive model clouds. At a maximum numerical refinement of six levels a finest numerical resolution of $\Delta x = 0.63$ pc (low-mass models) and $\Delta x = 1.02$ pc (massive models) is achieved per grid cell.

Moving CHVCs: The radial profiles of state variables for the moving clouds are obtained identically to the resting clouds. The setup differs in two ways from the resting cloud scenarios: First, the CHVCs are exposed to a potential flow $\vec{u}(\vec{x})$. Second, the cloud centre is shifted off the central cell of the computational domain. The CHVCs within this thesis are intended to move subsonically. Thus, regarding Landau & Lifshitz (2013) the potential flow around the initially spherical CHVCs is given by

$$\vec{u}(\vec{x}) = \begin{cases} \vec{v}_{\text{HIM}} + \frac{R_{\text{cloud}}^3}{2|\vec{d}|^3} [\vec{v}_{\text{HIM}} - 3\hat{n} \cdot (\vec{v}_{\text{HIM}} \cdot \hat{n})] & , \quad |\vec{d}| > R_{\text{cloud}} \\ 0 & , \quad |\vec{d}| \leq R_{\text{cloud}} \end{cases} . \quad (7.9)$$

The vector $\hat{n}$ is the unit vector of $\vec{d}$, i.e., $\vec{d} = |\vec{d}| \hat{n}$. Within the simulations the CHVCs move towards the $(-x)$ -direction. So, the HIM flows into the computational domain at lower x -boundary towards the positive x -direction, such that $\vec{u} = (u_x, 0, 0)$. Because the dot product $\vec{v}_{\text{HIM}} \cdot \hat{n} = 0$ for $\vec{d} \perp \vec{u}$, the streaming velocity (7.9) is greatest at the side edges of the clouds, and smallest at their front and rear sides. This velocity differences lead to pressure differences leading to an elongation of the CHVCs perpendicular to their motion according to Bernoulli's principle (section 3.3.4).

For distances $|\vec{d}| \gg R_{\text{cloud}}$ the potential flow just converges to the unperturbed streaming velocity $u_x = v_{\text{HIM}}$. However, if $|\vec{d}| \gtrsim R_{\text{cloud}}$ the vector components of the

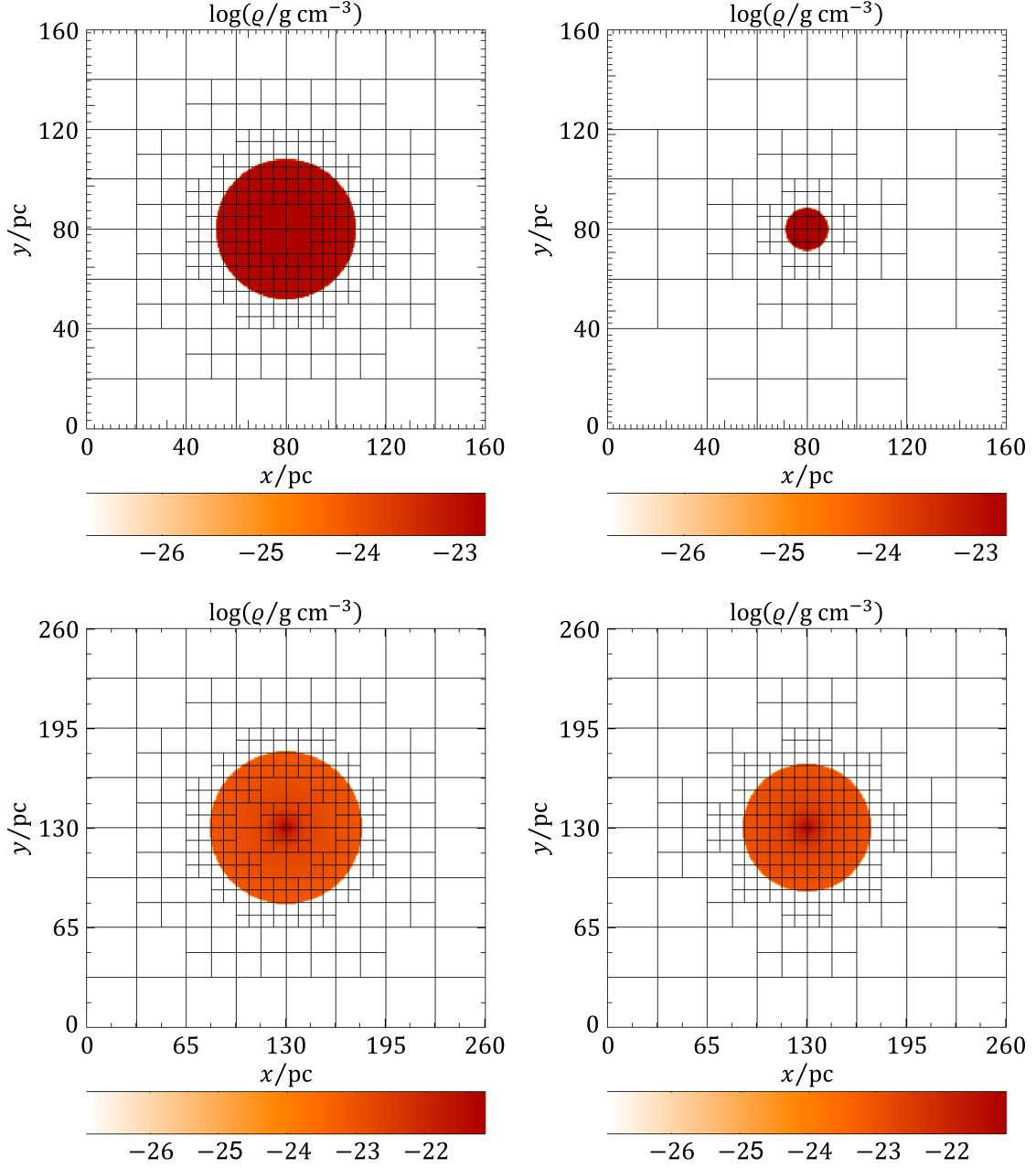


Figure 7.1: Initial density distributions of the four resting reference model clouds after being filled into the grid (*upper panel*: low-mass model clouds, *lower panel*: massive model clouds, *left column*: $Z = 0.1 Z_{\odot}$, *right column*: $Z = Z_{\odot}$). The adaptively refined blocks of the computational domain are shown up to six levels of refinement with one block consisting of $8 \times 8 \times 8$ grid cells. For the sake of clarity only two spatial dimensions of the computational domain are shown.

flow (7.9) outside the cloud are

$$\begin{aligned}
 u_x &= v_{\text{HIM}} + \frac{R_{\text{cloud}}^3}{2d_x^3} [v_{\text{HIM}} - 3v_{\text{HIM}}] = v_{\text{HIM}} - v_{\text{HIM}} \frac{R_{\text{cloud}}^3}{d_x^3}, \\
 u_y &= 0, \\
 u_z &= 0.
 \end{aligned} \tag{7.10}$$

The initial potential flow circulating around each model CHVC is illustrated in figures 5 and 6 in (Sander & Hensler, 2021).

In order to track any material stripped off the clouds as long as possible in the computational domain, the cloud centres are shifted upstream towards the lower x -boundary and the side lengths L of the computational domain are increased. For all simulated CHVCs it is thus $L = 340$ pc and $(i, j, k)_c = (100, 170, 170)$ pc. The finest numerical resolution is $\Delta x = 1.30$ pc at six levels of refinement. As is the case for the resting clouds, there are four moving reference model clouds (figure 7.2).

Galactic hot halo: Temperature and density for both scenarios of massive and low-mass clouds (table 7.2) all correspond to the super-virial phase (see section 4.2). The adapted solar metallicity is in line with solar and super-solar metallicities (e.g. Mg, C, Si) found in all three phases (Das et al., 2021). The values for T_{HIM} and n_{HIM} in the simulations are chosen to be constant, because the growth of filaments is suppressed above a critical length-scale (Nipoti & Binney, 2004; Sander & Hensler, 2021). The simulated CHVCs can be placed at $D_{\text{gal}} \sim 100$ kpc (as described in Sander & Hensler, 2021), which is roughly at $0.5R_{\text{vir}}$ of the halo of a Milky-Way like galaxy (cf. equation 1 in Fielding et al. (2020) for a total mass $\sim 10^{12} M_{\odot}$, and table 1 in Esmerian et al. (2021)).

Table 7.2: Inner and outer boundary conditions for integrating the radial temperature profile by means of equation (7.8). This yields the four reference models for resting clouds and CHVCs, respectively.

scenario	cloud values		halo values	
	$T_{\text{centre}}/\text{K}$	$Z/Z_{\odot}$	T_{HIM}/K	$n_{\text{HIM}}/\text{cm}^{-3}$
massive clouds	230	0.1	5.6×10^6	7×10^{-4}
		1.0		
low-mass clouds	1, 150	0.1	5.5×10^6	1×10^{-3}
		1.0		

Numeric boundary conditions: At all six sides of the three-dimensional computational domain the boundaries are of outflow type for advection of physical quantities in all simulations. Hence, matter and energy can both enter and leave the computational domain at all sides. The derivative of any quantity $\phi(\vec{r})$ perpendicular to a boundary $\partial\Omega$ of the computational domain is given by

$$\nabla\phi(\vec{r}) \cdot \vec{n}(\vec{r}) = \frac{\partial\phi}{\partial x}n_x + \frac{\partial\phi}{\partial y}n_y + \frac{\partial\phi}{\partial z}n_z = 0 \quad \forall \vec{r} \in \partial\Omega, \quad (7.11)$$

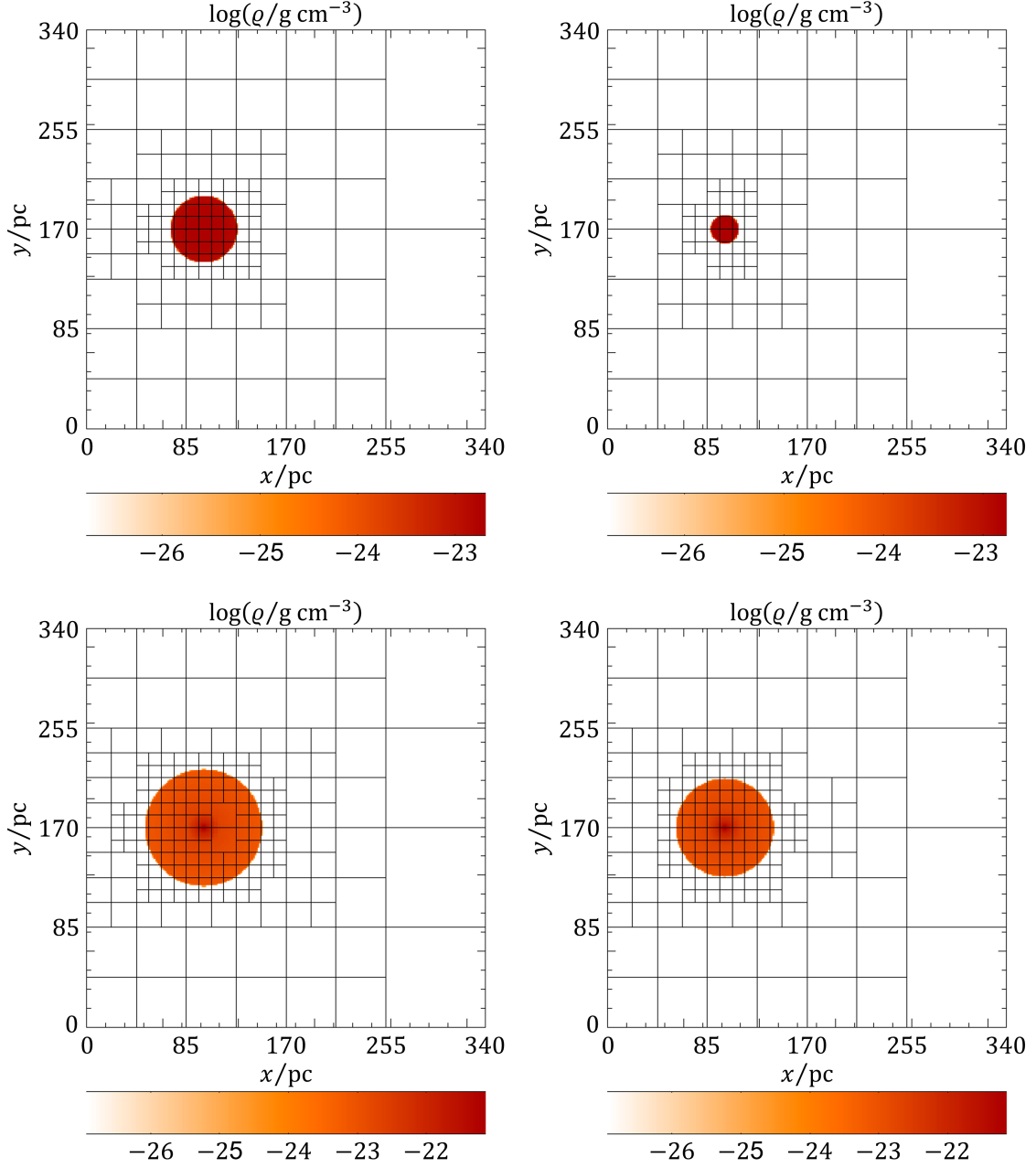


Figure 7.2: The same as in figure 7.1, but for CHVCs.

where $\vec{n} = \pm(n_x, n_y, n_z)$ is the exterior normal to the respective boundary. The boundaries are the outermost interior cells in the computational domain, which is surrounded by four layers of guard cells (see figure 6.1). By condition (7.11), the numerical solution is not extended into the wrapping guard cells (gradient of ϕ is zero). Instead, the values of $\phi(\vec{r})$ are copied from the boundary into the guard cells. The condition (7.11) is also known as von-Neumann or zero-gradient boundary condition.

While outflow boundary conditions are used for the gas, the multigrid solver imple-

mented in FLASH uses isolated boundary conditions to calculate the gravitational field Φ from Poisson's equation (Ricker, 2008, and section 6.2.7). How to obtain isolated boundary conditions by means of a multipole expansion of the gravitational potential Φ is shown in Tomida & Stone (2023, see their appendix A).

7.2 THE SIMULATION SUITE

The eight reference model clouds are diversified by different physical processes: heating and cooling processes, thermal conduction, self-gravity, different constant speeds or deceleration by drag (CHVCs only).

The subsonically streaming HIM exerts a dynamic counterforce

$$F_D = \frac{1}{2} A C_D \varrho_{\text{HIM}} v_{\text{HIM}}^2 \quad (7.12)$$

on a cloud of cross section $A = \pi R_{\text{cloud}}^2$ by transferring momentum to the cloud gas (see section 3.3.4). This drag is addressed in the simulations by an appropriate drag coefficient C_D (Henderson, 1976) that is described in Sander & Hensler (2021, see section 7.6). A Galilean transformation into the rest frame of a moving cloud is applied so that the simulated CHVCs are not pushed out of the computational domain due to drag. This step is performed in each evolution loop that FLASH runs through (see figure 6.3). First, the grid cell representing the cloud's center of mass $\vec{R}$ is identified. Thereafter, the velocity of the center of mass $\vec{V}$ is determined by figuring out the relative velocity between this particular grid cell and the ambient halo. Finally, $\vec{V}$ is subtracted from the fluid velocities in all grid cells in the computational domain. Thus, a relative motion between $\vec{R}$ and the grid is prevented.

All but one reference model clouds for CHVCs move at constant speed, which is maintained by keeping the velocity $\vec{v}_{\text{HIM}}$ constant at the inflow boundary of the computational domain. To account for a deceleration of one model CHVC the velocity decrement

$$dv = \frac{1}{2} \frac{A}{M_{\text{cloud}}} C_D \varrho_{\text{HIM}} v_{\text{HIM}}^2 \Delta t_{\text{int}} \quad (7.13)$$

caused by the force (7.12) is subtracted from $\vec{v}_{\text{HIM}}$ at inflow boundary in every timestep Δt_{int} .

The evolution of resting clouds and CHVCs is simulated for 100 Myr starting from one particular initial setup as explained in section 7.1. A plot file is output every 1 Myr. The entire suite of simulations is hence appropriate to address the key question in this

thesis:

Which physical processes suppress star formation in HVCs?

A natural explanation for prevented star formation arises from the simulations in terms of the interaction of decreasing dynamic pressure due to decelerated clouds and heat input via thermal conduction. The results are discussed in section 7.6.

In addition, some other important questions are answered.

1. Do mixing efficiencies allow for the chemical inhomogeneities the CHVCs constitute?
2. Will thermal conduction prevent from substantial mass ablation due to Kelvin-Helmholtz instability (KHI)?
3. Does thermal conduction lead to evaporation in a hot ambient medium or is HIM condensed onto the clouds?
4. Is time-independent analytical theory of thermal conduction combined with heating and cooling processes proper to describe the dynamics of interfaces between HVCs and the Galactic hot halo?
5. Do magnetic fields suppress thermally conducted energy onto clouds embedded in a highly ionized medium to a negligible level?
6. Is self-gravity pivotal even in those clouds that are below their Jeans mass?

The individual model clouds and setups are provided in Sander & Hensler (2019, three resting model clouds, table 1), Sander & Hensler (2021, 14 models of CHVCs, table 3), and Sander & Hensler (2023, four resting model clouds, table 1). The model clouds separate into massive clouds ($M \sim 6 \dots 9 \times 10^4 M_\odot$) and low-mass clouds ($M \sim 0.2 \dots 2 \times 10^4 M_\odot$). The mass ranges result from the two different initial cloud metallicities ($Z = 0.1 Z_\odot$ and $Z = Z_\odot$). The clouds thus bracket the mean values of UCHVCs observed in the GALFA-H I Survey (Saul et al., 2012) of $M \sim 2 \times 10^4 M_\odot$ and physical sizes ~ 100 pc at $D_{\text{sun}} \sim 100$ kpc.

7.3 ANALYSIS OF THE SIMULATION DATA

The raw simulation data are post-processed with the help of many IDL¹ scripts. FLASH provides an IDL script by which the adaptively gridded data are mapped to an equidistant mesh of the finest spatial resolution (6 levels of refinement) within the three-dimensional computational domain thus yielding $2^{D(3+j-1)} = 2^{24}$ fine-gridded data values. These mapped data are the basis for all subsequent analyses.

Colouring technique: FLASH offers the ability to include user-defined mass scalars. These are passive scalar quantities that do not exert forces to the dynamical system, but are advected due to hydrodynamics by solving individual equations of continuity. So,

$$\frac{\partial \varrho \Psi_j}{\partial t} + \nabla (\varrho \Psi_j \vec{v}) = 0 \quad (7.14)$$

with Ψ_j being the j -th mass scalar. The Ψ_j range between 0 and 1 and therefore represent a fraction of any quantity within a grid cell. The colouring intensity of a grid cell depends accordingly on this value and ranges from invisible ($\Psi_j = 0$) to full intensity ($\Psi_j = 1$). Within the simulations three mass scalars are used for (i) cloud gas Ψ_{cloud} , (ii) hot halo gas Ψ_{HIM} , and (iii) metal content Ψ_{metal} . The colour-intensity gradation of these three mass scalars in each grid cell makes it possible to visualize several dynamic processes during cloud evolution:

- mixing dynamics at the interface between cloud and HIM,
- spatial distribution of accreted metal-rich gas in the clouds,
- distinguish accreted ambient gas throughout the cloud and track its dynamics,
- evolution of cloud morphology during motion,
- dynamics of the metal distribution in the cloud can also be tracked if there is an initial difference between the metallicities of the cloud and the hot gas, and
- gas stripped off the cloud can be identified.

The initial values of the user-defined mass scalars within the computational domain are listed in table 7.3.

¹Interactive data language, see: <https://www.nv5geospatialsoftware.com/Products/IDL#language>

Table 7.3: User-defined mass scalars within the simulations and their initial values within a model cloud and in the ambient hot gas.

mass scalar	model cloud	hot gas
Ψ_{cloud}	1.0	0.0
Ψ_{HIM}	0.0	1.0
Ψ_{metal}	0.1 or 1.0	1.0

Cloud definition: Demarcating cloud gas from ambient HIM is pivotal for analyzing the evolution of clouds. The cloud is identified within the hot gas in two different ways. *Resting clouds* are distinguished by their density and temperature. Thus, all grid cells j belong to a resting cloud if they satisfy

$$\text{resting cloud} = \left\{ \text{all grid cells } j : \rho_j > 100\rho_{\text{HIM}} \text{ and } T_j < 5 \times 10^6 \text{ K} \right\}. \quad (7.15)$$

The above definition excludes dilute cloud gas that has already evaporated into the ambient hot gas while simultaneously including highly ionized interface gas.

Moving clouds are more tricky to distinguish from the rest of the gas within the computational domain. A grid cell j is counted in favour of a moving cloud if it is

1. gravitationally bound to the cloud and
2. spatially connected to the gravitational center of the cloud.

The above two conditions are checked in consecutive order. The first condition is examined by finding the gravitational center $\vec{R}$ of the cloud. This is simply the grid cell with the largest absolute value of gravitational potential within the computational domain. Its velocity $\dot{\vec{R}}$ is subtracted from the array storing the velocities $\dot{\vec{r}}_j$ of all grid cells. Hence,

$$\Delta v_j = \dot{\vec{R}} - \dot{\vec{r}}_j \quad (7.16)$$

provides the relative velocity between each cell j and $\vec{R}$. Consequently, the kinetic energy of grid cell j relative to the kinetic energy of the gravitational center of cloud is given by

$$\Delta e_{\text{kin}} = \frac{1}{2} m_j (\Delta v_j)^2 = \frac{1}{2} \rho_j V_j (\Delta v_j)^2 \stackrel{(*)}{=} \frac{1}{2} \rho_j V (\Delta v_j)^2. \quad (7.17)$$

The latter identity $(*)$ holds, because the data was previously mapped to an equidistant grid and thus $V_j = V = \text{const}$. The potential energy of grid cell j is calculated from

$$e_{\text{pot}} = -m_j \Phi_j = -\rho_j V_j \Phi_j = -\rho_j V \Phi_j, \quad (7.18)$$

where the gravitational potentials Φ_j of all grid cells are provided by FLASH in a

separate array. By comparing (7.17) and (7.18) a certain grid cell j is gravitationally bound to the cloud if

$$\Delta e_{\text{kin}} < |e_{\text{pot}}| \iff \Delta v_j < \sqrt{2\Phi_j}. \quad (7.19)$$

The second condition is checked for all cells identified by (7.19). So, the challenge is to figure out those gravitationally bound cells that are spatially connected to $\vec{R}$. This excludes cloud material, which is gravitationally bound, but has already been stripped off. To figure out spatially connected cells a space-filling algorithm is developed within this thesis. It loops through the following steps:

1. Start from the gravitational center $\vec{R} = (x_R, y_R, z_R)$ of the cloud and loop successively through all gravitationally bound cells in directions $(-x, y_R, z_R)$, $(+x, y_R, z_R)$, $(x_R, -y, z_R)$, and $(x_R, +y, z_R)$.
2. Check all adjacent grid cells whether they are gravitationally bound and not connected yet. Tag these cells as connected.
3. Loop through this new set of cloud cells and check adjacent grid cells again.
4. Repeat this procedure until all adjacent grid cells are either already bound and connected, or unbound. By this, all cloud cells are found, which lay within the (x, y) -plane that contains the gravitational center of the cloud.
5. Loop through this (x, y) -plane of cloud cells. Check all adjacent grid cells in directions $(x_R, y_R, -z)$ and $(x_R, y_R, +z)$ whether they are gravitationally bound and not connected yet. Tag these cells as connected.
6. Loop through this new set of cloud cells and check adjacent grid cells in all directions again.
7. Repeat this procedure until all adjacent grid cells are either already bound and connected, or unbound. By this, all cloud cells are found.
8. All identified cloud cells are stored in a separate array.

The space-filling algorithm must be extended to three dimensions as in only two dimensions filaments can be missed, which are connected to the cloud in one plane, but are not in another plane. Such a situation might appear if the cloud is irregularly shaped by ram pressure.

Collapsing regions: A region is a prospecting site of star formation if it collapses by means of gravity. This situation is fulfilled if the contained mass exceeds the Jeans mass within that region. The check for Jeans instability is applied to all cloud cells identified. So, a cloud cell j_c is gravitationally unstable if

$$m_{j_c} > M_J. \quad (7.20)$$

The Jeans mass M_J is given by equation (5.2). These unstable cloud cells are taken as the seeds of prospective collapsing regions. The procedure to identify an collapsing region is described in Sander & Hensler (2021, section 6.3).

7.4 WHY A CAREFUL MODELLING MATTERS

Within this section the effects of

- (a) self-gravity in sub-Jeans clouds, and
- (b) thermal conduction in presence of magnetic fields

are reviewed. It turns out that modelling simplifications do not prove to be necessarily accurate and can therefore lead to questionable results in numerical simulations of cool gas clouds residing in hot media.

The original idea for the manuscript was given by my supervisor Prof. Dr. G. Hensler, who also reviewed the manuscript iteratively. The manuscript is co-authored by Prof. Hensler and me.

My own contential contributions include:

- all simulations and calculations presented,
- all plots are based on self-written IDL scripts,
- line of proof in section 2.2,
- reasoning in chapter 3.

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The published paper is provided below.



Insights into the physics when modelling cold gas clouds in a hot plasma

Bastian Sander^{1,2★} and Gerhard Hensler¹

¹*Institut für Astrophysik, Universität Wien, Türkenschanzstraße 17, A-1180 Vienna, Austria*

²*Fraunhofer-Institut für Fabrikbetrieb und -automatisierung IFF, Sandtorstraße 22, D-39106 Magdeburg, Germany*

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ABSTRACT

This paper aims at studying the reliability of a few frequently raised, but not proven, arguments for the modelling of cold gas clouds embedded in or moving through a hot plasma and at sensitizing modellers to a more careful consideration of unavoidable acting physical processes and their relevance. At first, by numerical simulations we demonstrate the growing effect of self-gravity on interstellar clouds and, by this, moreover argue against their initial set-up as homogeneous. We apply the adaptive-mesh refinement code FLASH with extensions to metal-dependent radiative cooling and external heating of the gas, self-gravity, mass diffusion, and semi-analytic dissociation of molecules, and ionization of atoms. We show that the criterion of Jeans mass or Bonnor–Ebert mass, respectively, provides only a sufficient but not a necessary condition for self-gravity to be effective, because even low-mass clouds are affected on reasonable dynamical time-scales. The second part of this paper is dedicated to analytically study the reduction of heat conduction by a magnetic dipole field. We demonstrate that in this configuration, the effective heat flow, i.e. integrated over the cloud surface, is suppressed by only 32 per cent by magnetic fields in energy equipartition and still insignificantly for even higher field strengths.

Key words: conduction – diffusion – hydrodynamics – magnetic fields – methods: numerical – ISM: clouds.

1 INTRODUCTION

Today, computer simulations are a well-established approach for describing complex astrophysical systems and supplement observational and theoretical astrophysics by a third branch known as computational astrophysics. Simulations are essential for understanding systems whose evolution cannot be adequately described via analytical means. Here, galaxies are particularly well-suited objects of study, since they contain many interacting constituents. These include already existing stars and the interstellar medium (ISM), which comprises various co-existing physical states of gas (so-called ‘phases’) that can even lead to star formation. To add to the complexity, the ISM is determined by energetic processes and evolves dynamically in each of its different phases. Thus, for example, cold, interstellar (IS) gas clouds may be engulfed by hot gas – a process attracting particular research interest in the context of galactic matter cycles. When IS clouds are numerically investigated, the relevant physical processes must be modelled so as to jibe with observations, and the initial conditions must be established properly. In the literature on simulating IS clouds without star formation, there

are two striking simplifications in setting up cold gas clouds in a multiphase ISM:

(i) Self-gravity is neglected if the cloud mass is below its Jeans mass (e.g. Kwak, Shelton & Raley 2009) or the virial ratio in the cloud is much greater than one (e.g. Armillotta et al. 2017). As a consequence, IS clouds are considered single-phase, homogeneous, and isothermal. The assumption is that self-gravity has a negligible effect on the evolution of low-mass clouds.

(ii) Heat conduction is neglected if a magnetic field in the cloud is considered (e.g. Maller & Bullock 2004; Esquivel et al. 2006; Kwak et al. 2009). The assumption is that magnetic fields substantially hamper thermal conduction from the hot phase towards the cold cloud.

But these two simplifications lack a proof, or even an estimate from physical principles. In the first part of this paper (Section 2) we investigate whether self-gravity is a necessary initial condition when simulating the evolution of a low-mass cloud, i.e. one whose cloud mass is less than its Jeans mass. We argue that the Jeans criterion (or, more precisely, the Bonnor–Ebert criterion), however, is only a sufficient but not a necessary condition, because its derivation only describes the increase of perturbations. We are encouraged by studies such as that of Habe & Ohta (1992), who simulated supersonic, head-on collisions between non-identical clouds. They found that a gravitational instability in the cloud can be triggered by

★ E-mail: bastian.sander@univie.ac.at (BS), gerhard.hensler@univie.ac.at (GH)

such collisions even if the initial cloud mass is well below its Jeans mass. Furthermore, observations of real IS clouds reveal a core-halo density structure with radial profile $\propto r^{-2}$. This is a clear sign of the role played by self-gravity (Kaminski et al. 2014; Wyrowski et al. 2016), which was already deduced by Larson (1981) for virialized clouds. The observed head–tail density structure of the majority of high-velocity clouds (HVCs, Brüns et al. 2000; Ben Bekhti et al. 2006; For, Staveley-Smith & McClure-Griffiths 2013) teaches us that Rayleigh–Taylor instability is suppressed due to self-gravity (Murray et al. 1993; Hensler & Vieser 2002). We show that even a homogeneous low-mass cloud develops a radial density profile when self-gravity is considered initially.

In the second part of the paper (Section 3) we analytically address thermal conduction by electrons from a hot gas phase towards an embedded IS cloud that exhibits a strong magnetic dipole field, which is in equipartition with the energy density of the cloud. Within the model of the three-phase ISM (McKee & Ostriker 1977) cold gas clouds represent the cold neutral phase enclosed in a warm, slightly ionized phase with both being embedded in a hot phase. Since the phases are in mutual pressure equilibrium they must have different temperatures; thus, thermal conduction towards the cold clouds is a natural consequence. The hot phase is very dilute hence the mean free path of electrons (see e.g. Shu 1992) is of the order of ~ 10 pc, adopting $T \sim 10^6$ K and $n \sim 10^{-3} \text{ cm}^{-3}$. We go into the frequent argument that magnetic fields reduce the mean free path of electrons, by this, diminishing the effect of heat conducted by electrons to zero. We demonstrate that the effective heat flow, i.e. integrated over the cloud surface, is not suppressed substantially.

The studies of self-gravity and heat conduction presented in this paper are necessary to prepare simulations of cold gas clouds embedded in a hot tenuous plasma. We investigate both processes separated from each other in order to evaluate their respective effect on the evolution of cold gas clouds. A prime example of particular interest in astronomy are HVCs, which move through the hot halo of our Milky Way and contribute significantly to the Galactic gas-accretion rate (Richter 2016). HVCs have to be numerically investigated as it is not possible to disentangle the complex picture of the Galactic matter cycle from observations alone.

2 SELF-GRAVITY IN LOW-MASS CLOUDS

2.1 Initial set-ups and simulations

We carry out three-dimensional hydrodynamical simulations using the publicly available FLASH code (Fryxell et al. 2000), version 3.2¹ (see the reference paper by Dubey et al. 2009). It is an Eulerian hydrodynamics code, which integrates the (inviscid) Euler equations on a Cartesian, adaptively refined grid by means of the piecewise parabolic method (PPM, Colella & Woodward 1984; Woodward & Colella 1984). Our chosen computational domain is a cube with side length of 260 pc and a finest numerical resolution of $\Delta x = 1.0$ pc.

To study the effect of self-gravity on low-mass IS clouds in an isolated manner we put the clouds at rest with respect to the ambient gas. By that, internal cloud dynamics is not triggered by external forces. We focus on three models as indicated in Table 1. First, a fiducial cloud without both self-gravity and plasma cooling is simulated to check the numerics of the code (model H0). Second

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Table 1. Simulated model clouds. The respective physical processes are either considered (+) or disregarded (–).

Model	Self-gravity	Heating and cooling
H0	–	–
H1	–	+
H2	+	+

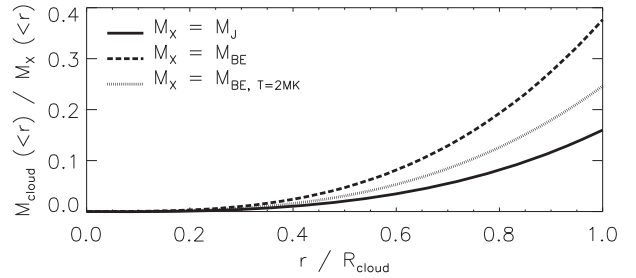


Figure 1. Radial distribution of cloud-to-Jeans-mass ratio (solid line) and cloud-to-Bonnor-Ebert mass ratio (dashed line for $T_{\text{ISM}} = 5.6 \times 10^6$ K, dotted line for $T_{\text{ISM}} = 2 \times 10^6$ K). The x-axis is normalized to the respective radius of each cloud.

and third: one cloud without (model H1) and another with self-gravity (model H2) are computed. We use FLASH’s Multigrid solver to integrate Poisson’s equation. This solver is adequate for handling arbitrary density distributions.

For comparison reasons, each cloud is initially homogeneous, isothermal, and isobaric and their masses and radii are the same: $M_{\text{cloud}} = 6.4 \times 10^4 M_{\odot}$ and $R_{\text{cloud}} = 41$ pc, respectively. The ambient gas has a temperature and density of $T_{\text{ISM}} = 5.6 \times 10^6$ K and $n_{\text{ISM}} = 7 \times 10^{-4} \text{ cm}^{-3}$, respectively. Temperature and density are chosen in such a way that the model clouds and the ambient hot phase are in pressure equilibrium at $\sim 4000 \text{ K cm}^{-3}$, which is in the favoured range of 10^3 to 10^4 K cm^{-3} for a three-phase ISM. According to Tumlinson, Peebles & Werk (2017) a value of $T_{\text{ISM}} = 5.6 \times 10^6$ K is consistent with observations of the Milky Way circumgalactic medium (CGM) ranging from 10^6 to 10^7 K. The density used in the models agrees with current measurements of the halo (10^{-5} to $\sim 0.002 \text{ cm}^{-3}$, Bland-Hawthorn & Gerhard 2016, and references therein). Internally, the clouds are in thermal equilibrium. Decreasing the temperature to the widely accepted value of $T_{\text{ISM}} \sim 2 \times 10^6$ K (Miller & Bregman 2015) would lead to an external pressure of $\sim 1400 \text{ K cm}^{-3}$. In order to get a model cloud at the same cloud temperature and with the same mass, the cloud density, n_{old} , must be lowered by the same factor at which T_{ISM} decreases (to maintain pressure equilibrium) yielding n_{new} . Accordingly, the cloud radius has to rise by $(n_{\text{old}}/n_{\text{new}})^{1/3}$ to obtain the same cloud mass. In our case, the radius increases to 59 pc. The decreased external pressure affects the Bonnor–Ebert mass of the cloud (cf. Section 2.2), such that the ratio of cloud-to-Bonnor–Ebert mass is lower as for $T_{\text{ISM}} = 5.6 \times 10^6$ K (cf. Fig. 1). The initial physical situation does not change and our reasoning still holds. We emphasize that we do not consider thermal conduction within the simulations.

The rates of heating and cooling of the plasma, and the degrees of dissociation and ionization are computed semi-analytically based on the local values of temperature, density, and metallicity. These local variables are calculated during runtime. Depending on temperature we have applied three different cooling laws from literature: for

¹ see <http://flash.uchicago.edu/site/flashcode/>

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$T < 900$ K we use the cooling curve from Falgarone & Puget (1985), which considers molecular line cooling. For $900 < T/K < 10^4$ we use the cooling function from Dalgarno & McCray (1972), and for temperatures above 10^4 K the plasma cools according to the rates in Boehringer & Hensler (1989). The plasma is heated owing to the photoelectric effect on dust particles (Weingartner, Draine & Barr 2006), ionization by UV radiation, by X-rays, and by cosmic rays (Wolfire et al. 2003), thermalization of turbulent motions, and condensation of molecular hydrogen on dust particles (both in Tielens 2010).

2.2 Simulation results

From linear perturbation analysis it is well known that self-gravity in clouds is important if at any radius $r \leq R_{\text{cloud}}$ the *Jeans mass* (e.g. Binney & Tremaine 1987)

$$M_J = \frac{\pi}{6} \left(\frac{\pi k \bar{T}}{G \mu} \right)^{3/2} \bar{\rho}^{-1/2} \quad (1)$$

is exceeded by the enclosed mass $M_r = M(\leq r)$. It is an indicator for a region of mass M_r , which is not confined by an external pressure, to be gravitationally stable ($M_r \leq M_J$) or unstable ($M_r > M_J$). However, if a non-vanishing ambient pressure $P_a = P(>r)$ takes effect on a cloud with radius r , the upper bound for its mass in order to resist gravitational instability is provided by its *Bonnor–Ebert mass* (Ebert 1955; Bonnor 1956)

$$M_{\text{BE}} = 1.18 \left(\frac{k \bar{T}}{\mu} \right)^2 G^{-3/2} P_a^{-1/2}. \quad (2)$$

The values $\bar{T}$, $\bar{\rho}$, μ in equations (1) and (2) denote mean values of temperature and density and the mean molecular weight, respectively, inside r and G is Newton’s constant of gravity. Thus, for any region,

$$M_r / M_{\text{BE}} > 1 \quad (3)$$

is a *sufficient* condition for self-gravity to be considered in the evolution of an IS cloud. However, it remains to prove the *necessity* of condition (3), namely,

$$\text{self-gravity is considered in an IS cloud} \Rightarrow M_r / M_{\text{BE}} > 1. \quad (4)$$

We approach this proof by means of hydrodynamical simulations. Condition (3) is not satisfied in any of the model clouds (Fig. 1), but models H1 and H2 differ by the consideration of self-gravity (Table 1). Model H1 is our reference cloud: it evolves a certain radial density profile not shaped by self-gravity. Oppositely, in model H2 self-gravity is considered. We now can use the method of proof called *reductio ad absurdum*: Let’s assume implication (4) to be true. Then self-gravity would not have any effect on the evolution of the density profile in cloud H2 *a priori*. Consequently, we would not see any (or only negligible) difference in the density profiles of models H1 and H2. By inspecting Figs 2 and 3, however, one clearly observes a strikingly different evolution of both the density distributions and the shapes of clouds H1 and H2. Both clouds condense hot ambient gas and thus increase their content of thermal energy. The expansion implied cannot be compensated in model H1. Obviously, cloud H2 is able to remain compact due to self-gravity despite its Bonnor–Ebert mass is far from being reached. Therefore, a contradiction to our assumption, that implication (4) is true, is shown and, by that, implication (4) is disproved. After 150 Myr, cloud H1 still oscillates entirely and with increasing elongation.

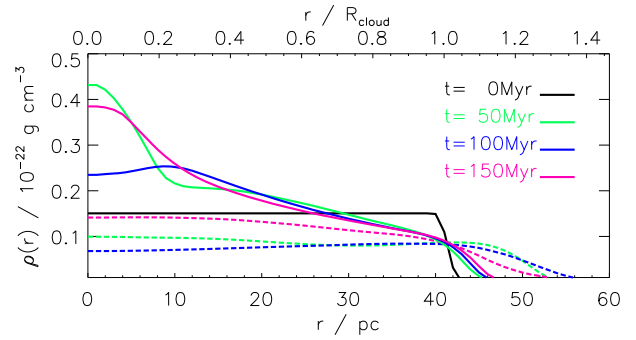


Figure 2. Evolution of the angle-averaged radial density profile in model H1 (dashed lines) and in model H2 (solid lines). The black line shows the initial density profile for both model clouds.

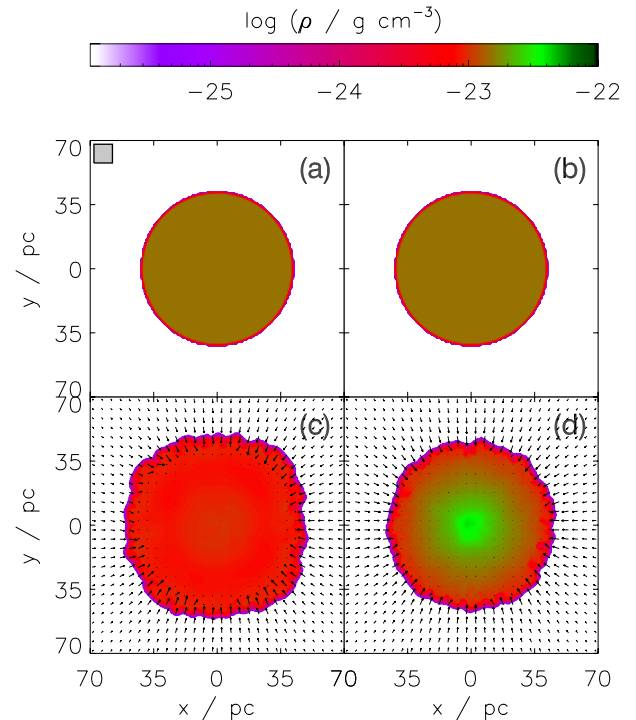


Figure 3. Evolution of density in the model clouds: (a) at initial stage, (b) fiducial model H0, (c) model H1, and (d) model H2 each after 150 Myr. The black arrows visualize the velocity field of gas. The grey square in the upper left corner is spanned by 10×10 cells of finest resolution ($\Delta x = 1.0$ pc).

However, cloud H2 remains almost constant further out but only inside the central 10 pc densities are changing. The main differences are:

- (i) The central density increases in cloud H2 by a factor $\gtrsim 3$ within 150 Myr. The cloud remains spherical and compact.
- (ii) Cloud H1 remains homogeneous over 150 Myr of evolution and expands significantly.

We conclude that self-gravity has a non-negligible effect on the evolution of even low-mass clouds and must in general be included to evolutionary models. Also a homogeneous distribution of density does not provide a realistic condition in the presence of self-gravity.

In plot (b) in Fig. 3 it is seen that the fiducial model H0 presents a stationary solution as expected for a homogeneous cloud, which does not consider any processes for energy exchange with the surrounding gas (i.e. plasma cooling and heating), where no gravitational forces are present, and which is in pressure equilibrium with its surroundings.

3 HEAT CONDUCTION IN A MAGNETIC DIPOLE FIELD

Since cool IS clouds are reasonably embedded in a warmer (up to hot) IS gas, thermal conduction determined by the temperature gradient is unavoidable. The path of charged particles such as electrons is hampered by magnetic fields. Vieser & Hensler (2007) have calculated that the mean free path of an electron perpendicular to a magnetic field in pressure equilibrium with the gas is reduced by a factor of 10^7 according to the Larmor radius with respect to collisions with neutral hydrogen atoms. This might lead to the general and frequently expressed, but nowhere documented, argument that thermal conduction cannot play a role for the energy budget of IS clouds.

A magnetic dipole field in spherical coordinates reads

$$\mathbf{B} = (B_r, B_\theta, B_\varphi) = \frac{m}{4\pi r^3} (2 \cos \theta, \sin \theta, 0), \quad (5)$$

with radius r , elevation angle θ , azimuthal angle φ , and magnetic moment m (i.e. strength of the dipole). In presence of large temperature gradients the heat flow becomes saturated and can be written as (Cowie & McKee 1977)

$$q_{\text{sat}} \equiv |\mathbf{q}_{\text{sat}}| = 0.4 n_e \left(\frac{2kT_e}{\pi m_e} \right)^{1/2} kT_e \approx 5 \Phi_s \varrho c_s^3, \quad (6)$$

where $v_e \equiv |\mathbf{v}_e| = [2kT_e/(\pi m_e)]^{1/2}$ is the thermal velocity of electrons, c_s is the speed of sound, ϱ is the density, and Φ_s covers both uncertainties owing to the treatment of flux-limited diffusion and the impact of potential magnetic fields. The saturation of heat flow being inherent in representation (6) accounts for a replenishment of the reservoir of heat-conducting electrons within the traveltime of sound waves, i.e. there is a maximum amount of thermal energy that can be transported by electrons moving at v_e .

Assuming a cloud of temperature T_{cl} embedded in an external medium of temperature $T_e > T_{\text{cl}}$, the temperature gradient across the cloud surface points radially inwards and hence the undisturbed flow lacks any angular component. So,

$$\mathbf{q}_{\text{sat}} \propto \mathbf{v}_e \parallel \mathbf{r}. \quad (7)$$

We may thus write for the electron velocity in spherical coordinates

$$\mathbf{v}_e = (v_r, 0, 0) \quad (8)$$

if no local forces are present, which may be due to, e.g. turbulence or density inhomogeneities. Assume the cloud to have a magnetic dipole field (5). The angle α between magnetic field (5) and heat flow (7) is given by the dot product

$$\cos \alpha = \frac{\mathbf{B} \cdot \mathbf{q}_{\text{sat}}}{|\mathbf{B}| \cdot |\mathbf{q}_{\text{sat}}|} \propto \frac{\mathbf{B} \cdot \mathbf{v}_e}{|\mathbf{B}| \cdot |\mathbf{v}_e|} = \left(1 + \frac{1}{4} \tan^2 \theta \right)^{-1/2}. \quad (9)$$

The radial velocity component (8) can be split into a component parallel, $v_{\parallel}$, and a component perpendicular, $v_{\perp}$, to the magnetic field lines. So,

$$v_{\parallel} = v_r \cos \alpha \quad (10)$$

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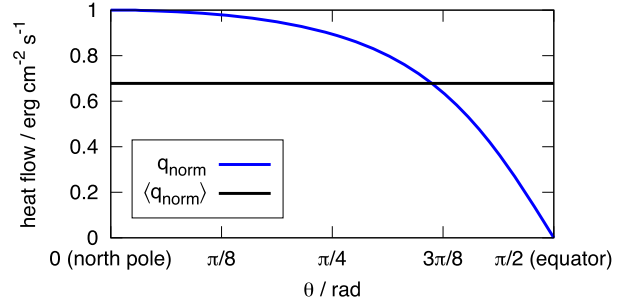


Figure 4. Normalized heat flow, q_{norm} (blue line), at a cloud-centric distance of $r = 40$ pc. The weighted mean of normalized heat flow, $\langle q_{\text{norm}} \rangle$ (black line), is averaged over all θ ranging from 0 to $\pi/2$.

$$v_{\perp} = \sqrt{v_r^2 - v_{\parallel}^2}. \quad (11)$$

By that, the heat flow (6) can be split in the same way. That is,

$$q_{\parallel} = 0.1 n_e v_{\parallel} \frac{3}{2} kT_e \quad (12)$$

$$q_{\perp} = 0.1 n_e v_{\perp} \frac{3}{2} kT_e, \quad (13)$$

such that $q_{\text{sat}} = \sqrt{q_{\parallel}^2 + q_{\perp}^2}$. The Lorentz force acts on $q_{\perp}$ only.

If no magnetic field is present, the electrons travel their respective mean free path for Coulomb collisions $\lambda_e = t_e v_e$, with an electron–electron equipartition time (Spitzer 1962)

$$t_e = \frac{3\sqrt{m_e}(kT_e)^{3/2}}{4\sqrt{\pi} n_e e^4 \ln(\Omega)}. \quad (14)$$

If a magnetic field is present, the electrons are deflected by the Lorentz force and hence gyrate around the field lines with Larmor radius

$$r_L = \frac{m_e v_{\perp}}{e|\mathbf{B}|}. \quad (15)$$

Therefore, the factor, by which $q_{\perp}$ is suppressed, is given by r_L/λ_e , since r_L measures the disturbing effect of the magnetic field on the straight line λ_e . So, the heat flow, which passes the magnetic field, reads

$$q_{\text{eff}} \equiv |\mathbf{q}_{\text{eff}}| = \sqrt{q_{\parallel}^2 + \left(\frac{r_L}{\lambda_e} q_{\perp} \right)^2}. \quad (16)$$

For sufficiently weak magnetic fields $r_L \sim \lambda_e$, and thus the heat flow perpendicular to $\mathbf{B}$ is not affected, but for stronger fields r_L can easily fall below $10^{-8} \lambda_e$ (see discussion in Vieser & Hensler 2007). The normalized heat flow $q_{\text{norm}} \equiv q_{\text{eff}}/q_{\text{sat}}$ measures the fraction of the flow (6) that is able to pass the magnetic field. The magnetic moment is chosen to be $m \approx 3 \times 10^{51}$ erg G $^{-1}$, such that $|\mathbf{B}| \sim 5 \mu\text{G}$ at cloud surface ($r = 40$ pc) at equator ($\theta = \pi/2$), which corresponds to equipartition with the energy density of the cloud. In Fig. 4 a clear dependence of q_{norm} on the elevation angle θ is visible. At equator a value of $q_{\text{norm}} \sim 10^{-9}$ is established.

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Averaging q_{norm} over all θ yields the weighted mean

$$\begin{aligned} \langle q_{\text{norm}} \rangle &= \frac{\int_{\theta=0}^{\pi/2} \int_{\varphi=0}^{2\pi} \sin \theta' q_{\text{norm}}(\theta') d\theta' d\varphi'}{\int_{\theta=0}^{\pi/2} \int_{\varphi=0}^{2\pi} d\theta' d\varphi'} \\ &= \frac{2}{\pi} \int_{\theta=0}^{\pi/2} \sin \theta' q_{\text{norm}}(\theta') d\theta'. \end{aligned} \quad (17)$$

As can be seen in Fig. 4, $\langle q_{\text{norm}} \rangle$ is suppressed to 68 per cent only. More extensive, but rather academic numerical studies on the impact of different magnetic field configurations on heat flow have been performed by Li, Frank & Blackman (2012). Their heat-transfer efficiency ζ has the same meaning like our q_{norm} . What they observe is a hampered conduction if the magnetic field is more tangled, i.e. if the contributions from field lines perpendicular to the heat flow are increased (their situation for $\zeta \rightarrow 0$). This corresponds to $q_{\text{norm}}(\theta = \pi/2) \sim 10^{-9}$ (Fig. 4). However, they figured out that if the global magnetic field dominates the local tangled field (in their terminology: $\zeta \rightarrow 1$), the impediment imposed by anisotropic heat conduction in the local tangled field is negligible. This result conforms with $q_{\text{norm}}(\theta = 0) = 1$. The findings of Li, Frank & Blackman (2012) hence substantiate our results.

4 SUMMARY AND CONCLUSIONS

The aim of our exploration is twofold. In the first part (Section 2) we present three-dimensional hydrodynamical simulations of three models of resting low-mass clouds (Fig. 1). After 150 Myr of evolution the morphologies of our model clouds H1 and H2 are strikingly different (Fig. 3). An initially homogeneous cloud without self-gravity (model H1) remains homogeneous and is shaped irregularly. The evolutionary process is totally different in the case of self-gravity (model H2), where the cloud remains compact, roundish, condenses centrally, and evolves a radial density gradient (Fig. 2). Hence we conclude that self-gravity has a major impact on even gravitationally stable clouds ($M < M_{\text{BE}}$), which disproves the necessity of condition (3).

We further calculated analytically the effective heat flow of electrons passing a magnetic dipole field, which contributes major to a multipole field (Section 3). It turned out that even in equipartition the heat flow, which is integrated over the cloud's surface, is reduced to 68 per cent only (Fig. 4). Similar results have been found for more complex configurations of magnetic fields. For example, Narayan & Medvedev (2001) have shown that the effective heat flow is suppressed to only ~ 20 per cent if the entanglement of the magnetic field lines extends over certain spatial scales. We thus deduce that even in the presence of strong magnetic fields the heat flow must not be neglected in simulations of cool clouds embedded in a hot plasma.

The study presented is thought as a preparatory effort for cloud models with heat conduction and the evolution of HVCs. In two forthcoming papers we show that the arguments we figured out in our analysis performed with resting clouds can be applied to fast-moving low-mass clouds, e.g. HVCs, as well.

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²see <http://vsc.ac.at/>

7.5 THERMAL CONDUCTION IN INTERSTELLAR CLOUDS

Within this section thermal conduction is analyzed in resting multi-phase clouds embedded in a hot, tenuous medium. The formation, evolution, and structure of a transition zone between cloud and ambient medium is analyzed in detail and is compared to theoretical considerations (Cowie & McKee, 1977). Condensation rates and mixing (efficiency, spatial distribution) of ambient gas with cloud material are discussed in terms of radial cloud-density gradient, self-gravity, and cloud mass.

The idea for the manuscript was developed together by my supervisor Prof. Dr. G. Hensler and me. Prof. Hensler reviewed the manuscript iteratively. The manuscript is co-authored by Prof. Hensler and me.

My own contential contributions include:

- all simulations and calculations presented,
- all plots are based on self-written IDL scripts,
- the correlation analysis in section 5.1.

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The published paper is provided below.

The effect of saturated thermal conduction on clouds in a hot plasma

Bastian Sander^{1,2★} and Gerhard Hensler^{1★}

¹*Department of Astrophysics, University of Vienna, Türkenschanzstraße 17, A-1180 Vienna, Austria*

²*Anhalt University of Applied Sciences, Strenzfelder Allee 28, 06406 Bernburg, Germany*

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ABSTRACT

We numerically investigate the internal evolution of multiphase clouds, which are at rest with respect to an ambient, highly ionized medium (HIM) representing the hot component of the circumgalactic medium. Time-dependent saturated thermal conduction and its implications like condensation rates and mixing efficiency are assessed in multiphase clouds. Our simulations are carried out by using the adaptive mesh refinement code FLASH. The model clouds are initially in both hydrostatic and thermal equilibrium and are in pressure balance with the HIM. Thus, they have steep gradients in both temperature and density at the interface to HIM leading to non-negligible thermal conduction. Several physical processes are considered numerically or semi-analytically: thermal conduction, radiative cooling and external heating of gas, self-gravity, mass diffusion, and dissociation of molecules and ionization of atoms. It turns out that saturated thermal conduction triggers a continuous condensation irrespective of cloud mass. Dynamical interactions with ambient HIM all relate to the radial density gradient in the clouds: (1) mass flux due to condensation is the higher the more homogeneous the clouds are; (2) mixing of condensed gas with cloud gas is easier in low-mass clouds, because of their shallower radial density gradient; and thus (3) accreted gas is distributed more efficiently. A distinct and sub-structured transition zone forms at the interface between cloud and HIM, which starts at smaller radii and is much narrower as deduced from analytical theory.

Key words: conduction – diffusion – hydrodynamics – methods: numerical – ISM: clouds.

1 INTRODUCTION

Astrophysical environments on and below galactic scales are inhomogeneous and not isothermal, which is observed in the interstellar medium (ISM, Lequeux 2005; Shelton 2009; Carraro 2021) and circumgalactic medium (CGM; Putman, Peek & Joung 2012; Tumlinson, Peebles & Werk 2017). Sample mechanisms that drive the Galactic matter cycle by maintaining these different physical states (so-called *phases*) on diverse spatial scales are galactic outflows (Norman & Ikeuchi 1989; McClure-Griffiths et al. 2003; Davé, Oppenheimer & Finlator 2011) and accretion of gas (Fraternali 2014; Richter 2016; Fox & Davé 2017), supernovae (Hensler 2011), and stellar winds (Zhang 2018). The coexistence of these phases in the ISM and CGM implies non-equilibrium situations being continuously triggered. These non-equilibrium situations lead to various dynamic processes like hot accretion flows (Faghei 2012), plasma cooling (Lehner, Savage & Wakker 2006), or thermal conduction (Nipoti & Binney 2004; Pittard et al. 2004; Brüggén & Scannapieco 2016), such that the phases in the ISM and CGM do not only contact and envelope each other, but also penetrate into each other in order to equilibrate by exchange of energy and mass. A static situation of the ISM is represented by the theory of a two-phase medium (Field, Goldsmith & Habing 1969), which is later extended by a third hot phase (McKee & Ostriker 1977). Recent studies and observations

reveal a more complex dynamic picture of the hot haloes around galaxies (Breitschwerdt et al. 2012; Falceta-Gonçalves et al. 2014; Lee & Lee 2019; Damle et al. 2022). High pressure and gravity in the cold neutral matter (CNM) balance with low pressure in the highly ionized medium (HIM) at their common interface, such that steep gradients in both density and temperatures are present. They in turn lead to a flow of heat and matter towards the CNM. Analytical approaches propose the evaporation of clouds in a few dynamical times (Balbus & McKee 1982; Giuliani 1984; Nagashima, Inutsuka & Koyama 2007), which contradicts observations of cool clouds embedded in diverse hot environments. Thus, the clouds are able to resist evaporation by combined physical mechanisms.

Previous numerical studies of multiphase clouds considering plasma heating and cooling, self-gravity, mass diffusion, and thermal conduction reveal that thermal conduction is saturated and leads to condensation thus providing a mechanism to mix ambient gas into clouds and, by this, enhancing their metallicity (two-dimensional simulations by Vieser & Hensler 2007b, hereafter *VH07b*). In a follow-up study Vieser & Hensler (2007a) show that the lifetime of moving clouds is distinctly extended by thermal conduction, which efficiently suppresses Kelvin-Helmholtz instabilities. Sander & Hensler (2019) work out that (1) self-gravity is important for even sub-Jeans clouds, and (2) thermal conduction may not be neglected for gaseous spheres even in presence of strong magnetic fields. Very recently, Sander & Hensler (2021, hereafter *SH21*) analyze the evolution of high-velocity clouds (HVCs) in a hot CGM and figure out that thermal conduction together with deceleration by drag naturally leads to the suppression of star formation in compact HVCs,

* E-mail: bastian.sander@hs-anhalt.de (BS); gerhard.hensler@univie.ac.at (GH)

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which conforms with observations (e.g. Davies et al. 2002; Hopp, Schulte-Ladbeck & Kerp 2007). In this work, we extend these investigations by analyzing saturated thermal conduction in an isolated manner without being superimposed by hydrodynamic interactions.

This paper is structured as follows: in Sections 2 and 3, we describe the physical processes being considered in our simulations, and Section 4 contains the applied numerical code and the model setup. The model clouds and the results of the simulations are discussed in Section 5. The paper is summarized in Section 6, where also the major conclusions are drawn.

2 THERMAL CONDUCTION

In the presence of large temperature gradients the classical description of the heat flux given by Spitzer (1962),

$$\mathbf{q}_{\text{class}} = -\kappa_{\text{class}} \nabla T, \quad (1)$$

with the coefficient for thermal conduction $\kappa_{\text{class}} = 1.84 \times 10^{-5} T^{5/2} / \ln(\Omega)$, and the Coulomb logarithm $\ln(\Omega) = 29.7 + \ln[T / (10^6 \sqrt{n_e})]$, yields a too high amount of energy being conducted by electrons (e.g. Campbell 1984). Instead, a saturated heat flux $\mathbf{q}_{\text{sat}}$ (see Cowie & McKee 1977; McKee & Cowie 1977) accounts for a finite reservoir of electrons and thus converges to a maximum value if all electrons are conducting heat irrespective of the steepness of the temperature gradient. It has been experimentally verified by Gray & Kilkenny (1980) that

$$\mathbf{q} = \begin{cases} \mathbf{q}_{\text{class}}, & \lambda_e / L_T \lesssim 2 \times 10^{-3} \\ \mathbf{q}_{\text{sat}}, & \text{else} \end{cases} \quad (2)$$

The first essential spatial scale is the mean free path of electrons in the hot, ambient medium,

$$\lambda_e = t_e \sqrt{\frac{3kT_{\text{amb}}}{m_e}}, \quad (3)$$

with temperature T_{amb} and electron-electron equipartition time (Spitzer 1962)

$$t_e = \frac{3\sqrt{m_e}(kT_{\text{amb}})^{3/2}}{4\sqrt{\pi}n_e e^4 \ln(\Omega)}. \quad (4)$$

The second spatial scale is the local temperature scale-height, $L_T = T/|\nabla T|$, which provides the typical spatial scale for temperature variation. If λ_e is significantly less than L_T (diffusion limit), temperature gradients are levelled off and the classical approach for thermal conduction is valid. In the opposite case, even strong temperature variations are preserved (free-streaming limit) and thermal conduction becomes saturated. Cowie & McKee (1977, hereafter CM77) suggest the expression

$$\mathbf{q}_{\text{sat}} = 5\Phi \varrho c_s^3 \quad (5)$$

for the saturated heat flux, which depends on the speed of sound c_s , local density ϱ , and a factor Φ accounting for the uncertainties in the definition of (5) regarding contributions in a plasma by magnetic fields, currents, instabilities etc. It holds $\Phi = 1.1$ if both ion and electron temperatures are equal. We use $\Phi = 1$ throughout all of our simulations analyzed in this work.

To ensure continuity at the transition between $\mathbf{q}_{\text{class}}$ and $\mathbf{q}_{\text{sat}}$ a proper representation for the heat flux must be chosen. Within this work the expression of Slavin & Cox (1992)

$$\mathbf{q}_{\text{eff}} = \mathbf{q}_{\text{sat}} \left(1 - \exp \left\{ -\frac{|\mathbf{q}_{\text{class}}|}{|\mathbf{q}_{\text{sat}}|} \right\} \right) \quad (6)$$

is used. It correctly converges to either of the heat fluxes in equation (2) for very small or very large temperature gradients, respectively.

The conductivity according to the effective heat flux (6) reads

$$\kappa_{\text{eff}} = \kappa_{\text{sat}} \left(1 - \exp \left\{ -\frac{\kappa_{\text{class}}}{\kappa_{\text{sat}}} \right\} \right), \quad (7)$$

with $\kappa_{\text{sat}} = 5\varrho c_s^3 / |\nabla T|$. We follow the procedure described in SH21 for calculating the temperature update according to saturated thermal conduction. Namely, we use the semi-implicit Crank-Nicolson method (Crank, Nicolson & Hartree 1947), which is applied in a directionally split manner, i.e. the temperature is updated along one spatial dimension while being kept constant along the other two dimensions. This procedure is repeated for the other dimensions, too. The order of direction sweeps is permuted analogously to the hydrodynamics (cf. SH21). From the new temperature distribution an adjusted heat flux

$$\mathbf{q}_{\text{eff}} = -\kappa_{\text{eff}} \nabla T \quad (8)$$

finally accounts for the change in energy by thermal conduction. It has to be noted that the temperature gradient can be very steep, especially across the cloud surface. The heat flux (8) can thus be very high in the vicinity of regions of nearly zero heat flux. We account for this special situation by a slope limiter (cf. SH21).

Begelman & McKee (1990) introduce the Field length

$$\lambda_F = \sqrt{\frac{\kappa_{\text{amb}} T_{\text{amb}}}{n_{\text{amb}}^2 \Lambda_{\text{max}}}}, \quad (9)$$

as a characteristic spatial scale to distinguish between growth and suppression of thermal instabilities. The Field length relates the thermal conductivity κ_{amb} with cooling efficiency $\Lambda_{\text{max}} = \max\{\Lambda_0, \Gamma_0/n_{\text{amb}}\}$ in the surrounding, hot phase that directly contacts the cloud. The rate coefficients for radiative cooling, Λ_0 [erg cm⁻³ s⁻¹], and heating, Γ_0 [erg s⁻¹], denote the respective density-independent powers being radiated (Raymond, Cox & Smith 1976; Boehringer & Hensler 1989). Thus, by equation (9) the Field length of ambient gas is calculated. If gas phases change, then λ_F changes accordingly. So, λ_F is a function of radius and time in cool clouds embedded in a hot plasma with time-dependent thermal conduction. If the cloud is interpreted as a cool density inhomogeneity embedded in the hot phase, its radius can be compared to λ_F and a direct consequence for the thermal structure of the cloud can be derived: if $R_{\text{cl}} \gg \lambda_F$, the cloud is dominated by external heating and radiative cooling while for $R_{\text{cl}} \ll \lambda_F$ classical thermal conduction governs. If, however, $R_{\text{cl}} \sim \lambda_F$, saturated thermal conduction is the dominating effect. The Field length must be resolved by at least three grid cells in numerical simulations (Koyama & Inutsuka 2004; Gressel 2009) to prevent results from being dependent on spatial resolution. We account for this in our simulations by resolving λ_F in equation (9) by 206 cells in the massive model clouds and 275 cells in the low-mass clouds.

CM77 introduce a global saturation parameter

$$\sigma_0 = \frac{(T_{\text{amb}} / 1.54 \times 10^7 [\text{K}])^2}{n_{\text{amb}} \Phi R_{\text{cl}} [\text{pc}]}, \quad (10)$$

which can be used to discriminate between classical ($\sigma_0 < 1$) and saturated ($\sigma_0 > 1$) thermal conduction. Here, R_{cl} is the cloud radius in terms of parsec, and Φ is defined in equation (5). The mass loss is considerably reduced with increasing σ_0 (see solid line in Fig. 1). For $\sigma_0 \gg 1$ also viscous heating becomes relevant and thermal conduction is best treated by a two-fluid approach (Balbus & McKee 1982). Based on equation (10), CM77 determine the rate of

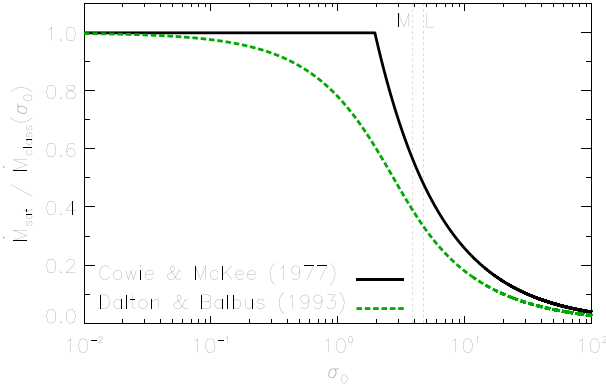


Figure 1. Ratios of saturated mass-loss rate to classical mass-loss rate analytically calculated by CM77 (black solid line) and Dalton & Balbus (1993) (green dashed line). The vertical thin dashed lines mark the values for σ_0 used for the simulated cloud models M and L (cf. Table 1).

evaporation due to saturated thermal conduction by

$$\dot{M}_{\text{CM}} = 3.25 \times 10^{18} n_{\text{amb}} \sqrt{T_{\text{amb}}} R_{\text{cl}}^2 [\text{pc}] \Phi F(\sigma_0) \text{ g s}^{-1} \quad (11)$$

for clouds without self-gravity, where

$$F(\sigma_0) = 2 \left[(\sigma_0 H)^{1+Ma^2} \exp \{-2.5 Ma^2\} \right]^{1/(6+Ma^2)}, \quad (12)$$

and

$$H = \begin{cases} h^h / (h-1)^{h-1}, & Ma \leq 1 \\ 11.5, & Ma > 1 \end{cases}, \quad h = \frac{11 + Ma^2}{1 + Ma^2}, \quad (13)$$

and Ma is the Mach number of the flow around the cloud or, more specific, within the zone of saturated heat flux. Dalton & Balbus (1993) extend these investigations by introducing a continuous transition between classical and saturated heat flux. By simultaneously solving their equations (25) and (30), they obtain the ratio $\omega(\sigma_0)$ (cf. dashed line in Fig. 1) between saturated and classical mass-loss rate. Hence,

$$\dot{M}_{\text{DB}} = \omega(\sigma_0) \dot{M}_{\text{class}} = \omega(\sigma_0) \frac{16\pi \mu_{\text{HIM}} \kappa_{\text{HIM}} R_{\text{cl}}}{25k_{\text{B}}}. \quad (14)$$

Theoretically, gas should evaporate from clouds for $\lambda_e \lesssim R_{\text{cl}} \ll \lambda_{\text{F}}$ and condense onto them if $R_{\text{cl}} \gg \lambda_e$ and $R_{\text{cl}} \gtrsim \lambda_{\text{F}}$. McKee & Begelman (1990) show that hot material condenses on isolated clouds if $R_{\text{cl}} \sim 0.24 \dots 0.36 \lambda_{\text{F}}$ (see also VH07b).

3 PHYSICAL PROCESSES

Beside thermal conduction we cover in our simulations different physical processes that most likely take place in interstellar clouds. They are explained in detail in SH21 thus we summarize only the most important facts here.

3.1 Self-gravity

The Poisson's Equation is solved in the entire computational domain. It is shown in Sander & Hensler (2019) that self-gravity may not be neglected even in clouds with masses below their Bonnor-Ebert mass. If self-gravity is stronger than acceleration of gas due to evaporation, the net acceleration at cloud surface points in the same direction like the density gradient. Hence, the condition for Rayleigh-Taylor (RT) instability is not given. So, it is expected that self-gravity significantly suppresses RT instabilities at the cloud surface (Murray et al. 1993).

3.2 Dissociation and ionization

The simulations contain cool clouds and hot plasma, thus temperatures range from a few hundreds up to several 10^6 K. While at lowest possible temperatures molecules can exist, they successively dissociate and finally become totally ionized with rising temperature. In our simulations we semi-analytically calculate both dissociation and ionization for hydrogen only. The fraction of H II is then related to He and metals. So, we have multiple phases (molecular, atomic, ionized) simultaneously present in our clouds. We note that we do not trace single species but rather have one fluid for which the composition is calculated based upon a given metallicity. The fractions of dissociation and ionization affect the magnitudes of cooling rates.

3.3 Plasma heating and cooling

The rates of heating, Γ , and cooling, Λ , of the plasma are treated semi-analytically. Depending on temperature we account for molecular line cooling (Falgarone & Puget 1985), atomic line cooling (Dalgarno & McCray 1972), and bremsstrahlung (Boehringer & Hensler 1989). The plasma is heated owing to the photoelectric effect on dust particles (Weingartner, Draine & Barr 2006), ionization by UV radiation, by X-rays, and by cosmic rays (Wolfire et al. 2003), thermalization of turbulent motions, and condensation of molecular hydrogen on dust particles (both in Tielens 2010).

Both Γ and Λ depend on local values of metallicity, temperature, density, and the fractions of dissociation and ionization. If metallicity is high, cooling rates are stronger and hence the temperature is lower in thermal equilibrium, which implies a lower pressure. Consequently, especially the mixing regions in the outskirts of clouds are expected to have a decreased pressure, which results in a flow of matter towards the outskirts.

3.4 Mass diffusion

We consider the transport of matter by diffusion using the formulation of the Chapman-Enskog expansion, which is valid at low gas densities (Bird et al. 1960).

3.5 Simplifications

Radiative transport is neglected, because we assume the model clouds to be optically thin. Magnetic fields are neglected as even strong dipole fields do not remarkably suppress the heat flux integrated over the entire cloud surface (Sander & Hensler 2019). It has been shown by Tribble (1989, cf. equation 11 there) that tangled magnetic field lines suppress the effective thermal conductivity between adjacent gas phases at different temperatures only by some small amount. A similar result is reported by Malyskin (2001) for a negligible homogeneous component of the magnetic field (see also Tao 1995).

We further consider chemodynamics on a simplified level, i.e. different species are not explicitly traced by their respective conservation laws. We rather calculate the fractions of H_2 , H I, H II, He, electrons, and metals in a semi-analytical fashion based on the local fractions of dissociation and ionization, respectively, for a single fluid.

3.6 Hydrodynamics

We solve the Euler Equations for clouds resting in the HIM and close them by an equation of state for the ideal gas. According to the source

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terms of self-gravity, heating and cooling, and thermal conduction, the hydrodynamics conservation law for total energy, e_{tot} , we solve numerically is given by

$$\frac{\partial \rho e_{\text{tot}}}{\partial t} + \nabla \cdot [\rho e_{\text{tot}} + P]v = \rho v g + \Gamma - \Lambda - \nabla \cdot q_{\text{eff}}. \quad (15)$$

4 THE SIMULATIONS

We conduct full 3D hydrodynamics simulations of cool, stratified, dense, multiphase clouds that are at rest with respect to an ambient, hot, rarefied medium resembling the HIM by means of the publicly available FLASH code¹ (Fryxell et al. 2000; Dubey et al. 2009). The (inviscid) Euler Equations are solved on a Cartesian, adaptively refined grid by applying the piecewise parabolic method (Colella & Woodward 1984; Woodward & Colella 1984). The Poisson's Equation is solved by a Multigrid solver, which is able to deal with arbitrary mass distributions. See SH21 for further details on the FLASH code and how we handle steep temperature gradients.

The present work investigates in more detail the effect of saturated thermal conduction on multiphase clouds and compares to previous numerical (VH07b) and analytical (CM77, McKee & Cowie 1977; Balbus & McKee 1982; Giuliani 1984; Begelman & McKee 1990; McKee & Begelman 1990; Nagashima et al. 2007) works. For this purpose cool, molecular to low-ionized (i.e. multiphase), dense, stratified clouds are set up at rest in a hot, tenuous, fully ionized plasma resembling the HIM ($T \gtrsim 10^{5.5} \text{ K}$) in the CGM. Within this paper we focus on the meaning of thermal conduction for the evolution of interstellar clouds that are investigated in SH21. The present study serves for investigating the non-dynamical processes that provide the initial conditions for HVCs. In particular, we are interested in the stabilizing effect thermal conduction has on HVCs by (1) formation of a transition zone between cloud and ambient gas, and (2) mass gain due to condensation of hot gas onto the cloud. For that purpose, our model clouds are setup at rest with respect to the ambient hot gas, which means they are neither in translational nor rotational motion. Dynamical effects like ram-pressure stripping or the Bernoulli effect hence do not superpose diffusion and we are able to study thermal conduction in the clouds in a fundamental manner.

The clouds divide into two mass categories: massive clouds (category *M*) and low-mass clouds (category *L*). Each of the model clouds considers all physical processes described in Section 3. The reference clouds (M, L) additionally account for thermal conduction while the model variations M_nc and L_nc do not, i.e. there is no conductive input of thermal energy from the ambient medium.

Initially, all four model clouds are in both hydrostatic and thermal equilibrium. By that, the radial temperature profile of the cloud is governed by

$$\frac{\partial T}{\partial r} = -\rho(T) \frac{GM(r < R_{\text{cl}})}{r^2} \left\{ \frac{\partial}{\partial T} \left[\frac{\rho(T)kT}{\mu(\rho, T)} \right] \right\}^{-1}, \quad (16)$$

where $\mu(\rho, T)$ is the mean molecular weight. From equation (16) the profiles of density and pressure are derived (see fig. 2 in SH21) and the radius R_{cl} is obtained at a cloud-centric distance, where $P_{\text{cl}}(r = R_{\text{cl}})$ equals the pressure of the ambient medium, P_{amb} . As being discussed in the present paper, model clouds of equal mass differ only in an additional heat input by thermal conduction, which destroys thermal balance over time. The Poisson's Equation is solved in the entire computational domain, hence HIM is accreted by the clouds due to gravity. The non-conductive model clouds thus serve as a

testbed of accretion by gravity compared to condensation by thermal conduction. Moreover, initial hydrostatic and thermal equilibrium must be maintained in clouds without thermal conduction, so they prove the implemented numerics to work correctly.

According to the HVC models in SH21 we still distinguish the metal content in the gas phases: All model clouds have an initial metallicity of $0.1 Z_{\odot}$ and the HIM has solar metallicity. By using different metallicities we are able to trace the enrichment of the cloud by heavy elements from HIM, which is triggered by gas mixing due to thermal conduction. As already mentioned in Section 3, our treatment of species implies that there is no effect of metallicity on gas dynamics. Temperature and density of the HIM in our simulations resemble those for the CGM (Tumlinson et al. 2017). By subsuming the parameter ranges for σ_0 , λ_F , λ_e , L_T (cf. Section 3), and the cloud mass, M_{cl} , the clouds we are interested in are located in the following subset of the relevant parameter space

$$\sigma_0 > 1, \quad R_{\text{cl}} \lesssim \lambda_F, \quad M_{\text{cl}} < M_{\text{BE}}, \quad (17)$$

whereas the HIM must satisfy (cf. equation (2))

$$\lambda_e \gtrsim L_T. \quad (18)$$

Table 1 contains the realisations of our setups in the parameter spaces (17) and (18).

The size of the computational volume is $(260 \text{ pc})^3$ for massive and $(160 \text{ pc})^3$ for low-mass clouds. Thus, the numerical domain is three times larger in diameter than the cloud, and boundary effects are not observed to have a feedback on the dynamics around the cloud. The chosen size of the integration domain implies a grid spacing of $\Delta x = 1.02 \text{ pc}$ ($\Delta x = 0.63 \text{ pc}$) on six levels of refinement for the simulations of massive (low-mass) clouds. In order to prove that obtained results are independent on resolution, we perform additional fine-grained simulations on seven levels of refinement yielding $\Delta x = 0.51 \text{ pc}$ ($\Delta x = 0.31 \text{ pc}$) for massive (low-mass) clouds (see Appendix A).

4.1 Massive model clouds

Massive clouds have masses of $9.0 \times 10^4 M_{\odot}$ that correspond to 31 per cent of their Bonnor-Ebert mass (Ebert 1955; Bonnor 1956)

$$M_{\text{BE}} = 1.18 \left(\frac{k \bar{T}_{\text{cl}}}{\bar{\mu}_{\text{cl}}} \right)^2 G^{-3/2} P_{\text{amb}}^{-1/2}, \quad (19)$$

which provides an upper bound for the mass of self-gravitating clouds embedded in a medium with pressure P_{amb} in order to be stable against gravitational instability. The identity (19) strictly holds for isothermal clouds only. Hence, we use the density-weighted average values of cloud temperature, T_{cl} , and mean molecular weight, μ_{cl} . The massive models show a distinct core-halo structure in density, temperature and, consequently, in pressure. The central region is dominated by the molecular phase which turns into an extended atomic phase at larger radii. Only a small rim at the boundary consists of warm, slightly ionized gas. Massive clouds are located in a hot plasma with temperature $T_{\text{HIM}} = 5.6 \times 10^6 \text{ K}$ and particle density $n_{\text{HIM}} = 0.7 \times 10^{-3} \text{ cm}^{-3}$. They exhibit a metallicity of 0.1 solar, which leads to a radius of 49 pc in consideration of equilibrium between cooling and heating in the cloud, and pressure balance between cloud and HIM. The central temperature is 265 K, which quickly raises to $\sim 970 \text{ K}$ at 10 pc and moderately increases to $\sim 1,650 \text{ K}$ at the boundary (Fig. 2, upper plot). That is, the molecular phase is confined to the inner 10 pc, beyond which dissociation is complete and only the atomic phase is present. Because temperature and hence conductivity are initially moderate inside the cloud, classical thermal conduction takes place. Since $\kappa_{\text{class}} \propto T^{5/2}$ a

¹ See <http://flash.uchicago.edu/site/flashcode/>

Table 1. Simulated model clouds with (models M, L) and without (models M_nc, L_nc) thermal conduction (also indicated in column 3 by + or –). Columns 4 to 9 consecutively list initial values of radius, mass, saturation parameter (cf. equation (10)), ratio of mean free path of electrons to scale height of temperature (cf. equation (2)), ratio of cloud radius to Field length, and the cloud mass in terms of Bonnor-Ebert mass. Columns 10 and 11 tabulate the side length of the computational cube and the finest numerical resolution, respectively, and column 12 shows the resolution of the Field length. The initial central values for temperature and gas density are given in columns 13 and 14. The metallicity in all model clouds is initially $0.1 Z_{\odot}$.

Model group (1)	Model name (2)	thermal cond. (3)	R_{cl} [pc] (4)	M_{cl} [$10^4 M_{\odot}$] (5)	$\sigma_0^{(*)}$ (6)	$\lambda_e/L_T^{(*)}$ (7)	$R_{cl}/\lambda_F^{(*)}$ (8)	M_{cl}/M_{BE} (9)	L [pc] (10)	Δx [pc] (11)	$\lambda_F/\Delta x^{(*)}$ (12)	T_{centre} [K] (13)	ρ_{centre} [$10^{-21} \text{ g cm}^{-3}$] (14)
massive clouds	M	+	49	9.0	3.8	1.22	0.23	0.31	260	1.02	206	265	0.8
	M_nc	–			–	–	–	–			–		
low-mass clouds	L	+	28	2.2	4.6	0.69	0.16	0.09	160	0.63	275	1,160	0.02
	L_nc	–			–	–	–	–			–		

Note. (*) This parameter relates to thermal conduction.

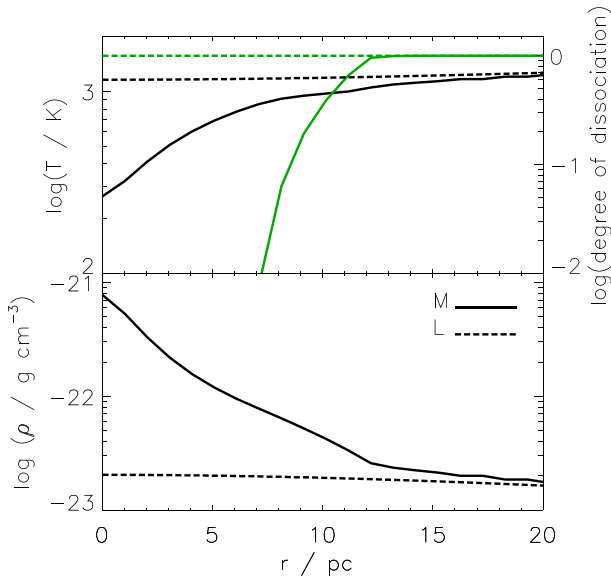


Figure 2. Upper plot: Initial temperature profile in the central cloud region for models M (black solid line) and L (black dashed line). The degrees of dissociation (green lines) are shown for each model cloud. Lower plot: Initial radial profiles of gas density for both models.

substantial heat transport towards the central cloud region is expected especially around 10 pc by pure thermodynamics. However, a present radial density gradient with the steepest slope inside of 15 pc (Fig. 2, lower plot) may hinder a mass flux into the centre.

In model M_nc thermal conduction is neglected but all other physical processes are identical to model M. Hence, the initial model is identical to the reference cloud in the distribution of each variable (like cloud mass, temperature, density, pressure).

4.2 Low-mass model clouds

In contrast, low-mass clouds are only weakly gravitationally bound with masses of $2.2 \times 10^4 M_{\odot}$ (corresponding to 9 per cent of their Bonnor-Ebert mass). They have a higher central temperature yielding a lower central density with respect to thermal equilibrium (Fig. 2). Consequently, the central pressure is low and a smaller radius is obtained. Low-mass clouds reside in a HIM with $T_{HIM} = 5.5 \times 10^6$ K and $n_{HIM} = 10^{-3} \text{ cm}^{-3}$. The reference model L takes into consideration the same physical processes as model M. It is nearly isothermal, homogeneous, and isobaric with $Z = 0.1 Z_{\odot}$, which leads to a radius

of 28 pc. The central temperature reads 1,160 K and smoothly increases to $\sim 1,340$ K at the boundary. Therefore, no molecular phase is present in the cloud as dissociation is complete even in the centre (Fig. 2, upper plot). Model L_nc is simulated equivalently to L, but without thermal conduction. Hence, the effect of an additional heat input on low-mass clouds can be studied.

5 DISCUSSION

5.1 Mass evolution

The mass-transfer rates, $\dot{M}$, are calculated by taking the difference in mass of the cloud at two subsequent time steps, i and $i + 1$. So,

$$\dot{M}_{i \rightarrow i+1} = \frac{M_{cl}(t_{i+1}) - M_{cl}(t_i)}{\Delta t}, \quad (20)$$

with $\Delta t = 1$ Myr. The mean mass-transfer rate is given by

$$\langle \dot{M} \rangle = \frac{1}{T} \sum_{i=0}^T \dot{M}(\tau_{sc} + i \Delta t), \quad (21)$$

where $T = (100 \text{ Myr} - \tau_{sc}[\text{Myr}])/\Delta t$. Thus, $\langle \dot{M} \rangle$ is based on the sound-crossing time

$$\tau_{sc} = \frac{R_{cl}}{\bar{c}_s}, \quad (22)$$

specific to each individual cloud. It is the typical time scale for a cloud to oscillate adiabatically (due to, e.g. local pressure enhancements). It is approximately the time a sound wave with mean speed $\bar{c}_s$ needs to travel into the centre and is therefore a measure for how fast a pressure balance can be established by pure hydrodynamics. We neglect cloud evolution before τ_{sc} in computing $\langle \dot{M} \rangle$ in order to exclude dynamical superimposing of repulsion resulting from initial onset of mass transfer. The cloud mass at a particular time, t_i , is given by

$$M_{cl}(t_i) = \sum_{j=1}^N \rho_j(t_i) V_j(t_i) = V \sum_{j=1}^N \rho_j(t_i), \quad (23)$$

i.e. the direct sum is taken over all grid cells j that are identified to belong to the cloud with constant volume V of a grid cell. A cell j is part of the cloud if both $\rho_j > 100 \rho_{amb}$ and $T_j < 5 \times 10^6$ K. The values for $\dot{M}$ calculated by equation (20) coincide with the values obtained by $\dot{M} = 4\pi r^2 \rho(r) v(r)$ at large distances r from cloud.

It is evident from Fig. 3 that ambient gas condenses onto clouds with thermal conduction. In general, the process of condensation applies to massive and low-mass clouds. By this, the initial mass of the massive (low-mass) cloud is increased by ≈ 1 per cent (≈ 2 per cent) after 100 Myr of evolution. The condensation rates increase over time

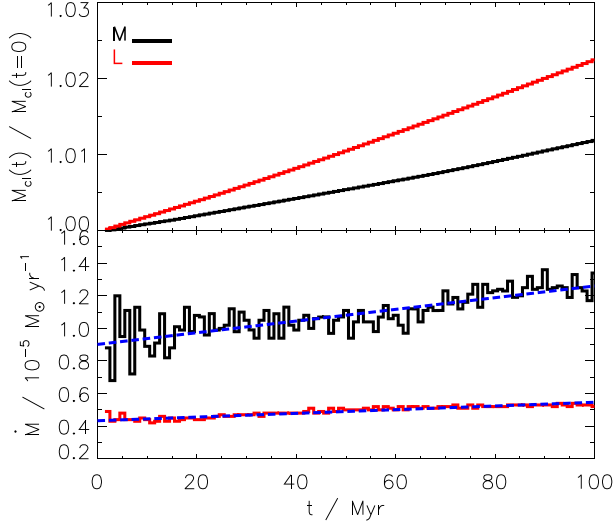
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Figure 3. Mass evolution for all clouds with thermal conduction. *Upper plot:* Evolution of cloud masses normalized to the respective initial cloud mass. *Lower plot:* Rates of condensation of ambient gas onto the clouds ($\dot{M} > 0$) with respective linear regression (blue dashed lines).

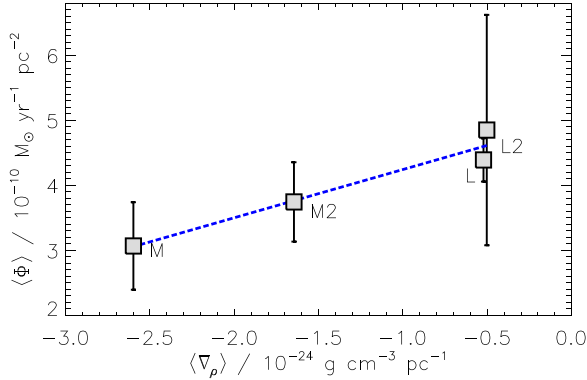


Figure 4. Mean mass flux due to thermal conduction onto clouds with respect to mean radial density gradient ($\langle \Phi \rangle > 0$ means condensation). The blue dashed line is the linear regression with a slope of $51 \text{ pc}^2 \text{ yr}^{-1}$. The errorbars give the standard deviation for each model cloud.

specific to cloud mass: $3.6 \times 10^{-8} M_{\odot} \text{ yr}^{-1} / \text{Myr}$ (massive cloud) and $1.1 \times 10^{-8} M_{\odot} \text{ yr}^{-1} / \text{Myr}$ (low-mass cloud) as can be seen in Fig. 3 (lower plot). The mean condensation rate for the massive cloud is more than twice the rate observed at the low-mass cloud at each and every time. However, the mean mass flux

$$\langle \Phi \rangle = \frac{1}{100 - \tau_{\text{sc}}} \sum_{i=\tau_{\text{sc}}}^{100} \frac{dM_{\text{cl}}/dt}{4\pi R_{\text{cl},i}^2} \quad (24)$$

is higher in model L as shown in Fig. 4. In models M and L the mean mass fluxes of condensing HIM are directed towards the clouds (i.e. $\langle \Phi \rangle > 0$). Since mass diffusion depends on density gradient, we study the density gradient between cloud centre and boundary, i.e.

$$\langle \nabla_{\rho} \rangle = \frac{1}{100 - \tau_{\text{sc}}} \sum_{i=\tau_{\text{sc}}}^{100} \frac{\rho_i(R_{\text{cl}}) - \rho_i(r=0)}{R_{\text{cl},i}}. \quad (25)$$

We extend the definition of $\langle \nabla_{\rho} \rangle$ until the cloud centre, because even in the massive cloud M there is a tiny amount of HIM mixing into the

Table 2. *Row 1:* Mean mass-transfer rate for each cloud as being computed between its specific τ_{sc} (cf. equation (22)) and 100 Myr. For comparison, the analytical rates $\dot{M}_{\text{CM}}$ of CM77 (equation (11)) and $\dot{M}_{\text{DB}}$ (Dalton & Balbus 1993, equation (14)) are provided in rows 2 and 3, respectively. Positive (negative) mass-transfer rates correspond to condensation (evaporation). *Row 4* contains the mean mass flux (mean mass-transfer rate per unit surface-area). Model clouds M2 and L2 only differ in metallicity from models M and L, respectively, which in turn implies a difference in radius. Thermal conduction is present in all four model clouds.

mean mass-transfer rate	Model			
	M	M2	L	L2
$\langle \dot{M} \rangle [10^{-6} M_{\odot} \text{ yr}^{-1}]$	11.1	8.8	4.9	1.5
$\dot{M}_{\text{CM}} [10^{-4} M_{\odot} \text{ yr}^{-1}]$	−10.3	−7.2	−4.9	−0.8
$\dot{M}_{\text{DB}} [10^{-4} M_{\odot} \text{ yr}^{-1}]$	−11.1	−7.8	−5.3	−0.9
$\langle \Phi \rangle [10^{-10} M_{\odot} \text{ yr}^{-1} \text{ pc}^{-2}]$	3.2	3.7	4.4	4.8

central region (cf. Section 5.2). From Fig. 4 we learn that a steeper radial density gradient in the cloud corresponds to a lower $\langle \Phi \rangle$.

Both $\langle \nabla_{\rho} \rangle$ and $\langle \Phi \rangle$ are the respective average values taken between τ_{sc} and 100 Myr of evolution. To quantify the $\langle \nabla_{\rho} \rangle$ - $\langle \Phi \rangle$ relation we add two further model clouds: M2 has a mean radius of 41 pc lying between models M and L and L2 has a mean radius of 11 pc, which is smaller than the radius of model L (M2 and L2 are identical to models M1.me and L1.me with solar metallicity in SH21). The mean mass fluxes result in $\langle \Phi \rangle = 3.7 \times 10^{-10} M_{\odot} \text{ yr}^{-1} \text{ pc}^{-2}$ (model M2) and $\langle \Phi \rangle = 4.8 \times 10^{-10} M_{\odot} \text{ yr}^{-1} \text{ pc}^{-2}$ (model L2). In order to be consistent with the models in SH21 the cooling efficiency of the plasma depends on Z . Metal lines dominate the line cooling for $T \gtrsim 2 \times 10^4 \text{ K}$ (Boehringer & Hensler 1989). Because the cloud boundary is at $T \lesssim 10^4 \text{ K}$ an enhanced metallicity is negligible for determining $\langle \Phi \rangle$. All four models are well approximated by a linear $\langle \nabla_{\rho} \rangle$ - $\langle \Phi \rangle$ relation with a slope of $51 \text{ pc}^2 \text{ yr}^{-1}$.

Density gradient $\langle \nabla_{\rho} \rangle$ and mass flux $\langle \Phi \rangle$ strongly correlate inversely with a Pearson's R-value of -0.977 . The correlation is significant on a level of 0.05 (p-value of 0.023). We conclude that homogeneity of clouds leads to an enhanced mass flow of ambient material onto clouds and this mass flow is due to thermal conduction only. The standard deviation for the smallest model cloud L2 is very high (Fig. 4), because its radius changes a lot during its evolution. So, the results for the nearly homogeneous clouds cannot be interpreted uniquely and we cannot propose for sure that $\langle \Phi \rangle$ converges to a constant value. Here we emphasize, that $\langle \Phi \rangle$ is due to thermal conduction only. If thermal conduction is not accounted for, there is no mass flow towards a hydrostatic cloud in thermal equilibrium including plasma cooling (cf. Section 5.2). Analytical mass-transfer rates predict evaporation for all clouds (Table 2) and are thus not appropriate to model time-dependent conditions for thermal conduction.

5.2 Mixing with ambient gas

According to the continuous mass gain for the heat conducting clouds M and L we suppose the velocity field of ambient gas to be directed towards the clouds over the entire evolution indicating condensation of ambient gas. In Fig. 5 condensation is illustrated with maximum gas velocities close to 43 km s^{-1} . As expected, in both models without thermal conduction there is only a tiny propagating net inflow of ambient HIM due to gravity. The inflow velocity does not exceed 0.1 km s^{-1} . So, their density structures are not changed at all irrespective of cloud mass. The result is not surprising as it proves the correct numerical treatment of hydrodynamics in an equilibrium

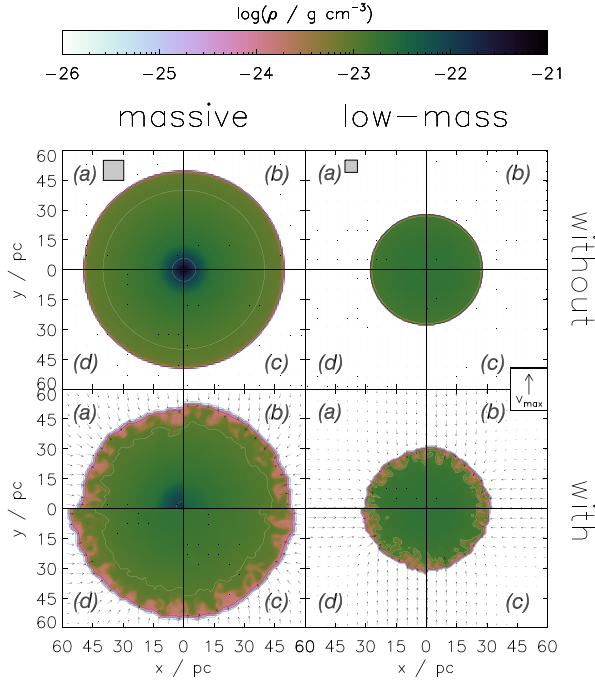


Figure 5. Evolution of density in the model clouds. The plots show the section plane (x, y) through cloud centre. The snapshots are taken after 25 Myr (a), 50 Myr (b), 75 Myr (c), and 100 Myr (d) of evolution for clouds without (upper panel) and with (lower panel) thermal conduction. Iso-density contours are shown for $\rho = 10^{-25} \dots 10^{-21} \text{ g cm}^{-3}$ separated by $\Delta(\log \rho) = 1$ dex (with the inner two contours being white just for a better contrast). Black arrows illustrate the motion of gas scaling linearly with respect to the maximum velocity $v_{\text{max}} = 43 \text{ km s}^{-1}$ as shown in the box to the right. The grey square in the upper left corner is spanned by 10×10 cells of finest numerical resolution ($\Delta x = 1.0 \text{ pc}$ for massive clouds and $\Delta x = 0.6 \text{ pc}$ for low-mass clouds).

setup. Thus, only clouds with thermal conduction condense material at a reasonable rate (Table 2). The central density peak in model M is smeared out by time and the cloud homogenizes before 50 Myr hence affected by thermal conduction. By this, the central cooling efficiency is reduced.

Condensation of ambient material and thus an efficient mixing with HIM is only possible in clouds with thermal conduction, which coincides with the findings of VH07b. In Fig. 6 the mixing of HIM into the cloud is shown by the fraction of HIM to cloud gas, f_{HIM} . In models M and L an outer mixing zone evolves in thickness where f_{HIM} can easily exceed 10 per cent. HIM is mixed with cloud material and is transported into deeper regions of the cloud while without thermal conduction the clouds are not enriched by HIM. The filling factor of HIM in model L is larger than in M indicating that accreted material can be distributed easier. After 100 Myr the entire cloud M is interspersed with at least traces of HIM.

We measure the time at each cloud radius, which is needed by the mixing front to enhance the HIM fraction at this particular radius to 1 per cent. In Fig. 7 it is obvious that in model M a HIM fraction of 1 per cent is reached later at every radius compared to model L. In model M, the mixing front only propagates until $0.3R_{\text{cl}}$ (corresponding to a radius of 15 pc) within 100 Myr. At exactly this radius the initial central density core profile flattens (see fig. 2 in SH21). We emphasize that in model M the central cloud region will be mixed with HIM on a much longer time scale than 100 Myr.

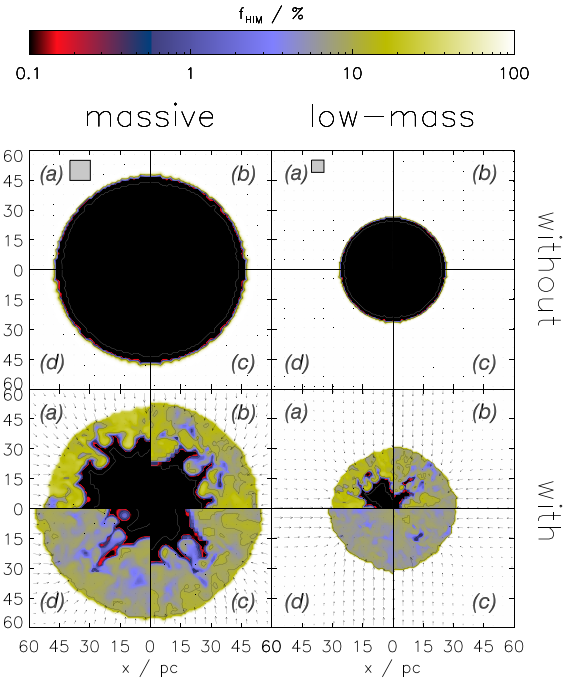


Figure 6. Same as in Fig. 5, but for the fraction f_{HIM} of ambient HIM mixed into the clouds. Isocontours are shown for $f_{\text{HIM}} = 0.1, 1$ and 10 per cent. Inside of the white contour there is no HIM at all.

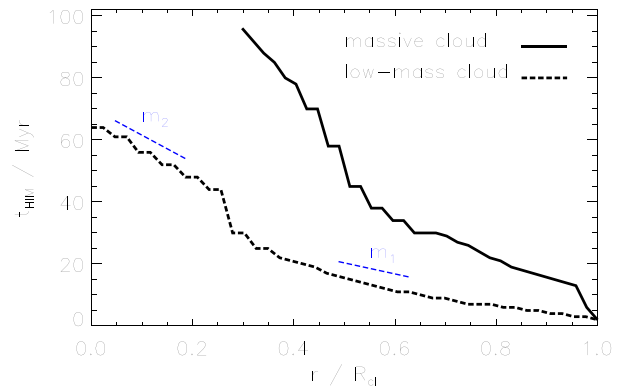


Figure 7. Time after which fraction of condensed HIM reaches 1 per cent in models M (solid line) and L (dashed line) at a given cloud-centric distance. The short blue lines represent the piecewise slopes $m_1 = -1.5 \text{ Myr pc}^{-1}$ and $m_2 = -2.9 \text{ Myr pc}^{-1}$ for the dashed line. The radii are normalized to the respective cloud radius.

Even though the steep density gradient gets shallower by time (cf. Fig. 5) the propagation velocity of the mixing front is substantially lowered. In contrast, in the low-mass cloud, which has an almost flat radial density profile, the mixing front approaches the cloud centre within 65 Myr. Around $0.28R_{\text{cl}}$ ($\approx 8 \text{ pc}$) the mixing front in model L decelerates and its velocity remains on a lower level inside this central cloud region. By taking the inverse of the piecewise slopes we can estimate the propagation velocity in model L from Fig. 7: for the outer 20 pc (from R_{cl} to $r = 8 \text{ pc}$) the mixing front needs 30 Myr yielding $m_1^{-1} \approx 0.67 \text{ pc Myr}^{-1}$. Likewise, $m_2^{-1} \approx 0.35 \text{ pc Myr}^{-1}$ for the central region inside $r = 8 \text{ pc}$ in model L.

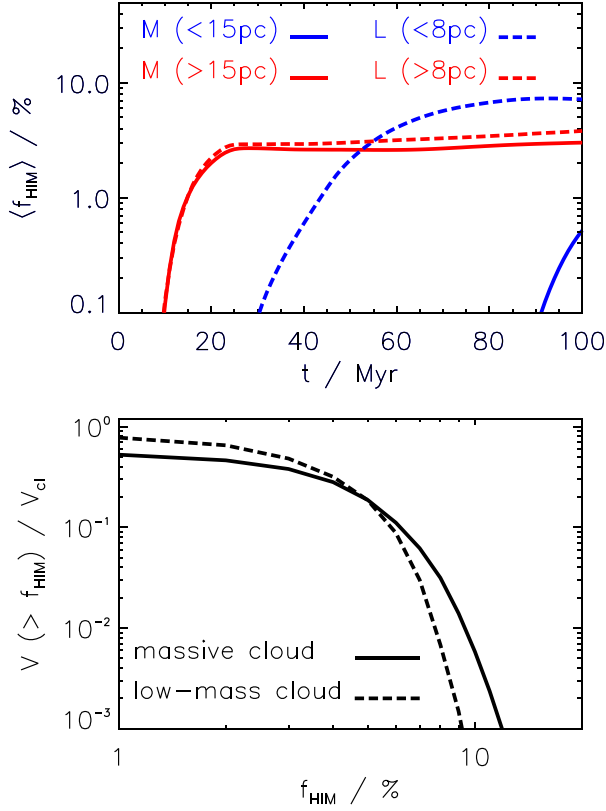
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Figure 8. *Upper plot:* Evolution of mean HIM fraction in model M inside and outside a radius of 15 pc (solid lines) and in model L inside and outside a radius of 8 pc (dashed lines). *Lower plot:* Normalized cumulative distribution of total cloud volume containing mixed HIM above a certain fraction after 100 Myr of evolution.

In order to analyze the change in propagation velocity of mixing fronts in models L at 8 pc and M at 15 pc we show in Fig. 8 (upper plot) the average HIM fractions in the respective inner and outer region in both clouds. As expected, in model M there is almost no HIM inside of 15 pc while radially outwards the fraction is roughly constant after 20 Myr with $f_{\text{HIM}} \approx 3$ per cent. In model L the mixing evolves distinctly. Accreted HIM is quickly transported into the inner region, because there is no density gradient retarding the mixing front. After 55 Myr the mean of f_{HIM} inside a radius of 8 pc exceeds the HIM fraction more outside. HIM accumulates centrally in cloud L, but does not turn back to the outer regions due to gravity. Both dashed lines in the upper plot in Fig. 8 may converge after 100 Myr. We thus conclude that the density profile of a cloud not only affects the spatial distribution of accreted gas, but also the direction of mixing: in model M the radial gradient is sufficiently steep to provide a reasonable obstacle to slow down propagation of accreted gas and the cloud is filled up with HIM from the outside to the inside. In contrast, if the cloud is nearly homogeneous, condensed HIM straightly gets through into the centre, accumulates there and the cloud fills up in the reverse direction.

It is readily apparent from Fig. 8 (lower plot) that only 50 per cent of cloud volume in model M contain more than 1 per cent of HIM after 100 Myr, while this fraction reaches 80 per cent in cloud L. However, the cloud regions with highest HIM concentrations are observed in model M: within 0.7 per cent of the total cloud volume the HIM fraction exceeds 10 per cent. We conclude again that the radial

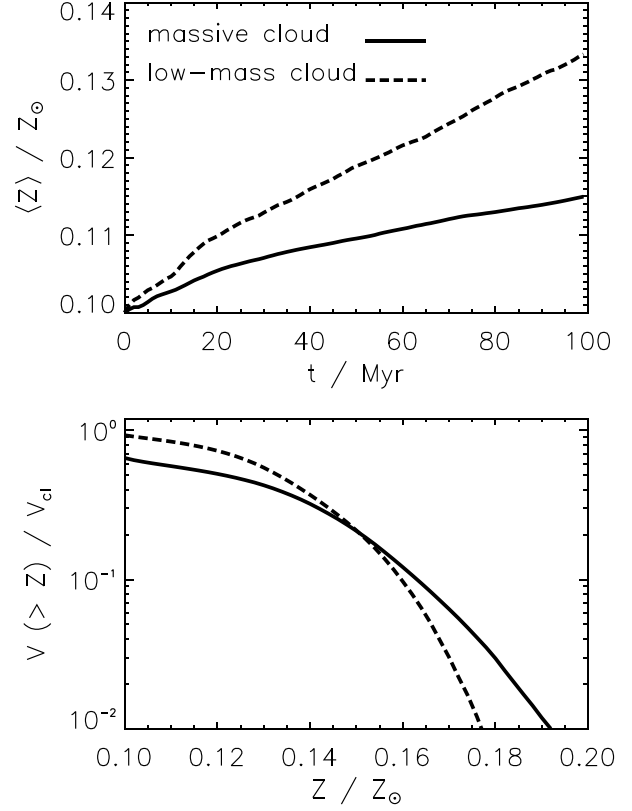


Figure 9. *Upper plot:* Evolution of density-weighted metallicity in model M within a radius of 40 pc (solid line) and in L within a radius of 20 pc (dashed line). *Lower plot:* Normalized cumulative distributions of metallicity in both model clouds after 100 Myr of evolution.

density gradient in model M accumulates in higher concentrations in the outer regions. Consequently, condensed HIM is much easier distributed in homogeneous clouds. Hence, the above result supports the findings in Sander & Hensler (2019), that self-gravity crucially affects the evolution of even sub-Jeans clouds. Mixing in moving HVCs without thermal conduction is analyzed by Gritton, Shelton & Kwak (2014). In their 3D simulations they observe fractions of halo gas in their model cloud B between 1 and 10 per cent for different ionization stages of oxygen and carbon after 100 Myr (see their figs 3 and 4).

A simple way to trace the HIM fraction that condenses onto the clouds is realized by studying the mean metallicity in each particular grid cell. The mean metallicity in both models M and L increases on average, while it is more enhanced in the low-mass cloud (upper plot in Fig. 9). The growth rates due to condensation of ambient gas are roughly $0.012 Z_{\odot} / 100 \text{ Myr}$ (model M) and $0.014 Z_{\odot} / 100 \text{ Myr}$ (model L) and are rather low. The fact that the clouds condense HIM at a rather low rate coincides with the findings of De Cia et al. (2021) that low-metallicity gas falling through the hot halo of the Milky Way towards the Galactic disk as HVCs sustains observed chemical inhomogeneities on scales of $10 \dots 100 \text{ pc}$. The HVCs thus would be different in metallicity for reasonable time scales (at least 100 Myr), because the cloud gas does not mix effectively with surrounding halo gas.

By inspecting the lower plot in Fig. 9 one finds that the regions of highest metallicity are found in the massive model cloud. The dynamics of metals is consistent with the mixing of condensed HIM

(see discussion above) hence supporting the reasoning, that it is easier for accreted gas to spread in a cloud with only a shallow radial density gradient. Vice versa, if a substantial density gradient is present in the central region then the accreted material accumulates preferentially in the outskirts. The metallicity in the outer regions in stratified clouds can exceed the metallicities found in homogeneous clouds. The distribution of HIM in the model clouds traces their enrichment by heavy elements if a metal contrast is given between cloud and HIM.

5.3 Structure of interface

During the evolution of both models M and L a transition zone forms at the interface between cloud and ambient, hot gas. From the upper plot in Fig. 10 it is obvious, that the condensation velocity of ambient gas is negative over the entire evolution thus condensation is continuous. The condensation velocity is between $-20 \dots -10 \text{ km s}^{-1}$ and is thus much smaller than the sound speed in the hot ambient gas and nearly one order of magnitude smaller than the local sound speed. But within the transition zone the condensation velocity is remarkably lowered by one order of magnitude to $\lesssim 2 \text{ km s}^{-1}$. This velocity is distinctly smaller than the sound speed in the local warm cloud gas. By this, material is accreted faster than being distributed within the clouds. The mean velocity is lowered, because the temperature in the clouds is much lower at a few 10^3 K . The inflow of hot ambient gas enhances the particle density thus increasing the cooling efficiency. By this, transfer of momentum from hot conducted electrons to cloud gas is compensated by an enhanced cooling. So, gas heating by conduction is not sufficient to overcome cooling in the outskirts of the clouds. The mixing of HIM with the cloud gas in deeper layers needs longer than the accretion of HIM: while HIM is accreted with a velocity of ~ 12 percent of local sound speed it is transported into deeper cloud regions with a local Mach number of 0.03 only. HIM is thus both condensed and distributed within the clouds only subsonically (according to the respective local sound speed).

Within the transition zone the temperature raises substantially from a few 10^3 K until it reaches T_{amb} . As the cooling function scales with ρ^2 and sharply increases towards $T \approx 10^4 \text{ K}$ the cooling efficiency has a maximum within the transition zone: for radii smaller than the inner boundary of the transition zone the temperature is too low. Radially outward of the transition zone the density is too low. Thus balance between maximum of cooling and energy input by thermal conduction is reached within the transition zone, i.e. here $\lambda_F = 1$, which is shown by red stars in Fig. 10. At $\lambda_F = 1$ an equilibrium temperature of $2 \dots 3 \times 10^5 \text{ K}$ is reached. Temperatures around 10^4 K can only be reached if the density is about 100 times higher, which is the case at smaller cloud-centric distances and can be seen by inspecting the density profile in Fig. 10. Outside of the transition zone λ_F reaches values of $\gtrsim 5 \times 10^2 R_{\text{cl}}$. Inside of the transition zone the cooling remains on a reasonable level thus λ_F immediately drops below $10^{-2} R_{\text{cl}}$ for $r < r(\lambda_F = 1)$ and thermal energy is lowered such that $E_{\text{therm}} \gtrsim E_{\text{grav}}$. By this, hot ambient gas condenses onto the clouds while theoretical mass-loss rates predict evaporation (Table 2).

We compare our findings to the theory described in CM77. They infer from the enthalpy flux that the interface between cloud and HIM is divided into three zones: (1) an *inner classical zone* extending from R_{cl} outward to a distance $R_{\text{sat}, 1}$, where the heat flux becomes saturated. The heat flux in the inner zone is classical, because the temperature and hence the conductivity are small. (2) A *saturation zone* extends between $R_{\text{sat}, 1}$ and $R_{\text{sat}, 2}$. The temperature increases steeply, such that the temperature gradient is large. Thus, the heat

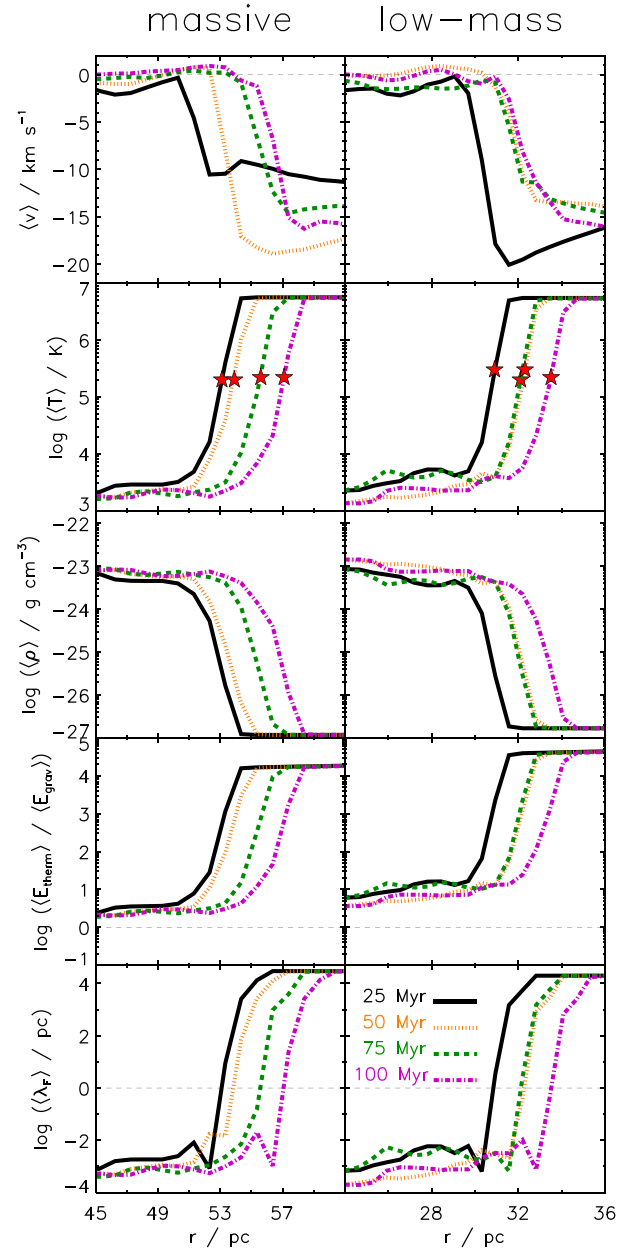


Figure 10. From top to bottom: Evolution of angle-averaged, radial profiles for inflow velocity, temperature, density, ratio of thermal to gravitational energy, and Field length around the boundaries of clouds with thermal conduction shown at four different snapshots. The red stars denote the radii at which $\lambda_F = 1$, i.e. the strength of thermal conduction equals that of radiative cooling.

flux is saturated. (3) Outside of $R_{\text{sat}, 2}$, in the so-called *outer classical zone*, the temperature gradient is small, because the temperature level of the HIM is reached. Consequently, the heat flux is classical again. At these high temperatures the cooling function dominates the input of conducted thermal energy (equilibrium is marked by red stars in Fig. 10).

CM77 analytically deduce the values for $R_{\text{sat}, 1}$ and $R_{\text{sat}, 2}$ based on both the Mach number M_s in the saturated zone and σ_0 (see their equations 48 and 58, respectively). The values for M_s and σ_0 do not

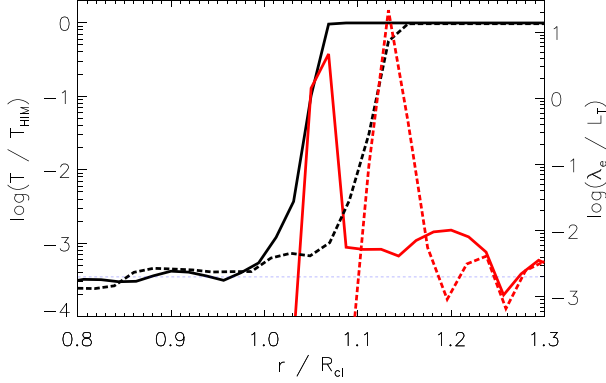
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Figure 11. Temperature profile normalized to ambient temperature over inner classical, saturation, and outer classical zone for models M (black solid line) and L (black dashed line) after 100 Myr of evolution. The ratio λ_e/L_T (red line for respective model cloud) distinguishes between classical and saturated thermal conduction (threshold $\lambda_e/L_T = 2 \times 10^{-3}$ (equation (2)) is indicated by the thin dashed line).

Table 3. Radial extension of inner classical zone (ranging from R_{cl} to $R_{sat,1}$), saturation zone (ranging from $R_{sat,1}$ to $R_{sat,2}$), and outer classical zone (above $R_{sat,2}$) according to CM77 (columns labeled ‘CM77’) and based on simulation results (columns labeled ‘Sim.’). The Mach numbers are averaged over inner classical zone (icz) or saturation zone (sz) while the ratios R_{cl}/λ_F are taken at particular zone transitions. All values are time averaged over 100 Myr. The uncertainties are standard deviations.

mean values	model M		model L	
	CM77	Sim.	CM77	Sim.
R_{cl} / pc	52 ± 3		30 ± 2	
$R_{sat,1}$ / pc	58 ± 3	54 ± 2	33 ± 2	31 ± 1
$R_{sat,2}$ / pc	77 ± 3	60 ± 4	46 ± 3	35 ± 2
$Ma(icz)$	0.09 ± 0.05	0.17 ± 0.08	0.2 ± 0.1	0.2 ± 0.1
$Ma(sz)$	0.04 ± 0.01	0.05 ± 0.02	0.04 ± 0.01	0.06 ± 0.02
$R_{cl}/\lambda_F(R_{cl})$		250		500
$R_{cl}/\lambda_F(R_{sat,1})$		0.38		0.26
$R_{cl}/\lambda_F(R_{sat,2})$		3×10^{-5}		5×10^{-5}

change in the analytically calculated saturation zone, because the temperature profile is constant in CM77. Based on the temperature profile in Fig. 11 we identify the transition zone that forms around our simulated clouds with the inner classical zone and the saturation zone from CM77.

We observe in our simulations that the saturation zone is located closer to the cloud boundary than in analytic calculations by CM77. Hence, the inner boundary of saturation zone $R_{sat,1}$ is on average closer to R_{cl} (Table 3). Likewise, the saturation zone $\Delta R_{sat} \equiv |R_{sat,2} - R_{sat,1}|$ is narrower in our simulations. From Table 3 we deduce $\Delta R_{sat} = 6 \pm 4$ pc in model M and $\Delta R_{sat} = 4 \pm 2$ pc in model L, or, in terms of the respective cloud radius, $\Delta R_{sat}/R_{cl} = 1.04 \dots 1.15$ (model M) and $\Delta R_{sat}/R_{cl} = 1.03 \dots 1.17$ (model L). This agrees with the results of VH07b. Our model clouds are distinct from the clouds considered in CM77 in terms of plasma cooling and heating, and self-gravity.

Within the saturation zone it holds $\lambda_e/L_T \gg 2 \times 10^{-3}$ (Fig. 11). Hence, according to condition (2) the heat flux is saturated. By inspecting the lower plot in Fig. 10 we find that λ_F reaches values $\gtrsim 10^4$ pc inside the saturation zone. Therefore, the condition for condensation is already fulfilled for values $R_{cl} \gtrsim 0.02 \dots 0.07 \lambda_F$, which decreases the threshold discussed in VH07b by a factor of at

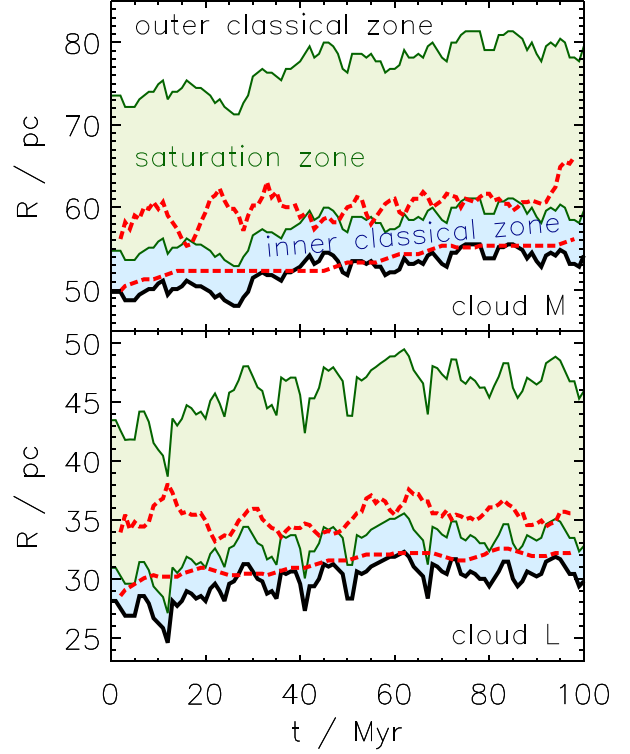


Figure 12. Evolution of structure of cloud interface in massive (upper plot) and low-mass clouds (lower plot). The colored areas regard the analytical calculations by CM77 for inner classical (blue), saturation (green), and outer classical zone (white). The red dashed lines enclose the saturation zone of our simulations.

least 3. As a consequence, condensation already sets in at larger λ_F , which correspond to cooling rates that are one order of magnitude lower as those found in VH07b.

We conclude that the condition for condensation is met within the saturation zones in our model clouds and thus ambient hot material condenses. Our results are consistent with the condensation condition $R_{cl}/\lambda_F \gtrsim 0.24 \dots 0.36$ discussed in VH07b.

5.4 Condensation rates

We already discussed the effects of thermal conduction, cloud mass, and radial density gradient on condensing hot ambient material onto the simulated clouds. Compared to analytic considerations of CM77 the interface structure in our simulated clouds is time-dependent and so are the condensation rates.

In Fig. 12 we show the analytically calculated radial interface structure based on CM77. The three zones are of constant thickness over the entire evolution of the clouds. The red lines in Fig. 12 show the respective saturation zone for our model clouds. It is visible to the eye that in model cloud L its thickness decreases by time. In model M the saturation zone gets slightly narrower (hardly noticeable in Fig. 12, upper plot). Simultaneously, the condensation rates $\dot{M}$ increase in both models (cf. lower plot in Fig. 3). Below we attempt to explain the relation between thickness of saturation zone and strength of condensation rate.

The interface between cloud and HIM structures radially into an inner classical zone, a saturation zone, and an outer classical zone due to saturated thermal conduction (cf. CM77). In the outer classical

zone it holds $R_{cl}/\lambda_F \ll 1$ (see the respective value at outer radius $R_{sat,2}$ in Table 3). Within the saturation zone λ_F drops such that $R_{cl}/\lambda_F \sim 1$. Hence, in the inner classical zone on average $R_{cl}/\lambda_F = 0.7$ (model M) and $R_{cl}/\lambda_F = 3.1$ (model L). Consequently, the condition for condensation of gas is fulfilled within both the saturation zone and the inner classical zone. By the process of condensation, hot ambient gas is transported into saturation zone and inner classical zone at a rate $\dot{M}$. Therefore, the particle density n grows within the saturation zone and the inner classical zone. Because the cooling efficiency is proportional to n^2 the cooling strength grows and hence according to equation (9) λ_F decreases. But if gas in the saturation zone cools the pressure drops and $\dot{M}$ increases. Actually, we find an anti-correlation between λ_F in the saturation zone and $\dot{M}$ with Pearson's R-values of -0.14 (model M) and -0.25 (model L). The anti-correlation is weakly significant for model M (p-value 0.16) and strongly significant for model L (p-value 0.01). We further find an anti-correlation between $\dot{M}$ and ΔR_{sat} for model M (R-value -0.17 , weakly significant with p-value 0.09) and model L (R-value -0.38 , strongly significant with p-value <0.01). From a physical point of view, the latter anti-correlation is attributed to the increased cooling efficiency by higher n due to a larger $\dot{M}$. A lower pressure leads to a narrower saturation zone. By this, the radial temperature gradient across saturation zone gets steeper. So, $\dot{M}$ increases again. We indeed find a positive correlation between the temperature gradient across the saturation zone and $\dot{M}$. For model M, the Pearson's R-value is 0.13 with a weak significance on level 0.05 (p-value 0.19). For model L, the R-value reads 0.23 and the correlation is strongly significant (p-value 0.02).

The above described 'condensation engine' does not stop if the reservoir of hot ambient gas is large compared to the cloud. As a result, we observe an increase of $\dot{M}$ and a decrease of ΔR_{sat} by time.

6 SUMMARY AND CONCLUSIONS

With our simulation suite we show both the limits of the analytical approach of thermal conduction by CM77 and extend the numerical analyses performed by VH07b. We conduct three-dimensional hydrodynamics and plasma-physics simulations of interstellar multiphase clouds at rest with respect to an ambient hot, tenuous, and highly ionized gas using the AMR code FLASH. Relevant physical processes considered numerically or semi-analytically are saturated thermal conduction, self-gravity, metal-dependent heating and cooling of the plasma by radiation and collisions, dissociation of molecules and ionization of atoms, and mass diffusion. In SH21 we simulate HVCs with thermal conduction while in the present work we analyze their resting counterparts. By comparing identical model clouds, which only differ in thermal conduction, we work out differences in evolution based solely on this physical process in presence of steep temperature gradients.

While analytical mass-transfer rates by CM77 only show evaporation, HIM condenses continuously on our model clouds during their entire evolution (Table 2). This originates from the fact that CM77 only integrate the time-independent equation for energy conservation in the heat conducting interface. But it turns out, that the rate of condensation is closely related to the effective heat flux crossing the transition zone at cloud surface: mixing of cool cloud gas with ambient hot plasma diminishes the temperature contrast between cloud and HIM and thus the effective heat flux reduces accordingly.

Regardless of mass, our model clouds are in the domain of condensation. This is in agreement with the two-dimensional results obtained by VH07b. Our simulated low-mass clouds are more easily heated up than the massive clouds, because their density is lower

and so is the cooling rate. The lack of a substantial negative radial density gradient leads to a faster distribution of condensed ambient gas in the clouds.

Condensation is found to be a valid process to mix HIM into the clouds. However, mixing is not equally efficient in all clouds and appears at a rather low rate. In model L condensed HIM is distributed more easily than in model M such that cloud and ambient gas homogenize faster. The mixing velocity in model M drops at a cloud-centric distance of ~ 15 pc where the density gradient increases. As a consequence, HIM accumulates outside of the steepest slope of the negative central density gradient. Because we use metallicity as a tracer of accreted HIM we draw two conclusions: (1) if a metallicity contrast is present, then metal enrichment of clouds is likely and (2) sub-regions with highest metallicity are found in clouds with steepest radial density profiles. The main results of our study are summarized as follows:

- (i) Saturated thermal conduction is able to disturb an initial thermal equilibrium in multiphase clouds such that condensation of ambient hot gas is initiated and maintained and a continuous mixing process is triggered.
- (ii) With thermal conduction considered, the central density peak in massive clouds levels off within 50 Myr.
- (iii) The velocity by which gas condenses onto clouds is higher than the mixing velocity inside the clouds. Hence, gas is accreted faster than being distributed.
- (iv) The shallower the radial density gradient the higher the mass flux due to condensation. We conclude that homogeneity in a cloud produces a higher material flow towards it. So as not to overestimate the condensation rate, it is important to consider self-gravity even in sub-Jeans clouds, which naturally implies a radial density gradient. This supports the results in Sander & Hensler (2019).
- (v) Self-gravity substantially affects the spatial distribution of hot ambient material condensed onto a cloud. The noticeable radial density gradient in stratified clouds prevents condensed HIM from being mixed into the central cloud region. Thus, HIM accumulates in the outer regions of stratified clouds thus the highest mixing fractions are found there. If cloud and HIM have different metal contents, the sub-regions with highest metallicity are observed in stratified clouds with a distinct radial density gradient. Consequently, a radial metallicity gradient forms. Condensed gas is more easily distributed in homogeneous clouds and hence mixing is more efficient.
- (vi) A transition zone forms at the interface between cloud and HIM due to saturated thermal conduction. The transition zone splits up into two adjacent zones with the same thermal characteristics described in CM77 for their inner classical zone and saturation zone. Compared to the analytic considerations in CM77 the saturation zones in our simulated clouds start at smaller radii, they are much narrower, and their thicknesses decrease by time. Within the transition zones the condition for condensation is met.

We conclude that plasma cooling, self-gravity, and thermal conduction play a key role in the evolution of multiphase clouds with a core-halo structure in density, temperature, and pressure. Self-gravity acts as a stabilizing process against RT instabilities, which can otherwise develop if the clouds are massive enough (Murray et al. 1993). The process of condensation provides a plausible, yet not efficient mechanism to mix ambient HIM into the cloud. If mixing is representative for HVCs, then their low metallicities (Wakker 2001) are presumably a consequence of reduced mixing. By this, the findings of De Cia et al. (2021) can be interpreted that HVCs are observed to keep their chemical inhomogeneities over dozens of Myr. As our most important result we deduce that clouds can resist

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thermal evaporation even in those parameter ranges, where hitherto analytical approaches predict evaporation.

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DATA AVAILABILITY

The data underlying this article (simulation results and analysis scripts) will be shared on reasonable request to the corresponding author.

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APPENDIX A: RESOLUTION STUDY

A comparison of the numerical to the analytical solution of pure thermal conduction in three spatial dimensions is already shown in Sander & Hensler (2021). Both solutions coincide on a reasonable level of accuracy.

In order to also proof the sufficiency of numerical resolution for the model clouds used in this work, we perform simulations with equal setup and halven the finest numerical resolution. To account for a finer grid spacing, we increase the level of refinement from 6 (massive: $\Delta x = 1.02$ pc, low-mass: $\Delta x = 0.63$ pc) to 7

²See <http://vsc.ac.at/>

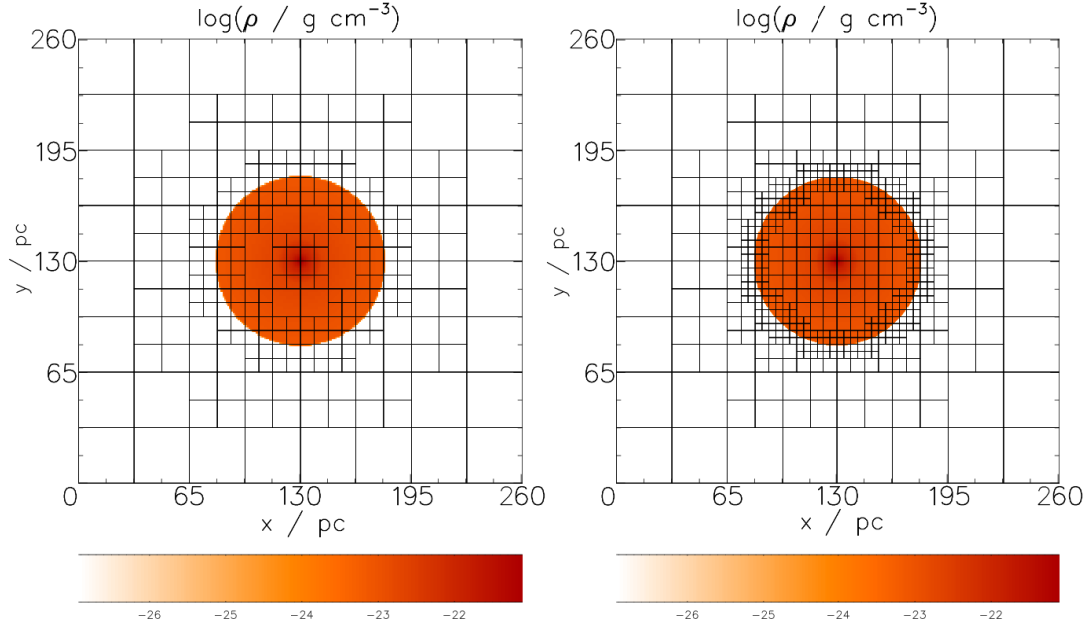


Figure A1. Density slice of a massive cloud with 6 levels of refinement ($\Delta x = 1.02$ pc, *left plot*) and 7 levels of refinement ($\Delta x = 0.51$ pc, *right plot*). The blocks of the computational mesh are shown with one block consisting of $8 \times 8 \times 8$ grid cells with respective side lengths of Δx .

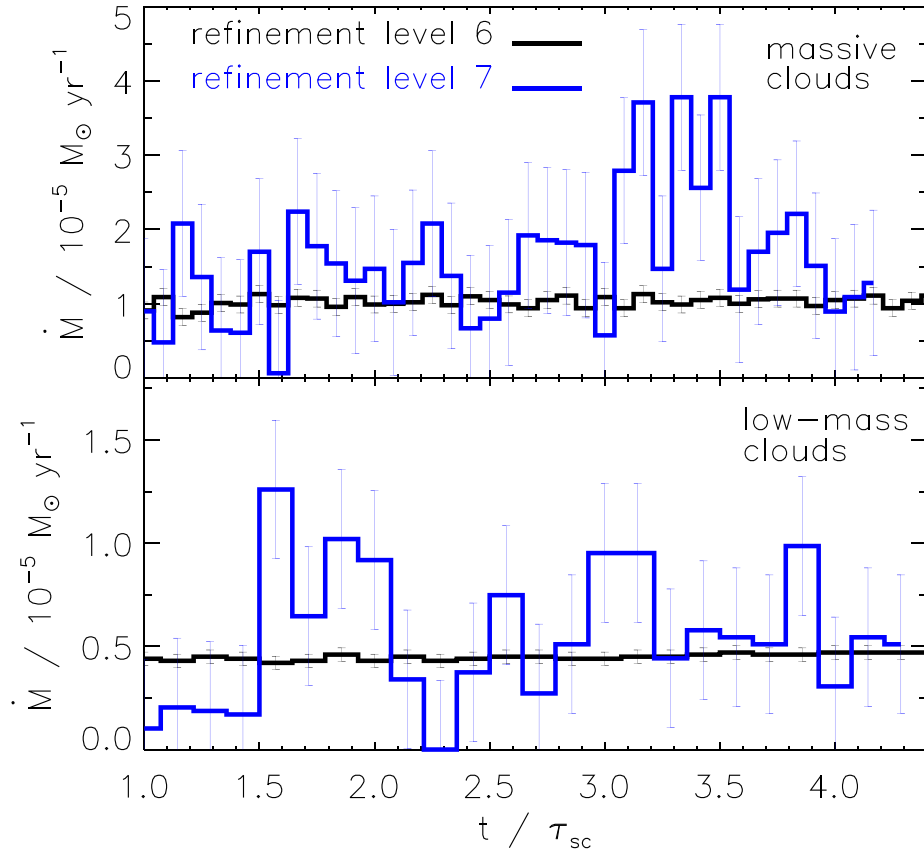


Figure A2. *Upper plot:* Mass-transfer rates of model M (black line) and its highly resolved variant (blue line). *Lower plot:* The same for model L.

1326 *B. Sander and G. Hensler***Table A1.** Model clouds M and L in comparison to their respective highly resolved (h.r.) variants. We compare the mean mass-transfer rates $\langle \dot{M} \rangle$ with standard deviations $\sigma(\langle \dot{M} \rangle)$ and the mean mass fluxes $\langle \Phi \rangle$.

values	Model			
	M	M (h.r.)	L	L (h.r.)
refinement level	6	7	6	7
Δx [pc]	1.02	0.51	0.63	0.31
$\langle \dot{M} \rangle$ [$10^{-6} M_{\odot} \text{ yr}^{-1}$]	11.1	11.2	4.9	5.0
$\sigma(\langle \dot{M} \rangle)$ [$10^{-6} M_{\odot} \text{ yr}^{-1}$]	1.1	9.8	0.3	3.4
$\langle \Phi \rangle$ [$10^{-10} M_{\odot} \text{ yr}^{-1} \text{ pc}^{-2}$]	3.2	3.7	4.4	5.1

(massive: $\Delta x = 0.51$ pc, low-mass: $\Delta x = 0.31$ pc). The mesh is adaptively refined based on the local density gradient, not on sound speed. This means, grid refinement can be coarse-grained in hot, but homogeneous regions (as is the situation far away from the cloud in the hot ambient medium). Furthermore, the mesh is conservatively refined, i.e. any two adjacent mesh regions (in terminology of the FLASH code they are called ‘blocks’) are only allowed to differ in one refinement level. Therefore, cold regions close to steep density gradients at the cloud boundary and the interface itself between cloud and ambient medium are at highest possible refinement level over the entire evolution. Conclusively, this interface, where condensation

takes place, is definitely resolved at level 7 in the highly resolved models, while the internal cloud part is equally refined in both resolution setups. To illustrate the above mentioned, we show the numerical grid for the two massive model clouds in Fig. A1.

We compare the model clouds with low and high resolution for about 4 sound-crossing times (cf. equation 22). The highly resolved models show larger variations in the magnitude of their mass-transfer rates. Consequently, they have much larger standard deviations (Fig. A2). Nonetheless, their mean rates $\langle \dot{M} \rangle$ do not differ by more than 2 per cent (cf. Table A1). The pairwise differences in the mean mass fluxes $\langle \Phi \rangle$ are a bit higher, because they relate to the surface of respective cloud, which scales with the square of cloud radius. A finer resolution leads to a slightly smaller cloud radius. In summary,

- (i) The mean mass-transfer rates of a low resolved and a highly resolved cloud deviate by less than one standard deviation.
- (ii) The Field length λ_F is already sufficiently resolved by 6 levels of refinement.

We hence conclude that the numerical resolution used in models M and L is sufficient to analyze processes related to thermal conduction.

This paper has been typeset from a $\text{\LaTeX}$ file prepared by the author.

7.6 STAR FORMATION IN COMPACT HIGH-VELOCITY CLOUDS

Within this section the evolution of modelled CHVCs moving through a hot halo is investigated in detail. Fourteen models are compared one with another as well as with observations and other simulations. Thermal conduction leads to a constant and filamentary stripping of cloud gas while otherwise large clumps are stripped off by KHI. Depending on the considered physical processes, some of the clouds are disrupted and some of the surviving clouds start to develop Jeans instabilities. In the most realistic model a natural explanation is found why no star formation is triggered. The $H\alpha$ intensities of the model clouds are similar to observed intensities in CHVCs.

The idea for the manuscript was developed together by my supervisor Prof. Dr. G. Hensler and me. Prof. Hensler reviewed the manuscript iteratively. The manuscript is co-authored by Prof. Hensler and me.

My own contential contributions include:

- all simulations and calculations presented,
- all plots are based on self-written IDL scripts.

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Physical effects on compact high-velocity clouds in the circumgalactic medium

Bastian Sander^{1,2★} and Gerhard Hensler^{1★}

¹*Department of Astrophysics, University of Vienna, Türkenschanzstraße 17, A-1180 Vienna, Austria*

²*Fraunhofer Institute for Factory Operation and Automation IFF, Sandtorstraße 22, D-39106 Magdeburg, Germany*

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ABSTRACT

We numerically investigate the evolution of compact high-velocity clouds (CHVCs) passing through a hot, tenuous gas representing the highly ionized circumgalactic medium (CGM) by applying the adaptive-mesh refinement code FLASH. The model clouds start from both hydrostatic and thermal equilibrium and are in pressure balance with the CGM. Here, we present 14 models, divided into two mass categories and two metallicities each and different velocities. We allow for self-gravity and thermal conduction or not. All models experience mass diffusion, radiative cooling, and external heating leading to dissociation and ionization. Our main findings are (1) self-gravity stabilizes clouds against Rayleigh–Taylor instability, which is disrupted within 10 sound-crossing times without; (2) clouds can develop Jeans-unstable regions internally even though they are initially below Jeans mass; (3) all clouds lose mass by ram pressure and Kelvin–Helmholtz instability; (4) thermal conduction substantially lowers mass-loss rates, by this, extending the clouds’ lifetimes, particularly, more than doubling the lifetime of low-mass clouds; (5) thermal conduction leads to continuous, filamentary stripping, while the removed gas is heated up quickly and mixes efficiently with the ambient CGM; (6) without thermal conduction the removed gas consists of dense, cool, clumpy fragments; (7) thermal conduction might prevent CHVCs from forming stars; and (8) clouds decelerated by means of drag from the ambient CGM form head-tail shapes and collapse after they reach velocities characteristic for intermediate-velocity clouds. Conclusively, only sophisticated modelling of CHVCs as non-homogeneous and non-isothermal clouds with thermal conduction and self-gravity explains observed morphologies and naturally leads to the suppression of star formation.

Key words: conduction – diffusion – hydrodynamics – methods: numerical – ISM: clouds.

1 INTRODUCTION

The coexistence of only gas phases in the interstellar medium (ISM) can theoretically be well described by the three-phase medium introduced by McKee & Ostriker (1977) (applied to the halo by Wolfire et al. 1995). Stratified, cool, dense, molecular gas, and cold neutral matter (CNM, $T \lesssim 100$ K) are surrounded by a warm, neutral to slightly ionized medium (WIM, $8000 \lesssim T/\text{K} \lesssim 10^{4.2}$). A hot, dilute, and highly ionized medium (HIM, $T \gtrsim 10^{5.5}$ K) provides the third thermally stable phase. Gas with temperatures in one of the two ranges between the three phases is thermally unstable and separates into either the cooler or the hotter phase by thermal instability (Field 1965). These two unstable phases provide continuous pressure transitions between the stable phases (Wolfire et al. 1995; Vazquez-Semadeni 2012). The physical conditions assumed for the HIM are observed in the circumgalactic medium (CGM) in the halo of our Milky Way (Tumlinson, Peebles & Werk 2017). Among the stable phases, the WIM has the highest volume-filling factor, which raises from ~ 0.4 in the galactic plane to ~ 0.6 at a galactic height of $z = 1$ kpc (Morfill & Scholer 1987) and contains more than 50 per cent of the total mass of the ISM (Tielens 2010). However, more recent

observations and simulations reveal a much more complex dynamical and thermal state of the ISM that cannot be explained within the simplified theoretical framework of McKee & Ostriker (1977): stars, electromagnetic radiation, cosmic rays, dust, magnetic fields, self-gravity, and turbulence contribute to the overall energy budget and dynamics and hence affect gas phases in various states (molecular, atomic, ionized, excited). These mutual interactions lead to many open questions (Cox 2005; McKee & Ostriker 2007). Especially for the CGM, it is not yet clear which small-scale physics can be deduced from the observed multiphase ionization structure (Tumlinson et al. 2017).

By connecting the high pressure of CNM, which manifests in the form of gravitationally bound clouds, with the low pressure of CGM, very steep gradients in both density and temperature are established in the WIM. A steep gradient in temperature between clouds and the hot gas results in the onset of a heat flux towards the cool cloud. Analytical approaches propose the evaporation of clouds in a few dynamical times (see, e.g. Balbus & McKee 1982; Giuliani 1984; Nagashima, Inutsuka & Koyama 2007). However, since interstellar clouds are observed ubiquitously, two scenarios are intuitive to assume (1) the accretion rate of intergalactic gas by the Milky Way and other galaxies in the local Universe is sufficient for replenishment (Bland-Hawthorn 2009; Richter 2012, 2016). (2) The clouds are stabilized against thermal (evaporation) and dynamical

* E-mail: bastian.sander@iff.fraunhofer.de (BS); gerhard.hensler@univie.ac.at (GH)

(e.g. ram-pressure stripping) interactions by processes that lead to accretion of mass, prevention of mass-loss, and suppression of hydrodynamical instabilities. Vieser & Hensler (2007a) figured out thermal conduction to play a key role in stabilizing interstellar clouds over reasonable evolutionary times. The numerical investigations of Pittard et al. (2004) have shown that evaporation is the most likely mechanism for mass loading of spherical accretion flows from embedded clouds if the ratio of accreted to ejected mass is high. Such a situation is observed in galactic winds from SNe II in the dwarf starburst galaxy NGC 1569 (Martin, Kobulnicky & Heckman 2002) and simulated by Recchi & Hensler (2007). Pioneering investigations on the effect of thermal conduction on cool clouds were performed by Cowie & McKee (1977) and McKee & Cowie (1977), who applied a saturation formalism to gaseous clouds in order to correctly describe the heat flux in the presence of steep temperature gradients. However, they did it in a stationary fashion, neglecting time-dependent contributions to the energy equation like mass diffusion, evaporation, condensation, metallicity (which, in turn, acts on the local cooling and heating functions of the plasma), and surface instabilities. Complementary to that, Begelman & McKee (1990) developed a time-dependent theory of two-phase media under consideration of gas heating and radiative cooling, and with mass exchange via condensation and evaporation.

The variety of interstellar clouds comprises the distinct class of high-velocity clouds (HVCs, see Wakker 2001; Richter et al. 2001; Faridani et al. 2014), which have line-of-sight velocities that are incompatible with the rotation speed of our Galaxy at their respective position. Seen in the local standard of rest (LSR), their velocities are mainly negative implying a net infall to the Galaxy. They are observed primarily at higher galactic latitudes ($|b| \gtrsim 25^\circ$, Westmeier, Brüns & Kerp 2005a) and throughout the entire electromagnetic spectrum. HVCs are multiphase clouds, which have a core-halo structure in temperature, density, and ionization fraction, and show a sub-solar metallicity. Due to a lack of suitable foreground and background sources, the determination of their distances is still a challenge. Distance brackets put HVC complexes in the Milky Way at distances between 1 and 30 kpc from the Sun. However, quite different from the HVC complexes are isolated and compact HVCs (CHVCs), which are not being resolved by a 0.5° beam and having angular sizes less than 2° of full width at half-maximum (Braun & Burton 1999). They distribute rather uniformly over the sky (Burton & Braun 1999; Braun & Burton 2000) and neither connect spatially nor kinematically to any other feature of neutral hydrogen down to a column density of $\sim 1.5 \times 10^{18} \text{ cm}^{-2}$ (van Woerden et al. 2004). This lack of association leads to the definition of a distinct class of HVCs that may have originated under the same environmental conditions, share a similar evolutionary history, and are in the same physical state. Among CHVCs one finds clouds with the highest LSR velocities. CHVCs can be located several 100 kpc away from the Sun (Braun & Burton 2000).

Another puzzle is the total lack of stellar objects associated with HVCs (e.g. Siegel et al. 2005). But Izumi et al. (2014) report on star formation in Cloud 1 (one of the Digel Clouds) triggered by the impact of the close-by HVC complex H on to the Galactic disc (see also the numerical study by Lepine & Duvert (1994) for star formation induced by infall of HVCs on the Galactic disc). Indeed, from observations it can be concluded that all necessary conditions for star formation in HVCs are fulfilled: there exist candidates with sufficient mass, molecular hydrogen and dust, and a variety of elements allowing them to cool radiatively over decades of temperatures. Many of them show emission in both $\text{H}\alpha$ and infrared, but however, no stellar counterpart could have ever been assigned

to these clouds (e.g. Hopp, Schulte-Ladbeck & Kerp 2003a,b; Stark, Baker & Kannappan 2015). However, a recent analysis of the *SECCO* survey assigns an identified stellar system (most likely a faint, nearly starless dwarf galaxy located in the Virgo Cluster) to the ultracompact HVC 274.68+74.70–123 (Bellazzini et al. 2015).

In this paper, we systematically explore the effects of various physical processes on the evolution of CHVCs moving through the CGM. We emphasize, that we do not model the large complexes of HVCs. The effects considered include self-gravity (Section 2.2), a variable metallicity affecting heating, and cooling of the plasma (Section 2.3), dissociation of molecules, and equilibrium ionization of atoms (Section 2.4), collisional heating and radiative cooling of plasma (Section 2.5), thermal conduction (Section 2.6), and diffusion of matter (Section 2.7). The individual processes are successively taken into account and the values for both cloud metallicity and mass are varied. We compare our results, on the one hand, to the findings of Vieser & Hensler (2007a), who conducted equal simulations in 2D on an equidistant grid under exploitation of cylindrical symmetry, in order to confirm their results with a more sophisticated 3D AMR numerical approach. Likewise, the simulations performed by Heitsch & Putman (2009), Armillotta et al. (2017), and Li et al. (2020) are discussed in the context of the conclusions drawn from simplified modelling of clouds as homogeneous and isothermal. On the other hand, we compare to the analytic calculations of Cowie & McKee (1977), whose approach solely focuses on classical and saturated mass-loss due to evaporation and neither accounts for self-gravity nor cooling. The latter is taken into account in McKee & Cowie (1977).

This paper is structured as follows: In Section 2, we describe the physical processes being considered in our simulations. In Section 3, we briefly explain the applied numerical code and Section 4 provides our approach for setting up the model clouds. The results of the simulations are presented in Section 5 and are finally discussed in Section 6. The paper is summarized in Section 7, where also the major conclusions are drawn.

2 PHYSICAL PROCESSES

In our simulation suite, we have covered physical processes that most likely take place in interstellar clouds. These processes are explained below. Radiative transport is neglected, because we assume the model clouds to be optically thin. Moreover, chemodynamics is taken into account on a simplified level, i.e. different species are not traced explicitly. Instead, the fractions of H_2 , H I , H II , He , electrons, and metals are calculated semi-analytically by using local fractions of dissociation and ionization, respectively.

2.1 Hydrodynamics

The dynamics of inviscid fluids in the absence of magnetic fields is determined by the Euler Equations

$$\frac{\partial \varrho}{\partial t} + \nabla \cdot (\varrho \mathbf{v}) = 0, \quad (1)$$

$$\frac{\partial \varrho \mathbf{v}}{\partial t} + \nabla \cdot (\varrho \mathbf{v} \mathbf{v}) + \nabla P = \varrho \mathbf{g}, \quad (2)$$

$$\frac{\partial \varrho e_{\text{tot}}}{\partial t} + \nabla \cdot [(\varrho e_{\text{tot}} + P) \mathbf{v}] = \sum_{j=1}^N S_j. \quad (3)$$

They yield the fluid's density ϱ , velocity $\mathbf{v}$, and total energy e_{tot} upon coupled integration. The S_j 's in equation (3) above are arbitrary

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source terms modifying the total energy of the fluid. By an equation of state,

$$P = (\gamma - 1)\varrho e_{\text{int}}, \quad (4)$$

the pressure P of the fluid is derived and the system of Euler Equations can be closed. Hence, the state of the fluid is determined for each and every time. The inner energy e_{int} is a function of the mean molecular weight μ thus equation (4) implicitly respects each species in molecular, atomic, and ionized form (see Section 2.4). For solar metallicity μ converges to ~ 2.4 in the CNM, to ~ 1.2 in the WIM, and to ~ 0.6 in the CGM. As we deal with only low-density environments in the simulations the assumption of an ideal gas is legitimated. Hence, $\gamma = 5/3$ in all simulations.

2.2 Self-gravity

The effect of self-gravity for an arbitrary matter distribution $\varrho(\mathbf{r})$ is described by solving the well-known Poisson's Equation

$$\nabla^2 \Phi(\mathbf{r}) = 4\pi G \varrho(\mathbf{r}), \quad (5)$$

for the gravitational potential $\Phi(\mathbf{r})$ (G denotes Newton's gravitational constant). Each mass element of $\varrho(\mathbf{r})$ contributes to $\Phi(\mathbf{r})$. If the total mass becomes a significant fraction of the (temperature-dependent) Jeans mass M_J for a particular $\varrho(\mathbf{r})$, self-gravity cannot be neglected as it relevantly affects the dynamics of all mass elements. However, it has been shown in Sander & Hensler (2019) that self-gravity has a measurable effect on the evolution of even clouds with masses well below their Jeans mass.

If self-gravity is stronger than the acceleration of gas in the opposite direction (e.g. due to evaporation or advection by a streaming CGM), the net acceleration at the cloud surface points in the same direction like the density gradient. Hence, the condition for Rayleigh–Taylor instability (RTI) is not given. So, it is expected that self-gravity significantly suppresses RTI at the cloud surface (see Murray et al. 1993) and the typical finger-like structures do not form.

According to Murray et al. (1993), the minimum mass of self-gravitating clouds above which they are stable against Kelvin–Helmholtz instability (KHI) when moving with a certain speed v_{rel} through an ambient medium with density ϱ_{CGM} reads

$$M_{\text{min}} \approx \frac{\sqrt{6\pi} v_{\text{rel}}^3}{\sqrt{\bar{\varrho}_{\text{cloud}}}} \left(\frac{\varrho_{\text{CGM}}}{G \bar{\varrho}_{\text{cloud}}} \right)^{3/2}, \quad (6)$$

with $\bar{\varrho}_{\text{cloud}}$ being the average mass density in the cloud.

To ensure that self-gravitating clouds, which are embedded in a medium with pressure P_{CGM} , are gravitationally stable initially, their total mass must be below their Bonnor–Ebert mass (Ebert 1955; Bonnor 1956)

$$M_{\text{BE}} = 1.18 \left(\frac{k_B \bar{T}_{\text{cloud}}}{G^3 \bar{\mu}_{\text{cloud}}} \right)^{1/2} P_{\text{CGM}}^{-1/2}. \quad (7)$$

Here, $\bar{T}_{\text{cloud}}$ and $\bar{\mu}_{\text{cloud}}$ are the average temperature and mean molecular weight in the cloud, respectively.

2.3 Metallicity

Observations indicate a common metallicity for CHVCs of 0.1 to 0.3 $Z_{\odot}$ (Wakker 2001). In order to study extreme cases, we consider two different initial cloud metallicities of $Z/Z_{\odot} = 0.1$ and 1.0. We note that the CGM has constant metallicity in all runs (see Section 4.3). The ratio of the abundances of hydrogen and helium is assumed solar and so are the ratios of metal abundances.

2.4 Dissociation and ionization of elements

The simulations contain cool clouds and hot plasma, so a temperature range from a few hundreds up to several 10^6 K is covered. For $T < 900$ K molecules exist, but they successively dissociate and finally become totally ionized with rising temperature. In our simulations, we semi-analytically calculate both dissociation and ionization for hydrogen in thermodynamic equilibrium. The fraction of H II is then related to He and metals. We note that we do not trace single species but rather have one fluid for which the composition is calculated based upon a given metallicity. The fractions of dissociation and ionization affect the magnitudes of cooling rates. See appendix A for further details.

2.5 Heating and cooling of plasma

The rates of heating $\Gamma = \Gamma_0 n$ and cooling $\Lambda = \Lambda_0 n^2$ of the plasma are calculated semi-analytically by applying cooling laws from literature: for $T < 900$ K the cooling curve from Falgarone & Puget (1985) is used, which considers molecular line cooling. For $900 < T/\text{K} < 10^4$, the cooling function from Dalgarno & McCray (1972) is applied, and for temperatures above 10^4 K the plasma cools according to the rates in Boehringer & Hensler (1989). Both Γ and Λ depend on the local values of metallicity, temperature, and density. These local variables are calculated in each time-step. The different cooling functions assume a solar ratio of metals.

The plasma is heated due to the photoelectric effect on dust particles (Weingartner, Draine & Barr 2006), ionization by UV radiation, by X-rays, and by cosmic rays (Wolfire et al. 2003), thermalization of turbulent motions, and condensation of molecular hydrogen on dust particles (both in Tielens 2010).

2.6 Thermal conduction

In the presence of large temperature gradients, the classical description of heat flux given by (Spitzer 1962)

$$\mathbf{q}_{\text{class}} = -\kappa_{\text{class}} \nabla T, \quad (8)$$

with the coefficient for thermal conduction $\kappa_{\text{class}} = 1.84 \times 10^{-5} T^{5/2} / \ln(\Omega)$, and the Coulomb logarithm $\ln(\Omega) = 29.7 + \ln[T/(10^6 \sqrt{n_e})]$, yields a too high amount of energy being conducted by electrons (e.g. Campbell 1984). Instead, a saturated heat flux $\mathbf{q}_{\text{sat}}$ (Cowie & McKee 1977; McKee & Cowie 1977) accounts for a finite reservoir of electrons and thus converges to a maximum value if all electrons are conducting heat and irrespective of the steepness of temperature gradient. The transition from classical to saturated thermal conduction depends on both the mean free path of electrons in the CGM,

$$\lambda_e = t_e \sqrt{\frac{3k_B T_{\text{CGM}}}{m_e}}, \quad (9)$$

where the electron–electron equipartition time (Spitzer 1962) reads

$$t_e = \frac{3\sqrt{m_e}(k_B T_{\text{CGM}})^{3/2}}{4\sqrt{\pi} n_e e^4 \ln(\Omega)}, \quad (10)$$

and the temperature scale height $L_T = T/|\nabla T|$ of the CGM, which provides the typical spatial scale of temperature variations. It has been experimentally verified by Gray & Kilkenny (1980), that

$$\mathbf{q} = \begin{cases} \mathbf{q}_{\text{class}}, & \lambda_e/L_T \lesssim 2 \times 10^{-3} \\ \mathbf{q}_{\text{sat}}, & \text{else} \end{cases}. \quad (11)$$

Cowie & McKee (1977) suggest the form

$$q_{\text{sat}} \equiv |q_{\text{sat}}| = 5\Phi q c_s^3, \quad (12)$$

with a factor Φ accounting for uncertainties specific to the approximation of flux-limited diffusion and the effect of magnetic fields. It holds $\Phi = 1.1$ if both ion and electron temperatures are equal. We use $\Phi = 1$ throughout all of the simulations.

To ensure continuity at the transition between q_{class} and q_{sat} , a proper representation for the heat flux must be chosen. Within this work the expression of Slavin & Cox (1992)

$$q_{\text{eff}} = q_{\text{sat}} \left(1 - \exp \left\{ -\frac{|q_{\text{class}}|}{|q_{\text{sat}}|} \right\} \right) \quad (13)$$

is used. It correctly converges to either of the heat fluxes in equation (11) for very small or very large temperature gradients, respectively.

The conductivity according to the effective heat flux (13) reads

$$\kappa_{\text{eff}} = \kappa_{\text{sat}} \left(1 - \exp \left\{ -\frac{\kappa_{\text{class}}}{\kappa_{\text{sat}}} \right\} \right), \quad (14)$$

with $\kappa_{\text{sat}} = 5q c_s^3 / |\nabla T|$ relating to the coefficient of thermal diffusion via $\alpha = \kappa_{\text{eff}} / (\rho c_V)$, where $c_V = \partial e_{\text{in}} / \partial T$ is the isochoric heat capacity.

We follow the procedure described in Vieser & Hensler (2007b) for calculating the temperature update according to saturated thermal conduction. From the new temperature distribution an adjusted heat flux

$$q_{\text{eff}} = -\kappa_{\text{eff}} \nabla T \quad (15)$$

finally accounts for the change in energy by thermal conduction. The code test for only thermal conduction given by equation (15) is shown in Appendix B.

If the wavelengths of temperature fluctuations in the ambient CGM are greater than the field length (Begelman & McKee 1990),

$$\lambda_F = \sqrt{\frac{\kappa_{\text{CGM}} T_{\text{CGM}}}{n^2 \Lambda_{\text{max}}}}, \quad (16)$$

these fluctuations are able to grow and lead to thermal instability (κ_{CGM} and T_{CGM} are the conductivity and temperature in the CGM, n is the local particle density and the maximum cooling rate reads $\Lambda_{\text{max}} = \max \{ \Lambda_0, \Gamma_0/n \}$). Thus, if $R_{\text{cloud}} \gg \lambda_F$, the thermal content of the cloud is dominated by external heating and radiative cooling whereas for $R_{\text{cloud}} \ll \lambda_F$ thermal conduction governs.

Cowie & McKee (1977) introduced a global saturation parameter

$$\sigma_0 = \frac{(T_{\text{CGM}} / 1.54 \times 10^7 [\text{K}])^2}{n_{\text{CGM}} \Phi R_{\text{cloud}} [\text{pc}]}, \quad (17)$$

which can be used equivalently to equation (11) to discriminate between classical ($\sigma_0 < 1$) and saturated ($\sigma_0 > 1$) thermal conduction. Here, R_{cloud} is given in terms of parsec, and Φ is defined in equation (12). The mass-loss rate by evaporation is considerably reduced with increasing σ_0 (Fig. 1). For $\sigma_0 \gg 1$ also viscous heating becomes relevant (Balbus & McKee 1982). Theoretically, gas should evaporate from clouds for $\lambda_e \lesssim R_{\text{cloud}} \ll \lambda_F$ and condense on to them if $R_{\text{cloud}} \gtrsim \lambda_F \gg \lambda_e$. Vieser & Hensler (2007b) have shown that the latter condition for condensation is already fulfilled for $R_{\text{cloud}} \gtrsim 0.24 \lambda_F$.

In order to correctly resolve λ_F in simulations, Koyama & Inutsuka (2004) and Gressel (2009) suggest a resolution of λ_F by at least three grid cells. Otherwise, results become resolution-dependent. In our simulations, λ_F is resolved by 158 cells in the massive model clouds and 130 cells in the low-mass clouds.

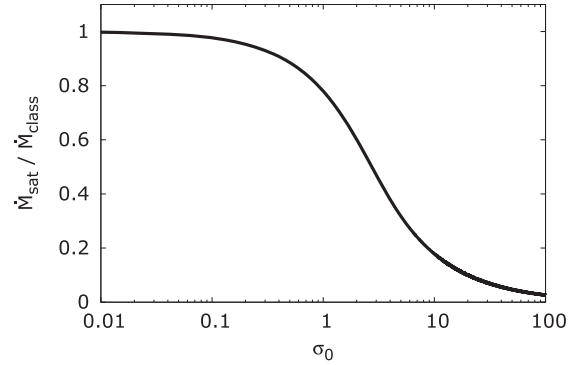


Figure 1. Ratio of saturated to classical mass-transportation rate (see Dalton & Balbus 1993).

2.7 Mass diffusion

The simulations respect the transport of matter by diffusion using the formulation of the Chapman–Enskog expansion. This is a valid description at low gas densities (Bird et al. 1960). The mass diffusion flow reads

$$\mathbf{j} = -D \nabla \varrho, \quad (18)$$

with D being the coefficient of mass diffusion. Assuming one species of gas,

$$D = 0.0018583 \frac{\sqrt{2T^3 / \mu_H}}{P \sigma_H^2 \Omega_{\text{int}}} \quad (19)$$

according to the Chapman–Enskog expansion. Here, σ_H and μ_H denote the cross-section and molecular weight, respectively, of hydrogen. The collision integral is given by

$$\Omega_{\text{int}} = \frac{1.06036}{T_*^{0.15610}} + \frac{0.19300}{\exp\{0.47635T_*\}} + \frac{1.03587}{\exp\{1.52996T_*\}} + \frac{1.76474}{\exp\{3.89411T_*\}}, \quad (20)$$

where $T_* = T / (E_L / k)$ and E_L is the depth of the Lennard–Jones potential of hydrogen.

2.8 Magnetic fields

We explicitly neglect magnetic fields, because so far only four rather imprecise measurements of magnetic field strengths in HVCs without any clue on orientation are available. Kazes, Troland & Crutcher (1991) searched for the Zeeman effect in four H I HVCs, where they measured a magnetic field $|\mathbf{B}| = (-11.4 \pm 2.4) \mu\text{G}$ in only HVC 132+23–212. More recently, McClure-Griffiths et al. (2010) reported on the strength for the line-of-sight component of the coherent magnetic field of $\gtrsim 6 \mu\text{G}$ in the HVC associated with the Leading Arm of the Magellanic System. Hill et al. (2013) and Betti et al. (2019) detected a magnetic field with strength $\gtrsim 8$ and $\gtrsim 5 \mu\text{G}$, respectively, associated with the Smith Cloud. Hence, magnetic fields of HVCs are weakly constrained by observations in orientation and strength. Another plausible cause for Faraday rotation can be the compressed magnetic field in the CGM in front of infalling HVCs (Jelínek & Hensler 2011). In this case, the source of Faraday rotation does not belong to the cloud. So, neglecting magnetic fields in our models does not contradict the observations.

It further has been shown theoretically by Sander & Hensler (2019) that the effective heat flux, i.e. integrated over the cloud surface,

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is reduced to only 68 per cent of its undisturbed value by a strong magnetic dipole field. Thus, neglecting magnetic fields in our models is theoretically justified with respect to thermal conduction.

3 NUMERICAL METHODS

3.1 The code

We carry out our simulations by using the publicly available code FLASH¹ (Fryxell et al. 2000; Dubey et al. 2009). It is a Eulerian hydrodynamics code, by which the (inviscid) Euler equations (1)–(3) are integrated on a Cartesian grid by means of the piecewise parabolic method (Colella & Woodward 1984; Woodward & Colella 1984).

Our simulations are performed in full three dimensions with outflow boundary conditions for the hydrodynamics at all sides of the computational domain, i.e. energy and mass can leave the domain, but are not allowed to return. We aim at investigating in detail the evolution of cloud structures thus the adaptive mesh refinement (AMR) property of FLASH is extensively utilized. The local numerical resolution is increased (decreased) based on the local density gradient by subsequently doubling (halvening) the number of blocks per dimension. One block consists of 8^3 grid cells. The side length of the computational cube in all of our models is 340 pc and the finest spatial resolution $\Delta x = 1.3$ pc with six levels of AMR.

We use the Multigrid solver provided by FLASH for solving Poisson’s Equation (for an introduction to multilevel methods see Zumbusch 2003) together with isolated boundary conditions. By this, the potential field in the cells at the boundary of the computational domain can be calculated by taking appropriate density values from so-called ghost cells surrounding the entire domain. The domain of integration is extended by four layers of ghost cells at each face, which get their physical values by copying them from the outermost cells of the domain. The same holds for each and every block the computational domain is subsequently divided into.

3.2 Time-steps

In multiphysics simulations, one is restricted by the time-step of the fastest process when integrating equations (1)–(3). Conservatively, integration is performed on the minimum value of all such time-steps in the entire computational domain. So, the wallclock times of simulations can increase enormously.

First, plasma cooling evolves on a time-scale

$$\Delta t_{\text{cool}} = \frac{\varrho e_{\text{int}}}{|\Gamma - \Lambda| + \varepsilon}, \quad (21)$$

where $\varepsilon > 0$ is a small number that prevents the code from dividing by zero at thermal equilibrium in a grid cell. In regions dominated by cooling Δt_{cool} can be very small.

Differently, thermal conduction is determined by the conduction time

$$\Delta t_{\text{cond}} = \frac{\min\{\Delta x^2, \Delta y^2, \Delta z^2\}}{\max\{\alpha\}}, \quad (22)$$

where $\alpha = \kappa_{\text{eff}}/(\varrho c_V)$ is the coefficient of thermal diffusion, $c_V = \partial e_{\text{int}}/\partial T$ is the isochoric heat capacity, and the Δ ’s are the spatial resolutions in x -, y - and z -direction. Near cloud boundaries, the numerical grid is fine-grained and α becomes very high thus Δt_{cond} is remarkably scaled down.

It thus appears to be very likely in our model setups that both Δt_{cool} and Δt_{cond} are much smaller than

$$\Delta t_{\text{CFL}} = C_{\text{CFL}} \frac{\min\{\Delta x, \Delta y, \Delta z\}}{\max\{|\mathbf{v}| + c_s\}}, \quad (23)$$

which is the numerically allowed maximum time-step according to the Courant–Friedrichs–Lewy condition (Courant, Friedrichs & Lewy 1928) and ensures a numerically stable solution of equations (1)–(3). Depending on the fluid speed $|\mathbf{v}|$, we set the safety factor $C_{\text{CFL}} = 0.1, 0.2$, or 0.4 in our simulations.

To circumvent this numerical restriction, we apply implicit integration methods. To update the energy based on heating and cooling (Section 2.5), we apply an implicit bisection method (cf. Engeln-Müllges, Niederdröck & Wodicka 2011). Thermal conduction (Section 2.6) is numerically addressed by the semi-implicit Crank–Nicolson method (Crank, Nicolson & Hartree 1947), which we have implemented in the FLASH code.

Hence, we are neither restricted by Δt_{cool} nor Δt_{cond} and the total energy density ϱe_{tot} can be updated according to

$$\frac{\Delta \varrho e_{\text{tot}}}{\Delta t} + \nabla(\varrho e_{\text{tot}} \mathbf{v}) = -\nabla(P \mathbf{v}) + \varrho \mathbf{v} \mathbf{g} + \Gamma - \Lambda - \nabla \mathbf{q}_{\text{eff}}, \quad (24)$$

where Δt is given by equation (23).

3.3 Slope limiter

Another numerical complicacy is due to very steep temperature gradients across cloud surfaces. The heat flux (15) can thus be very high in the vicinity of regions of nearly zero heat flux. By this, artificial numerical oscillations can be produced near steep temperature gradients when integrating equations (1)–(3). In order to subsequently limit these oscillations, we use a minmod (or: minbee) slope limiter to smooth temperature gradients. The minmod slope limiter preserves monotonicity of the employed upwind scheme thus ensuring diminishing of total variation (Toro 2009) of the solution even near discontinuities. By definition, the slope limiter becomes significant only around large temperature gradients. In regions of equal temperature, it is simply equal to zero.

4 MODEL SETUPS

4.1 Initial cloud profiles

The gas in the clouds is stratified hydrostatically, i.e. the condition for hydrostatic equilibrium

$$-\frac{\partial P}{\partial r} = \varrho \frac{GM(r)}{r^2} \quad (25)$$

must hold. In addition, the cloud is in thermal equilibrium at each and every point, i.e. the energy input by external heating is compensated by radiative cooling (cf. Section 2.5). We realize this condition by adapting the density to a certain temperature value, such that $\Lambda n^2 = \Gamma n$. By that, matter density becomes a function of local temperature, $\varrho = \varrho(T)$. By substituting P from equation (4) into equation (25), we derive the governing equation for the radial temperature profile of the cloud,

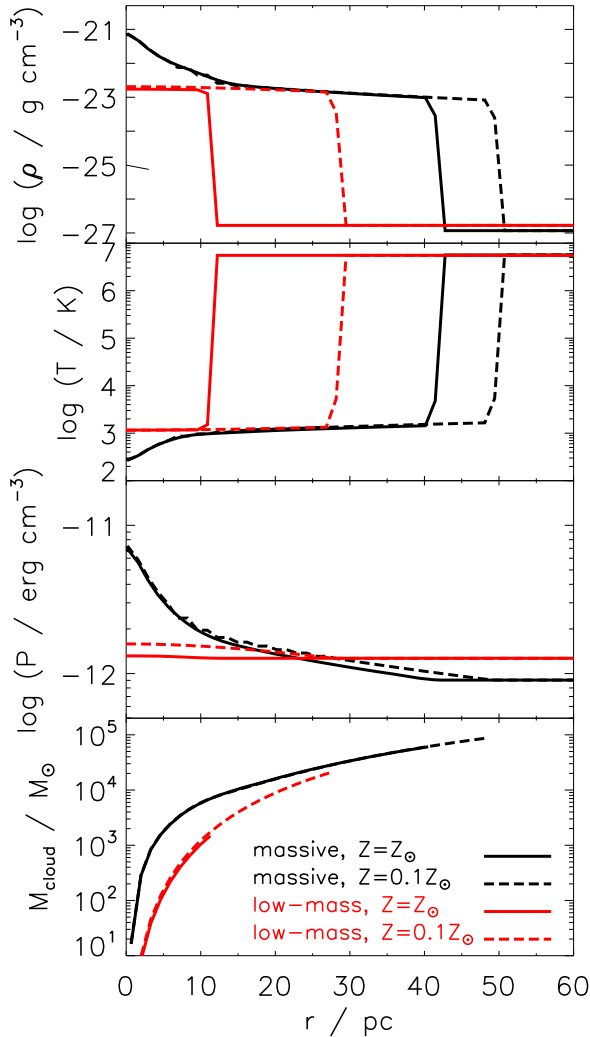
$$\frac{\partial T}{\partial r} = -\varrho(T(r)) \frac{GM(r)}{r^2} \left\{ \frac{\partial}{\partial T} \left[\frac{\varrho(T(r)) k_B T(r)}{\mu(\varrho, T(r))} \right] \right\}^{-1}, \quad (26)$$

where μ is the mean molecular weight. In order to solve equation (26), the central cloud temperature T_{centre} and the temperature and particle density of the CGM, T_{CGM} , and n_{CGM} , must be provided as boundary conditions (Table 1 shows our chosen values). Then,

¹ See <http://flash.uchicago.edu/site/flashcode/>

Table 1. Boundary conditions for obtaining the initial cloud models by integrating equation (26).

Parameter	Massive clouds (M)	Low-mass clouds (L)
$T_{\text{centre}}/\text{K}$	380	1150
$T_{\text{CGM}}/10^6 \text{ K}$	5.6	5.5
$n_{\text{CGM}}/10^{-3} \text{ cm}^{-3}$	0.7	1.0

**Figure 2.** From top to bottom panels: Initial radial profiles of density, temperature, pressure, and mass for all model clouds.

equation (26) is integrated outwards with the density being assigned according to thermal equilibrium. The radial integration is performed until the pressure of the cloud matches that of the CGM. By that, the cloud radius R_{cloud} is a natural consequence of pressure equilibrium. The radial profiles of the initial configurations for all model clouds are given in Fig. 2. At constant CGM temperature, the particle density of CGM is the only parameter affecting the ambient pressure. It thus controls the cloud radius when solving equation (26). By inspecting the third plot in Fig. 2, one observes a higher ambient pressure in the model setups for the low-mass clouds due to a higher ambient particle density. Thus, pressure balance is reached at a smaller radius.

Table 2. Initial parameters characterizing the simulation models with respect to constraints (27) and (28). The values are specific to clouds with $Z/Z_{\odot} = 0.1$ (in brackets for $Z/Z_{\odot} = 1.0$).

Parameter	Massive clouds	Low-mass clouds
σ_0	3.8 (4.6)	4.6 (11.2)
$M_{\text{cloud}}/M_{\text{min}}$	> 3	< 0.8
$M_{\text{cloud}}/M_{\text{BE}}$	0.25 (0.23)	0.08 ($\lesssim 0.01$)
$R_{\text{cloud}}/\lambda_{\text{F}}^a$	≤ 0.23 (0.20)	0.16 (0.07)
$\lambda_{\text{e}}/L_{\text{T}}$	1.05 (1.21)	0.59 (0.68)

^aWe point out that λ_{F} is not defined for clouds in which thermal conduction is not considered.

4.2 Cloud masses

The solution of equation (26) is mapped to the numerical grid of FLASH code, which consists of cells of equal size (adaptive mesh refinement is applied after the initial conditions are mapped to the grid). Due to the constraint of thermal equilibrium, the density is simultaneously provided, such that the mass of the cloud is given by adding together the masses of all grid cells inside the cloud radius (lower plot in Fig. 2).

In order to obtain a cloud that is dynamically stable initially and conducts thermal energy in a saturated manner, it must be constrained by the conditions

$$\sigma_0 > 1, \quad R_{\text{cloud}} \lesssim \lambda_{\text{F}}, \quad M_{\text{cloud}} \in [M_{\text{min}}, M_{\text{BE}}], \quad (27)$$

whereas the CGM must satisfy

$$\lambda_{\text{e}} \gtrsim L_{\text{T}}. \quad (28)$$

Table 2 contains the realizations of our setups in the parameter subspaces (27) and (28). The values for M_{cloud} are listed in Table 3.

4.3 Ambient CGM

Our current understanding of the CGM only allows hints to the range of density and temperature inhomogeneities, which are not fully understood (as e.g. heating of the intracluster medium (ICM) by the motion of galaxies, see Just et al. 1990) and must act on large scales. An example is the suppression of the growth of filaments in the ICM above a critical length-scale l_{crit} , derived by Nipoti & Binney (2004). Inserting our cloud parameters into their equation (6) yields $l_{\text{crit}} = 26 \text{ kpc}$ (massive clouds) and $l_{\text{crit}} = 15 \text{ kpc}$ (low-mass clouds). Our fastest model clouds are M1c_vel and L1b_vel. So, within 60 Myr they travel distances of 20 (M1c_vel) and 15 kpc (L1b_vel), which do not exceed l_{crit} in either case. Thus, we examine constant CGM conditions (cf. Table 1). Assuming a galactic stratification of ionized gas (see equation 6 in Gaensler et al. 2008) with a scale height of $2.6 \pm 0.4 \text{ kpc}$ (derived for the distribution of warm gas in the CGM based on observations of Si IV and O VI by Qu & Bregman 2019) our model clouds can be roughly placed at a galactic height of $\sim 100 \text{ kpc}$.

Although the CGM metallicity is uncertain around a range of 0.3–0.4 $Z_{\odot}$ we set the metallicity of the CGM to be solar in order to better distinguish the amount of CGM accreted by the cloud and the location of mixed CHVC gas. For densities and temperatures, we adapt for the CGM the metallicity does not play a crucial role, because cooling is mainly due to Bremsstrahlung. However, for mixing phases at $\lesssim 10^5 \text{ K}$ metal lines of various ionization stages determine the cooling strength such that metallicity matters there.

5336 *B. Sander and G. Hensler***Table 3.** Simulated model clouds with initial values of radius (Column 5), mass (Column 6), metallicity (Column 7), speed (Column 8), and Mach number (Column 9). Columns 3 and 4 show the physical processes that are either considered (+) or neglected (−).

Model group	Model	Self-gravity	Thermal conduction	$R_{\text{cloud}}/\text{pc}$	$M_{\text{cloud}}/10^4 M_{\odot}$	$Z/Z_{\odot}$	$v_{\text{cloud}}/\text{km s}^{-1}$	Ma
Massive clouds	M0 (reference model)	+	+	49	9.0	0.1	167	0.46
	M1_tc	+	−	49	9.0	0.1	167	0.46
	M1a_vel	+	+	49	9.0	0.1	83	0.23
	M1b_vel	+	+	49	9.0	0.1	250	0.69
	M1c_vel	+	+	49	9.0	0.1	333	0.93
	M1_dec	+	+	49	9.0	0.1	167 ^a	0.46 ^a
	M1_me	+	+	41	6.3	1.0	167	0.46
	M3_hom_tc_sg	−	−	41	6.4	1.0	167	0.46
Low-mass clouds	L0 (reference model)	+	+	28	2.2	0.1	165	0.46
	L1_tc	+	−	28	2.2	0.1	165	0.46
	L1a_vel	+	+	28	2.2	0.1	83	0.23
	L1b_vel	+	+	28	2.2	0.1	248	0.69
	L1_dec	+	+	28	2.2	0.1	165 ^a	0.46 ^a
	L1_me	+	+	11	0.15	1.0	165	0.46

^aCloud speed is a function of time according to the drag force of ambient CGM.

4.4 Subsonic flow

Our model CHVCs are subject to a subsonic flow of hot gas representing the motion of the clouds through the CGM of the Milky Way. By that, we cover the velocity range between 83 and 333 km s^{−1} (Table 3).

The flow is initially set up as a potential flow around the cloud, which is valid for subsonic flows (Landau & Lifshitz 1966). Since the cloud is initially a sphere, the velocity field of the flow at a distance $|\mathbf{d}|$ from the cloud is given by

$$\mathbf{v}(\mathbf{d}) = \begin{cases} \mathbf{v}_{\text{CGM}} + \frac{R_{\text{cloud}}^3}{2|\mathbf{d}|^3} [\mathbf{v}_{\text{CGM}} - 3\mathbf{n}(\mathbf{v}_{\text{CGM}} \cdot \mathbf{n})], & |\mathbf{d}| > R_{\text{cloud}} \\ 0, & |\mathbf{d}| \leq R_{\text{cloud}} \end{cases}, \quad (29)$$

with $\mathbf{n}$ being the unit vector pointing in the direction of $\mathbf{d}$. Keep in mind that a potential flow provides a laminar initial condition. In most of our simulation models, the velocity of the flow is constant. However, in two of the models the velocity of the CHVC relative to CGM decreases due to the drag force of the CGM. The decreased cloud velocity at time t_{i+1} is calculated by subtracting

$$\delta v = \frac{C_D}{2} \frac{A_{\text{cloud}}}{M_{\text{cloud}}} \rho_{\text{CGM}} v^2(t_i) \Delta t \quad (30)$$

from the cloud velocity at time t_i , where A_{cloud} and M_{cloud} are the effective area and the mass of the cloud, respectively, and Δt is the current integration time-step. The drag coefficient C_D is calculated according to Henderson (1976) and reads

$$C_D = C_1 + C_2 + C_3, \quad (31)$$

where

$$C_1 = 24 \left\{ Re + v_{\text{mol}} \left[4.33 + \frac{3.65 - 1.53 T_{\text{cloud}}/T_{\text{CGM}}}{1 + 0.353 Re + 0.48 \sqrt{Re}} \right] \right\}^{-1} \times \exp \left\{ -0.247 \frac{Re}{v_{\text{mol}}} \right\}, \quad (32)$$

$$C_2 = \exp \left\{ -\frac{Ma}{2\sqrt{Re}} \right\} \frac{4.5 + 0.38 (0.03 Re + 0.48 \sqrt{Re})}{1 + 0.03 Re + 0.48 \sqrt{Re}} + 0.1 Ma^2 + 0.2 Ma^8, \quad (33)$$

$$C_3 = 0.6 v_{\text{mol}} \left(1 - \exp \left\{ -\frac{Ma}{Re} \right\} \right). \quad (34)$$

The Reynolds number is $Re = \rho_{\text{CGM}} v(t_i) R_{\text{cloud}} / \eta$, with dynamic viscosity $\eta = 0.1$, and the molecular speed ratio reads $v_{\text{mol}} = Ma \sqrt{0.5\gamma}$, with γ being the ratio of specific heats and Ma is the Mach number of the flow. The drag coefficient is calculated for isothermal clouds in Henderson (1976) so we set T_{cloud} in equation (32) as the boundary temperature in our model clouds, since this is the temperature at interface to CGM. Equations (31)–(34) are valid expressions for subsonic flows.

4.5 Model clouds

We have simulated and analysed models differing in the considered physical processes or in model parameters, which are key to their evolution. In Table 3, the models are listed. There are 14 models in total allowing us to compare them pairwise and give clues on the effect of particular physics. All clouds are initially in both internal hydrostatic and thermal equilibrium as well as in pressure equilibrium with the ambient CGM. All models take into account heating and radiative cooling of plasma. The cloud speed is constant in all but the models with decelerated clouds. The clouds are divided into the two groups of massive (M) and low-mass (L) clouds.

(i) The massive clouds have a steep radial temperature profile which allows for a molecular phase in the centre and an extended atomic phase at larger radii. Only a small shell at the boundary consists of warm, slightly ionized gas. The density profile at all radii forms a distinct core-halo structure. All massive clouds are dynamically stable regarding $M_{\text{cloud}} > M_{\text{min}}$, but have masses below their respective Bonnor–Ebert mass ($< 0.3 M_{\text{BE}}$, see Fig. 3), which is a non-linear function of radius as the massive clouds are not isothermal. To study the effect of physical simplicity, we set up model M3_hom_tc_sg, which is isothermal with $T_{\text{cloud}} = 10^3$ K and homogeneous with $\rho_{\text{cloud}} = 1.5 \times 10^{-23}$ g cm^{−3}. The sound speed in the CGM has a value of 359 km s^{−1}.

(ii) The low-mass clouds have a higher central temperature yielding a lower central density with respect to thermal equilibrium. Consequently, the central pressure is low and together with a higher value of ambient particle density, a smaller radius is obtained. These clouds lack the molecular phase in their centres and they

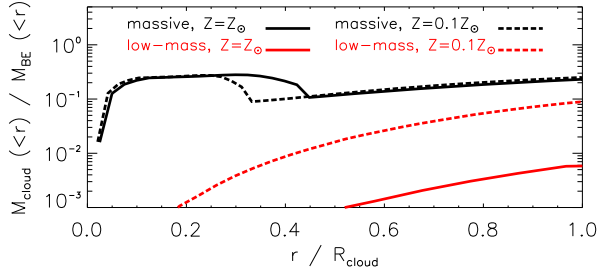


Figure 3. Radial cloud-to-Bonnor-Ebert mass ratios for all model clouds. The x-axis is normalized to the respective radius of each cloud.

are only weakly gravitationally bound. According to the mass ratio $M_{\text{cloud}}/M_{\text{min}}$ (Table 2), these clouds are expected to be disrupted. We can thus study if the lifetime of clouds, which are disrupted anyway, is increased by thermal conduction. Likewise, their masses are below the Bonnor-Ebert mass ($<0.1 M_{\text{BE}}$, see Fig. 3). As the low-mass clouds are almost isothermal, there is no such distinct feature in the radial profile of M_{BE} . The sound speed in the CGM has a value of 355 km s^{-1} .

For both massive and low-mass clouds, we assume two different metallicities to quantify the effect of metal-dependent radiative cooling: $Z/Z_{\odot} = 1.0$ is chosen in order to be comparable to the clouds investigated by Vieser & Hensler (2007b) and $Z/Z_{\odot} = 0.1$ is more likely for HVCs falling on to the Milky Way (Wakker 2001). Furthermore, the processes of thermal conduction and self-gravity are either considered in the models or not (cf. Table 3) in order to study their respective impact on cloud evolution. Different cloud velocities scan the range of actually observed LSR velocities.

4.6 Reference clouds

In order to study the effects of different parameters and processes, we define one respective reference cloud for each of the massive and low-mass models. Within the reference models, the physical processes of self-gravity and thermal conduction are considered, their simulated velocities are typical for CHVCs, and they have a metallicity of $0.1 Z_{\odot}$. These cloud models most resemble actually observed CHVCs (Williams et al. 2012).

5 MODEL COMPARISONS

In the following, we compare the simulated model clouds according to their differences as shown in Table 3. Hence, the respective clouds are compared in terms of their masses, the effect of thermal conduction, their speed together with the effect of drag by the surrounding medium, and the strength of radiative cooling determined by the metallicity in the clouds. Finally, the effect of physical simplicity is described.

For the purpose of tracking physical quantities (e.g. velocity fields) and locating physical processes (e.g. KHI, gravitational collapse, mixing of gas) in the clouds, we study slices through their density distributions. Integrated quantities, like column density, are discussed separately.

5.1 Cloud mass

Both massive and low-mass clouds effectively lose mass during their passage through the CGM with mean rates $\dot{M}$ of -3.6×10^{-4} and

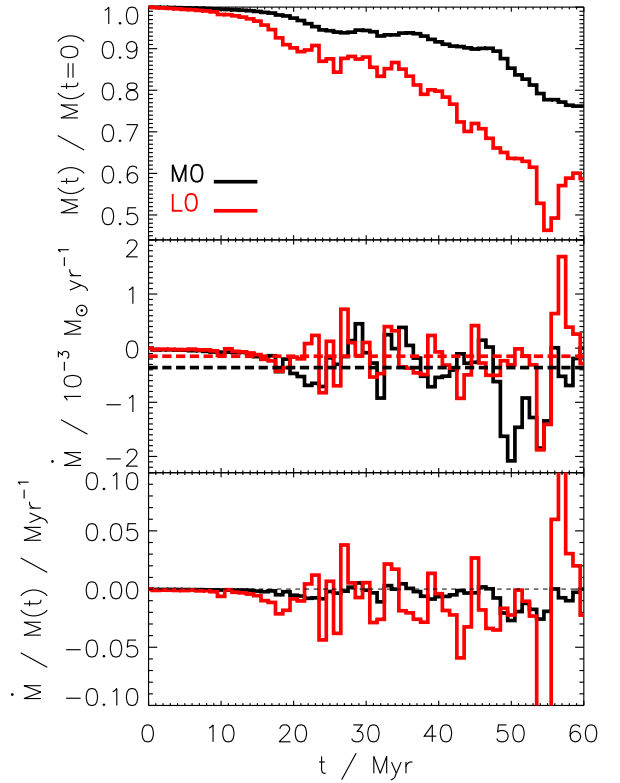


Figure 4. Evolution of mass (upper plot) for reference clouds. The corresponding mass-transportation rates (middle plot) show both loss ($\dot{M} < 0$) and gain ($\dot{M} > 0$) of mass with the average values (dashed lines) indicating a mass-loss. The relative mass-loss (lower plot) is higher in low-mass clouds.

$-1.6 \times 10^{-4} M_{\odot} \text{ yr}^{-1}$ (models M0 and L0, see Section 6.2). Low-mass clouds can lose up to 5 per cent of their current mass per Myr after 20 Myr and more than 10 per cent after 50 Myr (Fig. 4, lower plot) while for massive clouds their mass-loss is always less than 3 per cent per Myr. After 60 Myr of evolution, there is only about 60 per cent of initial mass left in the low-mass clouds while the massive clouds still contain 76 per cent.

In the middle plot of Fig. 4, it is striking that $\dot{M} > 0$ at some times. By inspecting the respective plots for models M0 and L0 in Figs 5 and 6, one observes the velocity field behind the cloud to be turbulent and hence directed towards the cloud. So, already stripped cloud material is advected back to the cloud by the turbulent rear wake. This effect is rather low, such that on average both models lose mass. The stripping is continuous, i.e. without any sudden peaks in mass-loss (Fig. 4), filamentary, and spatially extended (Figs 5 and 6).

5.2 Thermal conduction

By comparing models M0 with M1.tc (Fig. 5) and L0 with L1.tc (Fig. 6) one easily verifies the different evolution of clouds that differ in thermal conduction only. Material is stripped from clouds without thermal conduction in form of dense clumps that are confined and isolated from each other. From Fig. 7, we learn that these stripped clumps cool efficiently, because particle density is sufficiently high. The sharp cut around $2 \times 10^4 \text{ K}$ correlates with a maximum of the cooling strength, because hydrogen gets fully ionized. Thus, the budget of thermal energy in stripped fragments is regulated by

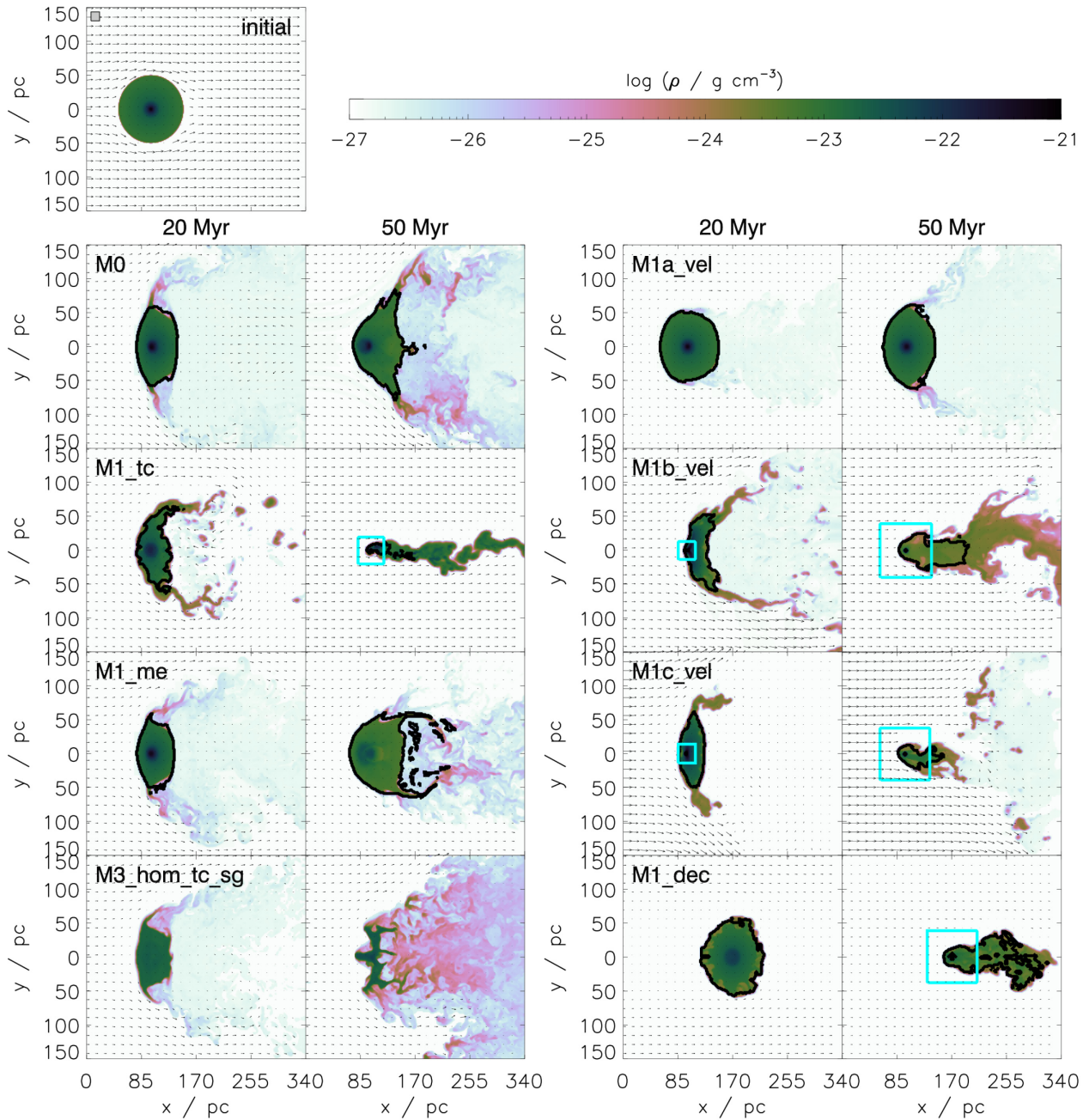
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Figure 5. Evolution of density in the massive model clouds. The plots show the section plane (x, y) through cloud centre. The uppermost plot shows the initial stage. The evolution is shown in snapshots after 20 (left-hand column) and after 50 Myr (right-hand column). Model M3_hom_tc_sg is shown at 55 Myr, because the finger-like structures, which are due to RTI, are substantially developed. Model M1_dec is shown at 100 Myr when the head-tail structure is more prominently developed. The black solid lines enclose cloud material (gravitationally bound and physically connected) and the light blue solid lines border gravitationally unstable regions. Black arrows illustrate the motion of gas. The arrow lengths are scaled to $v_{\text{ini}} = 333 \text{ km s}^{-1}$ (model M1c_vel). The grey square in the upper left-hand corner is spanned by 10×10 cells of finest numerical resolution ($\Delta x = 1.3 \text{ pc}$).

radiative cooling. In contrast, the material stripped from clouds with thermal conduction is too dilute so it is heated up quickly above 10^5 K , which conforms qualitatively with the findings of Armillotta et al. (2017). Irrespective of thermal conduction almost all of the stripped gas is above 10^3 K , i.e. it is fully dissociated.

Massive clouds without thermal conduction undergo a Jeans instability, while massive clouds with thermal conduction do not. In contrast, low-mass clouds do not experience any Jeans instability

irrespective of thermal conduction (models M0, M1_tc, L0, L1_tc in Figs 5 and 6). Low-mass clouds without thermal conduction are already disrupted after 30 Myr (Fig. 8), while those with thermal conduction still retain more than 50 per cent of their initial mass (Fig. 4). The mean mass-loss rates of massive clouds without thermal conduction are ~ 33 per cent higher than in their counterparts with thermal conduction. In low-mass clouds, this excess is even ~ 50 per cent (Section 6.2), which agrees qualitatively with the results of

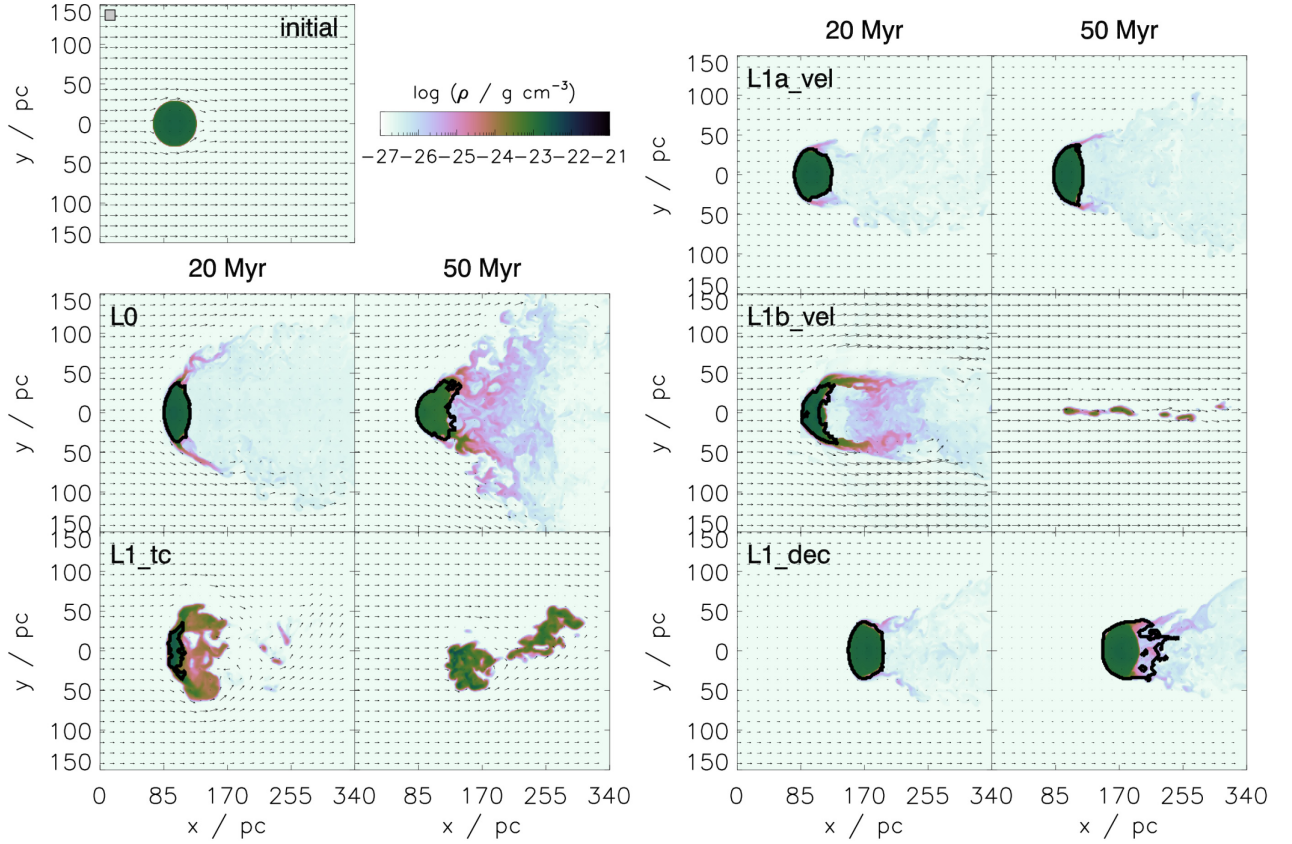


Figure 6. The same as Fig. 5, but for the low-mass model clouds. The arrow lengths are scaled to $v_{\text{ini}} = 248 \text{ km s}^{-1}$ (model L1b_vel).

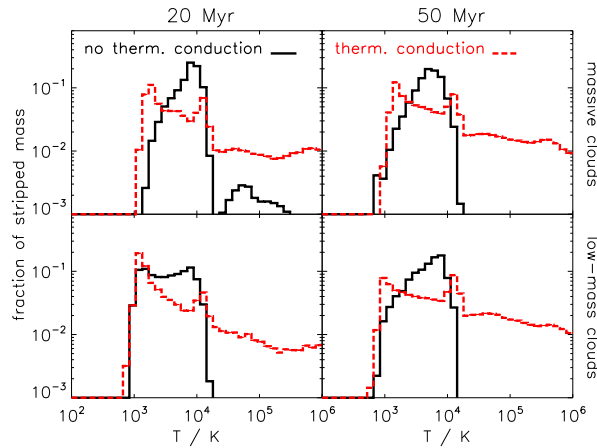


Figure 7. Distributions of temperature in the stripped material of massive clouds (upper panel) and low-mass clouds (lower panel) with thermal conduction (red dashed lines) and without thermal conduction (black solid lines) after 20 Myr (left-hand column) and after 50 Myr (right-hand column). The standard deviation in every bin is too low to be visible in the plots.

Vieser & Hensler (2007a), and implies a longer lifetime for clouds with thermal conduction. Thermal conduction thus acts in two ways: It prevents massive clouds from becoming gravitationally unstable and it prevents low-mass clouds from being tattered.

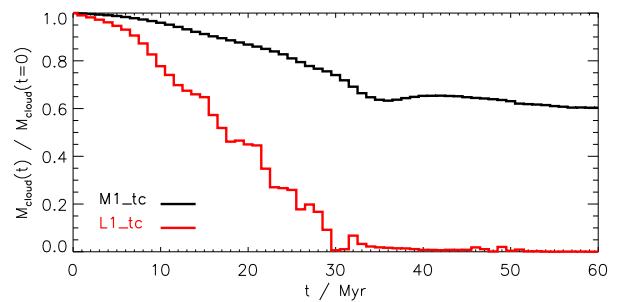


Figure 8. The same as Fig. 4, but for respective clouds without thermal conduction.

By comparing models M0 and M1_tc in Fig. 5, it is striking that the central region in the massive cloud without thermal conduction becomes gravitationally unstable. Because the density and hence the mass increase significantly in the central region, the mass stripped by ram-pressure can be compensated (Fig. 8). Consequently, the difference in mass-loss between models M0 and M1_tc starts to level out after ~ 35 Myr.

5.3 Cloud velocity

In addition to the amount of initial mass and thermal conduction also the relative speed between CHVC and CGM strongly influences the evolution of clouds. By inspecting models M0, M1a_vel, M1b_vel,

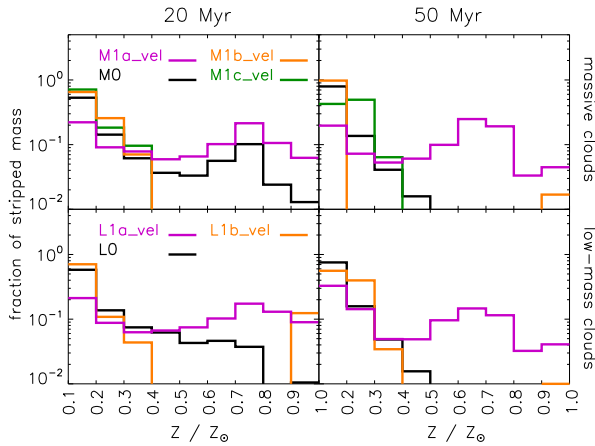
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Figure 9. Distributions of metallicity in the stripped material of massive clouds (upper panel) and low-mass clouds (lower panel) after 20 Myr (left-hand column) and after 50 Myr (right-hand column).

M1c_vel (Fig. 5) and L0, L1a_vel, L1b_vel (Fig. 6) a completely different morphology is seen after 50 Myr. In particular, low-mass clouds moving at 165 km s^{-1} (Mach 0.46) or below survive the first 50 Myr of evolution with 65 per cent of their initial mass, but when moving at 248 km s^{-1} (Mach 0.69) or above they are totally disrupted before 40 Myr (Section 6.2). However, massive clouds above 250 km s^{-1} (Mach 0.69) start to locally undergo a gravitational collapse in their central regions.

These very fast model clouds develop a so-called head-tail morphology, i.e. an asymmetric density distribution with a sharp gradient at the cloud head and a shallow gradient in the trailing part. The stripped trail of massive clouds is the more spatially confined the faster the clouds move (Fig. 5). Such head-tail structures are actually observed at many CHVCs (Putman, Saul & Mets 2011; Winkel et al. 2011). Two prime examples are HVC 125+41–207 (Brüns, Kerp & Pagels 2001) and HVC 297+09–253 (Ben Bekhti et al. 2006), two CHVCs that move at $v_{\text{LSR}} = -207$ and -253 km s^{-1} , respectively. Oppositely, the low-mass model clouds do not develop such a prominent tail.

The model clouds with a sub-solar metallicity of $0.1 Z_{\odot}$ can be used to study the impact of velocity on the spatial distribution of metals. From Fig. 9, we learn that in fast clouds (L1b_vel, M1b_vel, M1c_vel) the stripped material is mostly at initial cloud metallicity $0.1 Z_{\odot}$. Contrary, in slow clouds (L0, L1a_vel, M0, M1a_vel) stripped material has time to get enriched with CGM metals. Hence, a higher fraction of stripped material is at higher metallicity.

As can be seen in Fig. 9, the overall metallicity in the stripped gas of models M0 and L0 is lowered at 50 Myr compared to 20 Myr. The reduced metallicity is due to higher mass-loss rates around 50 Myr than at 20 Myr (middle plot in Fig. 4). Hence, more cloud material with $Z = 0.1 Z_{\odot}$ streams into the trailing parts of the clouds. For M0, a positive linear correlation of 0.36 between mass-loss rate and mean metallicity of the stripped gas is found at significance level of 0.05. For L0, there is no such correlation between mass-loss rate and metallicity in the tail.

The mixing with CGM appears dominantly at rear side of cloud. In Fig. 10, the rear side of clouds is at higher metallicity than the front side. Also in Fig. 10, we observe a lower metallicity behind high-speed clouds, because the region behind the cloud is refilled quickly with new cloud material. By inspecting Figs 5 and 6 another morphologic feature becomes apparent in model clouds M1b_vel and

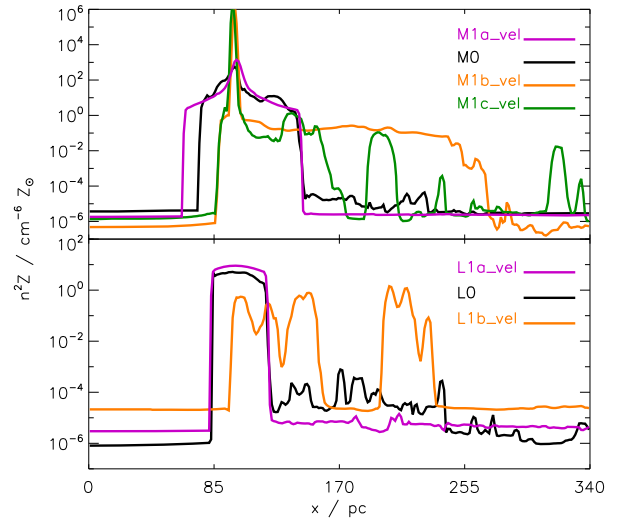


Figure 10. Metallicity profiles weighted with the square of particle density along the direction of motion for the massive clouds (upper plot) and low-mass clouds (lower plot) after 50 Myr.

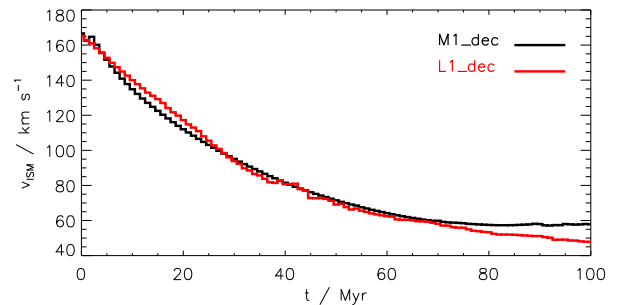


Figure 11. Deceleration of model clouds that experience drag.

L1b_vel, which is called a *bow-shock shape*, i.e. the cloud is shaped like a boomerang with a density gradient along the wings. Model M1c_vel develops this morphology later at ~ 35 Myr, but none of the slower clouds forms such a shape. Two observed examples with bow-shock shape are CHVC 172–60 and CHVC 157+03, which move at $v_{\text{LSR}} = -235 \text{ km s}^{-1}$ and -184 km s^{-1} (Westmeier, Brüns & Kerp 2004, 2005b). Their shapes are interpreted as a consequence of ram-pressure by an ambient medium.

With models M1_dec and L1_dec, we are able to study the effect of drag from ambient gas, which leads to a deceleration of CHVCs. In Figs 5 and 6, it is striking that decelerated, massive clouds can become gravitationally unstable. They develop a head-tail morphology much later than the respective clouds that move at a constant speed. The mass-loss is dramatically reduced, irrespective of initial mass (Section 6.2). In Fig. 11, we show that the terminal velocity of M1_dec is reached after ~ 80 Myr and reads $\sim 58 \text{ km s}^{-1}$. The terminal velocity of L1_dec is still not reached after 100 Myr, where the cloud is decelerated to $\sim 48 \text{ km s}^{-1}$. The mean deceleration rates are -1.1 and $-1.2 \text{ km s}^{-1} \text{ Myr}^{-1}$, respectively. So both models are decelerated well below the velocity threshold for HVCs of $v_{\text{LSR}} = 100 \text{ km s}^{-1}$ thus showing a transition to intermediate-velocity clouds (IVCs).

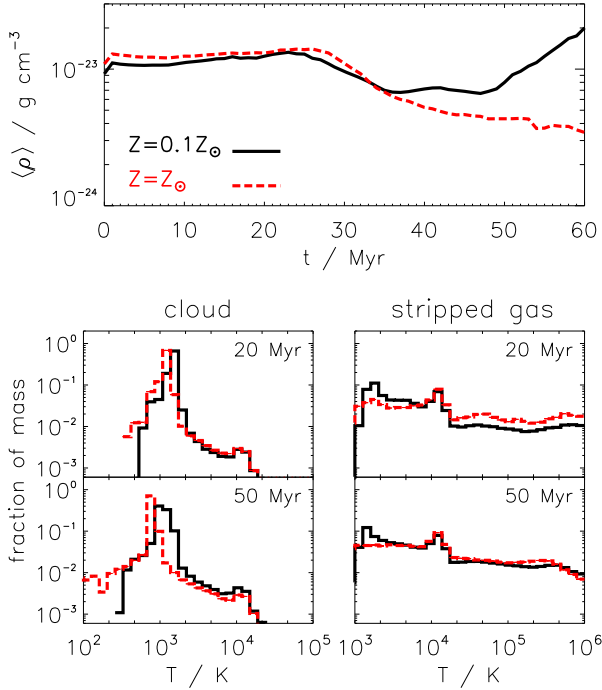


Figure 12. Upper panel: mean density of model clouds M0 (black solid line) and M1_me (red dashed line). Lower panel: distributions of temperature within the clouds (left-hand column) and in the stripped material (right-hand column) after 20 Myr (upper row) and after 50 Myr (lower row). The standard deviation in every bin is too low to be visible in the plots.

5.4 Radiative cooling

We explore the effect of cooling strength on the evolution of CHVCs by means of a different metallicity in the model clouds M0 ($0.1 Z_{\odot}$) and M1_me ($Z_{\odot}$). Because thermal conduction is active in both clouds, the stripping is filamentary and hence the stripped material bears the same temperature distribution irrespective of metallicity, hence irrespective of cooling strength (Fig. 12, right-hand column in lower panel). However, the cloud is cooler on average for $Z = Z_{\odot}$ at 50 Myr, even though the mean density inside cloud M1_me is significantly below that of M0 after 35 Myr (upper panel in Fig. 12). The left-hand in the lower panel in Fig. 12 shows that the peak of temperature distribution in M1_me is below and for M0 it is above 10^3 K. So, the metal-dependence of the cooling function is strong enough to have a measurable effect on cloud temperature.

5.5 Physical simplicity

Model M3_hom_tc_sg is intentionally physically simplified. We neither consider thermal conduction, nor self-gravity and the cloud is homogeneous and isothermal. In the respective plot in Fig. 5, it is observed that this cloud does neither develop a head-tail structure nor a bow-shock shape. It rather gets totally fragmented after ~ 60 Myr, because self-gravity does not stabilize the cloud against RTI thus the destructive modes can effectively grow. Such simplified cloud models do not maintain a density profile, which is a common feature being observed in CHVCs.

So, simplified clouds evolve completely different than hydrostatic cloud models with self-gravity and thermal conduction. They are not

Table 4. Observed CHVCs with high column densities.

Object	$\log (N_{\text{HI}} / \text{cm}^{-2})$	Reference
Six CHVCs	$\lesssim 21.0$	Braun & Burton (2000)
CHVC 125+41–207	20.1...21.0	Brüns et al. (2001)
HVC 291+26+195	$\lesssim 20.1$	Ben Bekhti et al. (2006)
HVC 297+09.253	$\lesssim 20.5$	Ben Bekhti et al. (2006)

realistic models in describing CHVCs, because they do not develop the same features as observations reveal.

6 DISCUSSION

6.1 Column densities

We have examined the column densities of neutral hydrogen, N_{HI} , as they would appear at an observers distance including beam dilution when using single-dish telescopes, which have spatial resolutions of several arcmin (e.g. Effelsberg, Green Banks, Parkes, Dwingeloo). By comparing column densities of observed CHVCs (Table 4) with our simulated massive clouds, we find reasonable agreement when accounting for beam dilution (Fig. 13, right-hand column). Exemplarily, we have chosen a beam with opening angle of 9 arcmin, like Effelsberg telescope, at a distance of 100 kpc, which is a typical distance for CHVCs. The diluted column densities of our simulated clouds are much lower (by ~ 2 orders of magnitude) depending on the filling factor of the cloud in the beam. By this, the very high peak column densities (Fig. 13, left-hand column) are no longer visible. With larger beams (e.g. Dwingeloo with 30 arcmin) or at greater distances (e.g. Braun & Burton 2000, find a distance bracket of 0.5–1 Mpc for CHVC 125+41–207) our simulated N_{HI} would appear even lower. Thus, by considering realistic beam dilutions the column densities of our model clouds do not contradict to observed CHVCs even though lying at the upper edge of consistency.

6.2 Survival of clouds

It has been demonstrated in Section 5 that it strongly depends on the combination of initial cloud mass, cloud velocity, and thermal conduction if CHVCs are being disintegrated or not during their passage through the CGM. There are several dynamical processes, which can lead to mass-loss of the clouds.

Several theoretical mass-loss rates have been determined to which we compare the mass-loss obtained in our simulations. Cowie & McKee (1977) determine the rate of mass-loss for a non-gravitating cloud by evaporation only (cooling effects are neglected) as

$$\dot{m}_{\text{CM}} = 3.25 \times 10^{18} n_{\text{CGM}} \sqrt{T_{\text{CGM}}} R_{\text{cloud}}^2 [\text{pc}] \Phi F(\sigma_0). \quad (35)$$

The strength of saturation is given by

$$F(\sigma_0) = 2 \left[(\sigma_0 H)^{1+Ma^2} \exp \{ -2.5 Ma^2 \} \right]^{1/(6+Ma^2)}, \quad (36)$$

where

$$H = \begin{cases} h^h / (h-1)^{h-1}, & Ma \leq 1 \\ 11.5, & Ma > 1 \end{cases}, \quad h = \frac{11 + Ma^2}{1 + Ma^2}, \quad (37)$$

and Ma is the Mach number of the flow around our clouds (cf. Table 3). Dalton & Balbus (1993) extended these investigations by introducing a continuous transition between classical and saturated heat flux. By simultaneously solving their equations (25) and (30) they obtain the ratio $\omega(\sigma_0)$ (cf. Fig. 1) between saturated and classical

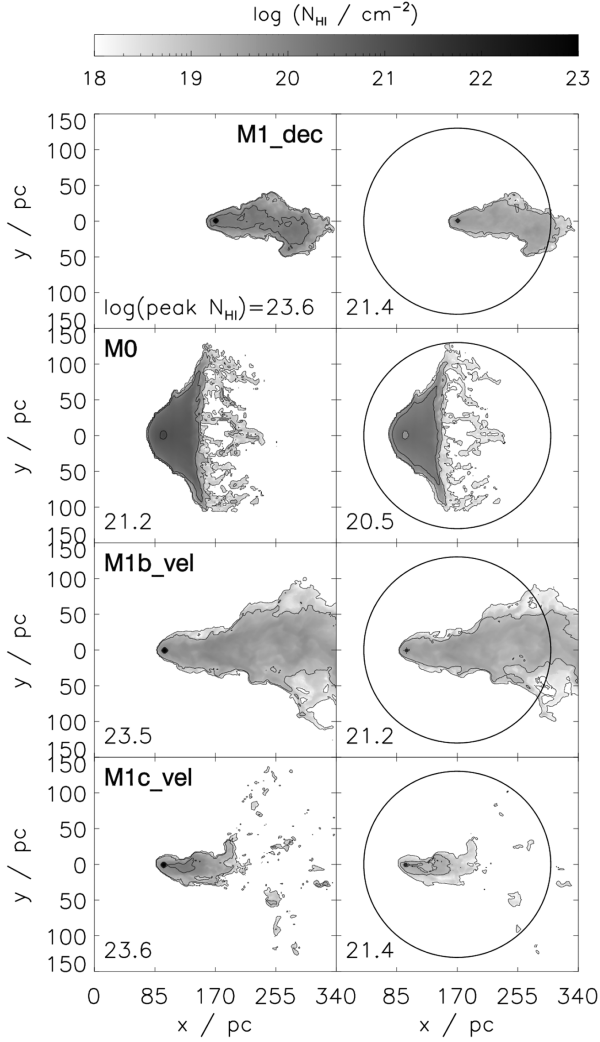
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Figure 13. Column densities of H I for selected massive clouds after 50 Myr (M1_dec after 100 Myr). The *black isodensity contours* represent $\log(N_{\text{H I}}/\text{cm}^{-2}) = 18, 19, \dots$ Left-hand column: actual $N_{\text{H I}}$ for the clouds. Right-hand column: diluted $N_{\text{H I}}$ as being obtained at a distance of 100 kpc by a beam of opening angle 9 arcmin. The *black circle* shows the area covered by this beam.

mass-loss rate. Hence,

$$\dot{m}_{\text{DB}} = \omega(\sigma_0) \dot{m}_{\text{class}} = \omega(\sigma_0) \frac{16\pi\mu_{\text{CGM}}\kappa_{\text{CGM}}R_{\text{cloud}}}{25k_{\text{B}}}, \quad (38)$$

with $\dot{m}_{\text{class}}$ given by Cowie & McKee (1977, see equation 16 therein) and our values for σ_0 are shown in Table 3. Hartquist et al. (1986) found

$$\dot{m}_{\text{H}} = [\varrho_{\text{CGM}}v_{\text{CGM}}(\bar{c}_{\text{cloud}}M_{\text{cloud}})^2]^{1/3} Ma^{4/3} \quad (39)$$

for the mass-loss of a cloud in a subsonic wind without considering self-gravity. As we consider self-gravity in all model clouds (except in model M3_hom_tc_sg), we also compare to the mass-loss rates analytically calculated by Arthur & Lizano (1997),

$$\dot{m}_{\text{AL}} \approx \frac{1}{24G} \frac{\bar{c}_{\text{cloud}}^4}{c_{\text{CGM}}} Ma, \quad (40)$$

Table 5. Mean mass-loss rates as obtained from the simulations (second column) compared to several analytical mass-loss rates given by equations (35), (38)–(40). All rates are given in units of $10^{-4} M_{\odot} \text{ yr}^{-1}$. Some analytical rates cannot be applied to models without thermal conduction or without self-gravity.

Model	$\langle \dot{m} \rangle_{\text{sim}}$	$\dot{m}_{\text{CM}}$	$\dot{m}_{\text{DB}}$	$\dot{m}_{\text{H}}$	$\dot{m}_{\text{AL}}$
M0	−3.6	−6.2	−7.7	−4.8	−2.7
M1_tc	−4.8	—	—	−2.7	−0.2
M1a_vel	−0.2	−8.6	−9.5	−1.9	−2.9
M1b_vel	−6.2	−1.9	−3.3	−4.9	−0.2
M1c_vel	−20.4	−1.7	−2.8	−9.6	−0.5
M1_dec	−0.2	−7.0	−8.2	−0.02	−0.06
M1_me	−0.7	−6.0	−6.7	−4.1	−5.6
M3_hom_tc_sg	−4.4	—	—	−4.7	—
L0	−1.6	−4.0	−4.4	−1.7	−1.4
L1_tc	−2.4	—	—	−0.3	−0.3
L1a_vel	−0.1	−4.3	−4.6	−0.8	−2.9
L1b_vel	−3.5	−4.1	−4.4	−2.9	−1.0
L1_dec	−1.9	−5.0	−4.9	> −0.01	−0.06
L1_me	−0.5	−1.5	−1.1	−0.03	−0.1

for self-gravitating clouds being subject to a subsonic wind. Here, $\bar{c}_{\text{cloud}}$ and c_{CGM} denote the average sound speed inside the cloud and in the CGM, respectively. The mass-loss in equation (40) accounts for the ablation by hydrodynamical instabilities.

Each of the definitions (35), (38)–(40) is based on only some of the physical processes we consider in our simulations. Thus, our mean mass-loss rates $\langle \dot{m} \rangle_{\text{sim}}$ are either underestimated or overestimated (Table 5). These analytical mass-loss rates are hence not sufficient to deduce a reliable mass-loss from multiphysics simulations of CHVCs. The clouds are subject to both KHI and RTI as a consequence of hydrodynamic interactions with the ambient gas. The KHI leads to the amplification of perturbations with wavenumbers

$$k > k_{\text{min}} \approx \frac{g\varrho_{\text{cloud}}}{\varrho_{\text{CGM}}v_{\text{rel}}^2}, \quad \text{if } \varrho_{\text{cloud}} \gg \varrho_{\text{CGM}}. \quad (41)$$

where $k = 2\pi/\lambda$, with λ being the wavelength of the perturbation, g is the gravitational acceleration at the cloud’s surface, ϱ_{cloud} and ϱ_{CGM} are the densities of CHVC and CGM, respectively, and v_{rel} denotes the relative velocity between them. The above relation (41) assumes discontinuities in both density and velocity between CHVC and CGM. Though this is valid initially (cf. Section 4.1), the dissipative effect of thermal conduction broadens the discontinuities over a finite domain (*transition zone*) hence converting them into very steep gradients. It has been shown by Vieser & Hensler (2007a, see their appendix B.2) that a finite region can be defined in k -space, which consists of all modes leading to an amplification of KHI. A cloud is stable against KHI for all modes outside of that instability strip, whose size inversely scales with the width of the transition zone. The thicker the transition zone, the thinner the instability strip. Thus, thermal conduction is able to suppress KHI in terms of decreasing the number of destructive modes. Our simulated model clouds M1_tc and L1_tc are subject to strong KHI that ablate big clumps of gas from the clouds. The disrupting KHI waves can propagate through the entire cloud hence leading to growing modes of instability everywhere in the cloud. Whereas the respective counterparts M0 and L0 are only mildly disturbed by KHI thus showing a continuous, filamentary, and low-density mass-loss (respective plots in Figs 5 and 6). One can deduce that thermal conduction extends the lifetime of even weakly gravitationally bound clouds by lowering the mass-loss.

RTI can develop at the front side of the CHVC when its initial density configuration is disturbed by the subsonically streaming CGM,

Table 6. Model clouds that survive their passage through the CGM for 50 Myr. The second column shows the remaining fraction of initial cloud mass after 50 Myr.

Model	$M_{\text{cloud}}(t = 50)/M_{\text{cloud}}(t = 0)$
M0	0.85
M1_tc	0.63
M1a_vel	0.99
M1b_vel	0.56
M1c_vel	0.75
M1_dec	0.97
M1_me	0.93
L0	0.65
L1a_vel	0.97
L1_dec	0.98

such that cloud gas is accelerated opposite to the direction of the density gradient in the cloud. If, like in model cloud M3_hom_tc_sg, self-gravity does not stabilize the cloud surface against RTI (Murray et al. 1993) these instabilities grow by time and the disturbed outskirts of the CHVC start to form typical finger-like structures (Fig. 5). These fingers of the dense cloud point into the dilute CGM in the opposite direction of motion. The wavelength of the dominant mode is $\lambda = 33$ pc (the separation between two adjacent 'fingers' in Fig. 5) and the growth rate is given by (Chandrasekhar 1961)

$$\omega^2 = (kv)^2 = ak \frac{\rho_{\text{CGM}} - \rho_{\text{cloud}}}{\rho_{\text{CGM}} + \rho_{\text{cloud}}}, \quad (42)$$

with $k = 2\pi/\lambda$ being the wavenumber of the perturbing mode and v is its phase velocity. By measuring $v \sim 2 \text{ km s}^{-1}$, the acceleration of the cloud surface is computed from equation (42) as $a \sim -3 \times 10^{-9} \text{ cm s}^{-2}$. According to Kull (1991) the e-folding time, i.e. the growing time for a certain mode, reads $\tau_e = (ka)^{-1/2}$. Hence the time, at which the spatial extension of the RT structures increases by a factor of e , is $\tau_e \sim 2.5$ Myr. The inverse denotes the growth rate $\tau_e^{-1} \sim 0.4 \text{ Myr}^{-1}$. Another dynamical effect the clouds are subject to is the stripping of material by ram pressure. Mori & Burkert (2000) determine the stripping radius for a galaxy of particular mass and moving at a certain speed. They balanced the thermal pressure of the cloud by the total pressure of the ambient medium, which in our case is

$$P_{\text{CGM}} = P_{\text{thermal}} + \rho_{\text{CGM}} v^2, \quad (43)$$

with P_{thermal} given by equation (4). The stripping radius for a specific cloud is the distance from cloud centre at which

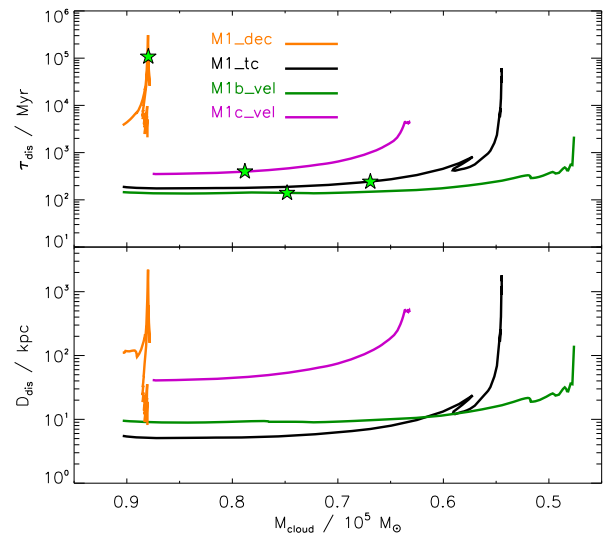
$$P_{\text{CGM}} > P_{\text{cloud}}(r) = \frac{k_B(r)T(r)}{\mu(r)}, \quad (44)$$

is true. Contrary to Mori & Burkert (2000), we consider self-gravity thus the cloud temperature is calculated by equation (26) while the corresponding density is constrained by thermal equilibrium. Equation (44) indicates disruption for all low-mass clouds, except model L1a_vel. However, according to Table 6 also the low-mass models L0 and L1_dec are not disrupted.

Analogously to Heitsch & Putman (2009), we calculate the disruption time $\tau_{\text{dis}} = M_{\text{cloud}}/(dM_{\text{cloud}}/dt)$ (time at which the cloud would be totally disrupted for a certain mass-loss rate) and disruption distance $D_{\text{dis}} = M_{\text{cloud}}/(dM_{\text{cloud}}/dz)$ (travelled distance after which the cloud would be totally disrupted for a certain mass-loss rate) for all our clouds. In opposite to the clouds simulated by Heitsch & Putman (2009), ours live much longer, because we assume much higher central densities and consider thermal conduction. Consequently, our mass-loss rates are much smaller than those in Heitsch & Putman

Table 7. Model clouds that experience a gravitational collapse with an onset at t_{collapse} (row 1). The following rows tabulate parameters specific to the instable regions after respective $t_{\text{collapse}} + 40$ Myr: mass in terms of cloud mass (row 2); size in terms of cloud size (row 3); mass in terms of Jeans mass (row 4); and gas surface-densities (row 5).

	M1_tc	M1b_vel	M1c_vel	M1_dec
$t_{\text{collapse}}/\text{Myr}$	29	17	19	47
M/M_{cloud}	1.00	0.98	1.00	0.90
V/V_{cloud}	1.00	0.58	0.64	0.51
M/M_J	140.4	8.0	25.7	1.6
$\Sigma/\text{g cm}^{-2}$	0.16	<0.01	0.02	<0.01

**Figure 14.** Both disruption time τ_{dis} and distance D_{dis} for clouds developing a gravitationally instable region. The light-green stars indicate the onsets of collapse.

(2009) and so our values of τ_{dis} and D_{dis} are greater than their values. In opposite to Heitsch & Putman (2009), four of our clouds undergo a gravitational collapse (Table 7). The values of τ_{dis} and D_{dis} for these clouds increase over time (Fig. 14) indicating that these clouds will not be disrupted. Such a result has not been observed in the clouds simulated by Heitsch & Putman (2009). In Sander & Hensler (2019), we have shown that self-gravity makes a distinct difference in the evolution of even homogeneous clouds with low densities and far below their Jeans masses.

Our massive clouds are comparable to model U in Vieser & Hensler (2007a), who already found that the mass-loss from a cloud with thermal conduction reduces after 60 Myr of evolution (see Section 5.2). This effect is even more pronounced in our simulations: For the entire simulation time, the mass-loss rate of model M0 does not exceed 75 per cent of the mass-loss rate of model M1_tc (Fig. 8 and Table 5). After 47 Myr, the mass-loss in model M1_tc reduces, because the cloud centre collapses, which is accompanied by an increasing density.

Armillotta et al. (2017) conducted simulations of CHVCs moving through a hot CGM. Their clouds with radii below 100 pc do not survive, which can be a consequence of a rather weak gravitational potential, because their clouds with radius of 50 pc contain only $160 M_{\odot}$, are homogeneous, isothermal, and do not consider self-gravity. Their clouds can most likely be compared to our model

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M3_hom_tc_sg even though their initial cloud is obtained by adopting the cloud density accordingly to environmental conditions for obtaining pressure equilibrium. Hence, model M3_hom_tc_sg is much more massive and compact as the clouds in Armillotta et al. (2017), but M3_hom_tc_sg is already disrupted by RTI. Here, physical simplicity leads to non-realistic results.

All of our model clouds, which are not disrupted after 50 Myr, are compiled in Table 6.

6.3 Gravitational collapse

In some of the CHVCs, which resist all disrupting effects (Table 6), local Jeans instabilities can develop. Within our simulations, all grid cells that belong to a cloud need to be identified. This is the case if a grid cell is

- (i) gravitationally bound, i.e. the potential energy at the cell's location is greater than its kinetic energy, and
- (ii) spatially connected to the cloud centre.

In that regard, we do not account for grid cells, which are gravitationally bound, but are already stripped off from the cloud. The cloud centre is defined to be the cell with the global minimum in potential energy. In contrast, for model M3_hom_tc_sg (no self-gravity) the above condition (i) is replaced by a threshold in density and temperature, i.e. $\rho_{\text{cell}} > 100 \rho_{\text{CGM}}$ and $T_{\text{cell}} < T_{\text{CGM}}$. To identify gravitationally instable regions, we adapt the procedure described in Federrath et al. (2010):

- (i) Find all instable grid cells in a cloud by checking if $M_{\text{cell}} > M_J$. Take these cells as the centres of prospective collapsing regions, because these cells will have the potential minimum of this region.
- (ii) Go radially outwards from each of the identified instable grid cells to their nearest neighbour cells and define this region to be the temporary control volume VC.
- (iii) Check for instability of the control volume: $M_{\text{VC}} > M_J$.
- (iv) Check for bounding of VC by applying the virial theorem: $|E_{\text{grav}}| \geq 2E_{\text{therm}}$.
- (v) If checks (iii) to (iv) are successful, go back to step (ii) and repeat the procedure for the new VC. If one of the checks (iii) to (iv) fails, the current extension of the control volume defines the size of the gravitationally instable cloud region.

By applying the above procedure to all surviving clouds (Table 6) we figured out four clouds (Table 7), whose central regions become gravitationally instable despite the fact that all models are initially below their Bonnor–Ebert mass (Fig. 3). Li et al. (2020) performed a variety of simulations of cool clouds in the CGM. For the sake of simplicity, their clouds are initially homogeneous and isothermal ($T \sim 10^4$ K) and are tuned to be in pressure balance with the ambient CGM. They find that self-gravity is only relevant for Jeans-instable clouds. However, this conclusion is drawn from too short evolution times of ≤ 1 Myr (see their fig. 2). We do the same comparison for our simplified model M3_hom_tc_sg by calculating it again with self-gravity. For a very early evolution ($\lesssim 2$ Myr), we find negligible differences between the clouds in both maximum particle density and mass-loss (Fig. 15), which supports the observations of Li et al. (2020). However, for later stages (> 40 Myr) differences become striking. Furthermore, in Sander & Hensler (2019), we have shown that self-gravity leads to a core-halo density profile even for resting (i.e. without dynamical interaction with the ambient medium) and initially homogeneous clouds with mass below their Bonnor–Ebert mass. In opposite to Li et al. (2020), the Jeans instabilities in our model clouds (Table 7) occur for a broad range of cloud velocities

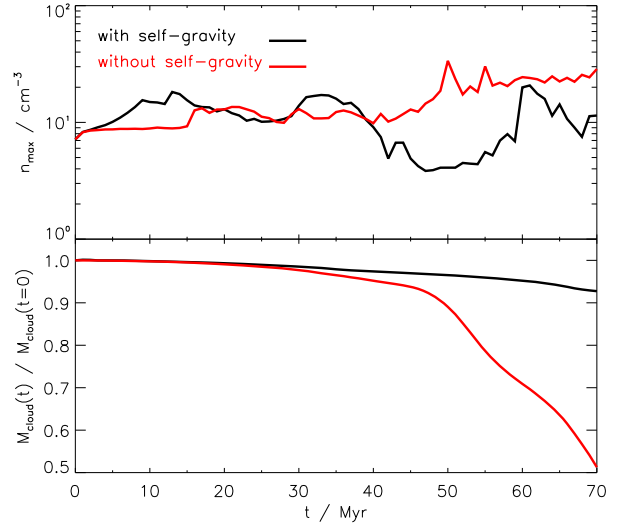


Figure 15. Evolution of maximum particle density (upper plot) and mass (lower plot) of homogeneous and isothermal clouds with and without self-gravity.

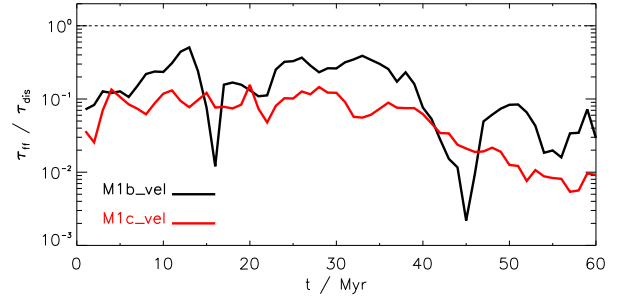


Figure 16. Free-fall time versus disruption time in fastest massive clouds.

and with or without thermal conduction. Hence, we do not find any evidence that supports the reasoning by Li et al. (2020) and we argue that self-gravity substantially affects the evolution of clouds with mass below their Bonnor–Ebert mass. If massive clouds are fast enough (M1b_vel, M1c_vel), neither thermal conduction nor efficient mass-loss can stop the gravitational collapse. For both models, the respective disruption time (Section 6.2) is always greater than the free-fall time (Fig. 16). We conclude that clouds with thermal conduction are able to collapse if they are very fast and sufficiently massive.

The Jeans-instable regions in the four models evolve differently, which we discuss in the following by means of Fig. 17 and Table 7. The instable region in the most realistic model, M1_dec, forms much later (47 Myr) than in the other three models, which is most likely due to both a decreasing dynamic pressure by deceleration and an enhanced thermal pressure in the cloud by thermal conduction. At this time, the cloud velocity has decreased below the threshold for HVCs (Fig. 11).

We observe that the size of the instable region increases to roughly half of the entire cloud size, but contains nearly all of cloud mass. The mean density is rather low and results in a gas surface-density below 0.01 g cm^{-2} . The thermal energy conducted into the instable region enhances the thermal pressure. At a certain point in evolution, this cannot be compensated for by the decreasing dynamic pressure,

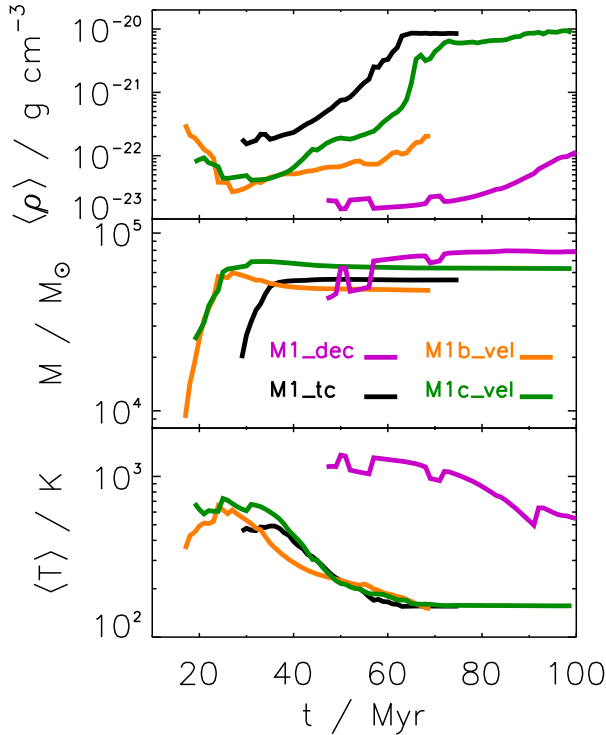


Figure 17. Evolution of mean density (upper plot), mass (middle plot), and mean temperature (lower plot) of gravitationally unstable regions.

which is due to cloud deceleration. The instable region thus expands. The density hence remains at low level and so is the cooling strength. Finally, the temperature in the instable region of model M1_dec is high compared to those in the other model clouds. As a consequence, its mass does not substantially exceed the Jeans mass.

In the fast models with thermal conduction, M1b_vel and M1c_vel, the ram-pressure is constantly high during the entire evolution and so the instable regions are compressed to much higher densities. This situation is more pronounced for the faster cloud M1c_vel, which is subject to a higher ram-pressure than model M1b_vel. The temperatures reach values ~ 150 K thus lowering the Jeans mass of the instable region. None the less, gas surface-densities are comparably low, too.

Model M1_tc is distinct from the other three models. The entire cloud becomes gravitationally instable. Without thermal conduction the thermal pressure in the cloud is low such that at a lower speed as in models M1b_vel and M1c_vel the ram-pressure is sufficient to compress the cloud such that the density can reach $10^{-20} \text{ g cm}^{-3}$. Being at a mean temperature of ~ 150 K the instable region greatly exceeds its Jeans mass. The gas surface-density reaches a value of 0.16 g cm^{-2} . According to Traficante et al. (2020, and references therein), the threshold of gas surface-density for a region to form high-mass stars is $0.05\text{--}0.1 \text{ g cm}^{-2}$ (stellar feedback would increase this threshold to 1 g cm^{-2} , Krumholz & McKee 2008). Thus, our cloud M1_tc would virtually be able to form even high-mass stars.

As has been discussed above, clouds with thermal conduction develop instable regions, but these do not get compact enough to become a star-formation site. For our studied cloud setups the instable regions are instead in some balance between gravity and thermal pressure. We conclude that thermal conduction might be a key process in preventing CHVCs to form stars. These findings

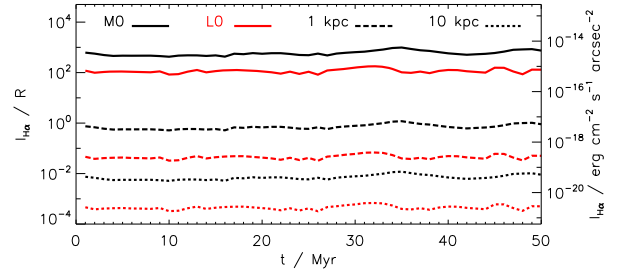


Figure 18. Evolution of $\text{H}\alpha$ intensity as measured for the simulated reference clouds M0 (black solid line) and L0 (red solid line). The dashed-lines and dotted lines show the respective intensities at distances of 1 and 10 kpc, in order to compare to Putman et al. (2003).

are in line with non-detections of stellar objects assigned to HVCs, in general, in all hitherto performed surveys (e.g. Burton, Braun & Chengalur 2001; Hopp et al. 2003a,b; Siegel et al. 2005; Hopp, Schulte-Ladbeck & Kerp 2007).

6.4 $\text{H}\alpha$ emission

A still puzzling feature of CHVCs is their partially bright emission in the $\text{H}\alpha$ disexcitation-line. In the following, we analyse how thermal conduction is able to contribute to the intensity of $\text{H}\alpha$ emission. We estimate the $\text{H}\alpha$ intensity emitted isotropically over 4π by any grid cell j of our model clouds by

$$I_{\text{H}\alpha}^{(j)} = A^{3 \rightarrow 2} S_{\text{H}\alpha}^{3 \rightarrow 2} \text{ erg cm}^{-2} \text{ s}^{-1}, \quad (45)$$

with $A^{3 \rightarrow 2} = 4.39 \times 10^7 \text{ s}^{-1}$ denoting the Einstein coefficient for spontaneous emission from $n = 3 \rightarrow 2$. By assuming ionization equilibrium (cf. Section 2.4 and appendix A), the source function $S_{\text{H}\alpha}^{3 \rightarrow 2}$ is given by Boltzmann's Equation for excitation temperature. As the clouds are assumed to be optically thin, each of the j contributions in equation (45) has to be summed up to the total amount of $I_{\text{H}\alpha}$. Thus, for N cloud cells

$$I_{\text{H}\alpha}^{\text{cloud}} = \sum_{j=1}^N I_{\text{H}\alpha}^{(j)} \quad (46)$$

yields the absolute value of $\text{H}\alpha$ intensity emitted by the cloud, which must be reduced by a dilution factor $\propto D^{-2}$ and corrected for the opening angle, Ω , of the cloud when it is observed from a certain distance D under a particular angle Ω . A correction for an optically thick cloud with column density above 10^{18} cm^{-2} is found in Barger et al. (2012, cf. their equation 7). Furthermore, the emitted $\text{H}\alpha$ radiation must be corrected for extinction in the line of sight, which is highly uncertain (Hill et al. 2013).

Observers mostly provide $\text{H}\alpha$ intensities in units of Rayleigh ($1\text{R} = 2.41 \times 10^{-7} \text{ erg cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$, van Woerden et al. 2004) or in units of $\text{erg cm}^{-2} \text{ s}^{-1} \text{ arcsec}^{-2}$ (where $\text{sr}^{-1} = 2.35 \times 10^{-11} \text{ arcsec}^{-2}$). In Fig. 18, the evolution of $\text{H}\alpha$ intensity (46) is shown for the reference clouds M0 and L0 directly at cloud surfaces and at distances of 1 and 10 kpc. For that, we only consider dilution by distance, but no extinction in the line of sight. It is striking that the intensity of massive clouds is always higher than that of low-mass clouds, which is due to a greater number of contributing cloud cells. Interestingly, the $\text{H}\alpha$ intensity is nearly constant for both clouds. This can be understood in terms of two counteracting processes: the cloud loses mass thus the number of grid cells contributing to the intensity decreases. But simultaneously the cloud is heated up and so the intensity increases. Both effects roughly cancel out.

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When comparing to observed emission measures of $H\alpha$ from CHVCs it turns out that the calculated intensities for our simulated clouds are at the same orders of magnitude for assumed distances of >1 kpc. Kutyrev & Reynolds (1989) report on $H\alpha$ emission from an HVC in constellation Cetus with a best-fitting mean value of $(8.1 \pm 1.9) \times 10^{-2}$ R, which consists of two fields with $(7.3 \pm 2.5) \times 10^{-2}$ R for field 1 and $(8.9 \pm 2.8) \times 10^{-2}$ R for field 2. They infer the speed of the HVC ($v_{\text{LSR}} = -310 \text{ km s}^{-1}$) to be supersonic and the ambient gas density to be $n \lesssim (4 - 6) \times 10^{-3} \text{ cm}^{-3}$ in order to produce the observed shock-excited optical emission in $H\alpha$. Both the cloud speed and CGM density in models M0 and L0 are below these values. Putman et al. (2003) performed deep $H\alpha$ spectroscopy towards several HVCs, among which there are five CHVCs with 20 to 220 mR at distances between 1.2 and 13.2 kpc. These clouds are exposed to a non-negligible ionizing UV background radiation in the Galactic halo due to 6 per cent of UV photons that escape normal from the Galactic disc (Bland-Hawthorn & Maloney 2002). The calculated $H\alpha$ intensities of our model clouds are thus similar to those of observed CHVCs. Note that we overestimate the intensities in our calculations, because we do not consider extinction in the line of sight. When comparing in detail to individual observed CHVCs it is likely that thermal conduction alone is not a proper mechanism to reproduce real $H\alpha$ intensities. An additional mechanism to enhance $H\alpha$ intensity by heating up the clouds might be provided by a diffuse, soft X-ray background as being stated by Bregman (2007). The cosmic UV background alone would account for only a negligible $H\alpha$ intensity of 5 mR and is exceeded by the Galactic UV radiation field for distances from the disc of $\lesssim 100$ kpc (Maloney & Bland-Hawthorn 1999).

7 SUMMARY AND CONCLUSIONS

We conducted 3D hydrodynamics simulations of multiphysical processes in HVCs by using the FLASH code. The capability of adaptive mesh refinement is extensively used with a finest numerical resolution of 1.3 pc. By this, we are able to sufficiently resolve hydrodynamical instabilities (Kelvin–Helmholtz, Rayleigh–Taylor) as well as the field length for thermal conduction. We tracked and compared the evolution of 14 different model clouds which differ in mass, size, metallicity, and speed. All models are initially in hydrostatic and thermal equilibrium internally and in pressure equilibrium with the ambient, hot medium. The physical processes of metal-dependent plasma cooling and heating, dissociation and ionization, and mass diffusion take place in all models. The model clouds either consider or exclude self-gravity and thermal conduction. Hence, we can directly compare the impact of cloud mass, speed, drag, metallicity, thermal conduction, and modelling simplicity on cloud evolution and discuss our findings in context of both their life-times and a potential gravitational collapse.

In total, 10 out of the 14 models survive, the first 50 Myr of their evolution (Table 6). These are all massive clouds with self-gravity (seven models) and three low-mass model clouds. To be not disrupted within 50 Myr, the low-mass clouds need thermal conduction, must not be too fast, and must not be too tiny in size (model L1_me). Finally, four out of these 10 models are not disrupted and in addition develop a local gravitational instability (Table 7). These findings contradict to other studies, where clouds are only resistive to disruption for much higher masses (Heitsch & Putman 2009), or larger radii and lower velocities (Armillotta et al. 2017).

By being too simplistic in the modelling of clouds they are easily disrupted by RTIs within 10 sound-crossing times. Wrong conclusions may be drawn if self-gravity, thermal conduction, and a core-halo density profile are not accounted for.

The metallicity of clouds affects the strength of radiative cooling. A high metallicity of $Z_{\odot}$ tends to produce cool clouds at mainly 10^2 – 10^3 K and with $0.1 Z_{\odot}$ the clouds are mainly at temperatures between 10^3 and 10^4 K. Since the observed core-halo structure in some CHVCs reveals core temperatures of only a few 10^2 K we conclude that metallicity is higher there, probably above $0.3 Z_{\odot}$.

It turns out that thermal conduction is a key process in the evolution of CHVCs. Even though all of our simulated low-mass clouds will be disrupted, thermal conduction delays disruption by substantially lowering the mass-loss rate. Clumpy fragments of high density, which can cool efficiently, are stripped from clouds without thermal conduction while thermal conduction causes a continuous, low-density, filamentary stripping of gas at high temperatures, which is more likely to resemble the smooth density gradients observed between head and tail in some CHVCs. We thus conclude that thermal conduction is at work in real CHVCs.

Massive clouds are gravitationally stable for sufficiently low velocities (Mach $\lesssim 0.5$), but for Mach $\gtrsim 0.7$ they indeed form Jeans-unstable regions. However, if thermal conduction is suppressed, the massive clouds can collapse even at Mach ~ 0.5 (model M1_tc) and are able to reach gas surface densities (0.16 g cm^{-2}) being high enough to account for a prospective star formation site. We strongly suppose that thermal conduction is a key process in preventing real CHVCs from forming stars.

It is not realistic, that clouds are not decelerated but keep constant speed, because the CGM transfers momentum to the clouds. Combining both physical effects of life-prolonging thermal conduction and stripping-reducing drag (cloud M1_dec) an unstable region can form. The collapse commences not before the clouds are already decelerated to speeds representative for IVCs. However, the decreasing ram-pressure is not sufficient to dominate the regions' thermal pressure maintained by thermal conduction. Hence, the unstable region cannot be compacted to surface densities high enough for star formation.

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8 DATA AVAILABILITY

The data underlying this paper (simulation results, analysis scripts) will be shared on a reasonable request to the corresponding author.

²see <http://vsc.ac.at/>

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APPENDIX A: DISSOCIATION AND IONIZATION IN THERMODYNAMIC EQUILIBRIUM

At sufficiently high temperatures the electronic bond of H_2 can be broken with

$$y := \frac{\varrho(H)}{\varrho(H) + \varrho(H_2)} \quad (A1)$$

being the fraction of dissociated H_2 , where ϱ denotes the respective mass density. In terms of the dissociation constant for H_2 (Black & Bodenheimer 1975), one obtains the conditional equation for y

$$\frac{y^2}{1-y} = \frac{2.11}{\varrho X} \exp\left\{\frac{-52,490}{T}\right\} =: F, \quad (A2)$$

where X is the mass fraction of hydrogen. Solving equation (A2) for y yields the degree of dissociation

$$y(\varrho X, T) = -\frac{F}{2} + \sqrt{\frac{F^2}{4} + F}. \quad (A3)$$

If temperatures increase above $T = E_{\text{ion}}/k_B$, with $E_{\text{ion}} = 13.595$ eV being the ionization level, hydrogen is going to be ionized. The necessary ionization energy can be provided by radiation (photons) or collisions (electrons). Vice versa, the ionized hydrogen can recombine again when it loses energy via radiation or collisions. Let n_0 and n_1 be the particle densities of neutral and ionized hydrogen, n_e be the electron number density, $R_{0 \rightarrow 1}$ and $C_{0 \rightarrow 1}$ be the rates of ionization by photons and collisions, and $R_{1 \rightarrow 0}$ and $C_{1 \rightarrow 0}$ be the respective recombination rates. We consider radiative ionization of hydrogen by soft X-rays with an ionization rate of $R_{0 \rightarrow 1}^x = 7 \times 10^{17} \text{ s}^{-1}$, and collisional ionization by both cosmic rays with an ionization rate of $C_{0 \rightarrow 1}^{\text{cr}} = 7 \times 10^{17} \text{ s}^{-1}$ and electrons with an ionization rate of $C_{0 \rightarrow 1}^{\text{el}} = n_e Q_{0 \rightarrow 1}^{\text{el}}$. The rate coefficient for collisional ionization by electrons

$$Q_{0 \rightarrow 1}^{\text{el}}(T) = 5.465 \times 10^{-11} \sqrt{T} \times \exp\left\{-\frac{13.543 \text{ eV}}{kT}\right\} \times (-0.435 + 0.3 \log T) \text{ cm}^3 \text{ s}^{-1}. \quad (A4)$$

is taken from Mihalas (1967).

On the other hand, we consider recombination by radiation with a rate $R_{1 \rightarrow 0}^{\text{rad}} = n_e \sum_{i=2}^{\infty} \alpha_i$. Captures to level $i = 1$ are excluded, since the released photons do not contribute to the net recombination ('on the spot approximation'). Hence, only recombination coefficient (e.g.

Spitzer 1998)

$$\alpha_2(T) = \frac{2.06 \times 10^{-11} Z^2}{\sqrt{T}} \phi_2\left(\frac{h\nu}{kT}\right) \text{ cm}^3 \text{ s}^{-1}, \quad (A5)$$

which includes all recombinations to level $i = 2$, is taken into account. The atomic number $Z = 1$ for hydrogen. The function ϕ_2 slowly varies with T and is tabulated in Spitzer (1998). So, in thermodynamic equilibrium,

$$n_0 (R_{0 \rightarrow 1}^x + C_{0 \rightarrow 1}^{\text{cr}} + n_e Q_{0 \rightarrow 1}^{\text{el}}) = n_1 n_e \alpha_2, \quad (A6)$$

with n_e being approximated by $n_e \approx n_1 + \chi_0(n_0 + n_1)$, where the second summand counts free electrons being released by species having lower ionization levels than hydrogen. The fraction of these species is given by $\chi_0 = 5 \times 10^{-4}$ (Dalgarno & McCray 1972). Let

$$f := \frac{n_1}{n_0 + n_1} \quad (A7)$$

be the fraction of ionized hydrogen. By substituting equation (A7) into equation (A6) one obtains

$$0 = f^2 + A f - B, \quad (A8)$$

with

$$A := \left[\chi_0 \alpha_2 + (\chi_0 - 1) Q_{0 \rightarrow 1}^{\text{el}} + \frac{R_{0 \rightarrow 1}^x + C_{0 \rightarrow 1}^{\text{cr}}}{n_0} \right] \left[\frac{1}{\alpha_2 + Q_{0 \rightarrow 1}^{\text{el}}} \right],$$

$$B := \left[\frac{R_{0 \rightarrow 1}^x + C_{0 \rightarrow 1}^{\text{cr}}}{n_0} + \chi_0 Q_{0 \rightarrow 1}^{\text{el}} \right] \left[\frac{1}{\alpha_2 + Q_{0 \rightarrow 1}^{\text{el}}} \right] \quad (A9)$$

are subsuming the rates of ionization and recombination. Solving equation (A8) for f yields the degree of ionization

$$f(n_0, T) = -\frac{A}{2} + \sqrt{\frac{A^2}{4} + B}. \quad (A10)$$

APPENDIX B: CODE TEST FOR IMPLEMENTATION OF SATURATED THERMAL CONDUCTION

We are going to show the evolution of a Gaussian temperature distribution being subject to diffusion only.³ For any twice differentiable, scalar function $f(x, y, z, t)$, $x, y, z, t \in \mathbb{R}$, $f: \mathbb{R}^3 \times \mathbb{R} \rightarrow \mathbb{R}^3 \times \mathbb{R}$, satisfying the heat equation

$$\frac{\partial f}{\partial t} - \left(\alpha_x \frac{\partial^2 f}{\partial x^2} + \alpha_y \frac{\partial^2 f}{\partial y^2} + \alpha_z \frac{\partial^2 f}{\partial z^2} \right) = 0, \quad (B1)$$

a delta distribution at any initial time t_0

$$f(x, y, z, t_0) = Q(t_0) \delta(x, y, z) \quad (B2)$$

spreads into a Gaussian distribution

$$f(x, y, z, t) = \frac{Q(t_0)}{[4\pi(t-t_0)]^{3/2} (\alpha_x \alpha_y \alpha_z)^{1/2}} \times \exp\left\{-\frac{1}{4(t-t_0)} \left(\frac{x^2}{\alpha_x} + \frac{y^2}{\alpha_y} + \frac{z^2}{\alpha_z} \right)\right\} \quad (B3)$$

and remains a Gaussian distribution for all times $t > t_0$. The analytical result (B3) is well known. By identifying $f(x, y, z, t)$ with the temperature field of a fluid and taking $\alpha = \alpha_x = \alpha_y = \alpha_z$ (isotropic diffusion), we obtain the evolution as being determined

³This test problem comes along with the FLASH code and is adapted accordingly by us.

by equation (B3). We point out that α is a function of the effective conductivity κ_{eff} (14). We have chosen $Q(t_0) = 10^4 \text{ K cm}^3$ at $t_0 = 0.5 \text{ s}$ (to avoid numerical difficulties implied by the Delta function at $t = 0$) and a both spatially and temporally constant density of $\varrho = 1 \text{ g cm}^{-3}$. The centre of the distribution is located at $(x_c, y_c, z_c) = (250, 250, 250) \text{ cm}$ and the side length of the computational domain is 500 cm. The domain is discretized by 128 cells per dimension yielding a grid spacing of $\Delta x = 3.9 \text{ cm}$. We use periodic boundary conditions. The background temperature is fixed at 10^{-3} K . We find $f(x_c, y_c, z_c, t_0) = 0.3 \text{ K}$ accounting for an initial temperature contrast of $\Delta T = 300$. In Fig. B1, the analytical solution (B3) is compared to the numerical one obtained by the FLASH code. Initially, the relative error is highest where the Gaussian temperature distribution passes into the background temperature: the change in slope cannot be approximated sufficiently by the grid, i.e. the numerical resolution is too coarse. However, the relative error is lowered to a reasonable level after 10 diffusion times ($\tau = 0.5 \text{ s}$). Boundary effects, which may superpose the numerical solution, are not observed. The numerical solution is hence sufficiently accurate.

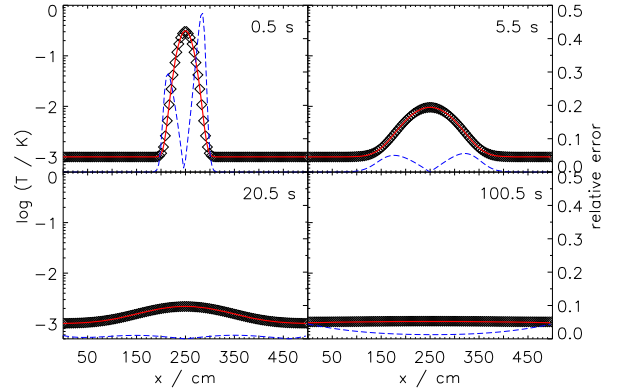


Figure B1. Numerically obtained evolution of a three-dimensional Gaussian temperature distribution being subject to diffusion only (black diamonds) compared with the analytical solution (red solid line) at different times. The blue dashed line shows the relative error between both solutions. The temperature profile is shown along the x -direction.

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7.7 DISCUSSION AND CONCLUSIONS

Within this section the main results of the simulations carried out in this thesis are discussed. The main conclusions are drawn to relate the simulation results to observations and to understand the role of CHVCs within the Galactic matter cycle. For detailed discussions see the published scientific articles presented in the respective sections 7.4, 7.5, and 7.6.

Self-gravity in low-mass clouds

A first major result regards self-gravity. By means of comparison of simulated resting model clouds it is shown that self-gravity governs the evolution of even homogeneous low-mass clouds that have only a few per cent of their respective Jeans mass. A low-mass, self-gravitating cloud, that is initially homogeneous, forms a radial density gradient thus leading to a hydrostatic stratification, and it remains compact and nearly spherical. Such core-halo structures are observed in many HVCs. Oppositely, without self-gravity considered an identical cloud evolves into a still homogeneous, irregularly shaped, and expanding cloud.

Common model simplifications may thus lead to spurious results. As an example, the disruption time of simulated HVCs may be drastically underestimated, the mass-loss rates may be overestimated, and a gravitational collapse may not happen if self-gravity is neglected (Heitsch & Putman, 2009). Moreover, without any radial density gradient the mixing process is completely different in terms of spatial distribution of accreted gas and mixing velocities (Sander & Hensler, 2021).

Effectivity of thermal conduction

As a second key result, an analytic calculation has shown that thermal conduction is effective even in presence of strong magnetic fields. Thermal conduction is examined in presence of a magnetic dipole field shielding a cloud. The effective heat flux towards the cool cloud is not suppressed significantly when integrated over the entire surface. It is only suppressed to $\sim 2/3$ of the unperturbed heat flux. In more complicated configurations of magnetic fields similar results are obtained (e.g., tangled fields, turbulent fields, multipole contributions, Narayan & Medvedev, 2001; Cho et al., 2003).

Hence, thermal conduction applies in much more astrophysical situations than assumed in some simulations of interstellar clouds in general and HVCs in special (e.g., Joung et al., 2012). However, other works support the necessity for taking thermal conduction into account.

Conductive transition zones

Owing to saturated thermal conduction a structured interface (transition zone) forms between multi-phase clouds and the ambient hot gas as predicted by analytical theory (Cowie & McKee, 1977). However, the transition zone that forms in the simulations within this thesis, which consider time-dependent thermal conduction, is much narrower and located closer around the cloud. The thickness of the transition zone effects the resilience of the cloud against KHI (Vieser & Hensler, 2007b).

The results in this thesis regarding conductive interfaces (transition zones) coincide with a most likely origin of observed Si IV/O VI, C IV/O VI, and N V/O VI ionic ratios in Complex C (Fox et al., 2004; Wakker et al., 2005), Complex M (Savage et al., 2005), the Magellanic Stream (MS) (Sembach et al., 2000), and other HVCs (Collins et al., 2007; Shull et al., 2011). The observed ratios are not consistent with collisional ionization equilibrium or turbulent mixing layers. Especially O VI is produced in transition-temperature plasma ($T \sim 10^5$ K, Narayanan et al., 2010), which is present in the transition zones formed in our simulations. In contrast, Gnat et al. (2010) apply non-equilibrium cooling and thermal conduction to calculate the ionization structure of transition zones. They conclude, that conductive interfaces increase the formation of highly ionized species, but the column densities obtained are not sufficient to account for the O VI columns observed in ionized HVCs.

The calculated $H\alpha$ intensities of the model clouds are similar to observed intensities in CHVCs (Putman et al., 2003a). However, thermal conduction alone is not sufficient to produce observed intensities hence other mechanisms likely contribute (Maloney & Bland-Hawthorn, 1999; Bregman, 2007).

Condensation of ambient hot gas

Within this thesis the findings of Vieser & Hensler (2007a) are confirmed that clouds with thermal conduction, which are embedded in a hot plasma, condensate external gas rather than being evaporated. For the first time in literature, condensation rates are quantified for a set of simulated multi-phase clouds differing in size, mass, radial density profile, and initial metallicity. A positive, linear correlation between radial density gradient and mean condensation rate is found. The condensation velocity is one order of magnitude below the sound speed of the ambient hot gas. Condensation rates may even be enhanced due to hydrodynamic drafting if the clouds are located in a multi-cloud environment (Williams & Shelton, 2022).

Mixing of cloud gas with ambient hot gas

The mixing dynamics of condensed hot gas with cloud gas is found to depend on the

radial cloud-density gradient and on thermal conduction. Without thermal conduction, mixing is literally not present. If saturated thermal conduction is accounted for, material is condensed with $10 \dots 20 \text{ km s}^{-1}$ it is decelerated to $\lesssim 2 \text{ km s}^{-1}$ in the transition zone. Thus, hot gas accumulates faster in the cloud outskirts than being transported into the central regions. Inside the cloud, the velocity of the mixing front and the spatial distribution of ambient material all relate to the radial density gradient. The more homogeneous the cloud is the easier and faster the condensed material gets distributed. Clouds with a steep density gradient accumulate most of the condensed gas in their outskirts thus establishing more persistent radial metallicity gradients. Our simulations of moving clouds confirm the results by Heitsch et al. (2022), who find that turbulent mixing does not enrich the HVCs substantially before $\sim 140 \text{ Myr}$ of evolution.

Gas removal from CHVCs

The simulations show that thermal conduction delays disruption of CHVCs by considerably lowering their mass-loss rates. In models without thermal conduction the removed gas consists of dense, cool, clumpy fragments, which is a clear indicator of substantial KHI (Barger et al., 2020). Oppositely, thermal conduction leads to continuous, filamentary stripping, while the removed gas is heated up quickly and mixes efficiently with the ambient HIM. The smooth density gradients are representative for observed head-tail CHVCs.

The difference in mass loss by thermal conduction confirms the findings by Vieser & Hensler (2007b). Armillotta et al. (2017) conduct simulations of the same kind as presented within this thesis, but in two dimensions. They find that KHI at the cloud surface are suppressed if thermal conduction is accounted for. However, all of their clouds are disrupted as a consequence of very low density and a lack of self-gravity. Thus, physical simplicity leads to completely different results.

Cloud morphologies

Based upon initial conditions the simulated CHVCs develop a head-tail, bow-shock, horseshoe-shaped, or completely irregular morphologies. All these morphologies are observed in many HVCs (Brüns & Mebold, 2004; Westmeier et al., 2004, 2005b,c). Within the simulations of moving clouds, horseshoe-shaped morphologies are generally observed at earlier stages of evolution while a head-tail structure forms later.

Star formation in HVCs

The initial conditions lead to either disruption, a gravitationally stable cloud, or a surviving cloud with onset of Jeans instabilities. Only four models develop gravitational

instabilities. By incorporating the physical processes of self-gravity, drag of ambient circumgalactic medium (CGM), heating and radiative cooling of plasma, and thermal conduction in one of the model clouds, we find a natural explanation for the suppression of star formation in CHVCs: the cloud starts to contract triggered by ram-pressure due to its motion relative to the ambient CGM. Simultaneously, the drag decelerates the cloud. Continuous heat input by thermal conduction increases the thermal pressure inside the cloud such that at some level the contraction ceases and finally stalls (figure 17 in Sander & Hensler, 2021). The densities do not reach values being necessary for star formation (Krumholz & McKee, 2008; Larson, 2010b; Clark & Glover, 2014; Traficante et al., 2020). Only the massive model cloud, which lacks the heat input by thermal conduction, would locally develop such high gas surface-densities to account for star formation. Without additional heat input by thermal conduction, the Jeans mass of the cloud remains low and is more easily exceeded as in the model cloud with thermal conduction. Additionally, much more gas is removed from simulated clouds without thermal conduction. It is thus much likely, that these model cloud do not live long enough to form stars.

To conclude, thermal conduction is a key process in preventing CHVCs from forming stars. This central result from the model cloud simulations is consistent with the observation that there are no stars assigned to HVCs. Alike the findings within this thesis, Afruni et al. (2023) find in their simulations that clouds accreted from the intergalactic medium (IGM) are not sufficient to sustain star formation in low-redshift disk galaxies. Werk et al. (2019) report that gas in form of intermediate-velocity clouds (IVCs), which crosses the disk-halo interface, may contain enough total mass to maintain the current star-formation rate of the Milky Way. Since both IVCs and HVCs do not form stars, they may serve to refuel gas exhausted by star formation in the disk.

The focus of interest within this thesis lies with a long-standing research question that has remained unsolved since the discovery of high-velocity clouds (HVCs) in 1963 by Muller et al. (1963). It is still completely unclear why no stars are observed in HVCs. At least within the densest candidates, which move at highest speeds through the most extreme regions of the Galactic hot halo, subregions could theoretically be compressed to the extent that their Jeans mass is exceeded. This distinct category of compact high-velocity clouds (CHVCs) has been investigated in particular in specific observational surveys and many hundreds of them are known, some of which show column densities far above 10^{21} cm^{-2} . Distinct density morphologies like head-tail, bow-shock, or horseshoe shapes imply strong interactions with the ambient hot gas. Core-halo structures in density, temperature, and ionization suggest that CHVCs are multi-phase and are able to radiatively cool over many orders of magnitude in temperature. Thus, an obvious contradiction exists between theoretical preconditions for star formation and observational non-detections of stellar counterparts assigned to CHVCs.

Within this thesis three-dimensional, multi-physics simulations of a suite of multi-phase clouds residing in a Galactic hot halo have been carried out by using the adaptive mesh-refinement code FLASH. The software used in this work was in part developed by the DOE NNSA-ASC OASCR Flash Center at the University of Chicago and the simulations were carried out using the Vienna Scientific Cluster. The FLASH code provides a strang-split method to solve for hydrodynamics, and a multi-grid solver to account for self-gravity. Within this thesis, the physical processes of saturated thermal conduction, plasma cooling and heating, mass diffusion, dissociation and ionization, a variable metallicity, and a temperature-dependent mean molecular weight have been added to FLASH. Semi-implicit and semi-analytical numerical schemes were implemented to treat these physical processes appropriately. The modelled Galactic hot halo resembles the super-virial phase observed in the circumgalactic medium (CGM) thus providing extreme conditions of the highly ionized medium (HIM).

The multi-phase model clouds separate into resting clouds and CHVCs. The clouds differ in mass, radius, metallicity, and radial profiles of density, temperature, and pressure. The low-mass clouds are nearly homogeneous and entirely dissociated while the massive clouds have distinct radial profile gradients and a molecular core. Self-gravity is found to be important even in the evolution of clouds with masses far below their respective Jeans mass.

The resting clouds are used to study saturated thermal conduction. In contrast to analytical theory the simulated clouds condensate ambient hot gas. The condensation rates are quantified and a linear correlation with the radial density gradient is found. The transition zone structures as analytically expected, but is much narrower, and closer around the cloud, which has consequences for the resilience against Kelvin-Helmholtz instability (KHI). The temperatures in the transition zones enable the formation of highly ionized elements (especially O VI) as being observed in CHVCs. It is shown analytically that thermal conduction is an essential mechanism to transport heat into interstellar clouds even in presence of strong magnetic fields.

The moving CHVCs are exposed to different subsonic potential flows, are infalling from an extragalactic origin, and are located at adapted galacto-centric distances of around 100 kpc. They are at the high end of observed column densities, are consistent with observed masses, multi-phase structures (core-halo), morphologies (head-tail, bow-shock, horseshoe-shaped). The modelled ambient medium they move through represents the super-virial phase of the Galactic hot halo at the highest observed halo densities. The modelled moving clouds are used to investigate the possibility of star formation within CHVCs. Regarding their multi-phase structure and high column densities, CHVCs are likely candidates for star formation, because they fulfill the necessary conditions. For the first time in scientific literature, a physical cause is revealed that confirms the observed lack of stars within CHVCs. It is validated within the simulations that saturated thermal conduction counteracts a time-dependent ram pressure. The ram pressure decreases by time, because CHVCs are decelerated, while thermal conduction continuously transfers thermal energy at a constant rate into the clouds. Even though a compression of cloud regions is initiated, a gravitational collapse is prevented by a maintained thermal pressure.

The research question addressed in this work is thus answered as follows:

Star-formation is not likely in HVCs even under extreme condition of the Galactic hot halo. A continuous heat input from the ambient hot halo by thermal conduction effectively prevents from a gravitational collapse.

This chapter comprises ideas for future research in order to extend the present work. Some promising questions arose during the preparation of this thesis.

New applications for already existing simulation data

Transition zones as test characteristics for the Galactic corona

Fox et al. (2004) observe ratios of highly ionized elements in Complex C, which are most likely produced in conductive interfaces (i.e., transition zones) and turbulent mixing layers. Such a transition zone thus provides an indirect evidence for a hot and dilute Galactic corona interacting with Complex C. Inspired by their work, an idea was developed during this thesis: the ambient hot gas shapes the spatial extension, ionization and temperature structure of the transition zones, and the condensation rate. Therefore, the characteristics of the transition zones, that are greatly analyzed in Sander & Hensler (2023), may provide a suited testbed to infer physical conditions of the ambient hot halo gas. The detailed simulations of the interior cloud structure may thus help to constrain the conditions of the circumgalactic medium (CGM) around particular high-velocity clouds (HVCs).

Relation between peak column density and cloud speed

Winkel et al. (2011) reveal an interesting result: five of their observed head-tail HVCs are consistent with a positive correlation of v_{GSR} and peak $N(\text{HI})$. A hypothesis is formulated stating that HVCs that experience drag for the longest time have lost more gas through ram pressure than HVCs that just entered the Galactic hot halo. Winkel et al. (2011) cannot prove a causal relationship, but it could be that the IVC feature they observe provides the oldest component among their cloud catalogue. The simulated model clouds in Sander & Hensler (2021) are predestined to be analyzed for such a peak $N(\text{HI}) - v_{\text{GSR}}$ relation. The hypothesis of Winkel et al. (2011) could thus be substantiated if such a relation is observed within simulated model clouds, which are either decelerated or move at constant speeds. The simulation data is immediately available for this analysis and thus forms the scientific basis for a next publication.

Metallicity gradients in head-tail HVCs

The prototypical head-tail cloud HVC 213.5+26.0+111 is identified within the catalogue of Brüns et al. (2000). This particular HVC is unique in the sense that two

quasar lines of sight pass through, which allows observers to probe the gas in front of the head and in the tail. The light from both background quasars passes two distinct regions with very different physical states in terms of column density, gas temperature, radial velocity, and probably metallicity. Accordingly, both the chemical and ionization structure can be studied spatially resolved.

Sander & Hensler (2021) show that the mixing process is not homogeneously distributed over the entire cloud surface, but is rather enhanced at the front side. This leads to a raising gradient in metallicity between head and tail. As the halo gas is expected to be of solar metallicity and above, this gradient would only be present if the cloud's metallicity is clearly sub-solar. Thus, the readily available simulation data of compact high-velocity clouds (CHVCs) with their metallicity gradients from head to tail can help to substantiate a prospective proposal for highly resolved optical absorption spectra of Ca II and Na I.

The observed optical data could be used to verify the heating processes being considered in the simulations. The derived column densities of Ca, together with already existing foreground extinction maps (Burstein & Heiles, 1982; Schlegel et al., 1998), constrain the amount of dust being present in HVC 213.5+26.0+111 (Kondo et al., 2006). In turn, the contributions by both the photoelectric effect and condensation of molecular hydrogen on dust particles to the total heating rate may be evaluated.

From the two quasar lines of sight the temperatures in the head as well as in the more diffuse outer gas of HVC 213.5+26.0+111 are derived. Taking ionization equilibrium as a basis we can derive the degree of ionization in the head-tail HVC. We then compare with the ionization structure thus obtained from our simulations and from earlier studies (Gnat et al., 2010). This will for the first time allow to disentangle whether thermal conduction or photoionization is the ruling process in governing the ionization structure of HVC 213.5+26.0+111, and to determine the extent by which thermal conduction leads to a metallicity enhancement of the clouds.

During the preparation of this thesis, such a proposal has been submitted to the Very Large Telescope (VLT) exploiting the UVES instrument by Jürgen Kerp, Gerhard Hensler, Benjamin Winkel, and Bastian Sander. The proposal was rejected as the probability for meeting the required seeing conditions have been rated as very small. However, the idea of coupling the proposed observations with appropriately simulated HVCs have been evaluated favorably. With regard to the extensively simulated model CHVCs within this thesis, it is thus advisable to re-submit this promising observation proposal.

Code extension with more sophisticated physics

Galactic UV background

The model of Bland-Hawthorn & Maloney (1999a,b, 2002) yields the fraction of ionizing ultraviolet (UV) photons escaping the Galactic disk. By this, the UV radiation field is obtained at every galacto-centric distance within the Galactic halo. This model is applied to HVCs to indirectly infer their distances from observed $H\alpha$ intensities (see appendix C). By implementing the model into FLASH the time-dependent contribution of photoionization to the heating of HVCs (section 6.2.6) can be taken into account. This would yield $H\alpha$ intensities for simulated clouds, which in turn could be used to put distance constraints to observed HVCs.

Non-equilibrium ionization

Within this thesis an equilibrium ionization is applied to the phases (section 6.2.1). The phases are thus assumed to be in thermodynamic equilibrium expressed via the rate equations (6.16). As described in Sander & Hensler (2023), the ratio of cloud radius to Field length in the hot ambient gas describes whether the thermal state of the cloud is dominated by heat conduction or radiative cooling. Gnat et al. (2010) report that non-equilibrium extends the spatial ranges at which thermal conduction governs flow dynamics. The equilibrium approach applied within this thesis is thus a lower bound for the spatial range of domination by thermal conduction. It is therefore worth to study the transition zones of model clouds in a non-equilibrium thermal state to analyze resulting $H\alpha$ intensities and produced ratios of highly ionized elements.

Galactic and local magnetic fields

Several simulations indicate an effect of Galactic magnetic fields on cloud evolution when they pass through the hot halo (cf. Grønnow et al., 2022; Jung et al., 2023). Both Galactic and local magnetic fields are not addressed within this thesis. The knowledge of the latter is currently very limited (see section 3.3.2). To extend the model setups with magnetic fields is interesting regarding two means: *(i)* the analytical estimate for the effective heat flux (Sander & Hensler, 2019) is reviewed in presence of a time-dependent magnetic field, and *(ii)* an additional magnetic pressure is exerted to the cloud thus influencing the formation of gravitationally unstable cloud regions.

Interaction of the Magellanic Stream with the Galactic hot halo

The physical processes that result in the variety of substructures in the Leading Arm of the Magellanic Stream (MS) can be explored by intertwining observations and simulations. First, observational data are extracted from the latest releases of the Parkes

Galactic All-Sky Survey (GASS) and the Effelsberg-Bonn H I Survey (EBHIS), respectively. The GASS/EBHIS data are combined with radio interferometric observations, e.g., performed with the Australia Telescope Compact Array (ATCA) or Westerbork Synthesis Radio Telescope (WRST). Therefore, a data base for the entire MS is generated providing spatial information down to the sub-parsec scale being invaluable for numerical studies.

Seond, the simulations as performed with FLASH disentangle the physical processes that lead to the various substructures in the Leading Arm, namely, self-gravity versus ram pressure and tidal forces, heating versus cooling, hydrodynamic instabilities, and possibly magnetic fields.

Appendices

A HYDRODYNAMIC INSTABILITIES

A.1 KELVIN-HELMHOLTZ INSTABILITY

If two gases of different densities ϱ_{high} and ϱ_{low} move relative to one another along a common interface and the surface tension of the dilute medium is too low, tangential perturbations grow and result in a Kelvin-Helmholtz instability (KHI). The amplified perturbations have wave numbers

$$k = \frac{2\pi}{\lambda} > k_{\text{min}} \approx \frac{g\varrho_{\text{high}}}{\varrho_{\text{low}}v_{\text{rel}}^2}, \text{ if } \varrho_{\text{high}} \gg \varrho_{\text{low}}, \quad (\text{A.1})$$

where λ is the wavelength of the perturbation, g is the gravitational acceleration, and v_{rel} is the relative tangential speed. Relation (A.1) assumes discontinuities in both density and velocity between cloud and ambient gas. Vieser & Hensler (2007a) and Sander & Hensler (2023) show numerically how the dissipative effect of thermal conduction broadens initial discontinuities over a finite domain (*transition zone*) hence converting them into very steep gradients. As has been worked out by Vieser & Hensler (2007b) a finite region can be defined in k -space, which consists of all modes leading to an amplification of KHI. A cloud is stable against KHI for all modes outside of that instability strip, whose size inversely scales with the extension of the transition zone. By this, KHI is inhibited by thermal conduction.

A.2 RAYLEIGH-TAYLOR INSTABILITY

A Rayleigh-Taylor instability (RTI) may develop if a dense gas is accelerated against a dilute gas. This situation may occur at the front side of a HVC when it moves towards the Galactic gravitational potential thus being accelerated in direction of the tenuous Galactic halo. The perturbed outskirts of the HVC start to form typical finger-like structures of the denser medium pointing into the dilute medium. The growth rate of the dominant perturbation mode is given by (Chandrasekhar, 1961)

$$\omega^2 = (kv)^2 = ak \frac{\varrho_{\text{low}} - \varrho_{\text{high}}}{\varrho_{\text{low}} + \varrho_{\text{high}}}, \quad (\text{A.2})$$

with a , $k = 2\pi/\lambda$, λ , and v being the acceleration, wave number, wavelength, and phase velocity of the dominant perturbation mode, respectively. According to Kull

(1991) the e-folding time, i.e., the timescale at which the spatial extension of a certain mode increases by a factor e , reads

$$\tau_e = \sqrt{\frac{\lambda}{2\pi a}}. \quad (\text{A.3})$$

The inverse of (A.3) denotes the growth rate τ_e^{-1} . Vieser (2001) extensively describes the formation of KHI and RTI.

B.1 THE LOCAL STANDARD OF REST

The reference frame, which is in rest relative to the sun, is called the *Local Standard of Rest* (LSR). It is a virtual point close to the sun that revolves in a perfect circular orbit around the galactic center. This orbit is only determined by gravitational forces acting on the fictitious point LSR.

Let l be the galactic longitude and b be the galactic latitude. Then, for each direction (l, b) the velocity of an object relative to the LSR is given by (van Woerden et al., 2004)

$$v_{\text{LSR}} = \left(\frac{R_{\odot}}{R} v(R) - v(R_{\odot}) \right) \sin l \cos b, \quad (\text{B.1})$$

where $R_{\odot} = 8.5 \text{ kpc}$ is the galacto-centric distance of the sun, $v(R_{\odot}) = 220 \text{ km s}^{-1}$ is the corresponding circular speed (e.g., Honma & Sofue, 1997), R is the galacto-centric radius of a point in the sightline, and $v(R)$ is the rotation velocity at distance R . Basically, $v(R)$ is that of a solid body for $R < 0.5 \text{ kpc}$ and $v(R) \approx v(R_{\odot})$ for all radii greater $R_{\odot}$.

B.2 THE GALACTIC STANDARD OF REST

The reference frame, which is in rest relative to the galactic centre, is called the *Galactic Standard of Rest* (GSR). If v_{LSR} is expressed relative to the GSR, it must be corrected for the component of galactic rotation:

$$v_{\text{GSR}} = v_{\text{LSR}} + v(R_{\odot}) \sin l \cos b. \quad (\text{B.2})$$

Hence, the velocity (B.2) is relative to the sun in a reference frame in which the Milky Way rotates.

B.3 THE DEVIATION VELOCITY

The deviation velocity was first introduced by Wakker (1991) to account for uncommon high-velocity gas irrespective of the galactic latitude. It is defined as

$$v_{\text{dev}} = \begin{cases} v_{\text{LSR}} - v_{\text{min}}(l, b) & , v_{\text{LSR}} < 0 \\ v_{\text{LSR}} - v_{\text{max}}(l, b) & , v_{\text{LSR}} > 0 \end{cases} , \quad (\text{B.3})$$

where $[v_{\text{min}}, v_{\text{max}}]$ is the range of velocities allowed by Galactic rotation for a given (l, b) direction. For clouds following Galactic rotation, $v_{\text{dev}} = 0$ (see also Begum et al., 2010). A usual operational threshold for defining gas to be at uncommonly high velocities is given by $|v_{\text{dev}}| > 60 \text{ km s}^{-1}$ and $|v_{\text{LSR}}| > 90 \text{ km s}^{-1}$. But the definition (B.1) of v_{LSR} indeed includes gas in normal rotation at low latitudes while at the same time many clouds at high latitudes, that move with a velocity being incompatible with differential Galactic rotation, are excluded. Between $l = 90^\circ$ and $l = 180^\circ$ the gas in the outer Galaxy appears at intermediate negative velocities because of the effects of differential rotation (Tamanaha, 1997).

C DISTANCE DETERMINATIONS

C.1 DERIVED QUANTITIES

Determination of distances to high-velocity clouds (HVCs) is of prime interest for understanding their physical states, because several fundamental properties are derived from the distance. The first key value is the size D_{HVC} of the cloud. By determining the

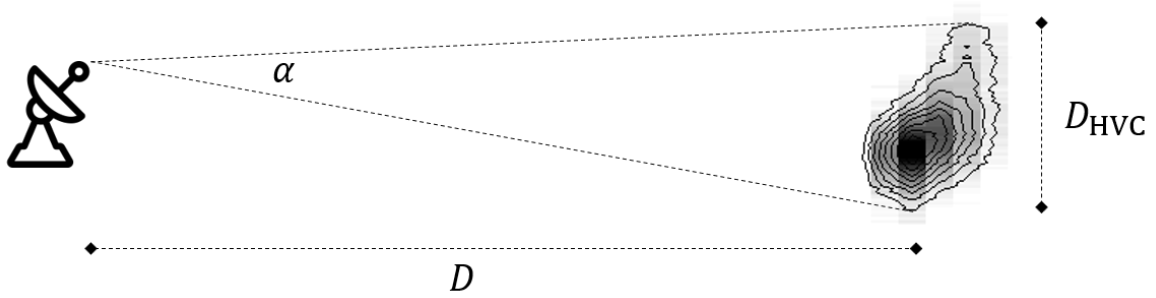


Figure C.1: Schematic diagram for deriving the size D_{HVC} of a high-velocity cloud, which appears under an angular size of α , from measured distance D . [HVC plot taken from Brüns et al. (2000), icon made by Rukanicon from www.flaticon.com]

opening angle α , under which the HVC appears in a telescope, a simple trigonometric method (Figure C.1) yields

$$R_{\text{HVC}} = 0.5D_{\text{HVC}} \approx \frac{\alpha}{2}D, \quad (\text{C.1})$$

for the radius of the cloud if D is the distance to the cloud and $\tan \alpha \approx \alpha$ for small opening angles. Hence, the scaling relation for the linear width reads

$$\boxed{D_{\text{HVC}} \propto D}. \quad (\text{C.2})$$

If the column density $N(\text{H I})$ is measured by a beam with opening angle α , then the particle density n_{HVC} in the cloud is derived from

$$n_{\text{HVC}} = \frac{N(\text{H I})}{D\alpha}. \quad (\text{C.3})$$

So,

$$\boxed{n_{\text{HVC}} \propto \varrho_{\text{HVC}} \propto D^{-1}}, \quad (\text{C.4})$$

where ρ_{HVC} is the matter density in the cloud. The volume V of a HVC scales with R_{HVC}^3 . Hence, the mass $M = \rho V$ scales like

$$\boxed{M_{\text{HVC}} \propto D^2}. \quad (\text{C.5})$$

Accordingly, the pressure $P_{\text{HVC}} = n_{\text{HVC}} k T_{\text{HVC}}$ scales like

$$\boxed{P_{\text{HVC}} \propto D^{-1}}. \quad (\text{C.6})$$

Even some typical timescales are derived from distances. Wakker & van Woerden (1997, see their Section 2.8) estimate the time a HVC needs to reach the Galactic plane by

$$t_{z=0} = \frac{z\text{-height}}{z\text{-velocity}} \sim \frac{D \sin b}{v_{\text{dev}} \sin b} \propto D. \quad (\text{C.7})$$

Distance determination provides also access to timescales relating to internal dynamics. Thus, the lifetimes of substructures at different length scales are implied. First, the timescale a dense core needs to shift considerably relative to each other core is given by

$$t_{\text{shift}} = \frac{\text{linear width}}{\text{velocity dispersion}} \sim \frac{\alpha D}{\Delta v} \propto D. \quad (\text{C.8})$$

Second, the time a core needs to cross over its own extension is estimated by

$$t_{\text{cross}} = \frac{R_{\text{core}}}{v_{\text{dev}}} \sim \frac{\Theta D}{v_{\text{dev}}} \propto D, \quad (\text{C.9})$$

where Θ and R_{core} are the half-brightness radius and the spatial radius of a core, respectively. Third, the timescale for substructures of a core to expand and dissolve is evaluated by

$$t_{\text{expand}} = \frac{R_{\text{sub}}}{\gamma} \sim \frac{\theta D}{\gamma} \propto D, \quad (\text{C.10})$$

with the half-brightness radius θ of such a substructure and the line width γ (full width at half maximum (FWHM)). By subsuming equations (C.7) to (C.10) one finds the general scaling relation

$$\boxed{\text{HVC timescale} \propto D}. \quad (\text{C.11})$$

That is why the distance is so essential to obtain. There are several ways of deriving the distance, which can roughly be separated into *direct* and *indirect* methods, which are described below.

C.2 DIRECT METHODS

With direct methods the distance of HVCs is directly inferred from observational results.

C.2.1 Absorption-line spectroscopy

Because no stars are assigned to HVCs yet, exact distance determinations are not feasible by absorption-line spectroscopy. Alternatively, at least one bright emission source is needed in front of and at least another one behind the HVC. The distances to these emission sources need to be known. Beers (1999) emphasizes the importance of field horizontal branch stars in determining the distances of HVCs, because these stars are numerous in the Galactic halo, luminous, and have an accuracy in distance determination of $\sim 10\%$. For an *upper* distance limit to be established it is necessary only to see absorption towards a background source of known distance. If 21-cm H I emission is observed along the same sightline at a similar velocity, then it is reasonable to assume that the emission and absorption are at least partially related. To define a *lower* distance limit it is necessary to see the IVC/HVC in 21-cm emission and *not* see it in absorption towards the selected probe star down to a level consistent with the H I emission (Ryans et al., 1997b). By this, a distance bracket can be derived constraining the HVC distance. The distance derived in this way is also called *spectroscopic distance*.

Typical lines used in this method are the optical Ca II H and K, and Na II doublet absorption. A huge number of halo stars is indexed by the Sloan Digital Sky Survey. These stars have been used to constrain the distances of almost all HVC complexes (e.g. van Woerden et al., 1999c; Putman et al., 2012). van Woerden et al. (2004) point out that the absorption-line method yields the most unique results. The distances of many isolated and compact HVCs cannot be determined by this technique due to a lack of suitable halo stars.

C.3 INDIRECT METHODS

Indirect methods are on the one hand applicable to a huge number of HVCs, which are not accessible by direct methods. On the other hand, indirect methods rely on theoretical models and require one or more key assumptions that are not unconditionally applicable to all HVCs (Ferrara & Field, 1994; van Woerden et al., 2004).

C.3.1 Emission of $H\alpha$

The de-excitation of excited hydrogen H II from energy level $n = 3 \rightarrow 2$ produces the optical emission line of $H\alpha$. In opposite to other emission lines the $H\alpha$ line is detected at many HVCs, because the clouds mainly consist of hydrogen.

By assuming that the brightness in $H\alpha$ is mainly due to an ionizing ultraviolet (UV) radiation field, the distance is derived from a model of the Galactic UV background (see Bland-Hawthorn & Maloney, 1999a,b, 2002, who have applied their detailed model to HVCs). If the clouds are subject to ionizing radiation they absorb photons with energies > 13.6 eV. Each ionization event produces on average 0.46 $H\alpha$ photons during the follow-up recombination (Pengelly, 1964; Martin, 1988). Thus, if the $H\alpha$ emission is mainly due to photoionization of the cloud, then the intensity is directly related to the incident Lyman continuum flux that depends on the location of the cloud (Tufté et al., 1998) and HVC distances can be constrained (Bland-Hawthorn et al., 1998; Putman et al., 2003a). Barger et al. (2012) used the photoionization model of the Milky Way (Bland-Hawthorn & Maloney, 1999a) and the model of the extragalactic background radiation (Lyman continuum model, Haardt & Madau, 2001) to infer distances based on $H\alpha$ emission.

Advantages of this indirect method are its application to many HVCs and that distances derived from $H\alpha$ emission tend to agree with directly measured distances (Putman et al., 2012). However, the Magellanic Stream is much brighter in $H\alpha$ as predicted by this model. Moreover, different $H\alpha$ intensities are measured at the same cloud, so the internal structure might further complicate the distance determination.

C.3.2 Emission of X-rays

Kerp et al. (1996a, 1998) correlate the *ROSAT* 1/4 keV all-sky survey with H I 21-cm line emission from the Leiden/Dwingeloo northern-sky survey for Complex C in the region $34 \leq l/^\circ \leq 86$ and $33 \leq b/^\circ \leq 79$. They find that large parts of Complex C are associated with soft X-ray signatures from which they conclude an impact onto the galactic disk. Their results put the HVCs of Complex C into a galactic height of $z = 1.5 \dots 2.5$ kpc. In a follow-up study Kerp et al. (1999) extend the observations to another region of Complex C at lower b and found large portions of Complex C to be spatially associated with an excess of soft X-ray emission. The excess in X-rays corresponds to a plasma temperature of $T \sim 10^{6.2}$ K.

C.3.3 Deep observations of neutral hydrogen

Deep H I observations of systems similar to the Milky Way or Local Group provide clues on distances to HVCs in our galaxy if a similar spatial distribution of HVCs is assumed. By this method, distance constraints are found especially for many compact high-velocity clouds (CHVCs) (Pisano et al., 2007; Westmeier et al., 2007).

C.3.4 The kinematic distance

Tamanaha (1997) derives kinematic distances D_{kin} to H I clouds by comparing the observed Local Standard of Rest (LSR) velocity at their longitude in the Galaxy with the velocity expected from a model galaxy with circular orbits and a flat rotation curve. It follows

$$D_{\text{kin}} \approx R_0 \left[\left(\frac{\Theta}{\Theta_0} \right) \left(1 + \frac{dv_{\text{LSR}}/dl + \Pi \sin l}{\Theta_0 \cos l} \right)^{-1} - 1 \right], \quad (\text{C.12})$$

where Θ and Π are the circular velocity component of the cloud and its radial velocity component parallel to the axis between Galactic centre and the Sun, respectively, l is the Galactic longitude of the cloud, and dv_{LSR}/dl its velocity gradient along the longitude. R_0 and Θ_0 are the galacto-centric distance of the Sun and its circular velocity component. Morras et al. (1998) obtain distance to Complex H on kinematic considerations and Wenger et al. (2018) derive kinematic distances based on Galactic rotation curves (Brand & Blitz, 1993; Reid et al., 2009) and updated values of the solar motion. Lockman et al. (2008) determine three different kinematic distances for the Smith Cloud.

Peek et al. (2007) derive a parameter Ξ with mean value $\langle \Xi \rangle = 2.7 \text{ cm}^{-3} \text{ pc}$, which by

$$D_{\text{kin}}[\text{pc}] = \frac{\Xi}{n_H} \quad (\text{C.13})$$

leads to distances consistent with the H α method, where n_H is the particle density in the halo at the location of the respective HVCs.

C.3.5 Thermal pressure equilibrium

By assuming thermal pressure equilibrium between HVCs and the Galactic halo a distance is obtained by combining cloud observations with a model of the stratification

of the halo. Namely,

$$\begin{aligned} P_{\text{HVC}} &= P_{\text{halo}} \\ \iff n_{\text{HVC}} T_{\text{HVC}} + P_{\text{turbulence}} &= n_{\text{halo}} T_{\text{halo}}. \end{aligned} \quad (\text{C.14})$$

The left-hand side of equation (C.14) is observationally determined by the column density (C.3) and by the $\text{FWHM} \propto \sqrt{T_{\text{HVC}} + P_{\text{turbulence}}/n_{\text{HVC}}}$ (van Woerden et al., 2004). The right-hand side must be confidentially modelled. Espresate et al. (2002) and Stanimirović et al. (2002) use this assumption to calibrate the density and temperature profiles of the Galactic halo with the known distances of Complex M and the core MS VI in the Magellanic Stream, respectively. In turn, unknown distances for other HVCs can be determined.

C.3.6 Ram-pressure equilibrium

Benjamin & Danly (1997) show that distances to infalling HVCs are obtained by balancing the ram pressure exerted to the cloud and the halo density. They assume that HVCs move through a gaseous halo with a height-dependent terminal velocity depending on local values of halo density and gravitational acceleration, the vertical scale-height, and on velocities and total column densities $N(\text{H I}) + N(\text{H II})$ of the clouds. The terminal velocity v_{T} can be used to infer a distance. The analytical model is given by equation (3) in Benjamin & Danly (1997) and some examples are illustrated in their figure 2. They accurately predict the distances to Complex M and some intermediate-velocity clouds (IVCs) by this model, which is found to be valid for $N(\text{H I}) \lesssim 10^{19} \text{ cm}^{-2}$.

C.3.7 Virial theorem

Blitz et al. (1999b) propose that HVCs most likely belong to the Local Group with a median distance of 1 Mpc and H I masses $\sim 10^7 M_{\odot}$. They assume the clouds to be in virial equilibrium, which requires a substantial amount of dark matter in HVCs in order to be stable and self-gravitating. By applying the virial theorem, Blitz et al. (1999b) derive a distance depending on the square of velocity dispersion σ^2 in the clouds (van Woerden et al., 2004). However, this model is ruled out for several reasons, e.g. a discrepancy to the H I mass function of field galaxies (Zwaan et al., 1997), an inconsistency with observed intra-group H I clouds or QSO absorption-line statistics (Zwaan, 2001a,b), or based on the fact that clouds embedded in a dark-matter halo would never develop head-tail structures (e.g. Plöckinger & Hensler, 2012).

D.1 THE 21-CM LINE OF NEUTRAL HYDROGEN

The 21-cm line results from a forbidden transition (magnetic dipole radiation) in the ground state $1s^2$ in atomic hydrogen. The two hyperfine structure levels at state $1s^2$ have total spins $S = 1$ (spins of nucleus and electron are parallel) and $S = 0$ (spins are anti-parallel). If the spin direction of the electron flips, a photon with energy 6×10^{-6} eV is released corresponding to an emission line in the radio domain with

$$\lambda_{21 \text{ cm}} = 21.1 \text{ cm} \quad \Leftrightarrow \quad \nu_{21 \text{ cm}} = 1420.4 \text{ MHz.} \quad (\text{D.1})$$

The transition probability (Einstein coefficient) is very low with $A = 2.87 \times 10^{-15} \text{ s}^{-1}$, which is equivalent to a lifetime of the upper hyperfine structure level of $A^{-1} = 11 \text{ Myr}$. Regarding the ubiquity of atomic hydrogen in the Universe, the 21-cm line is highly important to study large-scale structures and the dynamics of interstellar gas (e.g., Unsöld & Baschek, 2002; Tielens, 2005).

D.2 THERMAL LINE BROADENING

The thermal motion of particles can be converted into a line broadening. For only thermal broadening the relation between the Doppler temperature of a gas and an observed line width (Ben Bekhti et al., 2006; Saul et al., 2012; Faridani et al., 2014) is given by

$$T_D[\text{K}] = \frac{m_H (\Delta v_{\text{FWHM}})^2}{8k \ln 2} \approx 21.8 (\Delta v_{\text{FWHM}} [\text{km s}^{-1}])^2. \quad (\text{D.2})$$

The Doppler temperature assumes a pure Maxwellian velocity distribution of the gas and thus provides an upper limit for its kinetic temperature, i.e. $T_{\text{kin}} \leq T_D$. The kinetic temperature is presumably lower, because turbulent motions or unresolved small-scale structures may contribute to the line width.

D.3 COMMONLY USED SPECTRAL LINES

Several spectral lines help to explore the multi-phase nature of the circumgalactic medium (CGM). Table D.1 lists the transitions commonly used. An extensive compilation of resonant absorption lines is provided by Morton (1991, 2000, 2003).

Table D.1: Transitions at wavelengths λ that are commonly used to study the multi-phase CGM.

ionization mechanism	Ion	λ / Å	sample reference
photo-ionization (“low ions”)	Ly α	1216	Fox et al. (2005)
	H α	6565	Tufte et al. (1998)
	Na I-D	5890	Welsh et al. (1996)
	O I	6300, 1302	Barger et al. (2012)
	Mg I	2852	Wakker et al. (1998)
	Mg II	2796, 2803	Wakker et al. (1996)
	Si II	6716	Barger et al. (2012)
	Si II	1260, 1526	Danly et al. (1993)
	S II	6716	Barger et al. (2012)
	S II	1250, 1253, 1259	Gibson et al. (2001)
	N II	6584	Barger et al. (2012)
	C II	1334	Wakker et al. (1998)
	Ca II-K	3933	Welsh et al. (1996)
photo-ionization and collisions (“intermediate ions”)	Si III	1206.5	Ganguly et al. (2005)
	Si IV	1393.76, 1402.77	
	C III	977.02	
	C IV	1548.20, 1550.78	
collisions, thermal conduction, and mixing layers (“high ions”)	O VI	1031.93, 1037.62	Sembach et al. (2000)
	N V	1238.82, 1242.80	Ganguly et al. (2005)
	Ne VIII	770.41	Fox et al. (2005)

E COLUMN DENSITY

From 21-cm emission a H I brightness temperature T_b averaged over the main beam is obtained. If the 21-cm emission is sufficiently faint, the peak temperature of the spectral lines is usually low compared to the likely spin temperature¹ T_s of the H I gas (Martin et al., 2015). In this case, the 21-cm emission is optically thin and a conversion to H I column density $N(\text{H I})$ can be derived from the directly observed feature $W(\text{H I})$, which is the line integral of the H I emission spectrum over velocity. So,

$$W(\text{H I}) \equiv N(\text{H I}) [\text{cm}^{-2}] / C = \int_{v_1}^{v_2} T_b(v) dv \quad (\text{E.1})$$

with conversion factor $C = 1.823 \times 10^{18} \text{ cm}^{-2} (\text{K km s}^{-1})^{-1}$ (see equation 3 in Dickey & Lockman (1990) and equation 1 in Martin et al. (2015)). As Martin et al. (2015) mention, the integral (E.1) is evaluated separately for each H I emission spectrum, such that a column-density map is generated for each given velocity range $[v_1, v_2]$. Note that the 21-cm spin temperature T_s is identical to the kinetic gas temperature T_{kin} only if neutral hydrogen is in thermal equilibrium and level populations are described by Boltzmann statistics. There exist astrophysical environments where the 21-cm line is pushed away from thermal equilibrium especially by processes involving Ly α radiation (e.g. $T_s \gg T_{\text{kin}}$ in damped Ly α absorbers, Shaw et al., 2017).

¹The spin temperature of an H I region is the temperature that would produce the observed ratio of parallel to antiparallel spins if the H I gas were in thermal equilibrium (Ridpath, 2018).

F.1 STABILITY OF A NUMERICAL SCHEME

The stability of numerical schemes integrating non-linear, coupled partial differential fluid flow equations (e.g., the Euler equations as introduced in section 6.2.2) are commonly investigated by applying von Neumanns Fourier analysis (Crank et al., 1947). The method aims at introducing an error into a numerical solution and tracks its amplification in time. If the error (or, perturbation) remains bounded while evolving the solution, the numerical integration scheme is stable and unstable otherwise. The scheme is unconditionally stable if the perturbation does not grow for an arbitrarily large integration timestep, which is true for fully implicit methods. A numerical scheme is conditionally stable, if the integration timestep is constrained by any upper threshold. This holds for asynchronous methods like leapfrog, and semi-implicit methods like Crank-Nicolson. A numerical scheme is unconditionally unstable, if for any integration timestep the integration scheme produces spurious numerical solutions. An example is given by explicit methods like direct forward or backward Euler.

Below the stability condition for the generalization of the discretization (6.62) is explicitly given for one dimension. So,

$$T_j^{(n+1)} - T_j^{(n)} = \frac{\Delta t}{(\Delta x)^2} \left[(1 - \alpha) \left(T_{j+1}^{(n)} - 2T_j^{(n)} + T_{j-1}^{(n)} \right) + \alpha \left(T_{j+1}^{(n+1)} - 2T_j^{(n+1)} + T_{j-1}^{(n+1)} \right) \right] \quad (\text{F.1})$$

with $0 \leq \alpha \leq 1$. The value of α determines the integration scheme (F.1) to be explicit ($\alpha = 0$), semi-implicit ($\alpha = 0.5$, Crank-Nicolson type), or fully implicit ($\alpha = 1$). Let the wave

$$T_j^{(n)}(k) = \xi^n e^{ikj\Delta x} \quad (\text{F.2})$$

with amplitude ξ^n and phase $\exp\{ikj\Delta x\}$ be a solution to (F.1) at timestep n . At timestep $n + 1$ the wave (F.2) is evolved into

$$T_j^{(n+1)}(k) = \xi^{n+1} e^{ikj\Delta x} = a\xi^n e^{ikj\Delta x} = aT_j^{(n)}(k) \quad (\text{F.3})$$

with the amplification factor a for the perturbation ξ^n . A condition for stability of the numerical scheme (F.1) can be formulated for arbitrary α . By substituting (F.2) and

(F.3) into (F.1) one obtains

$$\begin{aligned}
(a-1)\xi^n e^{ikj\Delta x} &= \frac{\Delta t}{(\Delta x)^2} [(1-\alpha)(\xi^n e^{ik(j+1)\Delta x} - 2\xi^n e^{ikj\Delta x} + \xi^n e^{ik(j-1)\Delta x}) \\
&\quad + \alpha(\xi^{n+1} e^{ik(j+1)\Delta x} - 2\xi^{n+1} e^{ikj\Delta x} + \xi^{n+1} e^{ik(j-1)\Delta x})] \\
\iff a-1 &= \frac{\Delta t}{(\Delta x)^2} [(1-\alpha)(e^{ik\Delta x} - 2 + e^{-ik\Delta x}) + \alpha(\xi e^{ik\Delta x} - 2\xi + \xi e^{-ik\Delta x})] \\
\iff a &= 1 + \frac{2\Delta t}{(\Delta x)^2} \{(1-\alpha)[\cos(k\Delta x) - 1] + \alpha\xi[\cos(k\Delta x) - 1]\} \\
&= \frac{1 + \frac{2\Delta t}{(\Delta x)^2}(1-\alpha)[\cos(k\Delta x) - 1]}{1 - \frac{2\Delta t}{(\Delta x)^2}\alpha[\cos(k\Delta x) - 1]}. \tag{F.4}
\end{aligned}$$

The numerical scheme (F.1) is stable if and only if the amplification factor a in (F.4) remains bounded upon integration. Thus,

$$\text{stability} \iff |a| \leq 1 \iff -1 \leq \frac{1 + \frac{2\Delta t}{(\Delta x)^2}(1-\alpha)[\cos(k\Delta x) - 1]}{1 - \frac{2\Delta t}{(\Delta x)^2}\alpha[\cos(k\Delta x) - 1]} \leq +1. \tag{F.5}$$

Checking for the maximum of the cosine function yields

$$\max \{\cos(k\Delta x)\} = 1 \implies \frac{1 + \frac{2\Delta t}{(\Delta x)^2}(1-\alpha)[\max \{\cos(k\Delta x)\} - 1]}{1 - \frac{2\Delta t}{(\Delta x)^2}\alpha[\max \{\cos(k\Delta x)\} - 1]} = 1 \quad \forall \frac{2\Delta t}{(\Delta x)^2}. \tag{F.6}$$

Likewise, checking for the minimum of the cosine function yields

$$\min \{\cos(k\Delta x)\} = -1 \implies \frac{1 - 2\frac{2\Delta t}{(\Delta x)^2}(1-\alpha)}{1 + 2\frac{2\Delta t}{(\Delta x)^2}\alpha}. \tag{F.7}$$

The above expression (F.7) must be ≥ -1 to satisfy the stability condition (F.5). So, the condition for stability translates into a constraint for the integration timestep:

$$\Delta t \leq \frac{1}{1-2\alpha} \frac{(\Delta x)^2}{2} \begin{cases} = \frac{(\Delta x)^2}{2} & , \alpha = 0 \text{ (fully explicit)} \\ \rightarrow \infty & , \alpha \rightarrow 0.5 \text{ (Crank-Nicolson)} \\ < 0 & , \alpha > 0.5 \text{ (fully implicit)} \end{cases}. \tag{F.8}$$

Fully implicit methods are unconditionally stable.

F.2 ACCURACY OF A NUMERICAL SCHEME

The accuracy of a finite approximation can be checked by a Taylor expansion. A scheme is said to be accurate of n -th order if the n -th derivative is the highest discretized derivative considered in the scheme. The accuracy is separately determined for the spatial discretization Δx and the time discretization Δt .

The Crank-Nicolson scheme assumes an additional (virtual) timestep at time $n+1/2$, for which the time derivative is centrally discretized. The value of temperature does not exist at the virtual timestep $n+1/2$ so a discretization of higher order is obtained without including additional (real) timesteps. The left-hand side of equation (6.58) thus reads in discretized form

$$\left. \frac{\partial T}{\partial t} \right|_{n+1/2} = \frac{T^{(n+1)} - T^{(n)}}{2 \times \Delta t/2} + \mathcal{O}((\Delta t)^2). \quad (\text{F.9})$$

The right-hand side of equation (6.58) is discretized by using a mean value of temperature from old and new timestep

$$\kappa \nabla(\nabla T) = \kappa \left(\frac{T_{i-1}^{(n+1)} - 2T_i^{(n+1)} + T_{i+1}^{(n+1)}}{(\Delta x)^2} + \frac{T_{i-1}^{(n)} - 2T_i^{(n)} + T_{i+1}^{(n)}}{(\Delta x)^2} \right) + \mathcal{O}((\Delta x)^2). \quad (\text{F.10})$$

Hence, the Crank-Nicolson integration scheme is accurate to the second order in both time and space.

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