THE L1495-B218 FILAMENTS IN TAURUS SEEN IN NH$_3$ & CCS AND DYNAMICAL STABILITY OF FILAMENTS AND DENSE CORES

by

Youngmin Seo

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As members of the Dissertation Committee, we certify that we have read the dissertation prepared by Youngmin Seo entitled The L1495-B218 Filaments in Taurus Seen in NH$_3$ & CCS and Dynamical Stability of Filaments and Dense Cores and recommend that it be accepted as fulfilling the dissertation requirement for the Degree of Doctor of Philosophy.

Date: 22 July 2016

Yancy L. Shirley

Date: 22 July 2016

Christopher Walker

Date: 22 July 2016

John Bieging

Date: 22 July 2016

Kaitlin Kratter

Date: 22 July 2016

Paul F. Goldsmith

Date: 22 July 2016

Joan Najita

Final approval and acceptance of this dissertation is contingent upon the candidate’s submission of the final copies of the dissertation to the Graduate College.
I hereby certify that I have read this dissertation prepared under my direction and recommend that it be accepted as fulfilling the dissertation requirement.

Date: 22 July 2016

Dissertation Director: Yancy L. Shirley
STATEMENT BY AUTHOR

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SIGNED: Youngmin Seo
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DEDICATION

Dedicated to Mi Young Son and Yoon Seo
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We present deep NH$_3$ and CCS maps of L1495-B218 filaments and the dense cores embedded within the filaments in Taurus. The L1495-B218 filaments form an interconnected, nearby, large complex extending over 8 pc. We observed the filaments in NH$_3$ (1,1) & (2,2), CCS $N_J = 1_2-0_1$, and HC$_7$N $J = 21-20$ with spectral resolution of 0.038 km/s and spatial resolution of 31″. The CSAR algorithm, which is a hybrid of seeded-watershed and binary dendrogram algorithm, identifies 39 leaves and 16 branches in NH$_3$ (1,1). Applying a virial analysis for the 39 NH$_3$ leaves, we find only 9 out of 39 leaves are gravitationally bound, and 12 out of 30 gravitationally unbound leaves are pressure-confined. Our analysis suggests that a dense core may form as a pressure-confined structure, evolve to a gravitationally bound core, and then undergo collapse to form a protostar. We find that the L1495A, B213E, and B216 regions have strong CCS emission and the B211 and B218 regions have weak CCS emission. Analysis of CCS emission with NH$_3$ (1,1) and dust continuum emission shows that CCS is not a good tracer for starless core evolution. On the other hand, CCS appears to trace recently accreted gas in L1495A and L1521D.

We also present more realistic dynamic stability conditions for dense cores and filaments. In a new analysis of stability conditions we account for converging motions which have been modeled toward starless cores and take the effect of radiation fields. We find that the critical size of a dense core having a homologous converging motion with its peak speed being the sound speed is roughly half of the critical size of the Bonnor-Ebert sphere. We also find the critical mass/line density of a dense core/filament irradiated by radiation to be considerably smaller than that of the Bonnor-Ebert sphere/isothermal cylinder when the radiation pressure is stronger than the central gas pressure of dense core/isothermal cylinder. For regions in the inner Galaxy and near OB associations, the critical mass/line density of a dense
structure may be less than 20% of the critical mass/line density of Bonnor-Ebert sphere/isothermal cylinder.
CHAPTER 1

INTRODUCTION

The star formation process starts with the gravitational collapse of a molecular cloud that results in forming progressively smaller and denser structures such as molecular clumps and cores. If molecular cloud clumps and cores are gravitationally unstable, then they may further collapse to form stars. While this general framework is in place, the details of the star formation process are still poorly understood, particularly how a turbulent, magnetized molecular cloud fragments and forms dense cores, and how dense cores interact with their surroundings and evolve to form protostars. As the details of the star formation process determine the initial mass function of stars, galactic chemical evolution, and the initial conditions of planet formation, understanding the evolution of dense cores from their initial fragmentation from the molecular clouds to the stage of giving a birth to protostar is of fundamental importance to many areas in astronomy and astrophysics. In this thesis, I focus on the details of star forming process in dense filamentary structures in Taurus and the evolution of dense cores within the filamentary structures.

Molecular filaments are ubiquitous structures in molecular clouds throughout our Galaxy. Carbon Monoxide (CO) and dust continuum surveys revealed filamentary structures in nearby (D < 500 pc) molecular clouds such as Orion, Taurus, Ophiuchus, Musca-Camaeleon, Aquila, Serpens, Perseus, Ophiuchus, Cepheus, Polaris, Pipe nebula, etc. (e.g. Johnstone and Bally, 1999; Abergel et al., 1994; Cambrésy, 1999; Hatchell et al., 2005; Goldsmith et al., 2008; André et al., 2010, 2014). More distant and generally more massive “infrared dark clouds” (IRDCs), which are identified by absorption against the Galactic mid-infrared background, also have clear filamentary structures (e.g. Perault et al., 1996; Egan et al., 1998; Peretto and Fuller, 2009). The observed filamentary structures in nearby clouds and IRDCs span the range from low-mass star forming regions to high-mass star forming regions.
Molecular filaments are the most common birth places of young stars because most of dense cores are found to lie within filamentary structures. Polychroni et al. (2013) surveyed dense cores within L1641 of the Orion A cloud and they found 71% of the identified dense cores are within filamentary structures with extinction of $A_v > 5$ mag. Konyves et al. (2015) have also performed a similar study toward the Aquila cloud complex using the Herschel Space Observatory. They found 75% of dense cores lies within the filamentary structures with supercritical line densities (meaning gravitationally unstable) of $>16 \, M_\odot$/pc. In addition, more recent analysis toward L1495 in Taurus (Marsh et al., 2016) shows that all of the gravitationally bound starless cores are embedded within filaments with supercritical mass.

Young stellar associations or clusters seem to be correlated with “hubs”, which are typically dense clumps connecting less dense filamentary structures. Myers (2013) showed multiple examples of hub structures in nearby molecular clouds and IRDCs. He demonstrated that these hub structures are associated with stellar groups or clusters using the census of protostars from Spitzer Space Telescope infrared observations. Among these hub structures, Serpens-South has been studied extensively (Kirk et al., 2013a; Friesen et al., 2013). Serpens-South has a hub with at least 90 young stellar objects associated with the hub. Kinematics studies in HCN, HC$_7$N, C$^{18}$O, and $^{13}$CO toward Serpens-South show that the hub possibly formed through the collision of the two dense filaments and that the hub is also accreting the filaments and surrounding ambient gas. The accreted gas is expected to continue forming stars and form a stellar cluster. Other hubs may have similar kinematics as Serpens-South. However, these studies only focused on the large scale kinematics but not on the star forming processes in dense cores. Further studies are required toward hubs and filaments and their relation to stellar associations or cluster formation.

Many theoretical studies on the formation and evolution of molecular filaments and the formation of dense cores within filaments has been made through simulations (e.g. Bate and Bonnell, 2005; Vázquez-Semadeni et al., 2007; Heitsch et al., 2008; Chen and Ostriker, 2014; Moeckel and Burkert, 2015). Bate and Bonnell (2005)
show that a gravitationally unstable molecular cloud with supersonic turbulence develops filamentary structures as it contracts via gravoturbulence fragmentation. They also demonstrated the development of hubs. Star formation is considerably more active in the hub compared to within less dense filamentary structures connected to the hub. Chen and Ostriker (2014) studied the formation of filaments and dense cores by large scale colliding flows. They demonstrated that filamentary structures and dense cores simultaneously forms by colliding flows and some of dense cores become gravitationally bound. While these theoretical studies provide an overall evolutionary framework, detailed physical properties need to be compared with observations of filaments and dense cores to test the theoretical models.

The final evolutionary stage of forming a star is the collapse of dense cores. The study of the collapse of dense cores initially focused on isolated dense cores to facilitate a simpler analysis. Self-similar solutions and numerical solutions have been often used to investigate the details of the collapse dynamics (e.g. Ebert, 1955; Bonnor, 1956; Larson, 1969; Penston, 1969; Shu, 1977; Hunter, 1977; Foster and Chevalier, 1993; Gómez et al., 2007; Gong and Ostriker, 2009, 2011; Keto et al., 2015). In observations, collapse dynamics has been traced toward bright protostellar cores with infall signatures such as IRAS16293, B335, etc. (Pineda et al., 2012; Evans et al., 2015; Oya et al., 2016). and only a handful of starless cores such as L1544, L1498, L694-2, and L1197 (Ciolek and Basu, 2000; Lee et al., 2004, 2007). These theoretical and observational studies are limited to star formation in a quiescent isolated environment rather than the typical star formation environment in more complicated filamentary structures. The surrounding filaments introduce various perturbations to dense cores. Dense cores in filaments evolve in the gravitational potential of filaments, interact with flows surrounding filaments, and are under effects of radiation pressure and outflows originating from nearby young stellar objects. They are likely to show different dynamical evolution than that of isolated dense cores.

In this thesis, I address the following questions:

• How do filamentary structures evolve in a molecular cloud?
• How do dense cores form and evolve within molecular filaments?

• What are physical and chemical properties of molecular filaments and embedded dense cores?

• Do large scale flows affect the evolution of filaments and the formation of dense cores?

• Do large scale flows affect star forming processes in dense cores?

• How does a dense core in a filament collapse and form a star or multiple stars?

• How does a “hub” evolve and how is it different from the rest of filamentary structures?

To answer these questions, it is necessary to take both observational and theoretical approaches: In an observational approach, we surveyed the L1495-B218 filaments in Taurus and dense cores within the filaments in multiple molecular lines. The Taurus molecular cloud is an excellent testbed to study the star formation processes in filaments and dense cores because it is one of the closest molecular cloud to us, which means we have the highest spatial resolution for a given telescope beam size. The L1495-B218 filaments and surrounding molecular cloud have been previously studied in $^{12}$CO, $^{13}$CO, C$^{18}$O, and dust continuum (Goldsmith et al., 2008; Hacar et al., 2013; Palmeirim et al., 2013). In addition, young stellar objects are fully surveyed using the Spitzer Space Telescope, from which the evolutionary stages of regions in terms of star formation may be determined.

To study details of the dense core evolution, it is important to survey the complete population of dense cores from their early formation stage to protostellar stage and to understand their kinematics. In this thesis, we made high spatial and spectral resolution observations toward the L1495-B218 filaments in CCS and NH$_3$. CCS and NH$_3$ are often suggested to trace young and evolved dense cores (Suzuki et al., 1992), respectively, although some of studies show CCS does not trace young dense cores (e.g. Marka et al., 2012). NH$_3$ is also known for a excellent measure of kinetic
temperature. We observed the L1495-B218 filaments in CCS and NH$_3$ in order to survey a full population of dense cores within the filaments, obtain accurate kinetic temperatures of dense cores, and test which structure CCS traces. We also made deep observations in HC$_7$N, HCN, and HCO$^+$ because HC$_7$N is suggested to probe recently accreted gas and HCN & HCO$^+$ probe internal motions of dense cores.

In a theoretical approach, we established more realistic dynamical stability conditions of a filament and a dense core in order to correctly estimate which filaments or dense cores may form a star or multiple stars in a near future. A molecular cloud is typically turbulent with supersonic $rms$ velocity and with feedback from protostars such as radiation and outflows. Filaments and dense cores in molecular clouds are high likely to have been or to be perturbed by turbulence or feedback. Thus, it is important to consider those perturbations in understanding the dynamical stability of filaments and dense cores within molecular clouds. In this thesis, we mainly consider a large scale converging flow or an accretion flow, which is often found in filamentary star-forming regions (e.g. Taurus: Goldsmith et al. 2008; Palmeirim et al. 2013, Serpens-South: Kirk et al. 2013a), and non-ionizing radiation originating from nearby YSOs and the interstellar radiation field.

The structure of the thesis is as follows. In chapter 2, we present our deep, high resolution observations of the L1495-B218 filaments in the dense gas tracer NH$_3$. From the NH$_3$ observations, we calculate the kinetic temperature and the non-thermal velocity dispersion maps of the filaments. We also identified NH$_3$ dense cores along the filaments using a novel hybrid dendrogram/seeded-watershed algorithm. We made a virial analysis of the NH$_3$ dense cores and discussed the evolutionary path of dense cores embedded within molecular filaments.

In chapter 3, we present our high resolution observations in HC$_7$N and CCS toward the L1495-B218 filaments. We show that only three regions (two more evolved regions and one less evolved region) are bright in CCS and only two regions (more evolved regions) have a considerable emission in HC$_7$N. We also show that CCS does not trace core-like structures but instead traces recently condensed gas associated with dense cores. We discuss the future dynamical evolution of L1495A in B7 and
L1521D in B213E by studying CCS kinematics with respect to $^{12}$CO flows and by observing the two regions in the infall tracers HCN and HCO$^+$. We found that L1495A is at the center of a supersonic converging flow and likely to be a stellar association forming region while L1521D is a slowing contracting core similar to L1544.

In chapter 4, we extend the conventional stability condition of a dense core, which is the critical Bonnor-Ebert sphere (BES), by considering spherical symmetric homologous converging flows. We find that the critical mass and size of a dense core is considerably smaller when the peak inflow speed of a converging flow is same with the sound speed. We also study the limit of the peak infall velocities of the critical BES and uniform density cores with the condition of spontaneous gravitational collapse. We find that the two models have a specific range of the peak infall speed for a given density structure when a dense core undergoes pure gravitational collapse. If a dense core has an infall speed out of the range restricted by the two models, then it is likely perturbed by the surrounding environment. The stability criteria in this section are useful to test the stability of dense cores with converging flows, for example in L1495A-N.

In chapter 5, we also extend the conventional stability conditions of a filament and a dense core by considering the non-ionizing radiation pressure. Non-ionizing radiation can affect the dynamical stability of molecular dense structures by exerting pressure on dust grains, which are tightly coupled to the gas. We estimate the critical line densities of filaments and the critical masses and sizes of dense cores under various strength of the non-ionizing radiation field. We consider three dust grain models with different chemical composition of dust grains (silicate, carbonaceous, ice-mantle) as well as the two different types of radiation sources (ISRF and stellar radiation). We find that the stability thresholds (mass, size, line density, etc.) are considerably smaller than the conventional thresholds (the critical BES and Ostriker cylinder) when the radiation pressure is similar to the gas pressure at the filament or dense core center, which is easily achieved at an inner Galactic region (e.g. molecular ring and galactic center).
In chapter 6, we summarize our discussions and conclusions.
We present deep NH$_3$ observations of the L1495-B218 filaments in the Taurus molecular cloud covering over a 3 degree angular range using the K-band focal plane array on the 100m Green Bank Telescope. The L1495-B218 filaments form an interconnected, nearby, large complex extending over 8 pc. We observed NH$_3$ (1,1) and (2,2) with a spectral resolution of 0.038 km/s and a spatial resolution of 31". Most of the ammonia peaks coincide with intensity peaks in dust continuum maps at 350 µm and 500 µm. We deduce physical properties by fitting a model to the observed spectra. We find gas kinetic temperatures of 8 − 15 K, velocity dispersions of 0.05 − 0.25 km/s, and NH$_3$ column densities of 5×10$^{12}$ − 1×10$^{14}$ cm$^{-2}$. The CSAR algorithm, which is a hybrid of seeded-watershed and binary dendrogram algorithms, identifies a total of 55 NH$_3$ structures including 39 leaves and 16 branches. The masses of the NH$_3$ sources range from 0.05 M$_\odot$ to 9.5 M$_\odot$. The masses of NH$_3$ leaves are mostly smaller than their corresponding virial mass estimated from their internal and gravitational energies, which suggests these leaves are gravitationally unbound structures. 9 out of 39 NH$_3$ leaves are gravitationally bound and 7 out of 9 gravitationally bound NH$_3$ leaves are associated with star formation. We also find that 12 out of 30 gravitationally unbound leaves are pressure-confined. Our data suggest that a dense core may form as a pressure-confined structure, evolve to a gravitationally bound core, and undergo collapse to form a protostar.

2.1 INTRODUCTION

Filaments are common structures in molecular clouds (e.g. Ridge et al., 2006; André et al., 2010, 2014; Men’shchikov et al., 2010). They are important structures.
in star formation because most dense cores form within filaments and stars form within dense cores. Studies of isolated dense cores suggest that stars form mainly as a result of gravity within dense cores, but the formation of filaments and the evolution from filaments to dense cores are still unclear. Recent theoretical studies of molecular clouds suggest that a filament may form by colliding flows of turbulent media and some dense parts further undergo contraction by either self-gravity or ram pressure of colliding flows to make dense cores (e.g. Balsara et al., 2001; Padoan et al., 2001; Klessen et al., 2005; Vázquez-Semadeni et al., 2005; Gómez et al., 2007; Nakamura and Li, 2008; Gong and Ostriker, 2009, 2011; Heitsch et al., 2008, 2009; Chen and Ostriker, 2014). Since filaments are formed through colliding flows in this scenario, young and less massive dense cores within the filaments are gravitationally unbound, while evolved and massive dense cores tend to be gravitationally bound. Thus, observationally studying the physical state of dense cores within filaments with respect to their evolutionary stage is an essential step to understand how low-mass star formation proceeds within filaments.

The Taurus molecular cloud is one of the best testbeds for studies of both filaments and dense cores because it is nearby (D ∼ 140 pc, Loinard et al. 2007; Torres et al. 2009) and forming low-mass stars (no O and B stars) in a relatively isolated mode (with little spatially confusion, see review by Kenyon et al. 2008). The L1495-B218 filaments are prominent in the Taurus molecular cloud extending over 3 degrees on the sky. They are known to contain cores that range from chemically young starless cores (e.g. L1521B and L1521E; Hirota et al. 2004; Tafalla and Santiago 2004) to more evolved protostellar cores (in B7/L1495; see Rebull et al. 2010). In addition, this region is well mapped in dust extinction (Schmalzl et al., 2010), dust continuum emission (Herchel Space Observatory, Palmeirim et al. 2013), dust polarization (Chapman et al., 2011), low density molecular gas tracers (^{12}CO and ^{13}CO Goldsmith et al. 2008; C^{18}O Onishi et al. 1996; Hacar et al. 2013), and a dense gas tracer (N_{2}H^{+}, Hacar et al. 2013). These maps provide basic information on the structure of the filaments and the positions of some of the densest cores.
Hacar et al. (2013) observed the L1495-B218 filaments in C$^{18}$O $J = 1-0$ and N$_2$H$^+$ $J = 1-0$. C$^{18}$O $J = 1-0$ is a moderate density ($n_{\text{eff}} \leq 10^3 \text{ cm}^{-3}$ at 10 K)$^2$ gas tracer and is a good tracer with which to study velocity structure of moderate density regions since its optical depth is much smaller than those of $^{12}$CO and $^{13}$CO. Using C$^{18}$O, they estimated the LSR velocities and found that the L1495-B218 filaments are not a single coherent velocity structure but a complex group of at least 35 interwoven filaments. However, C$^{18}$O does not trace dense cores because CO is depleted onto dust grains at a high density and low temperature (e.g. Kuiper et al., 1996; Willacy et al., 1998; Kraemer and Jackson, 1999; Caselli et al., 1999; Bacmann et al., 2002; Bergin et al., 2002; Tafalla et al., 2002; Pontoppidan, 2006; Ford and Shirley, 2011). To find dense cores, the authors used N$_2$H$^+$ $J = 1-0$, which is a dense gas tracer with an effective excitation density above $10^4 \text{ cm}^{-3}$ (Reiter et al., 2011), but the densities that N$_2$H$^+$ traces are much higher than those traced by C$^{18}$O, which suggests that N$_2$H$^+$ does not trace the complete population of both young and evolved dense cores within the filaments. In order to identify the complete population of dense cores and probe the transition between filaments and dense cores, observations that trace gas with densities intermediate between those traced by C$^{18}$O and N$_2$H$^+$ are required.

Ammonia is an excellent probe that has an effective excitation density intermediate between C$^{18}$O and N$_2$H$^+$. Ammonia starts to form at an early stage of prestellar core evolution and becomes brightest during the later stage of prestellar core evolution (Suzuki et al., 1992). NH$_3$ traces densities $>10^3 \text{ cm}^{-3}$ (Flower et al., 2006; di Francesco et al., 2007). Owing to the NH$_3$ chemistry, the ammonia abundance steeply increases as a molecular cloud core becomes denser and it may be the most abundant molecule after CO that is easily observed within dense cores (Tafalla et al., 2004; Aikawa et al., 2005). Ammonia is also a good tracer for probing the physical condition of filaments. The inversion transitions of ammonia are known to provide an accurate measure of gas kinetic temperature along the line-of-

$^2n_{\text{eff}}$ is the effective excitation density, which is needed to produce $T_R = 1 \text{ K}$ line (Evans, 1999; Shirley, 2015). See Reiter et al. (2011) for a summary of effective excitation densities.
sight (e.g. Ho and Townes, 1983; Rosolowsky et al., 2008) since a transition between meta-stable energy levels is possible only through collisions. If the signal to noise ratio is large enough, the uncertainty of a kinetic temperature measured using the inversion transitions of ammonia may be as low as 0.1 K. An accurate measurement of the kinetic temperature also permits an accurate determination of non-thermal motions, making ammonia an excellent tracer for investigating kinematics.

We observed the L1495-B218 filaments in NH$_3$ (1,1) and (2,2). The complete maps spans over three degrees on the sky and is the largest NH$_3$ map published to date. In this paper, we present observations of filaments and dense cores traced by the ammonia emission. We analyze the physical properties of nested structure identified in the NH$_3$ emission. The kinematics will be explored in a subsequent paper. The layout of this paper is as follows: §2.2 briefly introduces our mapping of the L1495-B218 filaments in NH$_3$ and the data reduction procedures. In §2.3 we present the basic physical properties of the L1495-B218 filaments and the dense structures identified by the CSAR algorithm within the filaments. In §2.4 we discuss physical states of the dense structures identified by the CSAR algorithm. Finally, in §2.5 we summarize our results.

2.2 OBSERVATION & DATA REDUCTION

2.2.1 Mapping of the L1495-B218 Filaments in NH$_3$ (1,1) and (2,2)

We observed the L1495-B218 filaments in the Taurus cloud using the 100 meter Robert C. Byrd Green Bank Telescope (GBT12A295). Observations were carried out from April 2012 to February 2013 in 13 shifts totalling 66 hours. We used the K-band focal plane array (KFPA), which has 7 beams each with dual polarization in a hexagonal arrangement (for technical properties of the KFPA, see https://science.nrao.edu/facilities/gbt/observing/GBTog.pdf). We used the 7+1 mode of the KFPA, which allows the center beam to observe two frequencies si-

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3The National Radio Astronomy Observatory is a facility of the National Science Foundation operated under cooperative agreement by Associated Universities, Inc.
multaneously while the other 6 beams observe only a single frequency. The bandwidth of each frequency window is 50 MHz. For 7 beams, we set the rest frequency to be 23.702 GHz, which between the strongest hyperfine component of \( \text{NH}_3(1,1) \) (23.6944955 GHz; Lovas 1992) and \( \text{NH}_3(2,2) \) (23.7226333 GHz; Lovas 1992). For the central beam we selected a second frequency to include the CCS \( N_J = 1_2 \rightarrow 0_1 \) transition (22.344033 GHz; Yamamoto et al. 1990); these results is in the next chapter. The spectrometer is configured to have 16384 channels across 50 MHz, yielding a resolution of 3.05 KHz (0.0386 km s\(^{-1}\) at 23.705 GHz). Since the typical thermal broadening of starless cores is < 0.2 km s\(^{-1}\) (e.g. Jijina et al., 1999; Caselli et al., 2002b; Rosolowsky et al., 2008), the line profiles of \( \text{NH}_3 \) are well resolved in this setup.

We observed the L1495-B218 filaments in 20 rectangular regions that are arranged to cover high extinction regions (Schmalzl et al., 2010) while minimizing the total observation time. Each rectangular region is 20' by 6' or 6' by 20' oriented in Galactic coordinates. Data for each rectangular region were obtained by on-the-fly mapping of 25 rows 20' long with a row spacing of 15.12". The integration time per row was 150 seconds. The resulting integration time for a direction observed with all 7 beams is \( \sim 30 \) seconds, the total observation time for each rectangular region was 1 hour and 20 minutes. The reference OFF positions were picked to be where the extinction is low enough to suggest that they would have negligible \( \text{NH}_3 \) emission, and be close to each rectangular region. The average of two observations of the OFF position made before and after mapping each rectangular region was subtracted from each on-source observation. We also carried out “daisy” pattern observations to improve the signal-to-noise ratio (SNR) toward 29 bright \( \text{NH}_3 \) regions that we found from the mapping of rectangular regions. Each daisy pattern observation was made with a radius of 2.8' and 22 cycles following a rose curve. The daisy pattern observation completely covers the inner 2.8' radius with all 7 beams for an average integration time of 38 seconds per position. The total integration time for each daisy pattern is 20 minutes. We updated the pointing of the telescope using the quasar 0403+2600 between every mapping of a rectangular region or every
three daisy patterns. The pointing corrections were usually less than 10″, compared to the beam size of the GBT at 23 GHz of 31″.

Data reduction was performed using the GBT KFPA pipeline. The data were regridded into regular spaced pixels using AIPS and the gain of the 7 beams are calibrated through observing Venus and Jupiter in every observing shift. As a double check on the absolute flux calibration, we also observed the peak dust continuum position of the L1489PPC starless core (α = 04:04:47.6, δ = +26:19:17.9, J2000.0) every observing shift (Young et al., 2004; Ford and Shirley, 2011). The calibrations using Venus and Jupiter and the calibration using L1489PPC agreed within 25%.

2.3 RESULTS

2.3.1 L1495-B218 Filaments Seen in Ammonia

Integrated intensity of NH$_3$ (1,1) & (2,2) and Definition of Subregions

Figure 2.1 presents the integrated intensity of NH$_3$ (1,1) emission. The NH$_3$ (1,1) map clearly reveals dense cores associated with the L1495-B218 filaments. On the other hand, the filamentary structures between dense cores are barely detected, while those regions are very clearly visible in an extinction map (Schmalzl et al., 2010) and $^{13}$CO (Goldsmith et al., 2008; Pineda et al., 2010b). This suggests that the density of filamentary structures between dense cores may not be high enough to form NH$_3$ efficiently ($n_{H_2} < 10^3$ cm$^{-3}$; Flower et al. 2006) or that ammonia molecules exist but excitation conditions are insufficient to observe the (1,1) emission at our average rms level of 120 mK (Figure 2). The NH$_3$ (2,2) emission is detected only from dense cores and is much weaker than (1,1) emission, which implies that the kinetic temperature may be too low to excite to the $J$=2, $K$=2 state except in dense gas of $n_{H_2} > 10^4$ cm$^{-3}$ (Ho and Townes, 1983; Danby et al., 1988; Maret et al., 2009).

We compare our NH$_3$ (1,1) emission with 500 µm dust continuum emission seen by the SPIRE instrument of the Herschel Space Observatory in Figure 2.1 (Palmeirim et al., 2013; Marsh et al., 2014). The 500 µm dust continuum emission is thermal emission from dust grains and traces the column density structure of
the filaments. The bright regions of 500 µm and NH$_3$ (1,1) emission show a good agreement with each other. The 500 µm emission is more spatially extended than the NH$_3$ (1,1) emission. The lowest 500 µm intensity levels at which NH$_3$ (1,1) and (2,2) are detected at is 25 MJy/sr (N$_{H_2}$ = 5.6 × 10$^{21}$ cm$^{-2}$ at 10 K) and 177 MJy/sr (N$_{H_2}$ = 3.9 × 10$^{22}$ cm$^{-2}$ at 10K), respectively. This shows that NH$_3$ is a good tracer with which to identify intermediate to dense regions within filaments.

In Figure 2.3, we zoom into each region and show the location of protostars identified using IRAS and the Spitzer Space Telescope (Rebull et al., 2010, 2011; Kryukova et al., 2012). We divide the map into subregions based on Barnard’s identification (Barnard et al., 1927). We discuss each region and determine if it is more evolved (having a Class I/flat spectra protostar associated with NH$_3$ emission) or less evolved below.

B7/L1495 contains multiple Class II and III protostars, two flat spectra protostars, and one Class I protostar, none of which is embedded within NH$_3$ core but there are nearby NH$_3$ cores. The three dense cores seen in NH$_3$ are mostly commonly referred to as L1495A-S, L1495A-N, and L1495B (Lee et al., 2001). In B213 (also referred to as L1521D, Codella et al. 1997), there are two Class I protostars embedded in NH$_3$ cores in the eastern part of the region and another three Class I protostars, in the western part of the region. Therefore, B7 and B213 are the most evolved regions in the L1495-B218 filaments and B7 is likely to be more evolved than B213 (see Hacar et al., 2013).

B218 has only one Flat spectrum protostar embedded within a NH$_3$ core. Thus, we classify B218 as an evolved region with star formation, although the NH$_3$ emission is dominated by two bright starless cores. This region is also sometimes confusingly referred to as B217 (Benson and Myers, 1989). We use the notation B218 in this paper since it is closer to the original Barnard catalog position of B218 (Kenyon et al., 2008).

B10 also has only one protostar, which is Class II. In this work, we assume that B10 is a less evolved region because there is no other protostar except the one Class II protostar which is not directly associated with any dense cores in the region.
In contrast to the other regions, B211 and B216 show weak NH$_3$ (1,1) emission and do not possess any known protostars. Thus, both regions are less evolved. The B216 region is also named L1521B (Lee et al., 2001) and contains starless cores with large abundances of carbon chain molecules (Hirota et al., 2004).

Hacar et al. (2013) have also determined the relative ages of subregions using N$_2$H$^+$ 1−0 emission and the locations of protostars. Their estimates of relative age of subregions agrees well with ours except for B216. In their observation of N$_2$H$^+$ 1−0, they did not detect emission in B216. B216 has weaker NH$_3$ (1,1) emission and stronger CCS emission than B211 (Seo et al. in prep.) suggesting that B216 is a younger region. Thus, B211 and B216 are both relatively young and inactive parts of the L1495-B218 filaments with B216 likely being the most newly condensed part of the filaments.

**Spectral x'ine fitting**

To understand the physical properties of NH$_3$ structures in the L1495-B218 filaments, we fit a simple NH$_3$ spectra model to the observed spectra and deduce physical quantities including kinetic temperature $T_k$, excitation temperature $T_{ex}$, total optical depth of NH$_3$ (1,1) emission $\tau_1$, dispersion of emission line $\sigma_v$, and LSR velocity $v_{lsr}$ (See Appendix A for details of the fitting algorithm). The fitting algorithm includes only one velocity component, and this is satisfactory for all regions in the L1495-B218 filaments except B211 (see section 3.2.2 for the LSR velocity of NH$_3$). We only fitted spectra where the detection of NH$_3$ (1,1) and (2,2) emission is above 5$\sigma_{T_{mb}}$ and 3$\sigma_{T_{mb}}$ respectively, where $\sigma_{T_{mb}}$ is the baseline $rms$.

**Gas kinetic temperature and excitation temperature**

Maps of the gas kinetic temperature deduced from NH$_3$ (1,1) and (2,2) emission are presented in Figure 2.4. The gas kinetic temperature ranges from 8 K to 15 K in most of the region mapped. At the bright NH$_3$ peaks indicated by red arrows, the kinetic temperature decreases toward the core center. The drop in the kinetic temperature is about -1 to -3 K, or about 10% – 30% of the average kinetic temperature of
Figure 2.1 Left: 500 µm dust continuum emission (color) seen by the SPIRE instrument of the Herschel Space Observatory. Right: map of integrated intensity of NH$_3$ (1,1).
Figure 2.2 rms of the ammonia $T_{mb}$ map at a spectral resolution of 6.1 kHz (corresponding to 0.077 km s$^{-1}$), obtained by smoothing two spectral channels.
Figure 2.3 Maps of integrated NH$_3$ (1,1) emission and positions of known proto-stars. The x-axis is galactic longitude and Y-axis is galactic latitude. The names of previously studied cores are indicated in the image.
surroundings. This temperature drop may ultimately break the dynamical stability of these dense cores and lead them to collapse (see Khesali et al., 2013).

Figure 2.5 presents histograms of physical quantities to investigate statistical differences between subregions. In the first panel of Figure 2.5, we show histograms of the gas kinetic temperature. Medians of the gas kinetic temperature in subregions are all around 9.5 K with less than 1 K deviation (60% of all spectra have kinetic temperature within ±1 K of 9.5 K.). The two less evolved regions, B10 and B216, have lower kinetic temperature than the other more evolved regions (we exclude B211 from discussion because it has multiple velocity components and our NH\textsubscript{3} fitting model works only for a single velocity component.). The histograms of the gas kinetic temperature demonstrate a tendency that a less evolved region is cooler than a more evolved region and that a denser region is cooler than diffuse region if there are no embedded protostars. The gas kinetic temperature differences are, however, subtle.

Histograms of NH\textsubscript{3} excitation temperature are shown in the second panel of Figure 2.5. Median values of the excitation temperature are mostly around 5 K, which is only half of the median values of the kinetic temperature. This demonstrates that NH\textsubscript{3} in all regions is subthermally excited.

**LSR velocity**

NH\textsubscript{3} lines at 39 NH\textsubscript{3} integrated intensity peaks in the L1495-B218 filaments are shown in Figures 2.6 and 2.7 (for peak locations, see §3.2.1). The range of NH\textsubscript{3} \( v_{\text{lsr}} \) is from 5.0 km s\(^{-1}\) to 7.2 km s\(^{-1}\), which is the same range seen in the C\(^{18}\)O \( J = 1-0 \) spectra of Hacar et al. (2013). NH\textsubscript{3} structures are confined to a single velocity component except in B211. Comparing NH\textsubscript{3} to C\(^{18}\)O (Hacar et al., 2013) and \(^{13}\)CO (Goldsmith et al., 2008), we found that all NH\textsubscript{3} peaks have corresponding velocity components in \(^{13}\)CO, while only 28 out of 39 NH\textsubscript{3} peaks directly associated with the same velocity components in C\(^{18}\)O. Some NH\textsubscript{3} peaks where the NH\textsubscript{3} integrated intensity is relatively high (≥40 K km s\(^{-1}\), \textit{i.e.} four core-like structures neighboring in a row in B213) do not have corresponding velocity components in C\(^{18}\)O and \(^{13}\)CO.
Figure 2.4 Maps of the gas kinetic temperature in the L1495-B218 filaments. X-axis is galactic longitude and Y-axis is galactic latitude. Locations of protostars are marked with various symbols depending on the spectral classification of the protostars. Red arrows indicate dense cores having gas kinetic temperatures decreasing toward core centers. Regions with multiple velocity components in B211 are suppressed.
Figure 2.5 Histograms of physical properties for each region. Vertical lines denote the median values. High fractions at the very left edge of the first panel are due to low signal-to-noise level at the edge of map. Those are excluded in estimating median values. B211 is also excluded due to the presence of multiple velocity components. The dotted vertical line in the fourth panel denotes $R_p = 1$. 
emission at those places is also weak. We also found a few NH₃ peaks with both NH₃ and C¹⁸O emission but discrepancies in their LSR velocities, which may be a sign of chemical differentiation.

The LSR velocity structure of the L1495-B218 filaments was extensively studied by Hacar et al. (2013) using C¹⁸O J= 1−0. They identified 35 filaments using a “friend of friends” algorithm (FIVe; Hacar et al. 2013) and for each filament they deduced physical properties including mass, dispersion of \(v_{\text{lsr}}\), nonthermal motion, mean kinetic temperature. (See Table 3 in Hacar et al. 2013). Based on dense cores identified from their N₂H⁺ J = 1−0 observations, they suggested only 7 out of 35 filaments are “fertile” meaning they contain a N₂H⁺-bright dense core. However, N₂H⁺ traces gas with density higher than \(10^4\) cm\(^{-3}\), which is the central density of relatively evolved dense cores (Lada et al., 2008). It is likely that dense cores at early stages are not fully revealed by N₂H⁺. Ammonia molecules have an intermediate effective excitation density between those of C¹⁸O J = 1−0 and N₂H⁺ J = 1−0; therefore, our ammonia map may probe the velocity structure of both developing dense cores as well as later stages.

Matching the LSR velocity of NH₃ (1,1) and (2,2) structures to 35 filaments of Hacar et al. (2013), we found NH₃ (1,1) emission in filaments 6, 8, 10, 11, 15, 20, 28, 32 (Table 3 of Hacar et al. 2013). Hacar et al. (2013) suggested that filaments 6, 10, 11, 15, 20, 32 and 34 are fertile (i.e. having dense cores in their N₂H⁺ observation). Filament 34 is not within the field of view of our observation, so it is excluded in this study. Our NH₃ observations reveal that two more filaments, 8 and 28, also have NH₃ emission. Hacar et al. (2013) argued that a filament tends to be sterile when the line density of filament is less than 15 Mₜ pc\(^{-1}\), which is the critical density for an isothermal cylinder at 10 K (Stodólkiewicz, 1963; Ostriker, 1964). Filament 28 has a line density of 24.1 Mₜ pc\(^{-1}\) (Hacar et al., 2013) but there was no detection of dense cores in N₂H⁺ J = 1−0. With our NH₃ observations, we found that there are at least two dense cores in filament 28. On the other hand, filament 8 has a line density of 9.3 Mₜ pc\(^{-1}\) (Hacar et al., 2013), which is lower than the critical line density, but we see that there is a relatively bright NH₃ core. Since Hacar et al.
(2013) estimated the masses of filaments from \( \text{C}^{18}\text{O} \ J = 1–0 \), it may be possible that they underestimated the mass, due to CO depletion. The filament mass estimated from dust continuum may be larger than mass estimated from \( \text{C}^{18}\text{O} \ J = 1–0 \), but the filaments cannot be distinguished along the line of sight, which will result in an overestimation of filament masses. From the fact that filament 8 is relatively bright in \( \text{NH}_3 \ (1,1) \), it is likely that the line density of \( 9.3 \, \text{M}_\odot \, \text{pc}^{-1} \) is the minimum line density of filament 8 and the true line density may be higher than the critical value.

**Velocity dispersion**

Velocity dispersions range from 0.05 to 0.25 \( \text{km s}^{-1} \) except in B211 (B211 has multiple velocity components which are not pursued in this study). Overall, we found that velocity dispersions, which include both thermal and nonthermal components, tend to be wider in the less evolved regions than those in the more evolved regions. Histograms of the velocity dispersions (Figure 2.5) show this trend clearly. The median values of velocity dispersions for the least evolved regions, B216, is \( 0.118 \, \text{km s}^{-1} \), while the most evolved regions, B7, has a median dispersion equal to \( 0.104 \, \text{km s}^{-1} \) (the thermal velocity dispersion of \( \text{NH}_3 \) at 10 K is \( 0.069 \, \text{km s}^{-1} \)).

Using accurate gas kinetic temperatures deduced from \( \text{NH}_3 \ (1,1) \) and (2,2), we estimate the velocity dispersions of nonthermal kinetic motions \( \sigma_{\text{nt}} \) (Figure 2.8) as follows:

\[
\sigma_{\text{nt}} = \sqrt{\sigma_v^2 - \frac{kT_k}{\mu_{\text{NH}_3}m_h}},
\]

where \( k \) is Boltzmann constant, \( \mu_{\text{NH}_3} \) is the molecular weight of ammonia (17 a.m.u.), and \( m_h \) is the mass of a hydrogen atom. We find most regions of the filaments are supported mainly by thermal pressure (the fourth panel in Figure 2.5) while some edges of filaments are dominated by nonthermal support (\( \sigma_{\text{nt}} \geq \sigma_{\text{thermal}} = 0.188 \, \text{km s}^{-1} \) at \( T_k = 10 \, \text{K} \) and \( \mu = 2.33 \)). The median ratio of thermal support to nonthermal support \( (R_p \equiv c_s^2/\sigma_{\text{nt}}^2) \) is around five. In addition, the median ratios of thermal support to nonthermal support tends to be larger for more evolved regions (median \( R_p = 6.3 \) in more evolved region B7, while \( R_p = 3.7 \) in less evolved
Figure 2.6 $\text{NH}_3$ (1,1) lines (black histograms), the best-fit single component $\text{NH}_3$ (1,1) model (dashed red lines), C$^{18}$O 1−0 (green histograms; Hacar et al. 2013), and $^{13}\text{CO}$ 1−0 (blue histograms; Goldsmith et al. 2008; Narayanan et al. 2008) at the $\text{NH}_3$ peak positions of $\text{NH}_3$ leaves identified by CSAR. Black vertical lines denote locations and relative strength of $\text{NH}_3$ (1,1) hyperfine lines. Red-green arrows denote the LSR velocities of $\text{NH}_3$ leaves. $^{13}\text{CO}$ and C$^{18}$O have multiple velocity components in some of $\text{NH}_3$ leaves but $\text{NH}_3$ seems to be confined to one velocity component and is well fitted with a $\text{NH}_3$ line model having a single velocity component.
Figure 2.7 Continued from Fig. 2.6.
region B216). This suggests that nonthermal motions are considerably dissipated in more evolved regions and making it easier to initiate star formation, which has been seen toward dense cores in other molecular clouds (e.g. Goodman et al., 1998; Pineda et al., 2010a; Tatematsu et al., 2014a).

2.3.2 Dense NH₃ Structures Within the L1495-B218 Filaments

**Identifying dense structures in the L1495-B218 filaments**

The L1495-B218 filaments have many dense cores in different evolutionary stages of star formation, ranging from starless to protostellar cores. We use our ammonia observations to identify dense cores in the L1495-B218 filaments using the Cardiff Source-finding AlgoRithm (CSAR, Kirk et al. 2013b), which uses both a seeded watershed and dendrogram algorithm to find clumpy structures and their spatial relationships. We used the integrated intensity of the central hyperfine group (the strongest hyperfine group of nitrogen splitting) of NH₃ (1,1) to identify dense structures because the central hyperfine group gives the highest signal-to-noise ratio compared to the other four hyperfine groups. Parameters in the CSAR algorithm are set to find structures above a 7.5-σ integrated intensity level in the (1,1) transition (∼3σ in the (2,2) transition) and to identify a clump with a size larger than $\theta_{mb}/2 \sim 15''$ (fullbeam option) and a 2-σ enhancement in intensity from its background. We also used the removeedge option in order not to have overextended edges for a source. In this study, we use 2-σ as an intensity interval to clip intensity peaks instead of using typical value of 3-σ because with 3-σ intensity interval, the CSAR algorithm tends to combine two peaks into one elongated clumpy structure. These peaks typically lie on top of nested sources which are many σ brighter than their background. The CSAR algorithm also had trouble in excluding intensity spikes at edges of maps which have a high noise level ($\sigma_{T_b} \geq 0.3$ K). In order to exclude the edges of NH₃ maps, we made 12 patches and applied the CSAR algorithm to each patch.

The main reason we use NH₃ (1,1) instead of dust continuum is because NH₃ emission acts as a volume density filter. From the excitation curve of NH₃ (1,1),
Figure 2.8 Velocity dispersions of nonthermal components. The x-axis is the galactic longitude and Y-axis is the galactic latitude. Locations of protostars are marked with various symbols depending on spectral classification of protostars. Regions with multiple velocity components in B211 are suppressed.
the effective excitation density varies slowly over the range of gas kinetic temperature from 7 K to 15 K. The excitation temperature of dense structures identified by the CSAR algorithm is typically above 3.2 K, and the effective excitation density for the range of observed gas kinetic temperature and NH$_3$ abundance varies only from $7 \times 10^2$ cm$^{-3}$ to $1 \times 10^3$ cm$^{-3}$. This suggests that NH$_3$ (1,1) emission naturally filters out regions with a density lower than $\sim 7 \times 10^2$ cm$^{-3}$. This may make NH$_3$ structures appear smaller than corresponding dust continuum structures, but the NH$_3$ sources identified by the CSAR algorithm are true structures in volume density whereas analysis in dust continuum sees only an enhancement in total column density.

From the integrated intensity of the central hyperfine group of NH$_3$ (1,1), we found 39 leaves (a leaf: a structure containing an intensity peak) and 16 branches (a branch: a nested group of leaves). Figure 2.9 shows the locations of NH$_3$ leaves, and Figure 2.10 shows hierarchy among NH$_3$ leaves and branches in a dendrogram. The dendrogram has no lowest common branch because NH$_3$ (1,1) structures are highly fragmented due to the volume density and chemical filtering. In this study we refer to “active” leaves as those which have a Class I or Flat spectrum protostar within their lowest-level branches in the dendrogram (a.k.a. root) or within a distance of three times their radius. Note that this definition means that active leaves could be either starless or protostellar. Only 4 leaves have embedded protostars and 17 leaves are classified as active. Hacar et al. (2013) found 19 dense cores in N$_2$H$^+$ 1–0, while CSAR finds 39 leaves in NH$_3$ (1,1). This difference may be because NH$_3$ is an intermediate density gas tracer that reveals more cores than does the denser gas tracer N$_2$H$^+$ and our NH$_3$ observation ($\theta_{mb} \sim 31''$) have better angular resolution than the N$_2$H$^+$ observation ($\theta_{mb} \sim 50''$).

For the comparison between dust and NH$_3$, we also identified dust cores using CSAR. From the 500 $\mu$m dust continuum data of the Herschel Space Observatory ($\theta_{mb} \sim 36''$), we found 51 dust leaves. Figure 2.11 shows the locations and boundaries of dust leaves. 32 dust leaves coincide with 38 NH$_3$ leaves, while 19 dust leaves do not agree with any NH$_3$ leaf. Nine of the dust leaves contain at least two NH$_3$
leaves. This suggests that there may be multiple dense cores or fragments within a core since NH$_3$ is sensitive to volume density whereas the dust continuum is optically thin and depends on the total column density.

In order to compare how well dust continuum and NH$_3$ agree with each other, we investigate the separation of the dust peak position and the NH$_3$ peak position in a dust leaf associated with a NH$_3$ leaf. Most of NH$_3$ peak positions show agreement with corresponding dust peak positions with a median separation of 18″.7, which is slightly larger than half the FWHM beam size. However, there are four dust leaves having a separation between dust and NH$_3$ peaks larger than 30″ (the maximum separation is 66″.2), and three out of four are associated with a nearby protostar (Class I or Flat spectrum). This suggests a large separation may be due to chemical effect produced by the radiation from a protostar. The fact that most of the NH$_3$ sources agree with the dust sources demonstrates that NH$_3$ sources identified by CSAR may represent dense structures within the L1495-B218 filaments. In the following subsections all of the analysis are done with NH$_3$ sources determined from our NH$_3$ data.
Figure 2.9 Location of NH$_3$ leaves. Solid and dashed orange contours are NH$_3$ leaves and branches, respectively, identified by CSAR.
Figure 2.10 Dendrogram of NH$_3$ sources. Red solid lines denote active leaves (leaves that are embedding a ClassI/Flat spectra protostar within their lowest-level branches or closely associated with a ClassI/Flat spectra protostar within three times their leaf radiuses) and dashed red lines denote active leaves with an embedded protostar. Black circles indicate mass of NH$_3$ sources estimated from 500$\mu$m dust continuum. Red circles denote the virial mass. Purple dashed vertical lines indicates different subregions.
Size and mass

The physical properties of NH$_3$ leaves and branches are summarized in Tables 5.1, 5.2, and 2.3. The quoted kinetic temperature and velocity dispersion are the median values estimated from NH$_3$ fitting within NH$_3$ leaves and branches except for NH$_3$ leaf 11. This has a peak NH$_3$ (1,1) line intensity weaker than 5σ except at its peak location, so we fit the NH$_3$ line at the peak intensity location.

The column density of molecular hydrogen within each structure is estimated from the 500 µm continuum emission measured by Herschel Space Observatory using the following equation.

$$N_{\text{H}_2} = \frac{S_{\nu} \mu m_h B_\nu(T_{\text{dust}}) \kappa_\nu \Omega_{\text{ap}}}{m_{\text{gas}} m_{\text{dust}}}$$  \hspace{1cm} (2.2)

where $N_{\text{H}_2}$ is column density of H$_2$ molecules, $S_{\nu}$ is observed flux at frequency $\nu$, $\mu$ is the mean molecular weight, $m_h$ is the mass of hydrogen atom, $B_\nu$ is the black body brightness at frequency $\nu$, $T_{\text{dust}}$ is the temperature of a dust grain, $\kappa_\nu$ is opacity of dust grain at frequency $\nu$, $\Omega_{\text{ap}}$ is the solid angle used in the observation, and $m_{\text{dust}}/m_{\text{gas}}$ is dust-to-gas ratio, which we set to 0.01. In this work, we assume $T_{\text{dust}} = T_k$, where $T_k$ is the kinetic temperature deduced from our ammonia observations. $\mu$ is 2.8. We use $\kappa_\nu$ of 5.04 cm$^2$ g$^{-1}$ (OH5), which is column 5 of “MNR with thin ice mantle” model from Ossenkopf and Henning (1994). The peak H$_2$ column density of NH$_3$ sources in the L1495-B218 filaments ranges from $6.0 \times 10^{21}$ to $4.3 \times 10^{22}$ cm$^{-2}$, which is similar to the typical dense core column density in other molecular clouds (e.g. Jijina et al., 1999; Rosolowsky et al., 2008).

We define the size of a NH$_3$ source as the radius of a circle that has an area equal to that of a NH$_3$ CSAR leaf contour. In the L1495-B218 filaments, the FWHM from Gaussian fitting may not be a useful representation of the size of a dense core because the FWHM contour of a dense core often overlaps with the FWHM contour of its neighboring dense core. Since CSAR separates neighboring sources through hierarchical structure, we use the area identified by CSAR to estimate the size of a structure. Structures were found on scales ranging from 0.02 pc (4375 AU) to 0.1 pc (20000 AU).
Figure 2.11 Location of dust and NH$_3$ leaves. Blue contours are boundaries of dust leaves and green-black crosses are the peak positions of NH$_3$ leaves.
The masses of NH$_3$ sources are estimated by summing up the column density deduced from 500 $\mu$m continuum emission within each NH$_3$ leaf and branch (Table 5.2 and black circles in Figure 2.10). This is a typical method for estimating mass of a dense source since three dimensional morphology is usually unknown. If the three dimensional morphology is known or assumed, the mass of a local background may be subtracted, which will result a smaller dense source mass than the dense source mass estimated by simply summing up the column density. The resulting masses of NH$_3$ leaves range from 0.05 M$_\odot$ to 4.3 M$_\odot$, which are similar to typical dense core masses observed in other studies (0.1-10 M$_\odot$; Motte et al. 1998, 2001; Schmalzl et al. 2010; Onishi et al. 1996, 2002; Kirk et al. 2013b. The masses of NH$_3$ branches range from 0.78 M$_\odot$ to 9.5 M$_\odot$. In addition, we also compared three different methods of determining masses of NH$_3$ leaves in Table 5.2 and Figure 2.12 (mass estimated from dust continuum, virial mass shown in §4.1, and mass estimated from NH$_3$ column density with the fractional abundance of NH$_3$ relative to hydrogen molecules being $1.5 \times 10^{-9}$). The masses are strongly correlated with each other with a Spearmann’s rank correlation coefficient of $\rho = 0.93$. This strong linear correlation suggests that the assumptions (i.e. opacity and temperature) used to derive mass from the dust emission do not have strong systematic variations from the lowest mass to highest mass leaves. It also suggests that the properties derived from NH$_3$ sources identified by CSAR are representative of the properties of dense regions in the L1495-B218 filaments, regardless of whether those regions are traced by dust emission or NH$_3$.

We also compared nonthermal support to thermal support within NH$_3$ leaves identified by CSAR in Figure 2.12. All of the NH$_3$ sources are dominated by thermal support, with an average value of $R_p = c_s^2/\sigma_{nt}^2$ equal to 4.9, and there is no dependency on $R_p$ vs. mass. This shows that nonthermal motions are efficiently dissipated in dense regions and making it easier to initiate star formation.
Table 2.1. Properties of NH$_3$ sources

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<th>b $^b$</th>
<th>l $^c$</th>
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<th>$\delta$ $^e$</th>
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<th>Peak Dust$^e$</th>
<th>$N_d^{H_2}$</th>
<th>Embedded$^f$</th>
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<td>Td</td>
<td>Tnu</td>
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</table>

\(a\)IDs with decimal points are NH\(_3\) branches.

\(b\)Total integrated intensity of NH\(_3\) (1,1) emission.

\(c\)500 \(\mu\)m dust continuum of *Herschel Space Observatory*. Typical \(r_{ms}\) is 0.4 MJy/sr.

\(d\)Statistical uncertainty of peak column density is only due to uncertainties of the 500 \(\mu\)m dust continuum flux density and \(T_k\). \(T_d = T_k\) is assumed.

\(f\)Branches are omitted and noted as ‘-’. 
Table 2.2. Physical Properties of NH\(_3\) sources

<table>
<thead>
<tr>
<th>ID (^{a})</th>
<th>T(_k)(^{b}) [K]</th>
<th>(\sigma_v)(^{c}) [km s(^{-1})]</th>
<th>Mass from dust [M(_\odot)]</th>
<th>Virial Mass(^d) [M(_\odot)]</th>
<th>Mass from NH(<em>3) [M(</em>\odot)]</th>
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<tr>
<td>1</td>
<td>10.24(^{0.17}_{-0.42})</td>
<td>0.199(^{0.005}_{-0.005})</td>
<td>0.18</td>
<td>0.45</td>
<td>0.36</td>
</tr>
<tr>
<td>1.5</td>
<td>9.92(^{0.52}_{-0.86})</td>
<td>0.201(^{0.011}_{-0.005})</td>
<td>4.56</td>
<td>2.21</td>
<td>3.91</td>
</tr>
<tr>
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<td>9.66(^{0.09}_{-0.08})</td>
<td>0.197(^{0.005}_{-0.005})</td>
<td>0.06</td>
<td>0.19</td>
<td>0.09</td>
</tr>
<tr>
<td>3</td>
<td>10.23(^{0.93}_{-0.84})</td>
<td>0.204(^{0.014}_{-0.008})</td>
<td>4.38</td>
<td>1.45</td>
<td>1.03</td>
</tr>
<tr>
<td>4</td>
<td>8.49(^{0.99}_{-0.99})</td>
<td>0.210(^{0.011}_{-0.012})</td>
<td>0.31</td>
<td>0.34</td>
<td>0.05</td>
</tr>
<tr>
<td>5</td>
<td>9.65(^{1.52}_{-1.34})</td>
<td>0.252(^{0.018}_{-0.016})</td>
<td>1.74</td>
<td>1.93</td>
<td>0.71</td>
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<tr>
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<td>0.227(^{0.005}_{-0.020})</td>
<td>0.77</td>
<td>1.15</td>
<td>0.50</td>
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<tr>
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<td>0.201(^{0.028}_{-0.009})</td>
<td>1.85</td>
<td>1.61</td>
<td>1.39</td>
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<tr>
<td>7</td>
<td>9.48(^{0.43}_{-0.55})</td>
<td>0.194(^{0.006}_{-0.005})</td>
<td>0.88</td>
<td>0.98</td>
<td>0.77</td>
</tr>
<tr>
<td>7.5</td>
<td>9.29(^{0.78}_{-0.77})</td>
<td>0.202(^{0.022}_{-0.009})</td>
<td>9.52</td>
<td>3.89</td>
<td>7.58</td>
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<tr>
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<td>9.26(^{0.16}_{-0.21})</td>
<td>0.196(^{0.002}_{-0.003})</td>
<td>0.26</td>
<td>0.47</td>
<td>0.37</td>
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<td>0.198(^{0.008}_{-0.005})</td>
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<td>0.197(^{0.007}_{-0.005})</td>
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<td>0.89</td>
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<td>0.224(^{0.029}_{-0.012})</td>
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<td>0.61</td>
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<td>0.203(^{0.012}_{-0.011})</td>
<td>0.10</td>
<td>0.58</td>
<td>0.12</td>
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<td>0.213(^{0.026}_{-0.017})</td>
<td>0.95</td>
<td>0.98</td>
<td>0.98</td>
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<tr>
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<td>0.220(^{0.024}_{-0.021})</td>
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<td>1.47</td>
<td>1.24</td>
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<td>0.236(^{0.009}_{-0.013})</td>
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<td>0.31</td>
<td>0.03</td>
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<td>13.25</td>
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<td>0.227(^{0.036}_{-0.022})</td>
<td>2.79</td>
<td>2.67</td>
<td>2.12</td>
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<tr>
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<td>0.226(^{0.010}_{-0.008})</td>
<td>0.23</td>
<td>0.61</td>
<td>0.13</td>
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<tr>
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<td>0.217(^{0.035}_{-0.014})</td>
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<td>0.50</td>
<td>0.15</td>
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<tr>
<td>15.5</td>
<td>9.56(^{0.64}_{-1.06})</td>
<td>0.206(^{0.039}_{-0.012})</td>
<td>2.73</td>
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<td>0.196(^{0.007}_{-0.005})</td>
<td>1.62</td>
<td>1.22</td>
<td>1.97</td>
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Table 2.2 (cont’d)

|   | 18   | 18.5 | 19   | 19.25 | 20   | 21   | 21.925 | 22   | 22.85 | 23   | 23.75 | 24   | 24.5 | 25   | 26   | 27   | 27.75 | 28   | 28.5 | 29   | 30   | 30.5 | 31   | 32   | 33   | 34   |
|---|------|------|------|-------|------|------|--------|------|------|------|------|------|------|------|------|------|------|------|------|------|------|------|------|------|------|
|   | 0.43 | 0.54 | 0.71 | 0.99  | 1.22 | 1.36 | 0.87   | 0.58 | 0.82  | 0.51  | 0.94  | 0.15  | 0.05  | 0.48  | 1.31  | 1.01  | 0.99  | 0.95  | 0.99  | 0.67  | 1.08  | 1.14  | 0.49  | 1.96  | 0.53  | 0.57  |
|   | 0.40 | 0.47 | 0.47 | 1.24  | 1.29 | 1.30 | 0.82   | 0.75 | 0.77  | 0.54  | 0.80  | 0.16  | 0.23  | 0.78  | 1.22  | 1.09  | 1.31  | 1.32  | 1.32  | 0.97  | 1.18  | 1.09  | 0.52  | 1.32  | 0.58  | 0.91  |
|   | 0.06 | 0.06 | 0.06 | 0.013 | 0.012 | 0.016 | 0.017  | 0.014 | 0.017 | 0.005 | 0.018 | 0.003 | 0.013 | 0.009 | 0.012 | 0.004 | 0.010 | 0.010 | 0.014 | 0.008 | 0.022 | 0.014 | 0.008 | 0.013 | 0.027 |
|   | 0.09 | 0.010 | 0.006 | 0.010 | 0.007 | 0.011 | 0.012  | 0.011 | 0.012 | 0.011 | 0.012 | 0.003 | 0.003 | 0.009 | 0.012 | 0.011 | 0.014 | 0.014 | 0.005 | 0.007 | 0.012 | 0.011 | 0.005 | 0.011 | 0.015 |
|   | 12   | 9.75 | 0.71 | 4.06  | 0.30  | 0.92  | 8.97   | 1.89  | 6.84  | 0.95  | 4.77  | 0.88  | 1.81  | 0.27  | 0.24  | 0.08  | 0.85  | 0.15  | 0.77  | 0.19  | 0.19  | 0.19  | 1.08  | 0.26  | 2.49  | 0.07  |
|   | 40   | 0.40 | 0.22 | 2.15  | 0.66  | 1.30  | 3.82   | 1.39  | 3.55  | 1.11  | 7.04  | 0.77  | 1.38  | 0.53  | 0.77  | 0.45  | 1.65  | 0.70  | 0.68  | 0.68  | 0.68  | 1.39  | 0.40  | 1.87  | 0.69  |
Table 2.2 (cont’d)

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<td>0.209</td>
<td>0.018</td>
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<td>0.68</td>
<td>0.199</td>
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</tbody>
</table>

- **a** IDs with decimal points are branches in dendrograms.
- **b** Median $T_k$ of the NH$_3$ leaves and branches. Subscripts and superscripts are 15.9th %tile and 84.1th %tile of the $T_k$ distribution within NH$_3$ leaves and branches.
- **c** Median $\sigma_v$ of the NH$_3$ leaves and branches with the mean molecular weight $\mu$=2.33. Subscripts and superscripts are 15.9th %tile and 84.1th %tile of the $\sigma_v$ distribution within NH$_3$ leaves and branches.
- **d** The NH$_3$ leaves and branches are assumed to be homoeoidal ellipsoids for the virial mass calculation.
- **e** The fractional abundance of NH$_3$ is assumed to be $1.5 \times 10^{-9}$ (Typical fractional abundance ranges from $10^{-9}$ to $10^7$, e.g. Jijina et al. 1999; Tafalla et al. 2004).
**Shape and orientation**

The two dimensional morphology of the NH$_3$ sources is estimated using the moment of integrated intensity of NH$_3$ (1,1) with the center at the location of the intensity peak within a leaf or a branch. From the moment of integrated intensity we derive the principal axes, the axis ratios, and the angles with respect to galactic longitude, which are also summarized in Table 2.3. The distribution of angles of NH$_3$ leaves with respect to their lowest-level branches in the dendrogram is shown in the third panel of Figure 2.13. The minimum and the maximum phase angles of NH$_3$ leaves with respect to their lowest-level branches are -49° and 53°, respectively. Most of NH$_3$ leaves are well aligned with their lowest-level branches within a range of ±40°. Since a lowest-level branch corresponds to a filament, this suggests that NH$_3$ leaves are well aligned with filaments and were likely formed within the filaments.

The axis ratio of the NH$_3$ leaves is estimated as ratio of apparent short axis/long axis (Table 2.3). The distribution of axis ratios is shown in the fourth panel of Figure 2.13. The axis ratios range from 0.2 to 1.0 and are mostly around 0.5, which is close to the average value obtained from dense core observations by Myers et al. (1991a). In order to deduce the three dimensional morphology of NH$_3$ leaves, we compared the observed axis ratio distribution to the probability distribution of apparent axis ratios of spheroids using a Monte Carlo method (see Appendix B for more details). We found the best-fit axis ratio distribution when the spheroids are prolate and the orientation angles (the angle between the equatorial planes of spheroids and the observer’s line of sight) range from -80° to 80° (thick solid line in the fourth panel of Figure 2.13). The parameters of the prolate spheroids of the best-fit model are a mean axis ratio (short axis/ long axis) of 0.40 and a dispersion of 0.11 with $\chi^2 = 5.92$. We could not fit the distribution model using oblate spheroids within 3-σ uncertainty ($\chi^2 = 9$) unless we limit the orientation angles to a much narrower range than the observed distribution of the angles between the axes of NH$_3$ leaves and the axes of filaments (third panel of Figure 2.13). Thus, the NH$_3$ leaves are more likely prolate spheroids than oblate spheroids.
Figure 2.12 Relations between masses (the first, second, and third panels). The red line denotes $x=y$. Ratio of thermal support to nonthermal support (the fourth panel). The red line denotes $R_p = 1$. 
Table 2.3. Geometrical Properties of NH$_3$ sources

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\(^a\)IDs with decimal points are branches in dendrograms.

\(^b\)The phase angle of NH\(_3\) leaves with respect to the Galactic longitude. For the phase angle with respect to the equatorial plane, 62.6° should be added.

2.4 DISCUSSION

The physical and dynamical properties of dense structures depend on the history of their formation. Theoretical studies of molecular filament and dense core formation in a turbulent environment (the gravoturbulence fragmentation scenario) shows that a molecular cloud fragments and develops local density enhancements by colliding flows of turbulence and the density enhancements grow by self-gravity to evolve into filaments and dense cores. Dense structures formed following this scenario tends to be pressure-confined in their early evolution stage and proceed to a gravitationally bound state. On the other hand, dense cores formed in a quiescent, magnetized environment are typically gravitationally bound throughout their evolution since fragmentation typically takes place when a structure loses support against gravity. In addition, three dimensional morphologies and alignments of dense structures with magnetic fields also depend on the history of their formation. In a quiescent environment, magnetic fields play an important role in shaping and aligning dense structures, while the magnetic fields may not play a major role in determining the
Figure 2.13 Distribution of properties of the NH$_3$ leaves. Red histograms denote active leaves, and blue histograms inactive leaves. The first panel is the ratio of the virial mass to the observed mass. The red and blue vertical lines are the mean ratios of active and inactive leaves, respectively. The second panel is the ratio of the work done by pressure to the gravitational energy at the surface of NH$_3$ leaves. The red and blue vertical lines are the mean ratios of active and inactive leaves, respectively. The black dashed vertical line is the energy ratio of the critical Bonnor-Ebert sphere. The third panel is the distribution of the angle differences of the NH$_3$ leaves relative to their branches. The fourth panel is the axis ratio of the NH$_3$ leaves deduced from the principal axis analysis. The solid and dashed lines are the best-fit probability density of apparent axis ratio of prolate and oblate spheroids with a uniform orientation angle distribution within [-80°, 80°], respectively.
properties of dense structures in a turbulent environment. In this section, we discuss
the dynamical status, three dimensional morphologies, and statistical properties of
dense structures in the L1495-B218 filaments as well as uncertainties in our analysis.

2.4.1 Virialization

We investigate the dynamical states of NH$_3$ leaves using the virial theorem. An
equation of the virial theorem without magnetic contributions is

$$\frac{1}{2} \frac{d^2 I}{dt^2} = W + 2(T + \Pi) - \int p\vec{x} \cdot \hat{n}dS$$  \hspace{1cm} (2.3)

where $I$ is the moment of inertia, $W$ is the gravitational energy, $T$ is the kinetic
energy of systemic motions, $\Pi$ is the internal energy of thermal and nonthermal mo-
tions, and the integral term is the work done by the external pressure $p$. $S$ is the core
surface and $\hat{n}$ is the unit vector perpendicular to the core surface. We omitted the
terms related to magnetic fields because there is no measurement of magnetic fields
within NH$_3$ leaves. There are studies using the virial theorem including all terms
except magnetic fields (e.g. Hunter 1979; Seo et al. 2013) but only the gravitational
and the internal energies are conventionally used for studying dynamical stability
of a dense core in observational studies because it is hard to estimate other terms.
We also start with only the gravitational and the internal energies in estimating
the virial mass and further check whether or not the pressure term is important in
confining dense cores. We assume a time-independent steady state so that $d^2 I/dt^2$
= 0.

We assumed a NH$_3$ leaf as a homoeoidal ellipsoid with density falling as $\rho \sim r^{-\beta}$
(Bertoldi and McKee, 1992). The gravitational energy of a homoeoidal ellipsoid is

$$W = -\frac{3}{5} \frac{1 - \frac{\beta}{3}}{1 - \frac{2\beta}{5}} \frac{GM^2 \arcsin e}{R} \frac{\arcsin e}{e},$$  \hspace{1cm} (2.4)

where $e$ is the eccentricity of homoeoidal ellipsoid, and $R$ is the semi-major axis.
The internal energy is estimated from

$$\Pi = \frac{3}{2} M \sigma_v^2,$$  \hspace{1cm} (2.5)
where $\bar{\sigma}_v$ is a representative value of the velocity dispersion including thermal and nonthermal motions within a NH$_3$ leaf. It can be either the average, median, or mass-weighted values. In this study, we use median values of $\sigma_v$ within NH$_3$ leaves. Finally, the virial mass when $W = 2U$ is given as

$$M_{\text{vir}} = 5\frac{1 - \frac{2\beta}{5}}{1 - \frac{2}{3}} R\bar{\sigma}_v^2 \frac{e}{G \arcsin e}. \quad (2.6)$$

We first estimated a virial mass with $\rho \sim r^{-2} (\beta = 2)$, which is listed in Table 5.2. The eccentricities of leaves are estimated from the axis ratios of NH$_3$ leaves which are given in Table 2.3. The virial masses of NH$_3$ leaves are shown as red circles in Figure 2.10. The virial mass is calculated for only the NH$_3$ leaves because the NH$_3$ branches usually have either very elongated or irregular shapes, which makes it difficult to calculate the virial mass with conventional assumptions. We find that 30 out of 39 leaves have observed masses smaller than their corresponding virial masses (the first panel in Figure 2.13), which suggests that most of leaves are gravitationally unbound (the median ratio is about two). 7 out of 9 gravitationally bound leaves are active (red), while only two out of 22 inactive (blue) leaves are gravitationally bound. A virial mass of a homoeoidal ellipsoid with $\rho \sim r^{-1}$ is 50% larger than the corresponding virial mass of a homoeoidal ellipsoid with $\rho \sim r^{-2}$. For $\rho \sim r^{-1}$, there are only two gravitationally bound leaves, both of which are active leaves.

In order to assess the validity of the virial analysis, we must first carefully analyze the uncertainties. One uncertainty may come from a choice of $\bar{\sigma}_v$. In this study, we chose median values of $\sigma_v$ within NH$_3$ leaves but one may use a mean of $\sigma_v$ or a mass-weighted mean of $\sigma_v$ (mass weighted mean is $\bar{\sigma}_v^2 = \int N_{\text{H}_2} \mu m_h \sigma_v^2 dA / \int N_{\text{H}_2} \mu m_h dA$) within a NH$_3$ leaf. We found that the mean values of $\sigma_v$ are at most 10% larger (typically 5%) than the corresponding median values, which may result in at most 20% increase of corresponding virial mass. The mean values may be larger than the median values because $\sigma_v$ tends to be lower at the center of a NH$_3$ leaf and considerably increase toward edges of the NH$_3$ leaf. Mass weighted means on the other hand are in the range from 95% to 105% of the corresponding median $\sigma_v$ values. Thus, uncertainties originating from choices of $\bar{\sigma}_v$ do not strongly affect the
conclusion of virial mass analysis since ratios of viral masses to observed masses are quite large with respect to these uncertainties.

The largest uncertainties in estimating the observed mass from 500 µm continuum data are mainly due to assuming $T_d = T_k$ and uncertainty in dust opacity $κ$. Comparing the dust temperature, which is estimated by SED fitting of 70, 160, 250, 350, and 500 µm dust continuums (Palmeirim et al., 2013), to the gas kinetic temperature deduced from NH$_3$ (1,1) and (2,2), we found that the dust temperature is typically $3 - 5$ K higher than the gas kinetic temperature. If the dust temperature is 5 K higher than the gas kinetic temperature, the mass estimated from the dust temperature is 35% less than the mass estimated with the $T_d = T_k$ assumption and only four active NH$_3$ leaves are gravitationally bound.

Another uncertainty is from the dust-to-gas ratio. In this study, we use 0.01, which is a standard value assumed for molecular clouds (e.g. Beckwith et al., 1990; Lombardi et al., 2006; Schmalzl et al., 2010; Kirk et al., 2013b). However, detailed observations of interstellar dust grains in the diffuse ISM indicates that the dust-to-gas ratio is close to 0.0064 (Draine, 2011, and reference therein). This increases observed mass of the leaves about 55% from the current values. In this case, 20 out of 39 leaves are gravitationally bound and 19 leaves are still gravitationally unbound.

Not knowing three dimensional morphologies of leaves also brings uncertainty in estimating masses of leaves. We estimate the mass of a leaf by simply summing up the column mass of pixels that belong to the leaf, which corresponds to a cylindrical morphology rather than a spheroidal morphology. This has been a conventional method to estimate the masses of dense structures since we do not know three dimensional morphologies (e.g. Kirk et al., 2013b). In order to check the uncertainty from assuming different morphologies in estimating the mass of a dense structure, we estimated the masses of leaves by summing up the column density of pixels within leaves after subtracting the average column density of local backgrounds/branches. We found that the masses of leaves are typically $20 - 50$% less than the original values if we subtract a local background from the leaves in estimating leaf mass and that only two our of 39 leaves are gravitationally bound.
The typical uncertainty in the dust opacity is a factor of few (e.g. Shirley et al., 2011) and this results in quite a large uncertainty in estimating mass since $M \propto 1/\kappa$. We use OH5 dust opacity but this is not a direct measurement of dust opacity from molecular clouds but a simulated dust opacity based on a coagulated dust model. Shirley et al. (2011) made a direct measurement of the dust opacity at 450 $\mu$m and 850 $\mu$m in the outer envelope of a low-mass Class 0 core, B335, through a comparison between infrared extinction and submillimeter dust continuum. We compared the dust opacity of Shirley et al. (2011) to OH5 dust opacity at 450 $\mu$m and applied the same uncertainty range of dust opacity to OH5 dust opacity at 500 $\mu$m ($\kappa = 5.04^{+0.91}_{-1.10}$ cm$^2$ g$^{-1}$). The uncertainty in the observed mass ranges from 57% lower to 28% higher with respect to the observed mass given in Table 5.2. If we take the lower limit on the uncertainty in the observed mass, two active NH$_3$ leaves are gravitationally bound, while if we take the upper limit on the uncertainty in the observed mass 15 out of 39 are gravitationally bound.

Thus, uncertainties in $\sigma_v$, $T_d$, $m_d/m_g$, geometry, backgrounds, and dust opacity are unlikely to change our conclusion that most leaves are gravitationally unbound. This result suggests that a dense core may firstly form as a gravitationally unbound structure, evolve to a gravitationally bound core, and then undergo collapse to form a protostar. Indeed, 10 out of 15 bound leaves in the extreme opacity limit are in more evolved regions.

2.4.2 Pressure Confinement

External pressure may be a source to confine gravitationally unbound structure. Dib et al. (2007) carried out a detailed virial analysis for dense cores forming in a turbulent environment using a numerical method. They showed that the work done by the external pressure in the virial theorem is comparable to the other terms such as internal energy and gravitational energy (Dib et al., 2007). Since most of the NH$_3$ leaves are not gravitationally bound in our simple virial analysis, we checked whether or not work done by the external pressure is comparable to the internal and
gravitational energy. The work done by external pressure is estimated as follows
\[
\int p\vec{x} \cdot \hat{n} dS = \frac{4}{3\sqrt{1-e^2}} \min\{N_{\text{H}_2} \mu m_h \sigma_v^2\}_C A, \tag{2.7}
\]
where \(A\) is the projected area of NH\(_3\) leaf. \(N_{\text{H}_2} \mu m_h \sigma_v^2\) is the integration of pressure along the line of sight, where \(N_{\text{H}_2}\) is column density estimated from 500 \(\mu\)m dust continuum emission and \(\sigma_v\) is velocity dispersion of molecular gas with \(\mu = 2.33\). We choose the minimum value at the circumference of NH\(_3\) leaf, \(\min\{N_{\text{H}_2} \mu m_h \sigma_v^2\}_C\), because it corresponds to the minimum external pressure exerted on a NH\(_3\) leaf. The coefficient of \(4/(3\sqrt{1-e^2})\) is for spheroid geometry when one of the principal axes is aligned with the observer’s line of sight. We estimate the work done by external pressure using \(\min\{N_{\text{H}_2} \mu m_h \sigma_v^2\}_C\) because the external pressure is not well known. The estimated work done by external pressure ranges from \(5 \times 10^{40}\) erg to \(1 \times 10^{42}\) erg. In a majority of leaves, gravitational energies are larger than the corresponding work done by external pressure except for 12 leaves (the second panel of Figure 2.13), which suggests that those 12 leaves are pressure-confined. Active leaves tend to have a lower average ratio than that of inactive leaves, which indicates the gravity becomes more important than external pressure as a dense core evolves. Moreover, in 9 out of 12 leaves, the ratios of the work done by external pressure to the gravitational energy are larger than that of the critical Bonnor-Ebert sphere (black dashed vertical line). Thus, these results suggest that the youngest condensations are pressure confined.

A source of the external pressure may be investigated by comparing the internal pressure of the NH\(_3\) leaves and the pressure of the surrounding media. The mean internal pressure of NH\(_3\) leaves is \(<p/k> = 2 \times 10^6\) K cm\(^{-3}\), where \(p\) is roughly estimated as
\[
p = \frac{M \sigma_v^2}{V}, \tag{2.8}
\]
where \(V\) is the leaf volume and \(\sigma_v\) is the median value of \(\sigma_v\) within a NH\(_3\) leaf. We assumed that leaf volume is the same as the volume of a prolate spheroid having the observed axis ratio (see §4.2.3 for shapes of NH\(_3\) leaves). If we assume the density
structure of a leaf resembles the critical Bonnor-Ebert sphere, the surface pressure is about 40% of the mean internal pressure, \( < p_{\text{surface}}/k > = 8 \times 10^5 \text{ K cm}^{-3} \). This is significantly higher than the gas pressure of the interstellar medium \( < p_{\text{ISM}}/k > \approx 10^4 \text{ K cm}^{-3} \), Bertoldi and McKee 1992; Draine 2011 and still a bit higher than the interstellar turbulence ram pressure, \( < p_{\text{ram}}/k > \approx 5 \times 10^4 \text{ K cm}^{-3} \) (Lada et al., 2008).

Another source of external pressure may be the weight of the filament material outside of NH\(_3\) leaves since the self-gravity of the filaments produces a gravitational potential well and the weight of material outside of NH\(_3\) leaves exerts a confining pressure on the NH\(_3\) leaves. We estimated the pressure due to the weight of the filament material outside of the NH\(_3\) leaves as (Bertoldi and McKee, 1992; Lada et al., 2008):

\[
p_{\text{filaments}} = \frac{3\pi}{20} aG\Sigma^2,
\]

where \( \Sigma \) is the mean mass surface density of material outside of leaves, and \( a \) is a correction factor for the non-spherical geometry of the filaments. The values of \( a \) range from 1.0 to 3.3 depending on the eccentricities and density structures of the filaments (for \( a = 3.3 \), we took the maximum axis ratio of the NH\(_3\) branches, \( \max(Z/R) \sim 10 \), and a density structure of \( \rho \sim r^{-2} \)). The mean of the pressure on the leaves due to the weight of extended material is \( < p_{\text{filament}}/k > = 6 \times 10^5 \text{ K cm}^{-3} \), which is a bit smaller than the surface pressure estimated from Equation (2.8) but within a factor of two. This suggests that the weight of the filaments is one of main sources that confines gravitationally unbound leaves.

Ram pressure due to the inflow of material onto the filaments and dense cores from the molecular cloud may be another source of confining dense structures. Theoretical studies demonstrated that ram pressure due to mass accretion to dense structures may provide considerable confining pressure (Heitsch, 2013a,b) if a converging flow or an inflow is relatively isotropic (a non-isotropic converging flow may disrupt a dense structure, Dib et al. 2007). \(^{12}\text{CO}\) and \(^{13}\text{CO}\) observations of the L1495-B218 filaments show that there is a difference of 2 km s\(^{-1}\) in the LSR velocities from
the north-east to the south-west of B211 (Goldsmith et al., 2008; Palmeirim et al., 2013). This is similar to an infall velocity predicted in a similarity solution of gravitational infall onto a cylinder (Kawachi and Hanawa, 1998). If we accept that the difference in the LSR velocity is due to relatively isotropic inflow of materials from the molecular cloud to the filaments, the ram pressure due to an inflow motion is

\[ p_{\text{accretion}} = \rho_p(R)\delta V^2, \]  

(2.10)

where \( \rho_p \) is the density of the best-fit Plummer model to the filament (Palmeirim et al., 2013), and \( R \) is the outer radius of the filament which is 0.4 pc for the B211-B213 regions. \( \delta V \) is the difference in the LSR velocity across the B211 and B213 regions, which is about 2 km s\(^{-1}\). The ram pressure due to the accretion of materials from the molecular cloud to filaments are \( <p_{\text{filament}}/k> \simeq 4.5 \times 10^5 \) K cm\(^{-3}\), which is smaller than the surface pressure estimated from Equation (2.8) but still within a factor of two. Since there is projection effect in measuring the true inflow velocity, this may be a lower limit of ram pressure due to an inflow motion. Thus, ram pressure due to inflow of materials from the molecular cloud to the filaments may also be one of main sources of pressure that confine gravitationally unbound leaves.

2.4.3 Aligned Prolate Dense Structures

The formation of a filament or a dense core through the gravoturbulence fragmentation scenario predicts that a dense structure may not be strongly aligned with magnetic fields (Nakamura and Li, 2008) unless the magnetic field pressure is considerably stronger than the ram pressure of turbulence. On the other hand, formation through a quasi-static gravitational contraction may result in a strong alignment of the dense core with magnetic fields because contraction occurs more efficiently along magnetic field lines. Although magnetic field directions are not measured within the NH\(_3\) sources, magnetic field directions on filament scales (a few 0.1 pc) have been measured through dust polarization in optical and infrared wavelengths along the L1495-B218 filaments (Chapman et al. 2011 and references therein). In B7 and
B10 the magnetic fields are relatively parallel to the long axes of the filaments. For B213 to B216, the magnetic fields are perpendicular to the long axes of filaments. Dust polarization in B218 is not measured. Since we do not know the magnetic field directions within dense cores and magnetic field directions are aligned in only some regions, it is difficult to ascertain the role of magnetic field in aligning dense cores.

In the gravoturbulence fragmentation scenario, a dense core forms as an oblate spheroid and quickly proceeds to a prolate spheroid or triaxial shape (e.g. Gong and Ostriker, 2011; Kainulainen et al., 2014). The ratio of oblate dense cores to prolate dense cores depends on the Mach number of turbulence and the evolution time. A system with either a higher Mach number or a younger age shows a higher proportion of oblate dense cores. Since the NH$_3$ leaves are likely to be prolate (see §3.2.3), the L1495-B218 filaments may have been formed in a low Mach number turbulence ($M \leq 2$; Gong and Ostriker 2011) if the filaments are formed through the gravoturbulence fragmentation scenario. The velocity dispersions of filaments in the Taurus estimated from $^{13}$CO observations are indeed low Mach number dispersions that range from 0.15 km s$^{-1}$ to 0.65 km s$^{-1}$ (Panopoulou et al., 2014), corresponding to Mach numbers of 0.65 – 3.2, if we take the kinetic temperature to be 15 K (Goldsmith et al., 2008). Alternatively, magnetic fields may also make a dense core have a non-spherical shape because a dense core tends to contract along the magnetic field lines whereas a direction perpendicular to the field is supported by magnetic pressure. If magnetic pressure is the main force making a dense core non-spherical, the dominant shape is an oblate spheroid (Lizano and Shu, 1989). Since most of the NH$_3$ leaves are likely prolate spheroids, magnetic pressure does not seem to be the main cause of shaping a dense core into an elongated morphology, but gravoturbulence fragmentation may have played important roles in shaping the NH$_3$ leaves.

2.4.4 Comparing Size, Mass, and Velocity Dispersion

Larson proposed that there is an empirical relation between line width and size of a structure in a cloud scale (Larson, 1981). To see whether there is such a relation in
a sub-filament scale, we compare the mean velocity dispersion of NH$_3$ sources with the NH$_3$ source size in Figure 2.14. The velocity dispersion is almost flat on a scale from 0.005 pc to 0.1 pc (the mean velocity dispersion of molecular gas with $\mu = 2.33$ is 0.204 km s$^{-1}$), and it is close to the thermal velocity dispersion (0.188 km s$^{-1}$ at 10 K and $\mu = 2.33$). This demonstrates that nonthermal motions are significantly dissipated on scales less than 0.2 pc in the L1495-B218 filaments resulting in a breakdown in the cloud-scale supersonically turbulent scaling relation.

Observations toward molecular clouds shows that molecular clouds are supersonic turbulent (e.g. Larson, 1981; Scalo, 1984; Miesch and Bally, 1994; Padoan et al., 1999). In a smaller scale, observations toward molecular cloud cores show that dense cores are typically thermally supported and turbulence is significantly dissipated (e.g. Caselli and Myers, 1994; Goodman et al., 1998; Jijina et al., 1999; Caselli et al., 2002b). The dense NH$_3$ structures in this study also have the mean $R_p$ of 4.9 (§3.2.2). On the other hand, the $R_p$ values measured in C$^{18}$O $J = 1−0$ (Hacar et al., 2013) in the L1495-B218 filaments are close unity, and the mean $R_p$ value measured in $^{13}$CO $J = 1−0$ (Panopoulou et al., 2014) is 0.25. Since C$^{18}$O $J = 1−0$ traces less dense and larger structures than those seen by NH$_3$ (1,1) and $^{13}$CO $J = 1−0$ traces even more extended regions than those traced by C$^{18}$O $J = 1−0$, the dependency of $R_p$ values on the effective excitation density of the tracer suggests that supersonic turbulence cascades from cloud scale to core scale and is largely dissipated on a scale smaller than molecular filaments ($\leq$0.5 pc) in the Taurus.

Mass-size relations are often pursued to investigate the correlation between density structures of sources and their size. The mass and size distributions of leaves and branches in the L1495-B218 filaments are shown in Figure 2.14. Mass and size show a strong correlation of $M \sim r^{1.9}$ (red solid line for NH$_3$ leaves and branches) (power law indexes of leaves only and that of branches only are close to 2 as well). This is close to a global relation for molecular clouds where the power law index is 2 (Larson, 1981), which indicates a nearly constant mean column density. Kirk et al. (2013b) found a relation with a power law index of 2.35 in the Taurus from Herschel dust continuum observations. Their analysis includes structures larger than 1 pc,
Figure 2.14 Relations between physical properties of NH$_3$ sources. The left panel shows the relation of the velocity dispersion ($\mu=2.33$) and the size of the NH$_3$ leaves. The red dotted line denotes the thermal velocity dispersion when $T_k=10$ K and $\mu=2.33$. The right panel is the mass-size relation of the NH$_3$ sources. Black filled circles are the NH$_3$ leaves and blue boxes are the NH$_3$ branches. The red solid line is the best fit of $M \sim r^{1.9}$ to the NH$_3$ sources. The purple dashed line is $M \sim r^{1.5}$ from Kauffman et al. (2010) and the orange dashed-dotted line is $M \sim r^{2.35}$ from Kirk et al. (2013).

whereas the largest structure in this study is only 0.1 pc. If we only take structures smaller than 0.1 pc from Kirk et al. (2013b), the power law index is shallower than 2.35 and close to 2. Kauffmann et al. (2010) studied mass-size relations resolved in several nearby molecular clouds. They also found that the power law index is close to 2 in the Taurus on a size scale smaller than 0.1 pc, which agrees with our study.

A power index of 2 suggests that the dense structures in the L1495-B218 filaments are gravitationally unbound and have relatively constant column density while a gravitational bound Bonnor-Ebert-like core is expected to have a power law index closer to 1 since $\rho \sim r^{-2}$ for the outer profile of a Bonnor-Ebert-like structure. Our virial analysis also shows that most of NH$_3$ leaves in the L1495-B218 filaments are gravitationally unbound. If the NH$_3$ leaves are formed by colliding flows as shown in the gravoturbulence scenario (e.g. Gong and Ostriker, 2011), density structures are more likely to be shallower than $\rho \sim r^{-2}$. 
2.5 SUMMARY & CONCLUSIONS

In this chapter, we present extensive NH$_3$ (1,1) and (2,2) maps of the L1495-B218 filaments extending over 3 degrees on the sky in the Taurus molecular cloud with unprecedented depth (average $rms$ is 120 mK), angular resolution ($31''$), and velocity resolution (0.038 km s$^{-1}$). Using the maps, we study the physical properties of the L1495-B218 filaments by fitting 13109 NH$_3$ (1,1) and (2,2) spectra using an adaptive mesh refinement search of $\chi^2$ space. The main results are following.

1. From our ammonia observations and a protostar catalog (Rebull et al., 2010), we confirm that B211 and B216 are young, less evolved regions, while B213 and B7 are actively star-forming, older, more evolved regions. The young regions have NH$_3$ column density around $1 \times 10^{13}$ cm$^{-2}$ and do not have any protostars. On the other hand, the more evolved, older regions have NH$_3$ column density up to $1 \times 10^{14}$ cm$^{-2}$ and have multiple protostars. NH$_3$ emission is mostly confined to a single velocity component and the inversion levels are subthermally populated.

2. Gas kinetic temperatures in the L1495-B218 filaments deduced from NH$_3$ (1,1) and (2,2) lines reveal that the filaments to be very cold ($8 - 15$ K with median value equal to 9.5 K, 60% of spectra are in the range of $9.5 \pm 1$ K). We found that there is a small difference in gas kinetic temperature between the more evolved regions (B7, B213, and B218) and less evolved regions (B10, B211, and B216); with the more evolved regions having a higher median gas kinetic temperature (by at most 0.5 K) than the less evolved regions. Gas kinetic temperatures tend to decrease toward dense core centers at the level of a few Kelvins. This may ultimately affect the stability and dynamics of dense cores and filaments.

3. The nonthermal velocity dispersion of NH$_3$ (1,1) and (2,2) lines in the L1495-B218 filaments ($\sim 0.08$ km s$^{-1}$) is considerably narrower than those of C$^{18}$O $J = 1-0$ ($\sim 0.15$ km s$^{-1}$; Hacar et al. 2013) and $^{13}$CO $J = 1-0$ ($\sim 1$ km s$^{-1}$; Goldsmith et al. 2008) in the same region. This suggests that nonthermal motions in the regions traced by NH$_3$ have been significantly dissipated compared to those found in the more diffuse portions of the filaments. As a result, dense regions do not display a
Larson’s size-line width relationship.

4. Results from the CSAR clump-finding algorithm found 39 NH$_3$ peaks (leaves) and 16 nested groups (branches), which is about twice the number that Hacar et al. (2013) found in N$_2$H$^+$. This may be because NH$_3$ reveals lower density cores than does N$_2$H$^+$ and our NH$_3$ observations have 67% better spatial resolution than the N$_2$H$^+$ data of Hacar et al. (2013). The NH$_3$ leaves and branches are identified on a scale from 0.01 pc to 0.1 pc and have masses ranging from 0.05 M$_\odot$ to 9.5 M$_\odot$.

5. Masses of dense cores derived from dust continuum and NH$_3$ show a strong correlation in the L1495-B218 filaments. This suggests that NH$_3$ is an exceptional tracer for tracing dense cores in these filaments. A similar study across other molecular clouds will be interesting.

6. Most of the NH$_3$ leaves have observed masses smaller than the corresponding virial masses, which means they are likely to be gravitationally unbound structures. We found nine NH$_3$ leaves to be gravitationally bound and that 7 out of 9 leaves contain protostars or are within branches associated with star formation activity, while 30 NH$_3$ leaves are gravitational unbound and only 10 out of 30 unbound leaves contain protostars or are within branches associated with star formation. We also found that 12 out of 30 gravitationally unbound leaves are pressure-confined. This is a similar conclusion to that founds by Kainulainen et al. (2011) in larger scale clumps (>0.1 pc) in nearby clouds, while the dense structures in this study are from 0.005 pc to 0.1 pc.

7. Two sources may provide confining pressure to gravitationally unbound NH$_3$ leaves. The self-gravity of the filaments produces a gravitational potential well and the weight of filament materials outside of the NH$_3$ leaves may exert a confining pressure on the surface of the NH$_3$ leaves. The estimated mean pressure due to weight of filaments or external material is a bit smaller than the surface pressure needed to confine gravitationally unbound leaves but within factor of two. The other source may be ram pressure due to large scale inflows. $^{12}$CO and $^{13}$CO observations show that there may be inflows toward the L1495-B218 filaments. The minimum ram pressure due to the inflows is about half of the surface pressure needed to confine
gravitationally unbound leaves.

8. Most of the NH$_3$ leaves are well aligned with their lowest-level branches within the range of ±40 degrees. The distribution of the apparent axis ratios of NH$_3$ leaves is closer to that of randomly oriented prolate spheroids than that of randomly oriented oblate spheroids, which indicates most of the NH$_3$ leaves are likely to be prolate. Since magnetic pressure would shape a dense core to be an oblate spheroid, magnetic fields do not seem to play an important role in determining the observed axis ratio of NH$_3$ leaves.

9. The NH$_3$ leaves and branches each follow a mass radius relationship of $M \sim R^{-1.9\pm0.08}$.

Finally, both core-like and filamentary structures in the L1495-B218 filaments were successfully probed using NH$_3$, which suggests that NH$_3$ is an excellent tracer of $\geq 10^3$ cm$^{-3}$ gas for probing both dense cores and surrounding filaments in molecular clouds. We found that most of the NH$_3$ structures agree in the LSR velocity with the coherent filaments found by Hacar et al. (2013). We also found that most of dense NH$_3$ structures (the CSAR leaves and branches) coincide with the dynamically critical filaments ($\geq 15$ M$_\odot$ pc$^{-1}$) in Hacar et al. (2013). These results suggest that a dense core may form as a pressure-confined structure, evolve to a gravitationally bound state within critical filaments, and then undergo star formation.
We present deep CCS and HC$_7$N observations of the L1495-B218 filaments in the Taurus molecular cloud obtained using the K-band focal plane array on the 100m Green bank Telescope. The L1495-B218 filaments are a nearby, large complex extending over 8 pc and spanning various evolutionary stages. We observed the L1495-B218 filaments in CCS $N_J = 1_2 - 0_1$ and HC$_7$N $J = 21 - 20$ with a spectral resolution of 0.038 km/s and a spatial resolution of 31″. We also observed 24 out of 39 NH$_3$ cores in the infall tracers, HCN and HCO$^+$ 1–0. Strong CCS emission is found in two evolved regions, L1495A and B213E, and a young region, B216. We observed weak CCS emission in B218 and B211. No noticeable CCS emission is observed in B10 and B213W. HC$_7$N emission is only observed in L1495A and L1521D. We find that CCS and HC$_7$N intensity peaks do not coincide with NH$_3$ intensity peaks or dust continuum intensity peaks and that the the fractional abundance of CCS does not show a clear correlation with the dynamical evolutionary stage of dense cores. Our data suggest that CCS is not a good tracer to determine the dynamical evolutionary stage of dense cores. Kinematic analysis with low and intermediate density tracers suggests that L1495A may be a stellar cluster forming region with large scale inward flows and that L1521D is a slowly contracting core.

3.1 INTRODUCTION

CCS is a carbon-chain molecule which is frequently found in dense cores in nearby, low-mass star forming clouds (e.g. Suzuki et al., 1992; Wolkovitch et al., 1997; Rathborne et al., 2008; Roy et al., 2011; Marka et al., 2012; Tatematsu et al., 2014b). It is still debated which physical and chemical properties of star-forming

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1 Seo et al. in prep.
clouds are traced by CCS and how the chemical evolution of CCS fits into the scenario of star formation.

Suzuki et al. (1992) made the first CCS survey toward 49 dense cores in Taurus and Ophiuchus as well as a theoretical study on the chemical evolution of CCS with respect to the dense gas tracers NH$_3$ and N$_2$H$^+$. They found that carbon chain molecules including CCS and cyanopolyynes (HC$_X$N; X=3,5,7,...) are less abundant in star forming regions, where NH$_3$ emission is strong. Chemical modeling predicts that the abundance ratio of [CCS]/[NH$_3$] has an anti-correlation with time due to depletion (Suzuki et al., 1992; Aikawa et al., 2001). In the following years, surveys toward Bok globules and dense cores supported the results of Suzuki et al. (1992) by showing that protostellar cores have typically a lower fractional abundance of CCS compared to starless cores (Scappini et al., 1996; de Gregorio-Monsalvo et al., 2006; Foster et al., 2009; Tatematsu et al., 2014c). However, these studies only compared average values of [CCS]/[NH$_3$] between the starless and protostellar groups but have not studied a trend of [CCS]/[NH$_3$] or [CCS]/[N$_2$H$^+$] along the evolution from the starless to protostellar cores.

On the other hand, Benson et al. (1998) found four dense cores (two starless cores: L1512, L63, two protostellar cores: B5-N, and L1031A) that do not follow the chemical model of Suzuki et al. (1992) from the 20 dense cores with CCS emission in their survey of 60 dense cores in N$_2$H$^+$ and CCS. Observations toward isolated Bok globules in CCS and NH$_3$ also show that the isolated Bok globules do not have a clear anti-correlation in the column density ratio of N(CCS)/N(NH$_3$) from starless cores to protostellar cores (Marka et al., 2012). Furthermore, observations toward infrared dark clouds (IRDCs) in NH$_3$ and CCS using the GBT and the VLA show no clear evidence of a trend in N(CCS)/N(NH$_3$) from starless to protostellar clumps but rather show highly scattered values of N(CCS)/N(NH$_3$) (Dirienzo et al., 2015).

High resolution mapping observations have also show that chemical structures of some dense cores cannot be explained by the chemical models of Suzuki et al. (1992); Aikawa et al. (2001). A dense core in a static state is expected to form as a CCS-bright core, evolve to a NH$_3$ core with a CCS ring due to the central depletion...
of CCS and formation of NH$_3$, and further evolve to a NH$_3$-bright core without CCS emission in Suzuki’s model. Several evolved dense cores are found to have central depletions of CCS and ring or arc-shaped CCS emission surrounding NH$_3$/N$_2$H$^+$ cores (e.g. Kuiper et al., 1996; Williams et al., 1999; Ohashi et al., 1999; Tatematsu et al., 2014b), which qualitatively supports Suzuki’s chemical model. However, a quantitative analysis of the fractional abundances of CCS, NH$_3$, HCO$^+$, CO, etc. shows considerably inconsistencies with Suzuki’s prediction and suggests that either chemical processes in a dense core needs to be revised (Lee et al., 2003) or a dynamically evolving dense core (e.g. a core with infall motion) has considerably different chemical evolutionary path compared to that of a static core (Shematovich et al., 2003).

The CCS-bright “core” predicted by Suzuki is rarely found. In surveys toward the Taurus molecular cloud (Suzuki et al., 1992; Hirota et al., 2009), CCS is detected in 18 out of 29 cores but only L1521E is confirmed to have a CCS emission peak coinciding with the dust continuum peak (Hirota et al., 2002; Tafalla and Santiago, 2004), which indicates that the lifetime of the CCS peak is at most 0.1 Myrs (life time of a dense core is roughly 3 Myrs at the core mean density of $\bar{n} = 10^3$ cm$^{-3}$, Ward-Thompson et al. 2007). The frequency of finding CCS-bright cores is expected to be higher if Suzuki’s chemical model is correct since CCS is abundant up to $10^6$ years in Suzuki’s model. Thus, it is questionable that CCS is a good tracer with which to determine the evolutionary stages of dense cores.

Another possibility is that CCS traces gas recently condensed onto filaments and dense cores. The chemical evolution of CCS is similar to many other carbon-chain molecules (Friesen et al., 2013). Friesen et al. (2013) found that emission from cyanopolyynes (HC$_X$N; $X=5,7$) is bright but not spatially coincident with strong NH$_3$ emission in Serpens-South where large-scale inflows have been observed in a infall tracer such as HCN (Kirk et al., 2013a). They argued that cyanopolyyne emission is coming from gas newly accreted onto the filaments because it is expected to be weak due to depletion in a dense region that is old enough to have strong NH$_3$ emission. The depletion timescale of CCS is similar to that of cyanopolyynes but CCS is more abundant and typically brighter than HC$_5$N and HC$_7$N. This potentially
makes CCS a good candidate with which to trace accretion flows. In this study we investigate whether CCS emission is related with large-scale inflows onto filaments and infall motions in dense cores as well as search for CCS-bright young cores.

The L1495-B218 filaments in Taurus are good test beds in which to study physical and chemical properties of CCS in context of star formation in molecular clouds, because they are well observed from a large scale to dense core scale in multiple molecular tracers. Goldsmith et al. (2008) observed the Taurus molecular cloud including the L1495-B218 filaments in $^{12}$CO and $^{13}$CO 1-0. Their observations showed that there is a converging or shear flow toward the filaments (Narayanan et al., 2008; Palmeirim et al., 2013). Hacar et al. (2013) used maps of C$^{18}$O 1-0 emission to trace 35 velocity coherent filaments across the L1495-B218 regions. Thirty-nine dense cores within the filaments are identified and studied in NH$_3$ (Seo et al., 2015, hereafter Paper I) which have good spatial coincidences with bright dust cores. The dense cores show various evolutionary stages, including 35 starless cores and four protostellar cores (Rebull et al., 2010, 2011). The dense cores in the L1495-B218 filaments share a similar surrounding environment, making a consistent survey on the physical and chemical properties of CCS in dense cores possible.

We mapped the three-degree long L1495-B218 filaments in CCS $N_J = 1_2 - 0_1$ using the 100m Green Bank Telescope. We analyze the kinematics of CCS-bright regions and discuss what CCS emission tells us about dense core evolution. We also analyze the cyanopolyyn HC$_7$N $J = 21 - 20$ emission that was in the bandpass of previously published NH$_3$ observations in Chapter 2. The infall tracers HCN and HCO$^+$ $J = 1-0$ were observed toward 24 dense cores in the L1495-B218 filaments using the 12m radio telescope of the Arizona Radio Observatory. The layout of this chapter is as follows: §3.2 briefly introduces our mapping of the L1495-B218 filaments in CCS and the data reduction procedures. In §3.3 we present the basic physical properties of the L1495-B218 filaments seen in CCS and kinematics of three CCS-bright regions. In §3.4 we discuss what does CCS trace in a molecular cloud and physical states of CCS-bright regions by analyzing CCS with $^{12}$CO, $^{13}$CO, C$^{18}$O, NH$_3$, HCN, and HCO$^+$. Finally, in §3.5 we summarize our results.
3.2 OBSERVATION & DATA REDUCTION

3.2.1 Mapping in $N_J = 1_2 - 0_1$ and HC$_7$N $J = 21 - 20$

We observed the L1495-B218 filaments using the 100 meter Robert C. Byrd Green Bank Telescope in CCS (GBT12A-295 & GBT13A-126) in Spring of 2012 and 2013. In the first year, we mapped the filaments in NH$_3$ (1,1) & (2,2) and HC$_7$N $J = 21 - 20$ using 7 beams of K-band focal plane array (KFPA) and in CCS ($N_J = 1_2 - 0_1$) (22.344033 GHz; Yamamoto et al. 1990) using one beam of the KFPA (7+1 mode) from April 2012 to Feb 2013 (see Paper I). In the second year of observations, we mapped the L1495-B218 filaments in CCS using 7 beams. We mapped each region using position switching. Three OFF positions were selected having extinction low enough not to have CCS emission. Pointing of telescope was updated by pointing toward the quasar 0403+2600 every two hours. In order to calibrate the gain of seven KFPA beams, we observed either Venus or Jupiter in every observing shift and used the same calibration procedure described in Paper I (Seo et al., 2015) to place spectra in the $T_{mb}$ scale. We also observed the peak dust continuum position of the L1489PPC starless core ($\alpha = 04:04:47.6$, $\delta = +26:19:17.9$, J2000.0) in every observing shift (Young et al., 2004; Ford and Shirley, 2011) for a secondary check on the consistency of absolute flux calibration. The FWHM beam size of the GBT at 22 GHz is 31$''$.

The reduction of CCS data is done in the same way we did in Chapter 2. The integrated intensity map of CCS is shown in Figure 3.1 and the $rms$ of the CCS map is shown in Figure 3.2. The median noise level is 154 mK with the lowest noise level of 26 mK in B7.

The reduction of HC$_7$N is also done in the same way we did in Chapter 2. Since HC$_7$N $J = 21 - 20$ is in the same bandpass with NH$_3$ (1,1) and (2,2), the $rms$ is of the HC$_7$N map is the same with the $rms$ map shown in Figure 2 in Chapter 2.
Figure 3.1 Top: 500 $\mu$m dust continuum emission seen by the SPIRE instrument of the Herschel Space Observatory. Middle: map of integrated intensity of NH$_3$ (1,1). Bottom: map of integrated intensity of CCS $N_J = 1_2 - 0_1$
3.2.2 HCN & HCO$^+$ 1–0 Observations

We observed 24 out of 39 NH$_3$ peaks of Seo et al. (2015) using the 12m radio telescope of the Arizona Radio Observatory in HCN $J = 1$–0 (88.631847 GHz) and HCO$^+$ $J = 1$–0 (89.1885 GHz). The FWHM beam size was 66$''$. Observation were made from December 2012 to February 2013. For an accurate calibration, we observed Jupiter every one and half hour to adjust focus and pointing and to calibrate the spectra onto the T$_{mb}$ scale. A linear baseline was removed for the position switching observation. The integration for each core was 30 minutes with a typical rms of 150 mK at spectral resolution of 0.04 km s$^{-1}$. Data reduction was done using the CLASS program. For HCN 1–0, only 8 out of 24 cores only have a detection above 5$\sigma$, while we detected HCO$^+$ 1–0 in all dense cores. Only one dense core showed a strong blue infall asymmetry in both HCN and HCO$^+$ 1–0 and one other dense core showed a blue infall asymmetry in HCO$^+$, which we will discuss in section §3.4.

3.3 RESULTS

The L1495-B218 regions are a filamentary complex composed of 35 filaments (Hacar et al., 2013) with 39 NH$_3$ cores (see Seo et al., 2015). From previous sur-
veys in multiple molecular tracers and protostar surveys (Rebull et al., 2010, 2011; Hacar et al., 2013; Seo et al., 2015), the filaments show different evolutionary stages between the filaments and the Barnard regions. In terms of star forming activity, B7, B213, and B218 are more evolved regions while B10, B211, and B216 are less evolved (Hacar et al., 2013; Seo et al., 2015). In this section, we will analyze physical properties of CCS and HC$_7$N, and compare their properties based on the Barnard regions to see whether CCS is closely related with filament and dense core evolution.

3.3.1 Integrated Intensities of CCS and HC$_7$N

We present the integrated intensity of CCS $N_J = 1_2 - 0_1$ emission in Figures 3.1 and 3.3. The CCS $N_J = 1_2 - 0_1$ emission shows that CCS is bright ($>1$ K-km/s) in only three regions: north of B7/L1495 (L1495A), east of B213 (hereafter B213E), and B216. There is also weak emission in B211 and B218. No clear emission of CCS is seen in B10 and west of B213 (hereafter B213W) at our noise level (Figure 3.2). The lowest intensity of 500 $\mu$m dust continuum emission at which CCS is detected ($>3\sigma$) at is 15 MJy/sr while it is 25 MJy/sr for NH$_3$ (1,1); however, not all structures at 15 MJy/sr have CCS emission. In B10, CCS is not detected at our $rms$ level of 70 mK even though dust continuum emission spans from 250 MJy/sr to $<15$ MJy/sr. This demonstrates that CCS clearly traces lower density structures compared to NH$_3$ (1,1) but not every low density structure. CCS is bright in both more evolved (L1495A & B213E) and less evolved regions (B216), while it is weak or not detected in a couple of less evolved regions (B10 and B211).

We compare our CCS $N_J = 1_2 - 0_1$ map with an NH$_3$ (1,1) map (cyan contours, Seo et al. 2015) and a dust continuum map at 500 $\mu$m (orange contours, Palmeirim et al. 2013) in Figure 3.3. While the intensity peaks of the NH$_3$ and the 500 $\mu$m dust continuum show good spatial correlations, the peak intensities of CCS emission are typically offset from the intensity peaks of NH$_3$ and dust continuum. In L1521D and L1495A-N, the peak intensities of CCS are at the outskirts of NH$_3$ cores about 0.1 pc from the NH$_3$ peaks with CCS emission partially surrounding NH$_3$ cores. CCS also has emission in the west side of the L1521D region where there
are NH$_3$ dense cores no.30, no.31, and no.32. Among the three cores, CCS emission is brightest at the NH$_3$ core no.32, and its intensity peak has a relatively good agreement with the dust and NH$_3$ peaks within 0.05 pc (74 arcseconds). In B216, which is thought to be a less evolved region, the CCS peak is located about 0.1 pc west from the brightest NH$_3$ and dust continuum peaks and is outside of the NH$_3$ cores identified in Paper I. On the other hand, B216 has CCS emission that tends to be distributed over the 500µm intensity peaks which suggest CCS may trace a common envelope of multiple dense cores.

We compare HC$_7$N $J = 21−20$ with CCS $N_J = 1_2 − 0_1$ and NH$_3$ (1,1) emission in Figure 3.4. HC$_7$N $J = 21−20$ is only detected in L1495A and B213E at our rms level of 85 mK. HC$_7$N has a good correlation with strong CCS emission (>1 K) in both L1495A and B213E. As with the CCS emission, the HC$_7$N intensity peaks do not agree with NH$_3$ intensity peaks or the dust continuum intensity peaks. In addition, the HC$_7$N emission in B213E has an arc-like shape that surrounds an NH$_3$ dense core, which is similar with CCS emission. HC$_7$N $J = 21−20$ is not as spatially extended as CCS emission and the peak intensity of HC$_7$N is less than 0.7 K in all regions, which is much weaker than CCS intensity peaks (>2 K). Furthermore, the peak intensity positions of HC$_7$N do not exactly agree with CCS intensity peaks. This may be due to different excitation conditions ($E_{CCS}^u/k = 1.61$ K while $E_{HC7N}^u/k = 12.5$ K) and different chemical processes between CCS and HC$_7$N, even though they are both carbon-chain molecules.

3.3.2 Spectral Line Fitting

Often, CCS line profiles are not well characterized as a single Gaussian profile. Therefore, to understand the physical properties of observed CCS structures in the L1495-B218 filaments, we fit CCS lines with three different spectra models: single Gaussian, skewed Gaussian, and double Gaussian profiles. We use a least $\chi^2$ fit to
Figure 3.3 Maps of integrated CCS $^{12} - 0_1$ emission (color), integrated NH$_3$ (1,1) (cyan contour), 500 µm dust continuum (orange contour), and positions of known protostars. X-axis is galactic longitude and Y-axis is galactic latitude. NH$_3$ contours start at 1 K km/s and increase in steps of 1.5 K km/s. Dust continuum contours start at 50 MJy/sr and increase in steps of 25 MJy/sr. The names of previously-studied cores are indicated in the image.
Figure 3.4 Maps of integrated CCS $1_2 - 0_1$ emission (color), integrated NH$_3$ (1,1) (orange contour), HC$_7$N 21-20 (cyan contour), and positions of known protostars. The x-axis is galactic longitude and the y-axis is galactic latitude. NH$_3$ contours start at 1 K km/s and increase in steps of 2 K km/s. HC$_7$N contours start at 0.1 K km/s and increase in steps of 0.1 K km/s. The names of previously studied cores are indicated in the image.
the observed profile assuming the following model profiles:

\[
\phi_{\text{single}}(v) = A \exp \left( \frac{(v - v_{\text{lsr}})^2}{2\sigma^2} \right) \quad (3.1)
\]

\[
\phi_{\text{skewed}}(v) = A \exp \left( \frac{(v - v_{\text{lsr}})^2}{2\sigma^2} \right) \left[ 1 + \text{erf} \left( \frac{\alpha(v - v_{\text{lsr}})}{\sqrt{2}\sigma} \right) \right] \quad (3.2)
\]

\[
\phi_{\text{double}}(v) = A_1 \exp \left( \frac{(v - v_{\text{lsr},1})^2}{2\sigma_1^2} \right) + A_2 \exp \left( \frac{(v - v_{\text{lsr},2})^2}{2\sigma_2^2} \right) \quad (3.3)
\]

where \( v \) is the velocity, \( v_{\text{lsr}} \) is the LSR velocity of the line, \( \sigma \) is the velocity dispersion, \( A \) is the amplitude of line intensity, \( \alpha \) is the skewness, and the subscripts 1 & 2 denote velocity components in the double Gaussian model. We found the \( \chi^2 \) values are typically lower with fits using the skewed Gaussian model relative to the single Gaussian since the skewed Gaussian model has one more degree of freedom in the line fitting. We used the single Gaussian model rather than the skewed Gaussian model if \( \chi^2 \) values from both models are less than 3\( \sigma \). Figure 3.5 shows which model is used to fit CCS lines and examples of observed CCS line and best-fit models. CCS lines in all regions except part of B216 are fitted with single Gaussian profiles or skewed Gaussian profiles. In B216, we found there are multiple velocity components at 6.25 km/s and 6.5 km/s and the CCS lines in the region are better fit with double Gaussian profiles as indicated by reduced values of \( \chi^2 \).

Figure 3.6 shows the histograms of the LSR velocity in L1495A, B213E, and B216. In L1495A, the CCS velocity coincides with the \( \text{NH}_3 \) velocity but also has higher velocity components. In B216, the histograms show two velocity components at 6.25 km/s and 6.5 km/s. The \( \text{NH}_3 \) mostly overlaps with the CCS at 6.5 km/s but also has emission at a higher velocity than CCS from 6.7 km/s to 6.8 km/s, which suggest that B216 may be a complex region with multiple velocity components as similar to B211. In B213E, there are three peaks in the histogram and two of them are related with \( \text{NH}_3 \) velocity. The peak at 6.9 km/s is related with \( \text{NH}_3 \) core 32 in the west of B213E. The other two peaks at 6.4 km/s and 6.6 km/s are related with L1521D.

The velocity dispersion map deduced from CCS \( N_J = 1_2 - 0_1 \) emission is shown along with integrated \( \text{NH}_3 (1,1) \) emission (orange contours) in Figure 3.7. For B216
where there are two velocity components, we show the first component in the figure. The median velocity dispersion is 0.13 km/s. Considering the large molecular weight of CCS, the velocity dispersion is dominated by non-thermal motion (thermal dispersion is 0.038 km/s at 10 K). The median velocity dispersion deduced from CCS is similar to but slightly larger than the non-thermal component of NH$_3$ velocity dispersion (median value of 0.10 km/s, Seo et al. 2015). The velocity dispersion in the regions where CCS is more spatially extended than NH$_3$ emission is typically lower than the sound speed. The largest velocity dispersions in L1495A and B216 regions are 0.22 km/s and 0.21 km/s, respectively, which barely exceed the thermal sound speed (0.188 km/s). This indicates that NH$_3$ dense cores in L1495A as well as in B216 share common sub-thermal surroundings traced by CCS emission.

3.3.3 Kinematics

L1495A-N

Figures 3.8 and 3.9 present channel maps of CCS with respect to the integrated intensity of NH$_3$ (1,1) and 500 µm dust continuum in the L1495A region. The LSR velocities of NH$_3$ dense cores (dense cores no.1 and no.2 in Paper I) are 7.2 km/s and 7.25 km/s at their centers. The first panel in Figure 3.8 shows the blue-shifted CCS emission with respect to NH$_3$ in L1495A-N, whereas the following three panels presents red-shifted CCS emission with respect to NH$_3$ in L1495A-N. In the first panel, the blue-shifted CCS emission is bright in the northeastern and eastern parts of L1495A-N and elongated in the same direction of the long axis of L1495A-N. In the last panel, the red-shifted CCS emission is bright in the southwestern part of L1495A-N and also has elongated shape in the same direction as the blue-shifted CCS emission. The velocity change of the flow seen in CCS is three times the sound speed across a distance of 0.2 pc. These observations suggest that there is either a supersonic converging flow or a supersonic shear flow at L1495A-N from northeast to southwest direction.

To see whether the CCS flow is connected with a large scale flow, we examined
Figure 3.5 Maps of spectral fitting models (color) for CCS lines with the integrated intensity of NH$_3$ (1,1) (black contours) and examples of spectra model fitting (red lines) to observed lines (black lines). Blue, cyan, and yellow denote areas in which single, skewed, and double Gaussian models, respectively, are used to fit spectral lines. NH$_3$ contours start at 1 K km/s and increase in steps of 2 K km/s. Single Gaussian fitting is forced unless $\chi^2$ is greater than 9 in the single Gaussian fitting. In the line example, the top panel shows single Gaussian fitting in L1521D, the middle panel shows a skewed Gaussian fitting in L1495A, and the bottom panel shows double Gaussian fitting in B216.
Figure 3.6 Histograms of the LSR velocities in L1495A, B213E, and B216 (top) and CCS velocity vs. NH$_3$ velocity (bottom). Red histograms denote the LSR velocities of NH$_3$ emission, and black histograms denote the LSR velocities of CCS emission in the top figure. In the bottom figure, each point corresponds to each pixel that has both CCS and NH$_3$ emission.
Figure 3.7 Map of the velocity dispersion in L1495A, B213E, and B216 (color), and integrated intensity of NH$_3$ (1,1) (orange contours). The x-axis is galactic longitude and the y-axis is galactic latitude. NH$_3$ contours start at 1 K km/s and increase in steps of 2 K km/s. The color bar applies to all three sources.
$^{13}$CO channel maps with respect to the integrated intensity map of NH$_3$ (1,1) in Figure 3.10. The upper eight panels are zoomed-in toward L1495A and the bottom three panels show large scale $^{13}$CO channel maps. The channel maps show that the hub region may be a place where three groups of filamentary structures are colliding or converging. The large scale channel maps show that one group of filaments (Group I) stretches from the hub to B10, B211 and forms L1495-B218 filaments (6.23 km/s - 6.50 km/s). Another group of filaments (Group II) starts from the hub and extends to southwest direction (7.03 km/s - 7.29 km/s). The other group of filaments (Group III) stretches from the hub to the west direction (7.56 km/s - 7.82 km/s). In L1495A, the three groups of filaments seem to make two flows. One is a flow moving from west to east in the south region of L1495A-N (zoomed-in channel maps from 6.23 km/s to 7.29 km/s), which is related with two groups of filaments (Group I and II). The other flow is moving from north to south across L1495A-N (zoomed-in channel maps from 6.50 km/s to 8.09 km/s), which seems to be related with Group II and Group III. These 90-degree crossing flows can explain an odd "T"- or cross-shaped morphology of L1495A. In this kinematic picture, CCS seem to be connected with the flow from north to south crossing L1495A-N.

The LSR velocity maps of CCS and NH$_3$ are displayed side by side in Figure 3.11. There is a steep transition within L1495A-N in the LSR velocities of both CCS and NH$_3$ (pointed with black arrow). The transition is perpendicular to the gas flow direction of CCS and NH$_3$ and parallel to the long axis of the dense cores. The transition exhibits a velocity change of 0.3 km/s within a single beam size (0.022 pc) and a velocity gradient of $\sim$14.2 km s$^{-1}$ pc$^{-1}$, which is at least seven times larger than a typical velocity gradient measured for rotation ($\sim$2 km/s/pc, Goodman et al. 1993). This transition may be due to collision or shear of large scale flows.

The L1495A-N region is the only region to show blue asymmetric, self-absorbed profiles in both HCO$^+$ and HCN $J = 1-0$ (Figure 3.12). All hyperfine lines of HCN $J = 1-0$ show blue asymmetric line profiles. The infall velocities estimated using a simple two layer model (Equation 9 in Myers et al. 1996) are 0.41 km/s for F = 2-1, 0.17 km/s for F = 1-1, and 0.10 km/s for F = 0-1. The infall velocity measured
Figure 3.8 Channel map of CCS (contour) and integrated intensity of \( \text{NH}_3 \) (1,1) (color) in L1495. CCS contours start at 0.1 K km/s and increase in steps of 0.1 K km/s. The velocity range for each channel map is written in each panel in the unit of km/s. IDs of \( \text{NH}_3 \) dense cores are marked in the last panel.

in HCO\(^+\) 1-0 is 0.83 km/s. Among the lines, optical depths from the largest to the smallest are ordered as HCO\(^+\) 1-0, HCN 1-0 \( F = 2-1 \), HCN 1-0 \( F = 1-1 \), and HCN 1-0 \( F = 0-1 \), which means HCO\(^+\) is most sensitive motions at the outer layers of the core while HCN 1-0 \( F = 0-1 \) traces the kinematics of a more inner region of a core. These observations suggest that L1495A-N has a supersonic inflowing motion at the core outskirts and the speed of the inflow flow gradually decreases toward the center of L1495A-N.

**L1521D and \( \text{NH}_3 \) Core no.32**

Figures 3.13 and 3.14 present the channel maps of CCS with respect to \( \text{NH}_3 \) and dust structures in B213E. The \( \text{NH}_3 \) LSR velocity of L1521D is 6.75 km/s at its center. The blue-shifted CCS emission is bright and covers the outskirts of L1521D except the south side of the core (upper three panel of Figures 3.13 and 3.14). The red-shifted CCS emission is bright at the center and the west side of L1521D (the bottom three panels of Figures 3.13 and 3.14). The last two panels in Figures 3.13 and 3.14 show CCS emission near the west side of B213E where there are three \( \text{NH}_3 \) dense cores (Seo et al. 2015). It clearly shows that the CCS intensity peak coincides with the \( \text{NH}_3 \) intensity peak of the \( \text{NH}_3 \) core no.32 in the channel map.
Figure 3.9 Channel maps of CCS (contour) and 500 µm dust continuum (color) in L1495. CCS contours start at 0.1 K km/s and increase in steps of 0.1 K km/s. The velocity range for each channel map is written in each panel in the unit of km/s.

To see whether the CCS flows are connected with large scale flows, we investigated $^{12}$CO channel maps with respect to NH$_3$ and dust structures in B213E in Figure 3.15. The $^{12}$CO 1–0 emission is red-shifted in the northwest of L1521D and is blue-shifted in the south of L1521D. The CCS flows in L1521D do not have same direction with $^{12}$CO flows, thus the CCS flow is possibly a local flow at L1521D. In the west side of B213E where three NH$_3$ cores (no.30, no.31, no.32) are located, the LSR velocity of CCS changes from the southeast to the northwest direction, which is aligned with the $^{12}$CO flow.

We present the LSR velocity of B213E in Figure 3.16 in order to see whether the CCS flows are connected with NH$_3$. In L1521D, the LSR velocities in both CCS and NH$_3$ show a good agreement to each other where they have blue-shifted emission at the east-side of the core and smoothly change to red-shifted emission toward the west-side of the cores. The velocity gradient is 3 km/s/pc, which is similar to the average velocity gradient of fast rotating dense cores (Goodman et al., 1993). The LSR velocities of CCS and NH$_3$ also agree in the west of the B213E. In both NH$_3$ and CCS, the LSR velocity changes steeply from the dense core no.31 to the dense core no.32 with 13 km/s/pc, which is considerably steeper than a fast rotating core. The transition may be due to either the two cores are at different LSR velocities which overlaps at their outskirts or due to a recent collision with each others. Overall, in
Figure 3.10 Channel maps of $^{13}$CO (contours) and integrated intensity maps of NH$_3$ (1,1) (color) in the L1495-B218 filaments (top) and in L1495A (bottom). $^{13}$CO contours start at 0.5 K km/s and increase in steps of 0.13 K km/s. The velocity range for each channel is written in different color in the unit of km/s.
Figure 3.11 Maps of the NH$_3$ (top) and CCS (bottom) LSR velocities in L1495. The black arrow points to the steep transition in the LSR velocity within L1495A-N.
B213E, CCS and NH$_3$ trace the same structures, but only the west side of B213E is likely to be connected with large scale $^{12}$CO flows.

L1521D is the only other region among the L1495-B218 regions that shows an infall asymmetry. We found HCN $J=1-0$ emission is relatively weak compared to HCO$^+$ $J=1-0$ emission and could not use it to deduce an infall speed (Figure 3.12). On the other hand, HCO$^+$ $J=1-0$ has strong emission (>1 K) with a blue asymmetric profile. The estimated infall speed is 0.08 km/s, which is a bit less than half of the sound speed and substantially slower than motions measured toward L1495A-N. As we shall discuss in §4.3, these observations are consistent with a slowly contracting core.

B216

Figures 3.17 and 3.18 show the channel maps of CCS with respect to the NH$_3$ dense cores and dust structures in B216. The two NH$_3$ dense cores have LSR velocities of 6.66 km/s and 6.76 km/s. At the LSR velocities of the NH$_3$ structures, CCS has weak emission surrounding the NH$_3$ structures from 6.53 km/s to 6.78 km/s, which is a considerably narrow range of the LSR velocity compared to the velocity ranges.
Figure 3.13 Channel map of CCS (contours) and integrated intensity of NH$_3$ (1,1) (color) in B213E. CCS contours start from 0.1 K km/s and marks every 0.1 K km/s increase. The velocity range for each channel map is written in each panel and the unit is km/s. IDs of NH$_3$ dense cores are marked in the last panel.

Figure 3.14 Channel maps of CCS (contours) and 500 µm dust continuum (color) in B213E. CCS contours start at 0.1 K km/s and increase in steps of 0.1 K km/s. The velocity range for each channel map is written in each panel in the unit of km/s.
Figure 3.15 Channel maps of $^{12}$CO (color) and integrated intensity of NH$_3$ (1,1) (contours) in B213E. NH$_3$ contours start at 1 K km/s and increase in steps of 3 K km/s. The velocity range for each channel map is written in each panel in the unit of km/s.
Figure 3.16 Maps of the NH$_3$ (top) and CCS (bottom) LSR velocities in B213E.
Figure 3.17 Channel map of CCS (contours) and integrated intensity of NH$_3$ (1,1) (color) in B216. CCS contours start at 0.1 K km/s and increase in steps of 0.1 K km/s. The velocity range for each channel map is written in each panel and the unit is km/s. IDs of NH$_3$ dense cores and L1521B are marked in the last panel.

of CCS emission in B213E and L1495A ($\delta V \sim 1$ km/s). This suggest that an NH$_3$ dense core is likely to be a static and quiescent environment. The CCS emission in B216 spans from 6.10 km/s to 6.87 km/s and has two velocity components, one at 6.25 km/s and the other at 6.50 km/s (Figure 3.5). The CCS intensity peaks of both components do not coincide with NH$_3$ or dust continuum intensity peaks but instead is centered 0.08 pc in the southwest direction from the NH$_3$ peaks and dust continuum peaks. This regions (also known as L1521B) was previously observed by Hirota et al. (2004) as an example of a CCS-bright young core. Our observations clearly indicate that CCS is not tracing a dense core.
Figure 3.18 Channel maps of CCS (contours) and 500 µm dust continuum (color) in B216. CCS contours start at 0.1 K km/s and increase in steps of 0.1 K km/s. The velocity range for each channel map is written in each panel and the unit is km/s.
3.4 DISCUSSION

3.4.1 CCS as a Tracer of Recently Condensed Gas

Through our CCS survey, we found that CCS is rarely bright at the core center traced by NH$_3$ or dust continuum but rather peaks on the core outskirts or outside of cores. Only one core (no.32) out 39 NH$_3$ cores along the L1495-B218 filaments has a good agreement among the dust, NH$_3$, and CCS intensity peaks. Only two dense cores (no.32 & L1521E) in the Taurus molecular cloud have CCS and dust continuum intensity peaks that coincide. Particularly, CCS emission in B216 (L1521B), which only contains young dense cores with no star formation activity, clearly shows that CCS intensity peaks do not coincide with NH$_3$ or dust continuum peaks at an early evolutionary stages. This suggest that CCS is not a good tracer to trace core structure nor to use for a survey to find young starless cores.

We found that CCS emission is detected in both less evolved regions (B216 and B211) and in more evolved regions with star formation activity (L1495A and B213E). In particular, we found relatively strong CCS emission in L1495A-N and L1521D where NH$_3$ emission is strong and CCS is predicted to be depleted by chemical models (Suzuki et al., 1992; Aikawa et al., 2001). In these regions, NH$_3$ traces a core structure while CCS emission is bright at the outskirts of a NH$_3$ dense core with an arc-like or ring-like shape. Similar CCS geometries have been reported toward L1544 (Ohashi et al., 1999), L1498 (Hirota et al., 2004), and Orion A dense cores (Tatematsu et al., 2014b). Comparing the density concentrations among the starless cores in the L1495-B218 filaments, we found that L1495A-S, L1495A-N and L1521D are the most concentrated starless cores in the L1495-B218 filaments. Starless cores in the B213W and B10 regions, which are relatively less evolved than L1495A-S, L1495A-N, and L1521D, have no considerable CCS emission, while L1495A and L1521D have strong CCS emission. The Suzuki et al. (1992) chemical picture would have predicted strong CCS emission toward B213W and B10 and weaker emission toward L1495A and L1521D. These results coupled with the lack of CCS-peaked starless cores demonstrate that CCS is a not a good tracer to for determining the
dynamical evolutionary stage of a starless core.

One possibility of having bright CCS emission in an evolved dense core is that the core keeps accreting surrounding material. CCS forms rapidly, and this molecule becomes bright at the outskirts of the core. A similar mechanism for other carbon-chain molecules (e.g., cyanopolyyynes) has been suggested for filaments in Serpens-South (Friesen et al., 2013). Friesen et al. (2013) estimated chemical ages at HC7N-bright regions using the factional abundance ratio of [HC7N]/[NH3], since [HC7N]/[NH3] has an anti-correlation with time due to depletion of HC7N. They showed that HC7N-bright regions are considerably younger than NH3-bright regions and argued that HC7N bright regions are recently condensed gas due to accretion flows seen in HCN, HNC, and HCO+ (Kirk et al., 2013a). L1495A-N and L1521D have evidence for inward flows observed in HCN and HCO+ 1−0 (Figure 3.12), so we performed the same chemical age analysis toward L1495A-N and L1521D. The highest fractional abundance ratio of [CCS]/[NH3] in L1495A-N is 10 at the outskirts of L1495A-N. On the other hand, the lowest [CCS]/[NH3] is 0.2 at the NH3 intensity peak. Similarly, the highest [CCS]/[NH3] in L1521D is 10 at the outskirts of the core, whereas the lowest [CCS]/[NH3] is <0.05 at the NH3 intensity peak. An estimated chemical age for [CCS]/[NH3] = 10 using (Suzuki et al., 1992; Aikawa et al., 2001) is ≤0.1 Myr. For [CCS]/[NH3] = 0.2 and 0.05, estimated chemical ages are 0.7 Myr and 1.3 Myr, respectively, for gas with a density of 10^4 cm^-3 (0.5 Myr and 0.8 Myr, respectively for 10^5 cm^-3 density gas). The significantly younger chemical age for large [CCS]/[NH3] ratios suggest that the CCS-bright outskirts of L1495A-N and L1521D are the recently condensed gas at the outskirts of the cores due to inward flows rather than simple depletion in an isolated core. The exact ages will largely depend on the fidelity of chemical model and the assumptions in those models.

Our HC7N observations toward L1495A and L1521D also support the conclusion of L1495A and L1521D having accreted gas. The HC7N emission is bright at the outskirts of L1495A and L1521D, and the median fractional abundance ratio of [HC7N]/[NH3] at those outskirts are 0.05 and 0.03, respectively, with the highest [HC7N]/[NH3] = 0.16 and 0.1, respectively. These ratios are similar with the
\[ \text{[HC}_7\text{N]/[NH}_3\text{]} \text{ of the youngest region in Serpens-South (}[\text{HC}_7\text{N}/\text{NH}_3]\sim 0.03 \text{ in Clump no.2, Friesen et al. 2013)} \text{ with an estimated chemical age for those ratios being} \leq 0.1 \text{ Myrs old (McElroy et al., 2013; Friesen et al., 2013). Thus, the estimated age for HC}_7\text{N-bright gas in L1495A and L1521D is also expected to be} \leq 0.1 \text{ Myrs old. This agrees with the conclusion from the chemical age analysis using CCS.}

3.4.2 A Stellar Association/Cluster Forming Candidate: L1495A

L1495A is the hub of the L1495-B218 filaments and is the most evolved region. 17 out of 35 young stellar objects along the filaments are associated with L1495A. In addition, L1495A shows the highest column density along the filament as well as the most massive filament mass. Our NH$_3$ survey shows that L1495A is a gravitationally unstable clump with two dense cores in L1495A-S and two dense cores in L1495A-N (Seo et al., 2015). L1495A-N is particularly interesting because CCS and HC$_7$N are bright along with NH$_3$. We also observe a strong blue asymmetric infall profiles in HCN and HCO$^+$ 1-0 lines. In addition, gas flows in $^{12}$CO seems to be connected with converging flows seen in CCS. L1495A-N appears to be currently gravitationally accreting gas or is fed gas by a large scale flow.

We compared L1495A-N with theoretical models of gravitationally collapsing dense cores (Seo et al., 2013) to have an idea whether the inward flows toward L1495A-N are mainly due to gravitational accretion or to large scale colliding flows. The theoretical models calculate the minimum and the maximum infall speeds due to gravitational collapse for a given density concentration. We compared the infall speed and the dimensionless dense core radius $\xi_{\text{max}}$ of L1495A-N with the theoretical models. The $\xi_{\text{max}}$ of L1495A-N is approximately 8, which is estimated by taking the long axis of L1495A-N measured in Seo et al. (2015) and the mean density calculated by column density/long axis. The calculation shows that L1495A-N has an inward flow roughly seven times faster than the maximum infall speed that the gravitational collapse of a dense core can have with $\xi_{\text{max}}$ of 8. In addition, its morphology is the most elongated (axis ratio of 3.5:1) among the dense structures along the L1495-B218 filaments and its short axis is well aligned with the inflow direction of a large
scale flow seen in CO and CCS. This evidence suggests that L1495A-N is a dense structure likely formed by colliding large scale flow.

Since L1495A appears to be continuously fed by large scale inflows, we can estimate the star formation rate in L1495A. We use the inflow velocity of 0.8 km/s measured in HCO$^+$ 1−0. We took the lowest column density along the 5σ SNR edge of CCS emission in our map ($N = 5 \times 10^{21} \text{ cm}^{-2}$) and the long-axis of L1495A seen in CCS (0.7 pc). We assume a star formation efficiency of 30%, which is the average efficiency for dense structures in molecular clouds (André et al., 2014). Our estimate gives $16 \text{ M}_\odot/\text{Myr}$. L1495A is already associated with multiple young stellar objects and is likely to keep making stars, forming a stellar association/cluster, unless large scale inflows are disturbed by radiation pressure or outflows from protostars.

L1495A and Serpens-South share several similar characteristics such as bright carbon-chain molecules including CCS and cyaopolyynes along with bright NH$_3$ emission, association with multiple protostars, and the signatures of colliding flows and large scale accretion (Kirk et al., 2013a; Friesen et al., 2013; Nakamura et al., 2014). On the other hand, L1495A is much smaller in size, mass, and number of protostars compared to Serpens-South (L1495A: 0.7 pc, 60 M$_\odot$, 17 YSOs, Serpens-South: 10 pc, 360 M$_\odot$, 90 YSOs) (Kirk et al., 2013a; Seo et al., 2015). The velocity gradients of large scale colliding flows seen in $^{12}$CO seem to be similar between the two regions ($\sim$2 km/s/pc) (Goldsmith et al., 2008; Tanaka et al., 2013). It is no surprise that L1495A and Serpens-South have differences since their masses are considerably different. However, the similarities between L1495A and Serpens-South suggest that a stellar association or cluster may form in a relatively small space where collisions of gas flows occur.

3.4.3 A Slowly-Contracting Core: L1521D

L1521D is a dense core in an evolved region similar to L1495A. There are three Class I protostars within 0.2 pc of the core to the west. The L1521D core is gravitationally unstable based on a virial analysis (Seo et al., 2015). It shows relatively bright CCS and HC$_7$N emission as well as bright NH$_3$ emission and a signature of an infall.
motion in HCO$^+$. On the other hand, unlike L1495A-N, the CCS channel map shows a hint of a spherical flow (arc-like morphology in CCS emission) around the core and no evident connection with a large scale flow of $^{12}$CO. This suggests that the inward flow toward L1521D may be gravitational accretion rather than a large scale colliding flow.

We compared L1521D with theoretical models of gravitational collapsing dense cores (Seo et al., 2013) to see whether the inward motion in L1521D is mainly due to gravitational accretion. We compared the infall speed and the dimensionless dense core radius $\xi_{\text{max}}$ of L1521D with the theoretical models. The $\xi_{\text{max}}$ of L1521D is approximately 10, which is estimated by taking the long axis of L1521D measured in Seo et al. (2015) and the mean density calculated from column density/long axis. The calculation shows that L1521D has an inward flow slightly slower than the maximum infall speed that the gravitational collapse of a dense core can have with $\xi_{\text{max}}$ of 10. This suggests that L1521D is a collapsing core as a result of its self-gravity but is not likely to evolve to a stellar association since a large scale flow does not seem to related with the infall motion in L1521D.

L1521D shares similar characteristics with the slowly contracting starless core L1544 except that L1544 is an order of magnitude more centrally concentrated (Y. Shirley, priv. communication). The morphologies of CCS emission, the infall speed, and $\xi_{\text{max}}$ of L1521D are similar with those of L1544 (infall speed $\sim$0.1 km/s, Williams et al. 2006, $\xi_{\text{max}} \sim$10, Kirk et al. 2005). In addition, L1521D also has an extended inward motion (extending up to $\sim$0.2 pc from the core center) in a similar way with L1544 (Tafalla et al., 1998), and both are embedded within a filamentary structure. Thus, L1521D may be a dense slowly contracting core that is less evolved relative to L1544.

3.4.4 Search for the Transition-to-Coherence

The transition-to-coherence is a steep transition in the velocity dispersion from a subthermal dense core to superthermal surroundings (Goodman et al., 1998; Pineda et al., 2010a). It has been observed toward a few cores in Perseus, Tau-
rus, and Cepheus in NH$_3$, C$^{18}$O, and OH. It has been suggested as a boundary to distinguish a dense core from its surrounding. We carefully investigated any steep change in the velocity dispersion in the extended regions of CCS emission in L1495A and B216 to look for the transition-to-coherence which was not found in our NH$_3$ survey. In our CCS observations, we found that velocity dispersions of CCS lines are typically under the sound speed and do not observe any steep transitions. It is possible that the transition-to-coherence may be even further beyond the range of CCS emission, since the areas of CCS bright regions in L1495A and B216 are still smaller than the “coherent cores” found in Goodman et al. (1998); Pineda et al. (2010a).

Our findings are consistent with the physical model developed from the C$^{18}$O study by Hacar et al. (2013). The 35 velocity coherent filaments traced C$^{18}$O have typically velocity dispersion around the sound speed. We checked the nonthermal velocity dispersion of C$^{18}$O in L1495A, B213E, and B216, and it is very close to the sound speed ($<\sigma_{nt}>/c_s = 0.97, 1.15, \text{and } 1.03$, respectively). This suggest that the transition-to-coherence may exist at the edge of the C$^{18}$O traced filaments or even further beyond them. Considering that C$^{18}$O shows more extended emission than CCS in the L1495-B218 regions, it is not likely that we can observe the transition-to-coherence in our CCS map.

3.5 SUMMARY

In this chapter, we have presented CCS and HC$_7$N maps of the L1495-B218 filaments in the Taurus molecular cloud with an angular resolution of 31'' and a velocity resolution of 0.038 km/s. Using the maps, we investigated the chemistry and kinematics of the star-forming regions L1495A and B213E by comparing CCS and HC$_7$N to NH$_3$, 500 $\mu$m dust continuum, $^{12}$CO, and C$^{18}$O. The main results are as follow:

1. We carried out a complete survey in CCS and HC$_7$N along the L1495-B218 filament at noise levels of 120 mK for HC$_7$N $J = 21 - 20$ and 154 mK for CCS $N_J =1_2 - 0_1$. We found strong CCS emission (>1 K) toward two more evolved regions,
L1495A and B213E, and one less evolved region, B216. We observed weak CCS emission (<1 K) in a more evolved region, B218, and a less evolved region, B211. We found no noticeable CCS emission in B10 and B213W at our noise level. Our survey shows that CCS emission is not related with the global evolutionary stages of regions.

2. CCS intensity peaks do not agree with NH$_3$ and dust intensity peaks in most of the dense cores. Only NH$_3$ core no.32 has intensity peaks that are coincident in CCS, NH$_3$, and 500µm dust continuum emission. CCS emission in more evolved regions, L1495A and B213E, show an arc-like or a ring-like shape. L1521B (in B216), which was previously reported to be a young CCS-bright core also does not have CCS emission that coincides with NH$_3$ or dust continuum emission. These examples show that CCS rarely traces a core structure.

3. We found that CCS appears to trace recently accreted gas. CCS emission morphology and kinematics in more evolved regions agrees well with HC$_7$N emission morphology and kinematics. The chemical ages at NH$_3$ intensity peaks are relatively old (>500,000 years) compared to the chemical ages at the CCS intensity peaks (<100,000 years) using the chemical models of Suzuki et al. (1992) and Aikawa et al. (2001). For an evolved region, CCS may be a better tracer of recent accretion since it is more abundant than heavier carbon-chain molecules (e.g. HC$_X$N; x=5,7) and radiates emission from a state having lower energy at a given frequency.

4. Velocity dispersions of CCS lines range from 0.15 km/s to 0.3 km/s with the majority of CCS spectra (94%) having velocity dispersion less than 0.2 km/s. Considering the large molecular weight of CCS, this suggests that nonthermal motions in CCS bright regions are subthermal, which is the same conclusion we obtained from our NH$_3$ survey. The transition-to-coherence is not observed.

5. L1495A-N is likely a stellar association/cluster forming candidate. From the kinematics of NH$_3$, CCS, $^{12}$CO, HCN, and HCO$^+$, we found that L1495A-N has a supersonic converging flow and the flow is connected to a large scale flow (>1 pc) seen in $^{13}$CO. L1495A is already associated with multiple protostars and is likely to form more protostars due to the converging flow. This may lead the region to form
a stellar association or cluster.

6. L1521D is a slowly contracting dense core. The infall motion is not connected with a large scale flow seen in $^{12}$CO. Also, the infall speed is subthermal, which matches with a spontaneous collapse model due to self-gravity.
CHAPTER 4

STABILITY OF DENSE CORES WITH INWARD MOTIONS

In order to understand the collapse dynamics of observed low-mass starless cores, we revise the conventional stability condition of hydrostatic Bonnor-Ebert spheres to take internal motions into account. Because observed starless cores resemble Bonnor-Ebert density structures, the stability and dynamics of the starless cores are frequently analyzed by comparing to the conventional stability condition of hydrostatic Bonnor-Ebert sphere. However, starless cores are not hydrostatic but have observed internal motions. In this study, we take gaseous spheres with a homologous internal velocity field and derive stability conditions of the spheres utilizing a virial analysis. We propose two limiting models of spontaneous gravitational collapse: the collapse of marginally unstable Bonnor-Ebert spheres, and the collapse of critical uniform spheres. The collapse of these two limiting models are intended to provide the lower and the upper limits, respectively, of the infall speeds at a given density structure. The results of our study suggest that the stability condition depends sensitively on internal motions. A homologous inward motion with a transonic speed can reduce the critical size compared to the static Bonnor-Ebert sphere by more than a factor of two. As an application of the two limiting models of spontaneous gravitational collapse, we compare the density structures and infall speeds of the observed starless cores L63, L1544, L1689B, and L694-2 to the two limiting models. L1689B and L694-2 seem to have been perturbed, resulting in faster infall motions than for spontaneous gravitational collapse.

1Seo, Hong, & Shirley 2013, ApJ, 769, 50
4.1 INTRODUCTION

Low-mass star formation is a fundamental process in astrophysics because the majority of stars in our Milky Way are low-mass stars. Recent observations at (sub)millimeter and IR wavelengths show that starless cores are the initial stage of low-mass star formation (Myers and Benson, 1983; Benson and Myers, 1989; Ward-Thompson et al., 1994). Isolated low-mass starless cores represent the simplest environment of star formation compared to high-mass star forming region (e.g. Orion, Cygnus, Carina, etc.) because they are in less crowded regions with less contamination from strong radiation fields of neighboring stars and from feedbacks of stellar outflows. A starless core usually forms a single star or a few stars at most (e.g. Lada, 2006). Starless cores are therefore ideal objects in which to research the basic physics of star formation.

Star formation involves a series of complicated dynamical, chemical, and radiative processes. In one limit of theoretical studies, starless cores are formed through gravitational fragmentation and undergo hydrostatic collapse when they become gravitationally critical through loosing their thermal or magnetic pressure support (e.g. Mouschovias and Ciolek, 1999). In this limit, gravitational contraction is the dominant mechanism during collapse. Another limit of theoretical studies, which is still developing, proposes that turbulence is comparably important to self-gravity in forming stars (e.g. Mac Low and Klessen, 2004; Kudoh and Basu, 2008; Nakamura and Li, 2008; Gong and Ostriker, 2009). In this limit, colliding flows in supersonic turbulence create or disrupt starless cores, and some fraction of starless cores made from colliding flows may undergo collapse. The main differences between this limit and the former one are that the speed of colliding flows may be much faster than the speed of gravitational contraction and that the ram pressure of colliding flows may considerably shorten the time to form stars. Thus, studies of the internal velocity fields of starless cores are necessary to determine which limit is dominant.

The internal velocity fields of starless cores are probed using dense gas tracers such as HCO$^+$, HCN, CS, etc. (e.g. Lee et al., 1999, 2004; Gregersen et al., 2000;
Evans et al., 2005; Sohn et al., 2007; Stahler and Yen, 2009). Transitions of these tracers are usually optically thick lines which have an asymmetric profile compared to optically thin symmetric lines when there are internal motions. Surveys of the skewness of CS 2–1, CS 3–2, and HCN 1–0 lines were done by Lee et al. (1999, 2004), and Sohn et al. (2007). Their studies show that 25%, 27%, and 20% of their samples, respectively, have blue-skewed line profiles indicative of infall. In addition, studies of the internal velocity fields of L694-2 and L1197 with radiative transfer models and dynamical simulations (Lee et al., 2007; Seo et al., 2011; Gong and Ostriker, 2011) suggest that both starless cores have supersonic infall motions substantially faster and occurring in a narrower layer than predicted by quasi-static gravitational collapse. Such fast infall motions may appear in a scenario of colliding shocks where the turbulent surrounding may enhance infall speed in a narrow layer (Gómez et al., 2007) and considerably affect stability of the starless cores because the ram pressure $P_{\text{ram}} \sim \rho v^2$ becomes comparable to the surface pressure $P_{\text{surf}} \sim \rho c_s^2$. These studies demonstrate that internal motions of some starless cores may not be owing completely to gravitational contraction, but may also be under considerable influences from their surrounding environments. Their internal motions must be taken into account when testing their stability and understanding their internal dynamics.

Dynamical evolution of starless cores in the limit of negligible external dynamics (gravity dominant limit) are relatively well studied utilizing similarity solutions and numerical simulations (e.g. Larson, 1969; Penston, 1969; Shu, 1977; Hunter, 1977; Shu, 1977; Foster and Chevalier, 1993; Aikawa et al., 2005; Seo et al., 2011). The fiducial model in this limit is the collapse of a self-gravitating, isothermal, hydrostatic gas sphere bound by external pressure called the Bonnor-Ebert sphere (Ebert, 1955; Bonnor, 1956; Ebert, 1957, hereafter BE). This model has been widely used because it is a physically well established model with its column density profile similar to observed column densities in near-IR absorptions and dust continuum emissions (e.g. Alves et al., 2001; Evans et al., 2001; Tafalla et al., 2002; Harvey et al., 2003b; Kandori et al., 2005). In this model, the critical BE sphere collapses while maintaining a BE-like density profile (Foster and Chevalier, 1993; Aikawa et al., 2005;
Seo et al., 2011). However, the initial BE sphere is a hydrostatic sphere neglecting any internal motion, whereas real starless cores are likely contracting, expanding, or oscillating (e.g. Lee et al., 1999, 2004; Keto et al., 2006). The dynamics of starless cores in the limit of external dynamical influences (turbulence regulating limit) are also studied through simulations (e.g. Kudoh and Basu, 2008; Nakamura and Li, 2008), but do not have a simple fiducial model because of the chaotic nature of turbulence. Since only the first limit has a fiducial model, it is hard to compare observed starless cores to the simulations and interpret their dynamics with respect to the two limits. A dynamics model is required that provides a boundary between the two limits.

The main purpose of this work is to take an internal velocity field into account in the determination of the dynamical stability condition for starless cores. Our stability condition serves as a more realistic stability criterion for starless cores with internal motions. We also suggest a range of infall velocities at a given density structure that may classify observed starless cores with respect to the two dynamical evolution limits. The comparison of observed infall speeds of starless cores to this range will tell us whether the collapse of starless cores mostly depend on its internal properties or evolve together with its turbulent surroundings.

In this paper, first, we derive a new dynamical stability condition of the BE spheres with various dimensionless size and homologous infall motions. Hunter (1979) studies stability of a uniform field with homologous infall motions using the virial treatment. We apply Hunter’s analysis to the BE spheres and derive critical sizes of BE spheres for a given homologous infall motion. We define \( \xi_{\text{max}} \) as the dimensionless size

\[
\xi_{\text{max}} \equiv \frac{R_s(t) \sqrt{4\pi G \rho_c(t)}}{c_s},
\]

where \( t \) is time, and \( G \) is the gravitational constant, \( R_s \) and \( \rho_c \) are the outer radius and the central density of the core. We calculate dynamical domains in a two dimensional \( \xi_{\text{max}} \) vs. peak infall velocity space. The critical size \( \xi_{\text{crit}} \) for a given infall speed separates the stable and unstable regimes in the \( \xi_{\text{max}} \) vs. peak infall
velocity space.

Second, we search a range of infall velocities for isothermal spheres without any initial internal motion or any external perturbation. This is to find dynamical domain of the pure gravitational collapse in the $\xi_{\text{max}}$ vs. peak infall velocity space. Without any initial internal motion and external perturbation, the dynamics of the starless core depend only on the density distribution and there are two limits of the density distributions for starless cores: the BE sphere and a uniform sphere. The collapse of the critical BE sphere gives the slowest infall with respect to the density concentration (Seo et al., 2011) and provides the lower limit of the infall speed in spontaneous gravitational collapse (or pure gravitational collapse). A starless core that is collapsing slower than this model is likely to have been perturbed by external perturbations or is not isothermal. The uniform sphere is a starless core with the least pressure gradient against the self-gravity for its initial conditions. Inevitably, the sphere develops a faster infall motion than the collapse of the critical BE sphere at a given density structure (Seo et al., 2011). Since the uniform density field has the shallowest profile that a starless core can have, the collapse of the critical uniform sphere has the fastest infall speed with respect to a given density concentration (in the absence of external perturbations). Any starless core collapsing faster than the collapse of the critical uniform sphere is likely to have been subjected to converging flows, or have been perturbed to collapse faster than any pure gravitational collapse by an external source.

As an application of our study, four starless cores (L1544, L1689 B, L63, and L694-2) are compared with the new stability condition and the domains of two dynamical limits (gravity dominant and turbulence regulating limits). The four starless cores are the best four infall candidates with studies of both density structures and internal motions in dust continuum or extinction and molecular tracers of infall (e.g. HCN, CS, and HCO$^+$). In order to normalize the physical quantities of the four cores into dimensionless quantities of this study, we need to measure gas temperatures of the starless cores. We observed the four starless cores in NH$_3$ (1,1) and (2,2) inversion lines with the Robert C. Byrd Green Bank Telescope and
deduced their gas kinetic temperature. Their collapse dynamics are interpreted with the results of our dynamical calculations.

The layout of this chapter is as follows: §4.2 briefly introduces our observations of NH$_3$ lines measuring temperatures of the four starless cores. In §4.3 we introduce the result of Hunter (1979) and elaborate our application of his method to an arbitrary density distribution. In §4.4 we outline the two limiting models of the spontaneous gravitational collapse, and discuss the physical meaning of the regimes divided by the two models in the $\xi_{\text{max}}$ vs. peak infall speed space. In §4.5 we discuss limitations of our analysis and results. We also compare observations of the four starless cores with our stability conditions and discuss their internal dynamics. Finally, in §6, we will summarize our results and discussion.

4.2 GBT OBSERVATIONS

Observations of the NH$_3$ (1,1) and (2,2) inversion transitions were performed over three observing shifts in 2006 (GBT06A68) with the Robert C. Byrd Green Bank Telescope$^2$. Each observation consisted of a 4 minute frequency switched observation with a 4.11 MHz throw at 4 Hz centered on the 850 $\mu$m continuum peak position (see Table 1). The GBT spectrometer was set up with 50 MHz bandwidth and 6.1 kHz spectral resolution (0.08 km/s resolution). The spectra were folded and baselined using standard GBTIDL routines. The main beam efficiency was determined from observations of quasars (3C48 and 3C286) and planets to be $\eta_{mb} = 0.74 \pm 0.05$. In Section 4.5 we analyze the NH$_3$ observations for the four starless cores with measured infall motions (L1544, L1689B, L63, and L694-2).

4.3 GRAVITATIONALLY STABLE OR UNSTABLE: VIRIAL TREATMENT

To facilitate the virial analysis for a starless core, we assume the core is an isothermal gas sphere bounded by external pressure. We also assume that all internal
motions are spherically symmetric. The virial theorem for an isothermal medium with spherical symmetry is given by

\[ \frac{1}{2} \frac{d^2}{dt^2} \int_0^{R_s} \rho(r)r^2 \cdot 4\pi r^2 dr = \int_0^{R_s} 4\pi r^2 \rho(r)r^2 dr + 3 \int_0^{R_s} 4\pi r^2 \rho(r)c_s^2 dr - \int_0^{R_s} 4\pi r^2 \rho(r) \frac{G}{r} \int_0^r 4\pi r'^2 \rho(r')dr' dr - 4\pi R_s^3 c_s^2 \rho(R_s), \] (4.2)

where \( R_s \) is the radius of the core, and \( c_s \) is the sound speed. The above equation may be normalized with the following normalizing parameters: The sound speed in the core is

\[ c_{s,c} = 0.188 \, \text{km s}^{-1} \left( \frac{T}{10 \, \text{K}} \right)^{1/2} + \sigma_{nt}, \] (4.3)

where \( T \) is the temperature of the core, and \( \sigma_{nt} \) is the rms speed of non-thermal components. The normalizing parameter for length \( \alpha \) is defined as

\[ \alpha \equiv \frac{c_s}{\sqrt{4\pi G \rho_c}} = 0.044 \, \text{pc} \left( \frac{c_s}{0.188 \, \text{km s}^{-1}} \right) \left( \frac{10^4 \, \text{cm}^{-3}}{n_c} \right)^{1/2} \left( \frac{2.33}{\mu} \right)^{1/2}, \] (4.4)

where \( n_c \) is the number density of molecules at the core center, and \( \mu \) is the mean molecular weight. For the normalization of time, we take the sound crossing time \( t_0 \), which is the time taken by the sound wave to travel a unit length,

\[ t_0 \equiv \frac{1}{\sqrt{4\pi G \rho_c}} = 0.23 \times 10^6 \, \text{yr} \left( \frac{10^4 \, \text{cm}^{-3}}{n_c} \right)^{1/2} \left( \frac{2.33}{\mu} \right)^{1/2}. \] (4.5)

The normalized form of equation (4.2) is

\[ \frac{d^2}{d\tau^2} \int_0^{\xi_s} s(\xi)\xi^4 d\xi = 2 \int_0^{\xi_s} s(\xi)\xi^2 d\xi + 6 \int_0^{\xi_s} s(\xi)\xi^2 d\xi - \frac{2}{3} \int_0^{\xi_s} \xi s(\xi)\bar{M}(\xi)d\xi - 2s(\xi_s)\xi_s^3, \] (4.6)

where \( s \equiv \rho/\rho_c(t=0), \xi \equiv r/\alpha, \xi_s \equiv R_s/\alpha, \bar{M}(\xi) = \int_0^\xi \xi^2 s(\xi)d\xi \) and \( \tau \equiv t/t_0 \).

4.3.1 A Uniform Density Core with a Homologous Motion

The dynamical stability condition of a uniform density field with homologous infall motions is derived by Hunter (1979). He calculated how fast the converging motion
should be in order to make a stable core into a critical configuration after contraction. In this chapter, we show his results and briefly discuss its meaning in $\xi_{\text{max}}$ vs. peak infall speed space. For detailed derivation, please refer to his paper.

The relation between the initial speed of homologous motion and the ratio of initial and final size of the core (equation (4) in Hunter’s paper) in terms of our dimensional variables is as follow:

$$U_i^2 = 10 \left[ \ln R - \left( \frac{4}{3} - \frac{1}{R} \right) \left( 1 - \frac{1}{R^3} \right) \right],$$

(4.7)

where $U_i$ is the initial peak velocity of homologous motion divided by the sound speed. $R$ is the compression ratio defined as $\xi_{\text{si}}/\xi_{\text{sf}}$ where $\xi_{\text{si}}$ is the core radius at the initial moment and $\xi_{\text{sf}}$ is the core radius at the final moment. When there is no external pressure, equation (4.7) becomes

$$U_i^2 = 10 \left( \ln R + \frac{1}{R} - 1 \right).$$

(4.8)

The initial velocities $U_i$ are plotted in black lines as functions of the compression ratio $R$ in Figure 4.1. The solution with $U_i^2 < 0$ is not a true solution because the velocity becomes imaginary. The solution with $U_i^2 > 0$ and $R > 1$ is for a collapsing core, while the one of $R < 1$ is for a expanding core. The solutions for the cores bounded by the external pressure $P_{\text{ext}} = c_s^2 s(\xi_{\text{si}})$ and $P_{\text{ext}} = 0$ are marked with the solid and dashed lines, respectively. The solution of $P_{\text{ext}} = c_s^2 s(\xi_{\text{si}})$ has imaginary values from $R = 1$ to 2.25. This is because we assume that the density structure is kept as a uniform field during the contraction, which is unrealistic. Because of this assumption, the work done by the external pressure during the contraction cannot be correctly estimated. On the other hand, the solution without the external pressure has no imaginary values since the stability depends only on the gravitational potential energy and the kinetic energy.

Since the core mass is fixed and we assume that the uniform density structure is maintained, the radius and the central density of the core after contraction are $R_e/R$ and $\rho_c(t = t_f) = R^3 \rho_c(t = 0)$, where $t_f$ is the time epoch at the completion of contraction by inward motions. So, the dimensionless size $\xi_{\text{max}}$ becomes $\xi_{\text{max}}(t =$
\( t_f \) = \( \xi_{\text{max}}(t = 0) R^{1/2} \). The final dimensionless size \( \xi_{\text{max}}(t = t_f) \) of the core should be the same as the critical size of the static uniform density core. The critical size of the pressure-bound, static uniform density core is \( \xi_{\text{max}} = 2.25 \), which is numerically estimated using a Godunov-type hydrodynamics code within an error of 5% (Seo et al. 2011a, 2011b). For a given inward velocity \( U_i \) and the final critical size \( \xi_{\text{max}}(t = t_f) \), the compression ratio \( R \) is determined, and we can derive the corresponding \( \xi_{\text{max}}(t = 0) \) using equations (4.7) and (4.8). The stability diagram for uniform density cores with inward homologous motions is presented in Figure 4.2. The black solid line represents the core confined by the external pressure, and the black dashed line represents core with no external pressure. The regime on the left side of the black dashed line is where a core becomes gravitationally stable, while the one on the right side of the black solid line is the unstable domain for the critical uniform density core. The stability in the regime between the two lines depends on the external pressure.

4.3.2 The Bonnor-Ebert Sphere with Homologous Motion

Observed starless cores are not uniform density cores, but are reported to have BE-like density profiles. If the density is not constant, but a function of the radial distance, the moment of inertia \( I \) and the gravitational energy \( W \) take forms of \( C_0(\xi_s)M(\xi_s)\xi_s^2 \) and \( C_1(\xi_s)\tilde{M}(\xi_s)/\xi_s \), respectively, where \( C_0(\xi_s) \) and \( C_1(\xi_s) \) are functions of the radial distance. If the \( C_0 \) and \( C_1 \) do not sensitively depend on the radial distance and time during collapse, we may assume them to be constants. The \( C_0 \) value decreases from 0.2 to 0.14 for \( \xi_{\text{max}} \) from 0 to 7 and becomes almost constant for \( \xi_{\text{max}} \geq 10 \). Since the total variation of \( C_0 \) is about 30% for smaller cores of \( \xi_{\text{max}} < 6.5 \) while the observational uncertainty in \( \xi_{\text{max}} \) is at least a factor of two (Kandori et al., 2005), then assuming \( C_0 \) as a constant is a reasonable assumption unless the core is considerably disturbed and deviates from a BE-like structure (quasi-equilibrium contraction). Likewise, the \( C_1 \) value varies from 0.6 to 1 (40%), and thus we may also assume it is approximately a constant. In this study, we take \( C_0 \) and \( C_1 \) as \( C_0 = C_0(\xi_{\text{si}}) \) and \( C_1 = C_1(\xi_{\text{si}}) \), respectively.
The virial theorem for the BE sphere is given by
\[ \frac{d^2}{dt^2} C_0 \tilde{M}(\xi_s) \xi_s^2 = 2 \xi_s^2 C_0 \tilde{M}(\xi_s) + 2 \tilde{M}(\xi_s) - 2 C_1 \frac{\tilde{M}^2(\xi_s)}{\xi_s} - 2 \xi_s^3 s(\xi_s). \] (4.9)
We assume that the core does not gain any additional mass during contraction so the core mass \( \tilde{M}_0 \) is a constant \( \tilde{M}_0 = \tilde{M}(\xi_{si}) \). Assuming that the inward motion is a homologous motion, we have
\[ \ddot{\xi}_s = \frac{1}{C_0} \frac{1}{\xi_s} - \frac{C_1 \tilde{M}_0}{C_0 \xi_s^2} - \frac{\xi_s^2 s(\xi_s)}{C_0 \tilde{M}_0}. \] (4.10)
We integrate the above equation from \( \xi_{sf} \) to \( \xi_{si} \) and assume that the external pressure is constant \( s(\xi_{si}) = s(\xi_{sf}) \). Then,
\[ U_i^2 = U_f^2 + \frac{2}{C_0} \ln \left( \frac{\xi_{si}}{\xi_{sf}} \right) + 2 \frac{C_1 \tilde{M}_0}{C_0 \xi_s} \left( \frac{\xi_{sf}}{\xi_{si}} - 1 \right) - \frac{2}{3 C_0 C_2} s(\xi_{si}) \left( \xi_{si}^3 - \xi_{sf}^3 \right), \] (4.11)
where \( U_f \) is the velocity of the homologous motion at the final moment. Since the total mass of the core is fixed, we may write the core mass \( \tilde{M}_0 = C_2 (\xi_{si}) \xi_{si}^3 \). Equation (4.11) becomes
\[ U_i^2 = U_f^2 + \frac{2}{C_0} \ln \left( \frac{\xi_{si}}{\xi_{sf}} \right) + 2 \frac{C_1 \tilde{M}_0}{C_0 \xi_s} \left( \frac{\xi_{sf}}{\xi_{si}} - 1 \right) - \frac{2}{3 C_0 C_2} s(\xi_{si}) \left( 1 - \frac{\xi_{sf}^3}{\xi_{si}^3} \right). \] (4.12)
We assume that the kinetic energy of the homologous motion is expended at the final moment \( (U_f = 0) \) and the core becomes relaxed \( (\ddot{\xi}_{sf} = 0) \). Equation (4.10) at the final moment is
\[ \frac{1}{C_0} - \frac{s(\xi_{si})}{C_0 C_2} \frac{\xi_{sf}^3}{\xi_{si}^3} = \frac{C_1 \tilde{M}_0}{C_0}. \] (4.13)
Putting the above result into the equation (4.12), we get
\[ U_i^2 = \frac{2}{C_0} \left[ \ln R + \left( 1 - \frac{s(\xi_{si})}{C_2} \frac{1}{R^3} \right) \left( \frac{1}{R} - 1 \right) - \frac{s(\xi_{si})}{3 C_2} \left( 1 - \frac{1}{R^3} \right) \right]. \] (4.14)
When there is no external pressure, the above equation becomes
\[ U_i^2 = \frac{2}{C_0} \left( \ln R + \frac{1}{R} - 1 \right). \] (4.15)

In Figure 4.1 the square of initial velocity of equations (4.14) and (4.15) are plotted in blue lines as a function of the compression ratio \( R \). The solid blue line
represents the solution of $P_{\text{ext}} = c_s^2 s(\xi_{si})$, and the dashed blue line represents the solution without the external pressure. Only the solutions at $U_i^2 > 0$ and $R > 1$ correspond to collapsing cores. In this figure, the $C_0$ is fixed to be $C_0(\xi_{\text{max}} = 6.5)$. The solution of $P_{\text{ext}} = c_s^2 s(\xi_{si})$ also has imaginary velocity at the range of $R$ from 1 to 1.5 because the assumption of keeping a BE-like structure during the contraction results in incorrectly estimating the work done by the external pressure.

The critical size of BE cores with inward homologous motions can be calculated as follows: The final radius of the core after contraction by the inward motion, $R_{sf}$, is given by

$$R_{sf} = \alpha \xi_{sf} = \alpha_f \xi_{\text{max},f},$$

(4.16)

where $\alpha_f$ and $\xi_{\text{max},f}$ are the newly defined normalizing parameter and dimensionless size of the core, respectively, after the contraction. Since $\alpha \sim \rho^{-1/2}_c$, equation (4.16) may be reduced to

$$\xi_{\text{max},f} = \left(\frac{\rho_{cf}}{\rho_{ci}}\right)^{1/2} \frac{\xi_{si}}{R},$$

(4.17)

where the $\rho_{ci}$ and $\rho_{cf}$ are the central densities at the initial moment and after the contraction, respectively. The relation between the central densities is given by the conservation of the total mass of the core,

$$4\pi \alpha_i^2 \xi_{si}^2 \frac{d\psi}{d\xi} \bigg|_{\xi_{si}} = 4\pi \alpha_f^2 \xi_{\text{max},f}^2 \frac{d\psi}{d\xi} \bigg|_{\xi_{\text{max},f}},$$

(4.18)

where $\psi \equiv -\ln s$. Reducing the above equation and combining with the equation (4.17), we have

$$R = \frac{\xi_{\text{max},f} \frac{d\psi}{d\xi} \bigg|_{\xi_{\text{max},f}}}{\xi_{si} \frac{d\psi}{d\xi} \bigg|_{\xi_{si}}},$$

(4.19)

To be a critical core, the final size of the core $\xi_{\text{max},f}$ should be equal to the critical size of the static BE sphere, $\xi_{\text{max},f} = 6.5$. If the initial size is given, the compression ratio $R$ can be estimated for a critical core using the above equation. The corresponding
peak velocity of the homologous motion can be calculated from the equation (4.15). Thus, we may relate an internal motion that makes a core to be critical with an initial size of a BE sphere.

The critical sizes of BE cores with homologous internal motions are plotted in Figure 4.2. The blue solid and dashed lines are the critical sizes of the BE cores with the external pressure of $P_{\text{ext}} = c_s^2 s(\xi_{\text{si}})$ and without external pressure, respectively. Because of the imaginary velocity in the solution of $P_{\text{ext}} = c_s^2 s(\xi_{\text{si}})$, the $\xi_{\text{max}}$ value of the blue solid line starts from much smaller values than the critical size of the static BE core. The choice of $C_0$ and $C_1$ changes the starting points of lines for pressure bound cores, but not by more than factor of two. The critical sizes of smaller BE cores converge to those of uniform density cores because smaller BE cores have flatter density profiles. The critical size is very sensitive to internal motions. A homologous inward motion with a transonic speed can reduce the critical sizes of uniform density cores and BE spheres by more than half.
Figure 4.1 The relationship between the speed of inward homologous motion and the compression ratio \( R \) for a uniform core (black) and a Bonnor-Ebert core (blue). The pressure confined cores (solid lines) are easier to compress than the cores in a vacuum (dashed lines). Stronger compression is produced by a faster inward motion.

4.4 SPONTANEOUS COLLAPSE OF STARLESS CORES

The collapse of the critical BE spheres and uniform density spheres provide the lower and upper limits, respectively, of infall speeds in the spontaneous gravitational collapse of isothermal cores. Figures 3 and 4 of Seo et al. (2011) show density and velocity distributions of the collapse of the critical BE and uniform density spheres, respectively. Their density structures resemble each other at the later stage of evolution, but infall motions of the collapse of the critical uniform density sphere is always faster than the collapse of the critical BE sphere.

From the numerical simulations of the collapse of the critical BE and uniform density spheres, we measure the peak infall velocities of the two models as a function of \( \xi_{\text{max}} \). The results are plotted in Figure 4.3. The red solid line represents the collapse of the critical BE sphere, and the red dashed line represents the collapse of the critical uniform density sphere. The stability lines are also plotted in the same figure with the same colors and line styles used in Figure 4.2. The collapse of the
Figure 4.2 Critical size of the uniform core (black) and the BE core (blue) with respect to the peak speed of inward motion.
critical BE sphere shows a steady and slow growth of infall speed along with $\xi_{\text{max}}$. On the other hand, the collapse of the critical uniform density sphere results in a sudden increase of infall speed without a considerable density concentration because there is no pressure gradient at the initial moment. Owing to the sudden infall motion at the early evolution stage, the core bounces momentarily and generates a strong accretion shock, which travels outward and regulates the growth rate of infall motion. The accretion shock is shown as a kink or cusp in the velocity fields in Figures 3 and 4 of Seo et al. (2011b). The kinks in Figure 4.3 at $\xi_{\text{max}} = 8$ and 30 are when the accretion shock passes through the peak infall layer and when it escapes the core, respectively.

Since the collapse of the critical BE sphere provides the lower infall speed limit of the spontaneous gravitational collapse, the regime above the red solid line implies that the gravitational collapse is hindered by an external perturbation. The regime below the collapse of the critical uniform density sphere is for gravitational collapse enhanced by an external perturbation. For a starless core that evolves between the two lines, it is hard to tell which factors, either initial density distribution or external perturbation or both, determines its dynamical evolution. Both factors may be equally responsible.
4.5 DISCUSSIONS

4.5.1 Limitations of Our Analysis

The stability of isothermal gas spheres with homologous motions is derived utilizing the virial theorem. We impose homologous internal motions because they facilitate an analytic analysis. However, there is no observational evidence that starless cores have homologous internal motions when they are collapsing or perturbed. Moreover, a starless core does not maintain a homologous velocity field during its collapse even when a homologous velocity field is given as an initial perturbation. Nevertheless, our virial analysis provides a rough estimate of how much the conventional stability condition of BE sphere varies when there are inward velocity fields. Results in this paper suggest that the stability condition for starless cores is very sensitive to internal motions. A transonic homologous inward motion can reduce the critical size of the static uniform density and BE cores by about 67% and 54%, respectively.

Spherically symmetric starless cores are assumed in this study. For a spheroidal core in a hydrostatic state, its critical size is the same as the critical size of the BE sphere as long as its projected area is the same as that of the BE sphere (Lombardi and Bertin, 2001). Since the virial theorem is essentially an energy analysis, a stability analysis for spheroidal cores would be the same as this study except for the values of \( C_0 \) and \( C_1 \). Infall motions of the non-spherical cores were studied by Myers (2005) whom showed that the peak infall speeds are about 50% faster for a cylindrical core and about twice faster for an infinitely extended slab. Observed starless cores usually have \( \leq 2:1 \) axis ratios (Myers et al., 1991b), which is geometrically closer to the sphere than the cylinder. So, the evolution lines of the two models in Figure 4.3 may have slightly faster infall speeds if the cores are spheroids, but the difference should be quite a bit less than 50% of the plotted value.

Isothermality is another assumption in this study. Starless cores are observed to have temperature profiles decreasing inward because the interstellar radiation field heats the outskirts of the cores (e.g. Evans et al., 2001; Zucconi et al., 2001; Ward-Thompson et al., 2002; Pagani et al., 2004; Shirley et al., 2005; Crapsi et al.,
Figure 4.3 The collapse of the marginally unstable uniform density (red dashed line) and BE (red solid line) cores, and the critical sizes of the uniform density (black lines) and BE (blue lines) cores with respect to the peak speed of inward motion. The collapse of the BE core delivers the slowest infall motions, while the collapse of the uniform density core results in the fastest gravitational collapse. If a starless core goes through purely gravitational collapse, its peak speed of the inward motion is between values of the two red lines. If collapse of starless core is hindered by an external perturbation, it may collapse slower than the collapse of the BE core. If the peak speed of the inward motion is faster than the red dashed line, a core has been perturbed by an external perturbation in a way to enhance collapse.
Figure 4.4 Ammonia (1,1) and (2,2) lines observed toward four starless cores. The black solid lines are the observed lines and the red dashed lines are fitted models. Fitting is done in the five parameters space including $v_{\text{lsr}}$, line width, total optical depth of (1,1) transition, kinetic temperature, and filling factor. The best-fit parameters are listed in Table 4.2.
Launhardt et al., 2013). We assumed isothermality because it simplifies the dynamics calculation by obviating the need to perform radiative transfer calculations. The effects due to non-isothermal temperature profile have been studied by Sipilä et al. (2011). They demonstrated that the critical sizes of BE-like cores depend on their temperature profiles; however, the critical size varies only about 6% (from 6.38 to 6.76) from the original critical size (6.5), while the temperature varies 35% (from 7 K to 9.5 K). Considering that the observed measurements of $\xi_{\text{max}}$ has a typical error of 50%, our results with an isothermal assumption are still applicable to real starless cores.

Our analysis also ignores rotation of starless cores. Studies of molecular line widths toward starless cores show that cores slowly rotate (e.g. Heiles and Katz, 1976; Arquilla and Goldsmith, 1986; Goodman et al., 1993; Caselli et al., 2002b). If the rotational energy is comparable to the gravitational energy, rotation cannot be ignored in studying dynamical stability of starless cores. The typical observed ratio of the rotational energy to the gravitational energy is about 0.02 (Goodman et al., 1993). So, our assumption of ignoring rotation is practical for typical starless cores.

Magnetic fields may support starless cores against gravitational collapse. The magnetic field strength in starless cores from Zeeman splitting measurement and from polarization observations range from 10 $\mu$G to $\leq$ 200 $\mu$G (e.g. Ciolek and Basu, 2000; Levin et al., 2001; Crutcher et al., 2004; Kirk et al., 2006; Turner and Heiles, 2006). Magnetic pressure becomes comparable to the thermal pressure when the magnetic field strength reaches over 100 $\mu$G. Our stability analysis is not accurate for starless cores with magnetic fields of $\geq$ 100 $\mu$G. Inclusion of a magnetic term in the virial theorem is not practical for a dimensionless analysis because the inherent geometry of the field must be taken into account. Treatment of magnetic fields is beyond the scope of this current work.

4.5.2 Comparison to Observed Starless Cores with Infall

We compare the measured infall speeds and dimensionless sizes of observed starless cores which have clear infall signatures and well determined outer radii in Figure
4.6. The number of cores used in this study is limited because either only the density or the velocity structure is usually studied from observations of starless cores, but rarely both. The plotted lines are the same with those in Figure 4.3. The observed starless cores are marked with black filled circles. The $\xi_{\text{max}}$ value of L694-2 is from Harvey et al. (2003a). The peak infall velocity of L694-2 is measured by Lee et al. (2007) using the three hyperfine line of HCN J=1-0. The dimensionless size of L1544 is estimated by Kirk et al. (2005). The infall velocity of L1544 is studied by Williams et al. (2006) with interferometric observation of $\text{N}_2\text{H}^+$ J=1-0. For L63, we quote the dimensionless size of $\xi_{\text{max}} = 15$ from Kirk et al. (2005). There is no published infall speed study of L63, therefore, we estimate the infall velocity by applying the two-layer model (Equation (9) of Myers et al. 1996) to the HCN lines observed by Sohn et al. (2007). From the HCN $J=1-0 F'=1-1$ line, we estimate the infall speed of L63 to be about 0.1 km s$^{-1}$. The dimensionless size of L1689B is measured by Dapp and Basu (2009) and the infall speed is measured by Bacmann and Pagani (2008).

Internal temperature of those starless cores are required to normalize the infall speeds. Since internal temperatures of L63, L694-2 and L1689B are not studied with high precision, we deduced kinetic temperatures of the starless cores utilizing the hyperfine lines of $\text{NH}_3$ (1,1) and (2,2) inversion transitions. We fit ammonia spectra with five parameters including $v_{\text{lsr}}$, line width $\sigma_v$, kinetic temperature $T_k$, optical depth of (1,1) transition $\tau_1$, and filling fraction $\eta_f$ using the same LTE assumption and method in Rosolowsky et al. (2008). We fit signals with Signal-to-Noise Ratio $> 2$ and estimate the goodness of fit using the reduced $\chi^2$. We calculated the reduced $\chi^2$ in the full five parameter space instead of using a covariant matrix to more accurately determine 1σ uncertainties in the parameters. The results of the fittings are shown in Figure 4.4. The black solid lines represent observed ammonia lines and red dashed lines represent fitted line profiles. Reduced $\chi^2$ at the vicinity of the minimum value are plotted in Figure 4.5. The cross marks are the minimum points and the white lines represents 1σ uncertainty space in $T_k$ vs. $\tau_1$ space at the minima of $v_{\text{lsr}}$, $\sigma_v$, and $\eta_f$ (see Table 2). The temperatures of L1544, L63, L694-
2, L1689B from our study are 8.84±0.30, 9.50±0.41, 9.14±0.2, and 12.0±1.7 K, respectively. Crapsi et al. (2007) also deduced average temperature of L1544 using the same transition of NH$_3$ hyperfine lines, but with interferometric observations, to be 8.75K. Our results are consist within 1σ uncertainty. The temperature of L694-2 also agrees with the previous finding of Williams et al. (2006) which is 9.3K. The temperature of L1689B has a relatively large uncertainty. This is because L1689B appears to be optically thin in the NH$_3$ (1,1) line and determining its optical depth is degenerated with filling fraction.

Combining all physical quantities, four starless cores are plotted in Figure 4.6. We mark error bars only for L694-2 because the quoted references do not provide any error for ξ$_{max}$ and infall speeds for the other three cores. Generally, ξ$_{max}$ has about 50% uncertainty, and errors of infall speeds are hard to estimate unless a starless core is mapped with high angular resolution and with multiple molecular lines. Accepting that the observed physical quantities for the starless cores quoted are the best determination to date, we may say that all four cores are located in the regime below the collapse of the critical BE sphere. Two cores, L1689B and L694-2, seem to collapse faster than the collapse of the critical uniform density sphere. These two starless cores may be undergoing enhanced collapse by external perturbations, for example, turbulence, or a sudden increase of external pressure. Observations of their surrounding environments are needed to understand the nature of the perturbation. The other two starless cores L1544 and L63 seem to be in quasi-equilibrium collapse.
Table 4.1. Observed Starless Cores

<table>
<thead>
<tr>
<th>Source</th>
<th>RA (J2000.0)</th>
<th>DEC (J2000.0)</th>
<th>I(T_{mb}) [K km/s]</th>
<th>(\sigma_I) [K km/s]</th>
<th>(\sigma_{T_{mb}}) [K]</th>
</tr>
</thead>
<tbody>
<tr>
<td>L1544</td>
<td>05:04:17.2</td>
<td>+25:10:43.7</td>
<td>14.79</td>
<td>1.01</td>
<td>0.064</td>
</tr>
<tr>
<td>L63</td>
<td>16:50:14.9</td>
<td>-18:06:22.5</td>
<td>10.22</td>
<td>0.72</td>
<td>0.107</td>
</tr>
<tr>
<td>L694-2</td>
<td>19:41:04.3</td>
<td>+10:57:00.7</td>
<td>12.13</td>
<td>0.83</td>
<td>0.078</td>
</tr>
<tr>
<td>L1689B</td>
<td>16:34:48.3</td>
<td>-24:38:03.5</td>
<td>2.962</td>
<td>0.26</td>
<td>0.093</td>
</tr>
</tbody>
</table>

Table 4.2. Physical Quantities of Starless Cores Deduced from NH\(_3\) (1,1) and (2,2)

<table>
<thead>
<tr>
<th>Source</th>
<th>(v_{lsr}) [kms(^{-1})]</th>
<th>(\sigma_v) [kms(^{-1})]</th>
<th>(T_k) [k]</th>
<th>(\tau_1)</th>
<th>(\eta_f)</th>
</tr>
</thead>
<tbody>
<tr>
<td>L1544</td>
<td>7.196</td>
<td>0.134(^{+0.005}_{-0.005})</td>
<td>8.84(^{+0.30}_{-0.30})</td>
<td>3.17(^{+0.41}_{-0.35})</td>
<td>0.80(^{+0.06}_{-0.05})</td>
</tr>
<tr>
<td>L63</td>
<td>5.721</td>
<td>0.0825(^{+0.0053}_{-0.0026})</td>
<td>9.50(^{+0.41}_{-0.41})</td>
<td>2.62(^{+0.35}_{-0.35})</td>
<td>0.80(^{+0.06}_{-0.06})</td>
</tr>
<tr>
<td>L694-2</td>
<td>9.595</td>
<td>0.127(^{+0.003}_{-0.003})</td>
<td>9.14(^{+0.20}_{-0.20})</td>
<td>2.36(^{+0.15}_{-0.15})</td>
<td>0.77(^{+0.04}_{-0.04})</td>
</tr>
<tr>
<td>L1689B</td>
<td>3.591</td>
<td>0.172(^{+0.018}_{-0.032})</td>
<td>12.0(^{+1.7}_{-1.7})</td>
<td>0.30(^{+0.20}_{-0.20})</td>
<td>0.47(^{+0.52}_{-0.21})</td>
</tr>
</tbody>
</table>

\(^a\)Uncertainties of \(v_{lsr}\) are not estimated

4.6 SUMMARY & CONCLUSIONS

Starless cores are the precursors of low-mass star formation. Their stability and dynamics are important issues in developing a comprehensive picture of how low-mass stars form. However, owing to the difficulties of observing internal motions of starless cores through molecular line observations, only their density structures have been previously considered in investigating their stability and dynamics. In this study we find a stability condition for starless cores with internal motions. We first generalize the conventional stability condition of BE spheres by taking inward homologous motions into account. Even though homologous inward motions may not be a realistic internal motion that describes the observed internal motions of...
Figure 4.5 Reduced χ² values in T_k vs. τ₁ space. The other three parameters are fixed to be the best-fit values at the shown planes. The white crosses are the positions of the minimum reduced χ²_r and the white lines represents the 1σ uncertainty space. Uncertainties of each parameters are listed in Table 2.
Figure 4.6 Starless cores with clear infall signatures, L63, L1544, L1689B, and L694-2 are plotted with the generalized BE criterion and the limiting cases of the gravitational collapse. All cores are in unstable regime and gravitationally collapsing. The collapse of L694-2 and L1689B seem to be considerably enhanced by external perturbation.
starless cores, they provide a rough estimate how the stability conditions of starless cores vary with the speeds of inward motions. Second, we suggest two limiting cases of spontaneous gravitational collapse: the collapse of the critical BE sphere and the collapse of a uniform density sphere. The former model represents the slowest collapse by the self-gravity while the latter model sets the boundary for the fastest infall speed of spontaneous gravitational collapse at a given density. Any starless cores which considerably deviates from these two limiting models are likely to be strongly perturbed by the surrounding environment.

We applied our study to four starless cores which have strong infall signatures (large blue peak asymmetry). Some cores appear to collapse faster than infall speeds permitted by spontaneous gravitational collapse. This result suggests that external perturbations may be important in the collapse dynamics of low-mass starless cores. Further study of the density and velocity structures of starless cores and their surrounding environments are required to obtain more general conclusions concerning the collapse dynamics of low-mass star formation.
CHAPTER 5

THE ROLE OF NON-IONIZING RADIATION PRESSURE IN STAR FORMATION

Stars form when filaments and dense cores in molecular clouds fragment and collapse due to self-gravity. In the most basic analyses of gravitational stability, the competition between self-gravity and support from thermal pressure sets the critical (i.e. maximum stable) mass of spheres and the critical line density of cylinders. Previous work has considered additional support from magnetic fields and turbulence. Here, we consider the effects of non-ionizing radiation, specifically the inward radiation pressure force that acts on dense structures embedded in an isotropic radiation field. Using hydrostatic, isothermal models, we find that irradiation lowers the critical mass and line density for gravitational collapse, and can thus act as a trigger for star formation. For structures with moderate central densities, $\sim 10^3 \text{ cm}^{-3}$, the interstellar radiation field in the Solar vicinity has an order unity effect on stability thresholds. For more evolved objects with higher central densities, a significant lowering of stability thresholds requires stronger irradiation, as can be found closer to the Galactic center or near stellar associations. Even when strong sources of ionizing radiation are absent or extincted, our study shows that interstellar irradiation can significantly influence the star formation process.

5.1 INTRODUCTION

Star formation is a process by which interstellar gas becomes denser via a hierarchy of structures, with gravitational collapse playing a key role. Dense filaments are ubiquitous in molecular clouds, with 75% of denser starless cores residing in filaments (André et al., 2010). The formation of cores within filaments is explained

---

by gravitational instability (Inutsuka and Miyama, 1992). The formation of stars within cores is attributed to another gravitational collapse (Shu et al., 1987).

A basic understanding of gravitational collapse comes from the study of isothermal and pressure-confined gas in hydrostatic equilibrium. The classic solutions are the Bonnor-Ebert (BE) sphere (Ebert, 1955, 1957; Bonnor, 1956) and the isothermal Ostriker (1964) cylinder. These studies show that gravitational collapse occurs above the critical mass of the BE sphere and the critical line density (mass per length) of isothermal cylinders. The critical BE mass also applies to hydrostatic clouds of any geometry, provided the volume is finite (Lombardi and Bertin, 2001). Observed density profiles of filaments and cores are often well matched by these simple models (e.g. Bacmann et al., 2000; Alves et al., 2001; Kandori et al., 2005; Hacar and Tafalla, 2011).

While of fundamental importance, these classic solutions neglect many potentially significant effects. Magnetic fields, turbulence and detailed radiative transfer can alter the structure and stability of cores and filaments (McKee and Ostriker, 2007). This work focuses on a particular aspect of radiative transfer, the radiation pressure exerted on cores and filaments by ambient non-ionizing radiation. Ionizing radiation is known to have important effects on star formation in HII regions, i.e. near high-mass stars. Our focus on non-ionizing radiation applies not only to low-mass star forming regions, but also to regions of high column density into which non-ionizing radiation penetrates more deeply.

Non-ionizing radiation (from mid-UV to mid-IR) exerts radiation pressure on dust grains which are frictionally coupled to the gas (Draine, 2011). The radiation pressure force is weak in the diffuse interstellar medium because the interstellar radiation field (hereafter ISRF) is almost isotropic, with \( \sim 10\% \) anisotropy, (Weingartner and Draine, 2001). However, near and inside a dense structure, radiation becomes anisotropic due to shadowing by the structure itself. With this introduced anisotropy, radiation pressure becomes comparable to thermal gas pressure in the interstellar medium.

In this chapter we study how the radiation pressure force alters the structure
and gravitational stability of filaments and cores. Section 5.2 describes our model for hydrostatic irradiated structures. Section 5.3 presents our self-similar solutions in dimensionless coordinates. In §4, we apply our results to the physical conditions of star forming regions. In §5, we discuss limitations and future extensions of our model. We summarize our results in §6.

5.2 MODEL FOR IRRADIATED CYLINDERS AND SPHERES

5.2.1 Hydrostatic Structure with Radiation Pressure

We consider hydrostatic configurations of dense cores and filaments exposed to non-ionizing photons that exert radiation pressure. Our idealized model assumes spherical symmetry for dense cores and cylindrical symmetry for filaments. Radiation pressure acts on dust grains, which are uniformly mixed and perfectly coupled to the gas. Possible sedimentation of dust grains is addressed in §5.

Our models satisfy hydrostatic equilibrium and the Poisson equation:

\[
\nabla P_g = -\rho \nabla \psi + n_d \vec{f}_{\text{rad}} \tag{5.1}
\]

\[
\nabla^2 \psi = 4\pi G \rho \tag{5.2}
\]

where \( P_g \) is the gas pressure, \( \psi \) is the gravitational potential, \( \vec{f}_{\text{rad}} \) is the force exerted to a dust grain due to radiation pressure, \( n_d \) is the number density of dust grains per unit volume, and \( \rho = \rho_g + \rho_d \) is the total density of gas and dust, respectively.

The radiation pressure force on a dust grain is

\[
\vec{f}_{\text{rad}} = \pi a^2 \langle Q_{\text{pr}} \rangle \frac{\vec{F}_{\text{rad}}}{c}, \tag{5.3}
\]

where \( a \) is the effective spherical radius of dust grains, \( \vec{F}_{\text{rad}} \) is the energy flux of radiation field, \( c \) is the speed of light and \( \langle Q_{\text{pr}} \rangle \) is radiation pressure efficiency, described in detail in §4.1. The mass of a dust grain is

\[
m_d = \frac{\rho_d}{n_d} = \frac{4\pi}{3} \rho_m a^3 \tag{5.4}
\]

with \( \rho_m \) the internal, material density of dust grains.
The radiative flux is directed inwards, $\vec{F}_{\text{rad}} = F \hat{r}$ with $F < 0$ and $\hat{r}$ the unit vector along the spherical or cylindrical radial coordinate, $r$. We consider the extinction of radiation (but ignore scattering) by taking a two-ray approximation (see Figure 5.1). One ray, $F^-$, represents the flux entering the sphere/cylinder from the near surface, while $F^+$ represents the oppositely directed flux from the far surface, which has passed through the center. The total flux $F = -F^- + F^+$ with

$$F^- = F_0 \exp[-\tau(r)] \quad (5.5a)$$
$$F^+ = F_0 \exp[-\tau_{\text{tot}} + \tau(r)], \quad (5.5b)$$

The optical depth to the near surface, at $r = r_0$, is

$$\tau(r) = \pi \int_{r_0}^{r} a^2 \langle Q_{\text{ext}} \rangle n_d dr', \quad (5.6)$$

$\langle Q_{\text{ext}} \rangle$ is the spectrum-averaged extinction coefficient. The total extinction of a sphere/cylinder is $\tau_{\text{tot}} = \tau(-r_0)$. The normalization $F_0$ gives the one-sided surface flux, i.e. $F(r_0) \rightarrow -F_0$ for an opaque object with $\tau_{\text{tot}} \gg 1$. The validity of the two-ray approximation is addressed in §5.

We adopt an isothermal equation of state for the gas:

$$P_g = c_s^2 \rho_g, \quad (5.7)$$

where $c_s$ is the sound speed of gas, and $\rho_g$ is the gas density, consistent with the standard Bonnor-Ebert problem and in at least rough agreement with observed cores and filaments (Evans et al., 2001; Stamatellos et al., 2007; Seo et al., 2015).

To derive our version of the Lane-Emden equation, we take the divergence of $(1/\rho$ times) the equation (5.1), which combined with equation (5.2) gives

$$c_s^2 \nabla \cdot \left( \frac{1}{\rho} \nabla \frac{\rho}{1+Z} \right) = -4\pi G \rho + \nabla \cdot \left( Z \frac{\vec{f}_{\text{rad}}}{1+Z m_d} \right) \quad (5.8)$$

where the dust-to-gas ratio, $Z = \rho_d/\rho_g$. Solution of equation (5.8) also requires equations (5.3), (5.5) and (5.6) to specify the radiation force. To simplify the solution space, we also hold spatially constant both $Z$ and the dust opacities

$$\kappa_i = \frac{Z}{1+Z} \frac{3\langle Q_i \rangle}{4\rho_m a} \quad (5.9)$$
where the index $i = \text{“pr”}$ and “ext” labels the radiation pressure and extinction cases. In practice, we obtain solutions using the dimensionless equations described below.

### 5.2.2 Dimensionless Equations

Our dimensionless equations use the central density $\rho_c$, the dust-weighted sound speed $c_s' = c_s / \sqrt{1 + Z}$, and the characteristic scale height

$$\alpha \equiv \frac{c_s'}{\sqrt{4\pi G \rho_c}} \quad (5.10)$$

as scale factors. The dimensionless variables

$$\xi \equiv \frac{r}{\alpha}, \quad (5.11)$$
$$s \equiv \frac{\rho}{\rho_c}, \quad (5.12)$$

describe the radial distance and density, while $\tau$ is already dimensionless.

The inclusion of radiative effects adds two new dimensionless parameters. The dimensionless extinction,

$$\zeta \equiv \kappa_{\text{ext}} \alpha \rho_c, \quad (5.13)$$
gives a characteristic (but not the actual) the optical depth. The dimensionless radiation pressure strength,

$$\Upsilon \equiv \frac{\langle Q_{\text{pr}} \rangle}{\langle Q_{\text{ext}} \rangle} \frac{F_\circ / c}{P_c}. \quad (5.14)$$

normalizes the radiation pressure to the central pressure, $P_c = c_s'^2 \rho_c$, times the ratio of radiative efficiencies.

In these dimensionless units, the governing equations (5.6) and (5.8) read

$$\frac{d\tau}{d\xi} = \zeta s, \quad (5.15)$$
$$\nabla^2_\xi \ln(s) = -s + \zeta \Upsilon A(\xi; \tau) \quad (5.16)$$

where $\nabla^2_\xi$ is the standard Laplacian for the $\xi$ coordinate, i.e.

$$\nabla^2_\xi f = \xi^{1-D} d/d\xi (\xi^{D-1} df/d\xi)$$

with $D = 3$ for spheres or $D = 2$ for cylinders. The geometry of the radiation field is parameterized as $A = \nabla_\xi \cdot (\vec{F}_{\text{rad}}/F_\circ)$, a dimensionless
Figure 5.1 Schematic of our two-ray model for the flux of non-ionizing radiation inside a dense core or filament. The net flux at any internal radius \( r < r_0 \) arises from two competing rays: the inward directed flux, \( F^- \), and the outward directed flux, \( F^+ \), that has passed through the center of the object. An optically thin object in an isotropic radiation field will experience little net flux as the two contributions nearly cancel.
divergence. In the two-ray approximation,

\[ A = \left( -\frac{D-1}{\xi} + \zeta_s \right) e^{-\tau} + \left( \frac{D-1}{\xi} + \zeta_s \right) e^{-\tau_{\text{tot}} + \tau}. \]

(5.17)

We solve the equations (5.15), (5.16) and (5.17) subject to the boundary conditions \( s = 1 \) and \( ds/d\xi = 0 \) at the center, \( \xi = 0 \), \(^2\) and \( \tau = 0 \) at the outer boundary, \( \xi = \xi_0 \). In order to avoid matching conditions at the inner and outer boundaries, we set \( \tau(0) = \tau_{\text{tot}}/2 \) at the center. Not all choices of \( \tau_{\text{tot}} \) give valid solutions, but if \( \tau \) drops to zero within a finite radius then the solution is valid. We use a 4-th order Runge-Kutta integrator, and step slightly off \( \xi = 0 \) to avoid any singularity.

5.3 SELF-SIMILAR SOLUTIONS

5.3.1 Irradiated Spheres

The structure of irradiated, pressure-confined, isothermal spheres, is presented in the left panel of Figure 5.2. These dimensionless density profiles, \( s(\xi) \), are calculated as described in §5.2.2. In these sample curves, a large value of optical depth, \( \tau_{\text{tot}} \), is assumed so that the density at the outer radius is small, i.e. \( s(0) = s(\xi_0) < 10^{-3} \), which corresponds to a large density contrast \( \rho_c/\rho_0 = 1/s_0 \). The \( \Upsilon = 0 \) curve corresponds to the standard Bonnor-Ebert sphere. As radiation pressure increases (to larger \( \Upsilon \)), the outer density profile steepens and spheres become radially truncated. Physically, the outward pressure gradient force must increase to balance the inward radiation pressure.

Gravitational stability depends on the curve of dimensionless mass, \( m \), versus density contrast. In terms of the dimensional mass \( M \) and surface pressure \( P_0 \),

\[ m \equiv \frac{P_0^{1/2} G^{3/2} M}{c_s^4} = \sqrt{\frac{s_0}{4\pi}} \int_0^{\xi_0} s \xi^2 d\xi. \]

(5.18)

From the right hand side above, we note that as \( s_0 \) decreases (and thus the density contrast increases), the prefactor decreases \( \propto \sqrt{s_0} \), while the integral increases, due to the larger radius, \( \xi_0 \). This competition affects gravitational stability.

\(^2\)Solutions with \( s(0) \neq 1 \) simply correspond to a different normalization, and can be mapped onto equivalent \( s(0) = 1 \) solutions.
Pressure bounded spheres are gravitational unstable if (Bardeen, 1965; Stahler, 1983):

\[
\frac{\partial m}{\partial (\rho_c/\rho_0)} < 0.
\]  \hspace{1cm} (5.19)

This instability criterion is equivalent to the more intuitive Boyle’s law criterion, \( \partial P_0/\partial V_0 > 0 \), that gravitating spheres are unstable if an enhanced surface pressure induces expansion to a larger volume, \( V_0 \) (Bonnor, 1956; Lombardi and Bertin, 2001). In Appendix C we verify that this established correspondence also applies in the presence of other forces, such as radiation pressure.

The right panel of Figure 5.2 shows \( m \) versus density contrast for different values of radiation pressure, \( \Upsilon \). At low \( m \) solutions are gravitationally stable since \( m(\rho_c/\rho_0) \) has a positive slope. The local maximum where \( dm/d(\rho_c/\rho_0) = 0 \) defines marginal stability at the critical mass, \( m_{\text{crit}} \). For \( \Upsilon = 0 \), we reproduce the well known Bonnor-Ebert mass, \( m_{\text{crit}} = 1.18 \), and the maximum density contrast of 14.1. As radiation pressure increases (larger \( \Upsilon \)), both \( m_{\text{crit}} \) and the critical density contrast decrease. (But see below for a case where irradiation cause the critical density contrast to increase).

The properties of marginally unstable irradiated spheres are further explored in Figure 5.3, which also examines the effect of extinction, via \( \zeta \). All the panels in this figure correspond to the critical state with \( m = m_{\text{crit}} \), whose values are shown in the top panel. Both this critical mass and the corresponding radius (shown in the second panel) become smaller as either \( \Upsilon \) or \( \zeta \) increases. Either stronger irradiation or a greater opacity increases the surface radiation pressure force, which scales as \( \Upsilon \zeta \) (see equation 5.16).

The density contrast of critical spheres displays interesting behavior, shown in the third panel of Figure 5.3. For small extinctions, \( \zeta \lesssim 1.3 \), the density contract decreases gradually with increasing irradiation, as might be expected for smaller, lower mass spheres. For larger extinctions however, the density contrast develops a spike near \( \Upsilon = 1 \). We note that our normalization of the radiation pressure (to the central pressure) is clearly appropriate since interesting behavior occurs for \( \Upsilon \) near unity.
The origin of the spike in density contrast is explained by the total optical depth of the critical spheres, shown in bottom panel of Figure 5.3. The spike in the density contrast corresponds to the transition from high to low total optical depth. The plot shows that only weakly irradiated clouds (Υ < 1) can have a high total optical depth and remain stable. In this weakly irradiated regime, the optical depth scales simply with the opacity, via ζ, as radiative effects are not yet significantly affecting cloud structure or stability.

For stronger irradiation, as Υ approaches unity, the maximum optical depth of cores decreases, consistent with their smaller masses and sizes. The spike in density contrast occurs because strong radiation pressure forces, which steepen the density profile, are being felt throughout more of the sphere. For even stronger irradiation, with Υ exceeding unity, the core becomes so transparent that the radiative effects weaken, due to the flux cancellation depicted in Figure 5.1. In this highly irradiated regime, the critical density contrast decreases again and the behavior of critical spheres is surprisingly simple: the total optical depth is order unity value for all extinction values.

In summary, as irradiation increases, the marginally gravitationally stable state gradually transitions from the Bonnor-Ebert sphere to the sphere with optical depth near unity.

5.3.2 Irradiated Cylinders

Hydrostatic density profiles of irradiated cylinders are shown in Figure 5.4. Similar to the spherical case, irradiation steepens the density profile of cylinders at their outer edges. Non-irradiated isothermal cylinder already have very steep outer density profiles, with $s \propto \xi^{-4}(\rho \propto r^{-4})$. In the irradiated cases, even less mass resides at large radii.

An isothermal cylinder does not experience a Bonnor-type instability because $\partial P_0/\partial V_0$ is always negative (Lombardi and Bertin, 2001). Instead, the cylinder becomes unstable only when its line density exceeds the critical line density, which is the line density of a hydrostatic cylinder with infinite outer radius. The critical line
density in a dimensionless form is given as

\[ \lambda_{\text{crit}} \equiv \frac{G \Lambda_{\text{crit}}}{c_s^2} = \frac{1}{2} \int_0^\infty \xi s d\xi \]  

(5.20)

where \( \Lambda_{\text{crit}} \) is the dimensional critical line density of a hydrostatic cylinder. With formally infinite outer radius, the critical cylinder does not have a corresponding radius or density contrast. In reality of course cylinders are not infinite in radius, and they do have a finite density contrast set by the ambient medium. However the steepness of the outer density profile means that the mass of the infinite cylinder is a very good approximation for the stability boundary for finite radius cylinders. Because there is no Bonnor type instability, the external gas pressure does not affect the result. The optical depth of the critical cylinder is well defined, again due the steep density profile.

The critical line density and optical depth of irradiated cylinders are shown in Figure 5.5. The qualitatively behavior is the same as the spherical case. With low levels of irradiation, \( \Upsilon \ll 1 \), we recover the standard result for the critical line density of isothermal cylinders, \( \lambda_{\text{crit}} = 2 \). Increases in either the irradiation or the opacity (\( \Upsilon \) or \( \zeta \)) decrease the critical line density. When radiation effects are
Figure 5.3 Critical values of (from top to bottom) dimensionless mass $m_{\text{crit}}$, size $\xi_{\text{max}}$, density contrast $\rho_c/\rho_0$, and total optical depth $\tau_{\text{tot}}$ for marginally stable irradiated spheres as a function of the dimensionless radiation pressure $\Upsilon$. Different colored curves correspond to different dimensionless opacities, $\zeta$. 
strong, for $\Upsilon > 1$, critical cylinders converge to an optical depth of order of unity, independent of opacity.

5.4 RESULTS APPLIED TO STAR FORMING REGIONS

We now apply our dimensionless solutions to a wide range of interstellar conditions. We summarize the adopted physical parameters in section 5.4.1. We present our numerical solutions in section 5.4.2. In sections 5.4.3 and 5.4.4, we discuss implications for low and high mass star forming regions, respectively.
Figure 5.5 Critical line density (top) and total optical depth (bottom) of dusty gas cylinder irradiated by non-ionizing radiation as a function of the radiation pressure $\Upsilon/\zeta$. Different colors denote different dimensionless extinction, $\zeta$. 
5.4.1 Physical Parameters

Our solutions have a constant central number density, \( n_c \), of \( 10^3 \) or \( 10^4 \) cm\(^{-3} \). These values characterize quasi-spherical clumps and young cores, respectively (Bergin and Tafalla, 2007). Elongated filaments also span this range of central densities (Arzoumanian et al., 2011).

The effective sound speed

\[
 c_s = \sqrt{0.188^2 \left( \frac{T_k}{10 \text{ K}} \right) + \sigma_{nt}^2 \text{ km/s} = 0.188 \sqrt{\frac{T_{\text{eff}}}{10 \text{ K}}} \text{ km/s},
\]  

(5.21)

assumes a mean molecular weight of 2.33 proton masses. A kinetic temperature of \( T_k \sim 10 \text{ K} \) is a typical value for dense molecular gas (e.g. Leung, 1975; Hotzel et al., 2002; Tafalla et al., 2004). The non-thermal component of the velocity dispersion, \( \sigma_{nt} \), provides at most an order unity correction (e.g. Goodman et al., 1998; Pineda et al., 2010a; Hacar et al., 2013; Seo et al., 2015). Our numerical study considers a range of effective temperatures, \( T_{\text{eff}} = 7.5 - 25 \text{ K} \). The corresponding scale height is

\[
 \alpha = 0.034 \text{ pc} \left( \frac{c_s}{0.188 \text{ km/s}} \right) \left( \frac{10^4 \text{ cm}^{-3}}{n_c} \right)^{\frac{1}{2}}.
\]  

(5.22)

We fix the dust-to-gas ratio to be \( Z = 0.01 \), a typical interstellar medium value (Draine et al., 2007). For grain sizes, we adopt \( a = 0.05, 0.075, \) and \( 0.1 \mu m \), which are characteristic sizes in molecular regions (Köhler et al., 2015). Though not considered here, our analysis could be extended to accommodate particle size distributions.

We use three models for the chemical composition of dust grains: (1) astrosilicate grains (Weingartner and Draine, 2001), (2) carbonaceous grains (Draine, 2003) and (3) grains with an ice mantle covering a silicate-carbonaceous core. For the ice mantle grains, we assume a 1:1:1 ratio of astrosilicate, carbonaceous material and water ice by volume, in agreement with Li and Greenberg (1997). The material density of dust is assumed to be \( 1 \text{ g/cm}^3 \) (roughly 50% porosity) for ice grains and ice mantles, \( 1 \text{ g/cm}^3 \) for carbonaceous (about 40% porosity) and \( 2 \text{ g/cm}^3 \) (about 40% porosity) for astrosilicates.
For the radiation field we use the spectrum of the ISRF in the solar vicinity (Mezger et al., 1982; Mathis et al., 1983). Adopted values of the flux, $F_0$, range from $5 \times 10^{-4} - 0.4$ erg/cm$^2$/s. In the Solar vicinity (Galactocentric distance $D_G \sim 8$ kpc), the mean intensity of the non-ionizing ISRF may range from 0.015 erg/cm$^2$/s to 0.15 erg/cm$^2$/s (Keene, 1981; Mezger, 1990; Launhardt et al., 2013). These intensities are equivalent to a directed flux of $F_0 = 0.0025 - 0.025$ erg/cm$^2$/s).

We calculate the radiative efficiency parameters $\langle Q_{pr} \rangle$ and $\langle Q_{ext} \rangle$ with the Mie theory code, $miex$ (Wolf and Voshchinnikov, 2004) for the grain and radiation properties described above. Figure 5.6 shows results of the Mie calculations, both for the adopted ISRF and (for comparison) for main sequence stellar spectra of various effective temperatures (from Pickles, 1998).

With the above physical parameters, our dimensionless parameters scale as:

$$\Upsilon = 0.49 \left( \frac{\langle Q_{pr} \rangle}{\langle Q_{ext} \rangle} \right) \left( \frac{F_0}{0.02 \text{ erg/cm}^2/\text{s}} \right) \left( \frac{10^3 \text{ cm}^{-3}}{n_c} \right) \left( \frac{0.188 \text{ km/s}}{c_s} \right)^2$$

$$\zeta = 0.94 \left( \frac{Z}{0.01} \right) \left( \frac{\langle Q_{ext} \rangle}{1} \right) \left( \frac{n_c}{10^3 \text{ cm}^{-3}} \right)^{1/2} \left( \frac{c_s}{0.188 \text{ km/s}} \right) \left( \frac{0.1 \mu \text{m}}{\alpha} \right) \left( \frac{1 \text{ g/cm}^3}{\rho_m} \right)$$

for the radiation pressure strength and extinction, respectively. The fact that these parameters are order unity suggests (following the self-similar analysis of section 5.3) that irradiation will indeed be important in moderately dense regions.
5.4.2 Gravitational Instability Criteria with Irradiation

Molecular clumps and cores

Figure 5.7 shows the critical mass and size of irradiated molecular clumps \((n_c = 10^3 \text{ cm}^{-3})\) and young cores \((n_c = 10^4 \text{ cm}^{-3})\), modeled as 10 K spheres. In the left panels, we use ice-mantle dust grains with three different dust sizes. With weak irradiation, lower density clumps have a higher critical mass and larger critical radius. This behavior agrees with the standard density scaling of isothermal Bonnor-Ebert spheres, \(M_{\text{crit}} \propto R_{\text{crit}} \propto 1/\sqrt{n_c}\).

As with the self-similar solutions, increased external irradiation decreases critical masses and radii. A given level of irradiation more strongly affects lower density clumps, because radiation pressure is larger relative to the central pressure. Indeed the irradiation parameter, \(\Upsilon \propto F_0/n_c\) at constant temperature, quantifies this pressure ratio. Thus while the critical state of clumps is significantly downsized by the radiation fields characteristic of the Solar neighborhood (the grey bands in Figure 5.7), denser cores are significantly affected only when the radiation field approaches levels seen in Galactic center.

The left panels of Figure 5.7 also show that the choice of grain size only modestly affect the critical state of spheres. This result is consistent with the fact (seen in Figure 5.6) that efficiency parameters, \(\langle Q \rangle\), increase with the size of small grains so that the relevant combination \(\langle Q \rangle/a\) varies weakly with size in this regime. In other words, optical properties are more sensitive to the grain mass than the grain surface area in the long wavelength regime.

The effects of grain composition are shown in the right panels of Figure 5.7. For silicate grains, the effects of irradiation are weaker – as measured by the reduction in critical mass and radius at a given flux – than for carbonaceous or ice-mantle grains. This effect is explained by the lower radiative efficiencies of small silicate grains, as shown in Figure 5.6. In a dense molecular region, temperatures are low enough that we do not expect bare silicates. Either the ice-mantle case or carbonaceous cases should be more realistic, depending on the fractional abundance of carbon species.
Figure 5.6 Spectrum-averaged radiation pressure efficiency (left) and spectrum-averaged extinction (right) as a function of dust grain size. *Silicate* denotes that dust grains are composed of astro-silicates, *carbonaceous* denotes carbonaceous dust grains, and dust grains of *ice mantle* are assumed to be composed of 1:1:1 of astro-silicates, carbonaceous, and water ice.
(Greenberg and Li, 1999).

Table 5.1 presents critical masses and radii of irradiated spheres for a range of effective temperatures. Critical masses and radii are larger at higher temperatures. Also, at higher temperatures, and thus higher central pressures, a correspondingly stronger external radiation pressure is needed to lower the critical masses and radii.
Figure 5.7 Critical mass (top), critical size (bottom) of molecular clumps and cores as a function of radiation field strength. The left panels show the critical mass and size of clumps and cores with ice-mantle dust grains. The right panels show the critical mass and size for different chemical compositions of dust grains, while dust size is fixed to be $a = 0.05 \mu m$. The effective temperature is fixed to be 10 K in all solutions. The central gas densities $n_c$ are written in the panels and marked with different colors. Different line styles denote different sizes of dust grain (left panels) and different chemical composition of dust grain (right panels). The grayed area denotes a range of the ISRF strength in the Solar vicinity. The dashed-dotted and the dashed-double dotted gray lines denote the average ISRF at the molecular ring ($D_G = 4$ kpc) and the Galactic center, respectively.
## Table 5.1: Critical Sizes and Masses of Dense Clumps and Cores under Radiation

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<th>$F_0^6$ (erg/cm(^2)/s)</th>
<th>$\alpha^c$ (µm)</th>
<th>$T_{\text{eff}}^d$ (K)</th>
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<td>0.141, 3.04</td>
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<td>0.138, 3.76</td>
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\(a\) The central density of gas
\(b\) The radiation flux normal to the dense clump/core surface
\(c\) The radius of dust grain. The ice-mantle dust grain model is used.
\(d\) The effective temperature of the internal supports including thermal and non-thermal components
Molecular filaments

The critical line density of irradiated filaments at 10 K is shown in Figure 5.8. In the low-irradiation limit, the critical line density, $\Lambda_{\text{crit}} = 2c_s^2/G$ is independent of central density, unlike the $M_{\text{crit}}$ of spheres. Analogously to the spherical case, the critical line density of cylinders decreases for increasing levels of irradiation. Also like the spherical case, irradiation is felt more strongly by lower density cylinders, grain size dependence is weak and silicate grains are less affected by irradiation. The different geometry of spheres and cylinders gives rise to some subtle differences. For instance $\Lambda_{\text{crit}}$ is slightly less sensitive to increasing irradiation than is $M_{\text{crit}}$. Table 5.2 presents critical line densities of irradiated cylinders for a range of effective temperatures.

5.4.3 Low-Mass Star Forming Regions

In low-mass star forming regions, both nearby young stellar objects (YSOs) and the diffuse ISRF contribute to the radiation field. We first consider quiescent regions with no YSOs. Under the ISRF in the Solar vicinity (shaded gray in Figure 5.7), the critical mass of a molecular clump ($n_c = 10^3$ cm$^{-3}$) is significantly reduced, to values as small as 5.5 M$_\odot$, i.e. 40% of the non-irradiated value (see red curves in Figure 5.7). On the other hand, the stability properties of dense cores ($n_c > 10^4$ cm$^{-3}$) are only modestly affected by the Solar ISRF (see blue curves in Figure 5.7). Thus while the Solar ISRF does not affect the collapse of dense cores, this level of irradiation can affect the formation of dense cores within less dense clumps.

Closer to the Galactic center, the ISRF is more intense, around $\sim$9 times stronger in the molecular ring ($D_G = 4$ kpc) compared to the Solar vicinity, and $\sim$50 times stronger at the Galactic center (Mezger, 1990). (These levels of irradiation are marked by dashed-dotted and dashed-triple dotted lines, respectively, in Figure 5.7.) In the inner Galactic regions, the critical sizes and masses of both dense clumps and dense cores may be significantly smaller ($<20\%$) than those in the Solar vicinity. At face value, these results favor the formation of lower mass stars in the
Figure 5.8 Critical line density of molecular filaments as a function of radiation strength. Left panel shows the critical line densities with ice-mantle dust grains. Right panel show the critical line densities with different chemical composition of dust grains, while dust size is fixed to be $a = 0.05 \, \mu m$. The effective temperature is fixed to be 10 K in all solutions. The central gas densities $n_c$ are written in the panels and marked with different colors. Different line styles denote different size of dust grain (left) and different chemical composition of dust grain (right). The grayed area denotes a range of the ISRF strength in the Solar vicinity. The dashed-dotted and the dashed-double dotted gray lines denote the average ISRF at the molecular ring ($D_G = 4 \, kpc$) and the Galactic center, respectively.
Table 5.2 Critical Line Densities of Filaments under Radiation

<table>
<thead>
<tr>
<th>$n_c^a$ (cm$^{-3}$)</th>
<th>$F_0^b$ (erg/cm$^2$/s)</th>
<th>$\alpha^c$ (µm)</th>
<th>$T_{\text{eff}}^d$ = 7.5 K</th>
<th>10 K</th>
<th>15 K</th>
<th>20 K</th>
<th>25 K</th>
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<td>$10^4$</td>
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</table>

$a$ The central density of gas

$b$ The radiation flux normal to the filament surface

$c$ The radius of dust grain. The ice-mantle dust grain model is used.

$d$ The effective temperature of thermal and non-thermal support.
more intensely irradiated inner Galaxy. On the other hand, competing effects such as higher temperatures (e.g. Walmsley and Ungerechts, 1983; Ao et al., 2013) and extra non-thermal (e.g. magnetic) support favor higher mass star formation towards the Galactic center. More detailed models are needed to include these effects self-consistently.

In more evolved regions, molecular clouds receive enhanced irradiation from nearby YSOs, stellar associations or clusters (e.g. NGC 1333, B1688 & B1689 in ρ Ophiuchus, L1495A in Taurus, see Volgenau et al., 2006; Bontemps et al., 2001; Maruta et al., 2010; Seo et al., 2015). Consider a stellar association following Kroupa’s IMF (initial mass function) with ~600 stars and with the earliest stellar type being B9V. For a typical cluster diameter of 1 pc (Nilakshi et al., 2002), we roughly and conservatively estimate the radiation field by placing all stars in a 0.5 pc radius shell around a central reference point. A non-ionizing radiation flux of 0.05 erg/cm$^2$/s would be incident upon an opaque cloud at this reference point. The actual radiation field will of course depend on more realistic stellar locations and will be larger closer to the brightest stars. Figure 5.7 shows that this nominal flux of 0.05 erg/cm$^2$/s is enough to reduce the critical mass of dense cores from 4.4 $M_\odot$ to ~ 3$M_\odot$. This order unity change might not seem significant, given our idealized model. However the idea that radiation from nearby YSOs could help regulate the IMF merits further discussion, which we begin in the summary.

Irradiation also helps to trigger the formation of starless cores within filamentary structures. A non-irradiated young filament with $n_c \sim 10^3$ cm$^{-3}$ at 10 K should not fragment until it reaches a critical line density of 16.4 $M_\odot$/pc. However irradiation by the ISRF in the Solar vicinity can lower the fragmentation threshold to 9 $M_\odot$/pc (see gray shaded area in the left panel of Figure 5.8). Irradiation can thus explain why some filaments contain dense cores, despite having line densities below the standard (non-irradiated) critical value (Benedettini et al., 2015). For example, a cold filament in L1495-B218 (no.28 in Hacar et al., 2013) has a line density of 9.3 $M_\odot$/pc but embeds NH$_3$ starless cores (Seo et al., 2015). On the other hand, filaments with high line densities, e.g. in Aquila (Arzoumanian et al., 2011), seem
to require extra internal support – e.g. magnetic – to explain their existence beyond nominal stability thresholds. Thus a detailed understanding of the physical conditions of a filament – including the radiation environment – is required to determine gravitational stability.

5.4.4 High-Mass Star Forming Regions

High mass star forming regions are strongly affected by radiation from OB associations. Observations preferentially find YSOs near the rims of HII regions and near globules within HII regions (Jose et al., 2013). Massive stars can trigger star formation in different ways. In the collect-and-collapse scenario, the stellar winds of massive stars are the trigger (Elmegreen and Lada, 1977; Whitworth et al., 1994; Hosokawa and Inutsuka, 2005; Dale et al., 2007). In the radiation driven implosion (RDI) scenario, radiation pressure from ionizing photons is responsible (Bertoldi, 1989; Lefloch and Lazareff, 1994; White et al., 1997; Kessel-Deynet and Burkert, 2003; Gritschneder et al., 2009; Chauhan et al., 2009; Bisbas et al., 2011; Walch et al., 2013).

However, triggered star formation is harder to understand far from the direct influence of stellar winds and ionizing photons. For instance, enhanced YSO populations are observed deep within globules in the HII regions G028.83-0.25 and G041.92+0.04 (Dirienzo et al., 2012). Large scale compression could play a role. Alternately, we propose that the prodigious non-ionizing radiation from OB associations helps trigger star formation in more embedded regions. Due to a greater penetration depth, non-ionizing radiation can influence embedded regions (quantitatively $N_H \approx 10^{18} \, \text{cm}^{-2}$ for $\tau = 1$ at the Lyman limit of 912Å, while $N_H \approx 10^{21} \, \text{cm}^{-2}$ for a visual extinction $A_V = 1$, Cardelli et al., 1989; Gay et al., 2012). Moreover, triggered YSOs can subsequently promote sequential star formation via non-ionizing radiation pressure that extends into embedded regions.

Consider, for example, the Elephant’s trunk nebula (IC1396A), often cited as an example of triggered star formation. IC1396A is a globule within an HII region containing more than 50 embedded YSOs (Reach et al., 2004; Sicilia-Aguilar et al.,
2005, 2006b,a, 2014; Morales-Calderón et al., 2009; Getman et al., 2012) and a total mass of $\sim 200M_\odot$ (Morgan et al., 2010). YSOs at the bright rim of the globule – illuminated by an O6.5 star at 4.5 pc – are believed to be formed through the RDI (Sicilia-Aguilar et al., 2014). We propose that YSOs in the interior of the globule were triggered by non-ionizing radiation. In particular many YSOs formed in-between the bright outer rim of IC1396A and its central cavity, which is being cleared by a bright stellar association that includes the intermediate mass variable V390 Cep. With irradiation from both sides, the YSOs embedded in this in-between region start to resemble our idealized model of isotropic irradiation. The flux of the OB association at the surface of globule is 0.4 erg/cm$^2$/s and the flux from the YSOs at the globule center may be 0.05 erg/cm$^2$/s if we assume that the combined spectral type of YSOs are a single A0 type star. From figure 5.7 these estimated flux levels can strongly affect star formation. While our simple model can not make detailed predictions for such a complicated environment, we hope to motivate more study of the effects of non-ionizing radiation pressure environments like IC1396A.

When HII regions are illuminated by later type stars, i.e. OB associations with spectral types later than O6V, non-ionizing radiation will have a greater relative importance compared to ionizing radiation. For instance in the Gum nebula, cometary globules such as the CG30/31 complex (Nielsen et al., 1998; Kim et al., 2005) and the CG4/Sa101 complex (Rebull et al., 2011) are far from OB associations, but they are near to a stellar association with multiple B and A type stars. Similarly, the Cone nebula is a cometary dense cloud embedded within an HII region that is illuminated by Allen’s source (i.e. the B star NGS2264IRS Thompson et al., 1998). The YSOs in these regions are assumed to be a result of triggered star formation, even though the effects of stellar winds and ionizing radiation are relatively weak. On the other hand, we find that the non-ionizing radiation from nearby stellar associations or late-type OB associations is strong enough to trigger gravitational collapse by reducing the critical core mass (Figures 5.7 and 5.8).
5.5 MODEL ASSUMPTIONS AND EXTENSIONS

Our study focuses on the role of non-ionizing radiation pressure on dust grains in molecular filaments and cores. Our neglect of ionizing radiation is appropriate not only for low mass star forming regions but also where ionizing sources are strongly extincted. Similarly, we neglect some radiative forces on dust grains, specifically the photoelectric and photodesorption forces (Weingartner and Draine, 2001). These forces – which arise when UV photons remove electrons or atoms, respectively, from the surface of a dust grain – are more important in the diffuse interstellar medium, e.g. in the cold neutral medium, than in the dense molecular regions we consider.

We assume an isotropic background radiation field, where the anisotropy required for a radiation pressure force is introduced by an opaque cylinder or sphere. In reality, the background radiation will be anisotropic, though the ISRF in the Solar vicinity is only asymmetric at the 10% level (Weingartner and Draine, 2001). For highly anisotropic irradiation from a nearby star or association, our results will not apply directly. However, anisotropic non-ionizing irradiation may still promote star formation in a similar way as ionizing irradiation from an OB association can trigger star formation (Bisbas et al., 2011). By considering the isotropic component of radiation fields here, future studies can determine the effects of more complex irradiation.

Observed cores and filaments are not the perfect spheres or cylinders used in our study. Observed dense cores are spheroids with a typical axis ratio of $\leq 2:1$ (Myers et al., 1991b). The BE stability of general spheroids is similar to that of perfect spheres. In particular, Lombardi and Bertin (2001) found that the maximum density contrast of a stable, isothermal cloud is independent of shape. While filaments are not perfect cylinders, their length-to-width ratios are typically $>10:1$ (André et al., 2010, 2014; Men’shchikov et al., 2010; Hacar et al., 2013). Our idealized geometries are consistent with the isotropic irradiation and provide a reference for more complex geometries.

Our two-ray approximation to the extinction of radiation simplifies the detailed
angular dependence of the radiation field, i.e. we neglect rays that obliquely intersect the radius vector. In reality, a radiation field that is isotropic for all incoming angles at the surface would become more radially directed with depth (due to the increased slant optical depth of oblique rays). Figure 5.9 tests our approximation by comparing the two-ray approximation to a detailed extinction calculation that includes the angular dependences. The detailed calculation shows only a minor reduction in flux, of at most a few percent. Thus the two-ray approximation is acceptable, especially in the context of other, more severe, approximations.

The isothermal structure in our models is only roughly consistent with observations, which find a factor of two variation in temperature within filaments and dense cores. Central temperatures (∼8 K) are typically lower than outer regions (∼20 K) due to interstellar radiation heating (Evans et al., 2001; Zacconi et al., 2001; Pagani et al., 2004; Shirley et al., 2005; Crapsi et al., 2007; Launhardt et al., 2013; Palmeirim et al., 2013; Seo et al., 2015). With a more realistic temperature structure, the critical line density of filaments increases by 20 – 30% (Recchi et al., 2013). Non-isothermal spheres show a similar correction Sipilä et al. (2011). A more complex model with detailed heating and cooling could self-consistently determine both the temperature structure and the radiation pressure profile for a given level of irradiation.

We neglect rotation, which is probably a weak dynamical effect. Observations of velocity variations across a filament or a dense core indicate that rotational energy is typically less than 5% of the internal or gravitational energy (e.g. Arquilla and Goldsmith, 1986; Goodman et al., 1993; Caselli et al., 2002a).

We assume perfect coupling between dust grains and gas, neglecting grain sedimentation. Since radiation pressure dominates gravity at the surface of our objects (for most adopted irradiation levels), we estimate the terminal speed of dust grains as

\[ v_{\text{terminal}} = \frac{3(Q_{\text{pr}})F_0}{4\rho g c_s c}, \]  

using the Epstein drag law appropriate for dilute gases (Youdin, 2010). The terminal velocity is independent of grain size since radiation pressure and drag forces both
scale with the grain cross section. Figure 5.10 shows the terminal velocity and the settling time (across a typical scale of 0.1 pc) versus level of irradiation. The red lines in the plots correspond to the sound speed and sound crossing time. With high levels of irradiation and low gas densities (appropriate near the surface) significant sedimentation could occur. This tendency to settle is counteracted by mixing due to turbulent diffusion or other large scale motions (Youdin and Lithwick, 2007). Sedimentation will also be reduced for more optically thin objects that experience flux cancellation. Sedimentation of grains could alter dynamical stability by piling dust grains and concentrating the radiation force in a thin shell, and could also affect observed extinction profiles (Whitworth and Bate, 2002). Time-dependent numerical simulations are likely required to explore the full effects of grain sedimentation and its dynamical effect on dense structures.

5.6 SUMMARY & CONCLUSIONS

We study the gravitational stability of hydrostatic cylinders and spheres bathed in an isotropic, non-ionizing radiation field. We find that the radially-inward radiation pressure force promotes – and can trigger – the gravitational collapse of cores and filaments. The classic stability thresholds – the critical mass of Bonnor-Ebert spheres and the critical line density of isothermal Ostriker cylinders – are significantly lowered once the surface radiation pressure reaches the magnitude of the central gas pressure. Thus a given level of irradiation more strongly affects objects with a lower central density, provided the column density is large enough that the optical depth is at least an order unity. Our analysis shows that the critical state of highly irradiated spheres or filaments is characterized by an order unity optical depth. Physically, these highly irradiated objects must become partially transparent to incident irradiation to avoid implosion.

Standard interstellar radiation fields are strong enough to influence gravitational stability. For instance, consider a spherical molecular clump with a central density of $\sim 10^3$ cm$^{-3}$ at 10 K. The maximum stable mass of such an object is 14 M$_\odot$ without
Figure 5.9 Comparison of a two-ray approximation to a full radiative transfer calculation. A flux of radiation field is shown as a function of radius within the critical Bonnor-Ebert sphere. Total extinction, $A_V$, is 10 for this example. Only absorption by dust grains is considered while scattering is neglected. The black line denotes a full radiative transfer calculation, and the red line denotes a two-ray approximation. The difference between the two fluxes within the sphere is <3% of the flux at the sphere surface.
Figure 5.10 Plot of the terminal speed (top) and settling time across 0.1 pc (bottom) for dust grains accelerated by radiation pressure (flux level on the x-axis) and damped by gas drag for different densities. Horizontal reference lines (in red) show the sound speed (for $T = 10$ K) and sound crossing time (over 0.1 pc). For high fluxes and low gas densities, short settling times allow grain sedimentation. However, internal motions could counteract sedimentation.
radiation (the Bonnor-Ebert mass) and only 5 Mₜₜ if subject to interstellar irradiation of the Solar vicinity. For more evolved, i.e. denser, cores with central densities \( \sim 10^3 \text{ cm}^{-3} \) a stronger radiation field is needed to affect gravitational stability. These stronger radiation fields can be found towards the Galactic center or in active star forming regions near YSOs or stellar associations.

The gravitational stability of molecular filaments is similarly affected by irradiation. A young filament with a central density of \( \sim 10^3 \text{ cm}^{-3} \) at 10 K has a maximum line density of 16 Mₜₜ/pc without radiation, the Ostriker value. Interstellar irradiation in the Solar vicinity can lower the critical line density to only 9 Mₜₜ/pc.

At face value, our results imply that the initial mass function (IMF) of stars could vary with the radiation environment. Alternatively, we propose a mechanism to maintain a universal IMF with non-ionizing radiation pressure. Imagine that in some star forming region there is a temporary over- (or under-) production of massive stars, for either physical or stochastic reasons. Our results show that the resulting enhancement (or reduction, respectively) in irradiation can then reduce (or increase, respectively) subsequent fragmentation masses. The triggering of star formation by the evolving radiation pressure could thus help regulate the IMF.

The full implications of non-ionizing radiation pressure are not yet clear due to neglected effects in our model. Additional support from magnetic fields or strong turbulence will counteract the destabilizing effects of radiation pressure. Furthermore, our static models neglect crucial evolutionary and dynamical effects, notably the processes that set the mass spectrum of prestellar cores. Nevertheless, our results demonstrate that non-ionizing radiation pressure is strong enough to influence both star formation and filamentary structure, so its full implications should be explored in future work.
6.1 Summary & Conclusions

In this thesis, I have used both observational and theoretical approaches to study the evolution of dense starless cores. I shall first summarize the observational results followed by the theoretical results.

We present extensive NH$_3$ (1,1) & (2,2), CCS, and HC$_7$N maps of the L1495-B218 filaments extending over 3 degrees on the sky in the Taurus molecular cloud with unprecedented depth (median rms of 154 mK with the lowest rms of 26 mK for CCS, median rms of 120 mK for HC$_7$N), angular resolution (31$''$), and velocity resolution (0.038 km/s). Using the maps, we studied the physical and chemical properties of the L1495-B218 filaments and compared with theoretical star forming scenarios in filamentary structures in molecular clouds. The main results are as follows:

1. Using our ammonia observations and the protostar catalog from Spitzer Space Telescope observations (Rebull et al. 2010), we confirm that the L1495-B218 filaments are at various evolutionary stages: B10, B211, and B216 are relatively young regions, while L1495/B7, B213, and B218 are actively star-forming, older, more evolved regions.

2. Gas kinetic temperatures in the L1495-B218 filaments deduced from NH$_3$ (1,1) and (2,2) lines reveal that the filaments to be very cold (8 – 15 K with median value equal to 9.5 K, 60% of spectra are in the range of 9.5 ± 1 K). The nonthermal velocity dispersions of NH$_3$ (1,1) & (2,2) and CCS lines in the L1495-B218 filaments (<0.15 km s$^{-1}$) are typically less than the thermal velocity and are considerably narrower than those of C$^{18}$O $J = 1–0$ (∼0.15 km s$^{-1}$; Hacar et al. 2013) and $^{13}$CO $J = 1–0$ (∼1 km/s; Goldsmith et al. 2008) in the same region.

3. 39 NH$_3$ peaks (leaves) and 16 nested groups (branches) are identified using
the CSAR clump-finding algorithm. The NH$_3$ leaves and branches are identified on a scale from 0.01 pc to 0.1 pc and have masses ranging from 0.05 M$_\odot$ to 9.5 M$_\odot$. NH$_3$ peaks have good spatial agreement with dust continuum peaks, which suggests that NH$_3$ is an excellent tracer for finding dense cores in the L1495-B218 filaments.

4. Most of NH$_3$ leaves are gravitationally unbound. Nine NH$_3$ leaves are found to be gravitationally bound and that 7 out of 9 leaves are either protostellar or within branches associated with star formation activity, while 30 NH$_3$ leaves are gravitational unbound and only 10 out of 30 unbound leaves are either protostellar or within branches associated with star formation. We also found that 12 out of 30 gravitationally unbound leaves are pressure-confined.

5. Our CCS survey shows that CCS intensity peaks do not spatially agree with NH$_3$ and dust intensity peaks in most of dense cores. Only one out of 10 dense cores that have both NH$_3$ and CCS emission have intensity peak that coincide in NH$_3$ and CCS. CCS emission in more evolved regions, L1495A and B213E, show an arc-like or a ring-like geometry rather than a core-like shape, while only NH$_3$ core no.32 have core-like CCS emission. These examples show CCS emission rarely traces a core structure.

6. CCS does not trace dynamical evolutionary stages of dense cores. CCS is bright in the most concentrated dense cores in L1495A and L1521D, while less evolved starless cores in B213W and B10 has no noticeable CCS emission. The youngest NH$_3$ cores in B216 have CCS emission 0.1 pc to the south.

7. The hub (L1495A) of the L1495-B218 filaments is likely to be a stellar association or cluster forming region. The hub region contains multiple young stellar objects and it is currently accreting gas from large scale flows. Particularly, we found that the dense core L1495A-N is likely formed by a supersonic large scale converging flow.

We studied the dynamical stability of a dense core and a filament under external perturbations to have more realistic dynamical stability conditions, which can be used in a complex environment (e.g. L1495-B218 filaments). In this thesis, we mainly considered two external perturbations: a converging flow, which may be
originated large scale turbulent flows, and non-ionizing radiation including stellar radiation from nearby YSOs and the interstellar radiation field. The main results are as follows:

8. A dense core with an inward (converging) motion is more unstable than a dense core in a quiescent environment. We found that a homologous inward motion with the peak inward motion speed of the sound speed has the critical size smaller than half of the critical Bonnor-Ebert sphere.

9. The infall speed of a spontaneously collapsing core by self-gravity without any external perturbation has a unique peak speed for a given density structure. The gravitational collapse of the critical Bonnor-Ebert sphere shows the lowest collapse speed at a given density concentration, and the gravitational collapse of a uniform sphere shows the fastest collapse speed at a given density concentration. As an example, we found that L694-2 is collapsing considerably faster than the fastest infall speed of gravitational collapse, which suggests that it is perturbed to enhance its infall speed by an external perturbation.

10. A dense core or a filament bathed in an isotropic, non-ionizing radiation field is more unstable than a structure without any irradiation. We find that the radially-inward radiation pressure force promotes – and can trigger – the gravitational collapse of cores and filaments. For structures with moderate central densities, \( \sim 10^3 \text{ cm}^{-3} \), the interstellar radiation field in the Solar vicinity has an order unity effect on stability thresholds.

11. For more evolved objects with higher central densities, a significant lowering of stability thresholds requires stronger irradiation, as can be found closer to the Galactic center or near stellar associations. Even when strong sources of ionizing radiation are absent or extincted, our study shows that interstellar irradiation can significantly influence the star formation process.
6.2 Future Work

Through our survey in NH$_3$, CCS, and HC$_7$N emission toward the L1495-B218 filaments in the Taurus molecular cloud, we studied the star-forming processes from a molecular cloud to a dense core. The star-forming process from a dense core to the formation of a first hydrostatic core, where the fragmentation of a dense core and disk formation are expected in a ≤10,000AU scale, is not studied through our survey due to lack of a high enough spatial resolution to resolve 10,000 AU in the nearest star-forming regions. We proposed a pilot survey in $J = 1−0$ transition of HCN, HCO$^+$, and N$_2$H+ using ARGUS on 100m GBT toward dense cores in L1495-B218 filaments. The GBT ARGUS will provide a high spatial resolution (1000 AU at Taurus), which is one-tenth of a typical wide binary separation (Tobin et al., 2016). HCN and HCO$^+$ 1−0 trace an infall motion in a dense core and N$_2$H+ traces dense gas for possible fragmentation.

Our studies on the dynamical stability conditions made the first step to understand the collapse dynamics of a dense core or a filament affected by non-ionizing radiation and a converging flow. We hope to follow up on the non-linear processes of triggered star formation by non-ionizing radiation using a numerical method. We already have developed a three dimensional hydrodynamics code describing the motions of dust grains as super-particles. We will add a simple radiative transfer description to the code and simulate radiation-dust-gas interactions. This future study will help to understand triggered star formation in a region with considerable radiation strength such as near HII region, PDR, in the molecular ring, or the Galactic center.
We derived physical quantities through finding the minimum reduced \( \chi^2_r \) value between the observed \( \text{NH}_3 \) (1,1) and (2,2) spectra and the spectra model. Our method of finding the minimum \( \chi^2_r \) value is a bit different from previous work in that we do not use a \( \chi^2_r \) minimization method that utilizes a covariance matrix to calculate the uncertainties but instead explore a wide range of parameter space in adaptive grids to obtain the global minimum and determine the shape of \( \chi^2_r \) volume as functions of the parameters. The common method of minimizing \( \chi^2_r \) value utilizes a covariance matrix because the method is relatively easy and computationally fast, but parameters should be random variables. If a parameter has a specific non-random distribution, the 1-\( \sigma \) uncertainty of that parameter may be not symmetric and may differ from the standard deviation estimated by the covariance matrix.

Exploring parameter space gives proper estimation of uncertainties, but it is usually computationally expensive if the parameter space is large and is searched at high resolution. To reduce computational time, we use the adaptive mesh refinement method (AMR method). First, we set up ranges of parameter space to explore \([x_{0,\text{lower}}, x_{0,\text{upper}}]\), where the subscript 0 denotes zeroth layer, the subscript lower means the lower limit of the parameter space which we are exploring, and the subscript upper means the upper limit. The number of grids in each parameter can be chosen arbitrarily, but for efficiency we used fewer than 20 grids for each parameter. Second, we calculate the minimum \( \chi^2_r \) at given grid points and find its location \( x_{0,\text{min}} \). Third, we set up new range of parameter space and explore as the first layer \([x_{1,\text{lower}}, x_{1,\text{upper}}]\). The range can be given as

\[
x_{1,\text{lower}} = x_{0,\text{min}} - \frac{x_{0,\text{upper}} - x_{0,\text{lower}}}{2} \quad \text{when} \quad x_{1,\text{lower}} \geq x_{0,\text{lower}}
\]  

(A.1)
and

\[ x_{1,\text{upper}} = x_{0,\text{min}} + \frac{x_{0,\text{upper}} - x_{0,\text{lower}}}{2} \]

when \( x_{1,\text{upper}} \leq x_{0,\text{upper}} \). 

(A.2)

If \( x_{1,\text{lower}} \) is smaller than \( x_{0,\text{lower}} \), then \( x_{1,\text{lower}} = x_{0,\text{lower}} \). Likewise, if \( x_{1,\text{upper}} \) is larger than \( x_{0,\text{upper}} \), then \( x_{1,\text{upper}} = x_{0,\text{upper}} \). Fourth, we estimate the \( \chi^2_r \) value within the new range of parameter space with the same number of grids we used in the first step. At this step, the resolution is double that of the previous layer. Fifth, we repeat the second to the forth steps until we achieve the desired resolution. If there is the global minimum of \( \chi^2_r \) value within the range of the zeroth parameter space, this code can find the global minimum with an uncertainty of the resolution of the final layer unless the widths of local minima and global minimum are much narrower than the resolution of the final layer. In our calculations, we set the finest velocity resolution of the model to be at least quarter of the velocity resolution of our observation, so we may not miss the global minimum.

For the model of the NH\(_3\) line profiles, we use the model described in Rosolowsky et al. (2008).

Figure A.1 shows two examples of fitting the observed NH\(_3\) spectra and Figure A.2 shows the \( \chi^2_r \) of the two examples in \( T_k \) vs. \( \tau_1 \) space. The top panel of Figure A.1 is the best-fit of the simple emission model to an observed strong NH\(_3\) (1,1) emission of \( \sim 32 \sigma_{T_{\text{mb}}} \) and the left panel of Figure A.2 is the reduced \( \chi^2 \) of fitting in \( T_k \) vs. \( \tau_1 \) space with the other three parameters fixed at the best-fit values. The bottom panel of Figure A.1 shows an example of a weaker emission spectra of NH\(_3\) (1,1) with the peak signal-to-noise ratio (SNR) being \( \sim 7.5 \) and the right panel of Figure A.2 is the reduced \( \chi^2 \) of fitting in \( T_k \) vs. \( \tau_1 \) space with the other three parameters fixed at their best-fit values. Uncertainties of physical parameters are about three times larger in the weak line than those in the strong line.

We fit 13,109 spectra which satisfy the detection of NH\(_3\) (1,1) and (2,2) emission is above \( 5 \sigma_{T_{\text{mb}}} \) and \( 3 \sigma_{T_{\text{mb}}} \), respectively, where \( \sigma_{T_{\text{mb}}} \) is the baseline \( \text{rms} \). The parameter ranges for searching for the best-fit are \([7.5 \text{ K}, 30.0 \text{ K}] \) for \( T_k \), \([2.75 \text{ K}, 30 \text{ K}] \) for \( T_{\text{ex}} \), \([5 \times 10^{-3}, 20] \) for \( \tau_1 \), and \([0.05 \text{ km s}^{-1}, 0.3 \text{ km s}^{-1}] \) for \( \sigma_v \). The \( v_{\text{lsr}} \) is
preliminarily estimated from the peak intensities of NH$_3$ (1,1) and (2,2) lines and a more accurate value is searched for in a range of ±0.1 km s$^{-1}$ from the guessed value. Typical 1-σ uncertainties are ±0.77 K for $T_k$, ±0.85 K for $T_{ex}$, ±0.26 for $\tau_1$, ±0.013 km s$^{-1}$ for $\sigma_v$, and ± 0.020 km s$^{-1}$ for $v_{lsr}$. The overall uncertainty distributions of 13,109 fits to the spectra are shown in Figure A.3.
Figure A.1 Examples of spectra model fitting (red dashed lines) to observed lines (black solid lines). The top panel shows NH$_3$ (1,1) lines having a peak signal to noise ratio (SNR) equal to 32, and the bottom panel shows NH$_3$ (1,1) lines having a peak signal to noise ratio equal to 7.5.
Figure A.2 Examples of reduced the $\chi^2$ space in $T_k$ vs. $\tau_1$. The left panel shows the $\chi^2$ space of fitting the line with a peak SNR of 32, and the right panel shows $\chi^2$ space of fitting the line with a peak SNR of 7.5. The white contours are 1-$\sigma$ uncertainty limits.
Figure A.3 Distributions of 1-σ uncertainties in the spectra fittings. Vertical lines represent the median values.
APPENDIX B

MODELING APPARENT AXIS RATIO DISTRIBUTIONS

We analyzed morphology of dense structures through estimating the minimum $\chi^2$ value (not reduced $\chi^2$) between the observed distribution of apparent axis ratios and distribution models. The distribution models are made using a Monte Carlo method. For each model, we imposed ten million spheroids and assumed all of the spheroids are either in oblate or prolate shapes. The apparent axis ratios of the spheroids are determined by the intrinsic axis ratio of the spheroid and the angle between observer’s line of sight and the equatorial plane of the spheroid as follow:

\[ p(q, \theta) = \left[1 - (1 - q^2) \cos \theta\right]^\frac{1}{2} \text{ for oblate,} \]  
\[ p(q, \theta) = q \left[1 - (1 - q^2) \sin \theta\right]^\frac{1}{2} \text{ for prolate,} \]  

where $p$ is the apparent axis ratio of the spheroid, $q$ is the intrinsic axis ratio of the spheroid (short axis/long axis), and $\theta$ is the angle between observer’s line of sight and the equatorial plane of the spheroid. We assumed that the intrinsic axis ratios of the spheroids follow a Gaussian distribution and the angles between the observer’s line of sight and the equatorial planes of spheroids have random values with a uniform distribution within $[-\theta_{\text{limit}}, \theta_{\text{limit}}]$. The mean value and width of a Gaussian distribution and $\theta_{\text{limit}}$ serves as input parameters of the model. Using above equations with given input parameters, we modeled a distribution of apparent axis ratios of the spheroids. We made 400,000 models with different mean values $\langle q \rangle$ and widths $\sigma_q$ of Gaussian distributions of the intrinsic axis ratios and $\theta_{\text{limit}}$. The range of $\langle q \rangle$ that is exploited in our models is from 0.01 to 1 with an interval of 0.01 and the range of $\sigma_q$ is from 0.01 to 1 with interval of 0.01. The range of $\theta_{\text{limit}}$ is from 0 degree to 90 degree with an interval of 5 degree. We compared all 400,000 distribution models to the observed distribution and found the best-fit model by minimizing the $\chi^2$ value.
APPENDIX C

THE DIMENSIONLESS BONNOR STABILITY

We now show that the dimensionless stability criteria of equation (5.19) applies even in the presence of radiation pressure, or for that matter any additional spherically symmetric force. The existence of a stability boundary at \( \frac{dm}{ds_0} = 0 \) can be derived, for instance, from an analysis of the free energy of the system (Stahler, 1983). We choose a simpler path and show that the dimensional Bonner instability condition in physical units, \( \frac{dP_0}{dV_0} > 0 \) at fixed mass, remains equivalent to \( \frac{dm}{ds_0} > 0 \) and thus equation (5.19).

First, we express the dimensional outer radius, \( r_0 = r(s_0) = \alpha \xi_0 \), in terms of the dimensionless mass, \( m \), as

\[
    r_0 = \frac{GM}{c_s' s_0} \sqrt{\frac{\xi_0}{m}} \quad (C.1)
\]

using equation (5.18).

Next we consider perturbations to hydrostatic solutions at fixed dimensional mass, \( M \), and temperature, i.e. \( c_s' \). Equations (5.18) and (C.1) give

\[
    d \ln(P_0) = 2d \ln(m), \quad (C.2a)
\]

\[
    d \ln(r_0) = d \ln(\xi_0) + \frac{1}{2}d \ln(s_0) - d \ln(m). \quad (C.2b)
\]

We can now relate the Bonnor stability criterion to dimensionless variables as

\[
    \frac{d \ln(P_0)}{d \ln(V_0)} = 3 \frac{d \ln(P_0)}{d \ln(r_0)} = 6 \frac{d \ln(m)/d \ln(s_0)}{d \ln(\xi_0)/d \ln(s_0) + \frac{1}{2} - \frac{d \ln(m)}{d \ln(s_0)}}. \quad (C.3)
\]

Thus a stability transition at \( \frac{dP_0}{dV_0} = 0 \) also gives a transition at \( \frac{dm}{ds_0} = 0 \). The desired sign of the stability criterion for \( \frac{dm}{ds_0} \) requires that the denominator in the final term of equation (C.3) be positive.

We verified numerically that this denominator remains positive for our solutions, but a more general argument proceeds as follows. First, note that this denominator
is positive for the classic Bonnor-Ebert solution, as it must be since the correspondence between the stability criteria holds in this case. Next, consider that any finite radiation pressure (or other force) can be increased incrementally from zero to produce a continuous family of solutions that starts with the Bonnor-Ebert solution (as visualized in Figure 5.2). Since none of these incremental steps can produce an infinite divergence in $dP_0/dr_0$, the denominator in question cannot change sign, and the desired correspondence between dimensionless and dimensional stability criteria holds.
REFERENCES


