

MAGISTERARBEIT

Titel der Magisterarbeit

**THE CONNECTION BETWEEN THE LOCAL BUBBLE
AND THE ^{60}Fe ANOMALY IN THE DEEP SEA
HYDROGENETIC FERROMANGANESE CRUST**

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Abstract

The deep ocean's crust 237KD, analyzed by Knie et al. 2004 [24], shows a significant increase of the abundance of the radioisotope ^{60}Fe 2.2 Myr ago. ^{60}Fe is mainly produced in supernova explosions. This indicates that one or more supernovae must have gone off in the solar vicinity to eject enough material to be deposited on Earth.

The Local Bubble is an X-ray emitting, HI deficient, cavity in the local interstellar medium (ISM). By analyzing young stars in the radius of 200 pc around the Sun, Fuchs et al. 2006 [17] found that the Local Bubble was formed by 14 - 20 supernova explosions starting around 13 Myr ago. The remaining stars of the corresponding stellar moving group belong today to the subgroups Upper Centaurus Lupus and Lower Centaurus Crux of the Scorpius OB2 association. Following the trajectories of this cluster back in time, we find a minimal distance to Earth of about 65 pc around 2.2 Myr ago, which coincides with the appearance of the ^{60}Fe peak.

In this work we also test a supernova model developed by Kahn in 1998 [22]. Due to the occurrence of not only one but many supernova explosions the ISM does not have a homogeneous density distribution around the explosion center, because the medium has already been disturbed by previous supernova explosions. With Kahn's method we compute analytically the time the remnant takes to hit the Earth and how much ^{60}Fe is deposited.

Thus for the first time an ab initio model for the ^{60}Fe deposition has been compared to the observations. Our results are in very good agreement with the measurements.

Zusammenfassung

In der Tiefseemangankruste 237KD im Pazifischen Ozean wurde 2004 von Knie et al. [24] ein vermehrtes Vorkommen des Radionuklids ^{60}Fe vor 2.2 Millionen Jahren gemessen. Dieses Element wird hauptsächlich in der Endphase von massereichen Sternen und während einer Supernovaexplosion erzeugt. Deshalb wird angenommen, dass eine oder mehrere Supernovaexplosionen in der Umgebung des Sonnensystems stattgefunden haben, um diese Isotopenanomalie auf der Erde zu verursachen.

Die Lokale Blase ist durch ein geringes Auftreten von neutralem Wasserstoff im lokalen interstellaren Medium gekennzeichnet und emittiert im weichen Röntgenbereich. Junge, massereiche Sterne in einem Radius von 200 pc um die Sonne wurden von Fuchs et al. 2006 [17] analysiert. Dabei stellte sich heraus, dass sich in der heutigen Region der Lokalen Blase etwa 14 - 20 Supernovae in den letzten 13 Millionen Jahren ereignet haben. Verbleibende Sterne der Stellarstromgruppe gehören heute zu den Untergruppen Upper Centaurus Lupus und Lower Centaurus Crux des Sternhaufens Scorpius OB2. Die Bahnen dieser Sterne wurde zeitlich zurückgerechnet, und es ergab sich, dass der Minimalabstand der Bahnen zur Sonne etwa 65 pc vor 2.2 Millionen Jahren betrug, was zeitlich gut mit der Erhöhung der ^{60}Fe Konzentration übereinstimmt.

In der vorliegenden Arbeit wird ein Supernovamodell getestet, welches 1998 von Kahn [22] erstellt wurde. Da es sich nicht nur um eine, sondern um mehrere Supernovae handelt, können wir nicht annehmen, dass das Interstellare Medium um das Explosionszentrum herum eine homogene Dichteverteilung aufweist. Die Materie sollte durch vorangegangene Explosionen schon in einer Schale aufgesammelt worden sein. Mit Kahns Methode wird in dieser Arbeit berechnet, wie lange der Supernovaüberrest

braucht, um die Erde zu erreichen, und wieviel ^{60}Fe sich dabei im Ozean ablagert.

In dieser Arbeit wird zum ersten Mal ein ab initio Modell für die ^{60}Fe - Verteilung mit den Beobachtungen verglichen. Es gelingt uns, die gefundene Isotopenverteilung analytisch zu reproduzieren; die Berechnungen stimmen sehr gut mit den Messungen überein.

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CHAPTER 1

Introduction

This chapter will answer the questions that might have arisen by reading the title of this thesis. It will explain the keywords and the purpose of this work.

1.1 The Local Bubble

The Local Bubble (LB) is a cavity of thin hot gas in the local interstellar medium (ISM). It extends about 600 pc in the galactic plane and 200 pc perpendicular to it. The LB is about 13 Myr old and the medium inside has an average density of 5×10^{-3} atoms cm^{-3} and a temperature of 10^6 K ([4],[17]). It is deficient of HI and emits in the soft X-ray spectrum. Our Sun is located inside the LB, but in a less hot and less dense region called the Local Fluff. This is an interstellar cloud with a density of about 0.1 atoms cm^{-3} and a temperature of 7000 K. In about 70000 yr the Sun will leave this cloud [7]. The first hints towards the existence of the LB were discovered by Stuart Boyer in the 60s. He found that soft X-rays are emitted from all directions in space. This was later confirmed by the satellites ROSAT and XMM-Newton. By measuring the absorption spectra of background stars and molecular clouds, it was later observed that there is also a deficiency of HI in this region [7]. Other evidence was found by Sfeir et al. in 1999 [43] and Lallement et al. in 2003 [27]. They analyzed absorption lines of NaI from background stars and found that in the solar vicinity the absorption is very low until a distance of about 100 pc and then the absorption increases considerably. They concluded that in this distance a shell is located where the density increases abruptly. Towards the galactic north pole no definite

border was found, so it was suggested by Welsh et al. [48] that the LB opened and a chimney was formed.

The LB was excavated by multiple supernova (SN) explosions by a stellar moving group [17]. They heated up the medium, pushed the ISM away and compressed it into a shell. The remaining stars of this moving group belong today to the Scorpius OB2 association, a young star cluster. Further information is provided in Chapter 3.

There are other superbubbles in our galaxy, for example the Loop I, which is our neighboring bubble. Both bubbles formed approximately at the same time and they are separated by a wall (see Fig. 1.1). It is estimated that 14 -

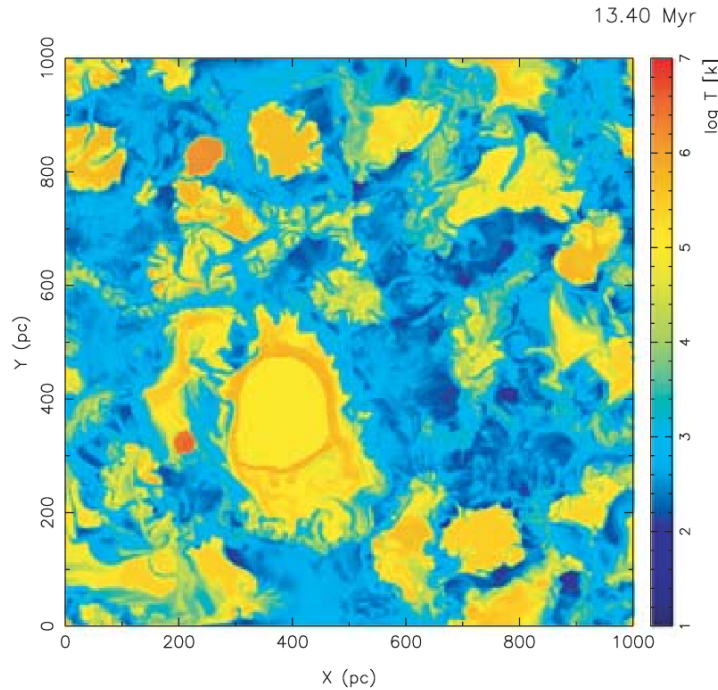


Figure 1.1: A temperature distribution of the ISM today in the galactic midplane. We see the Local Bubble and the Loop I bubble that were both caused by multiple supernova explosions and are separated by a shell. This numerical simulation was carried out by Breitschwerdt and de Avillez in 2006 [8]. They found that both bubbles will probably merge in the future.

20 SN exploded in the LB [17] and 38 in Loop I. Unlike in the Loop I, where there are still ongoing supernova explosions (about 39 expected), no active cluster with early type stars was found in the region that is occupied by the LB today. This probably leads to a merging of the two bubbles in the future [8].

1.2 A hydrogenetic ferromanganese crust

Ferromanganese crusts are found on the bottom of the ocean, primarily on sea-mountains and plateaus, on deep-sea volcanoes and the mid-ocean ridges in a depth of 400 - 5000 m. As the name suggests, these crusts contain large amounts of Mn and Fe, but also other elements like Ti, Co, Ni etc.

We distinguish between three types of manganese crusts, which differ in their formation processes. These types are the hydrothermal, the hydrogenetic and the diagenetic manganese crust.

The hydrothermal manganese crust forms by hot fluids rich on minerals that escape from thermal springs on the ocean's floor. We know these phenomena as black or white smokers. The elements settle on the surface and build a vent with a growth rate of some m/Myr.

The hydrogenetic ferromanganese crust is the type that is relevant for this thesis, as we will see in Chapter 2. It grows by obtaining its composition of elements out of the ambient water and contains around 20 % or more of Mn and around 15 % of Fe. It has a very low growth rate with less than 10 mm/Myr and is the final product of sedimentation processes. This is the reason why they are much more dense (about 1.3 g cm^{-3}) than a marine sediment with an average growth rate of 2 - 3 mm/Myr [32]. Fig. 1.2 shows such a hydrogenetic crust.



Figure 1.2: A hydrogenetic ferromanganese crust. Courtesy of oceanexplorer.noaa.gov.

The third type is the diagenetic manganese crust, which grows by gaining the elements from the surrounding sediment and water. It grows faster

than a hydrogenetic crust with a growth rate of 10 - 200 mm/Myr.

The information presented in this section was taken from M. Poutivtsev [36], who measured the ^{53}Mn concentration in different manganese crusts and confirmed the half-life of this radionuclide in his PhD thesis.

1.3 The radionuclide ^{60}Fe

The isotope ^{60}Fe is an unstable radionuclide with the recently measured half-life of 2.62 Myr [39]. It is produced by nucleosynthesis in massive stars, mainly by the r- and s-process (f. ex. [5]) and during the supernova explosion. The dependence of the ^{60}Fe yield on the mass of the stars is discussed in the subsequent section 1.4. Following Limongi and Chieffi 2006 [29] we will describe the production in more detail.

The major contribution occurs from the slow neutron capture on ^{59}Fe , which is an unstable isotope with a half-life of about 44 days. Since it thus undergoes a β^- -decay, the temperature must be high enough to produce enough neutrons for the $^{59}\text{Fe}(n,\gamma)$ process. On the one hand there is a temperature limit of 2×10^9 K. Above this value, ^{60}Fe will be destroyed by the (γ,n) and the (γ,p) photodisintegrations. On the other hand the temperature in the central-burning phases of He and C is too low, less than 4×10^8 K, to produce enough neutrons. A neutron density higher than $n = 10^7 \text{ cm}^{-3}$ is needed for an efficient production of ^{60}Fe . This is why it is primarily synthesized in the shell-burning of He and C. Only a little amount arises from Ne burning.

In the shell-burning phase the temperature is of the order of 4×10^8 K and higher for He and $\geq 10^9$ K for C. The neutrons are mainly produced by the $^{22}\text{Ne}(\alpha,n)^{25}\text{Mg}$ reaction and reach densities of $n \geq 3 \times 10^{10} \text{ cm}^{-3}$ in the He burning shell and $6 \times 10^{11} \leq n \leq 2 \times 10^{12} \text{ cm}^{-3}$ in the C burning shell (for stars with masses of 20 - 120 $M_{\odot}$).

The $^{23}\text{Na}(\alpha,n)^{26}\text{Al}$ reaction in the Ne burning phase produces large amounts of neutrons, but since there is no stable convective shell that lasts until the explosion, no significant amount of ^{60}Fe is synthesized.

The final contribution to the synthesis of ^{60}Fe emerges from the supernova explosion. The blast wave moves through the star and produces ^{60}Fe in the regions where the temperature is of the order 2.2×10^9 K. This occurs

mostly in the base or in the C convective shell.

1.4 The ^{60}Fe yields in massive stars

As mentioned before, the abundance of ^{60}Fe depends on the mass of the star. Numerical simulations to calculate the ^{60}Fe yield for different types of stars have been carried out by several research groups, for example Woosley & Weaver 1995 [51], Rauscher et al. 2002 [38], Limongi & Chieffi 2006 [29], Karakas & Lattanzio 2007 [23] and Woosley & Heger 2007 [50]. Most of them cover the massive stars in the mass range of 11 - 120 $M_{\odot}$; Karakas & Lattanzio computed the yields for AGB stars in the mass range of 1 - 6 $M_{\odot}$. Fig. 1.3 shows some selected publications to see how the results differ from each other.

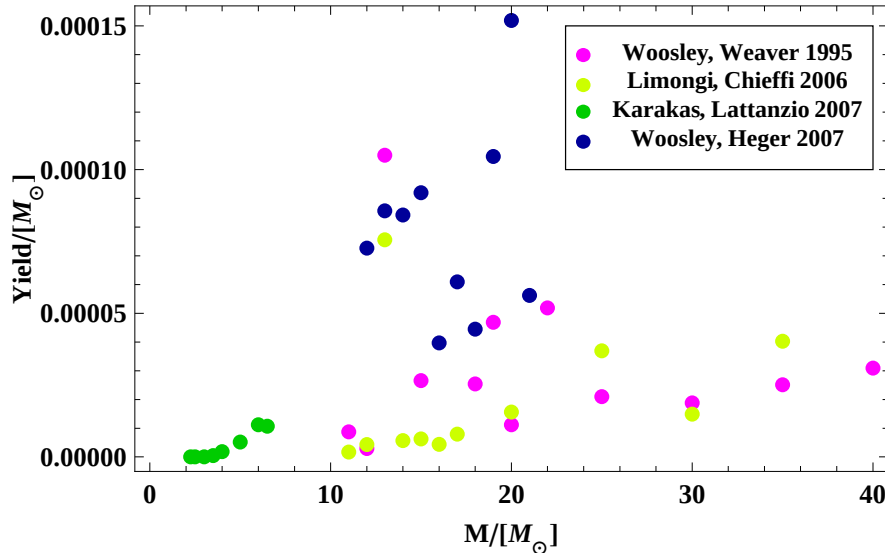


Figure 1.3: Dependence of ^{60}Fe yield on the stellar mass based on different computations.

The large range in ^{60}Fe yields for different calculations results from the sensitivity of nucleosynthesis to a lot of factors. The amount of ^{60}Fe produced depends on the stellar structure and its initial composition of elements, the location of the convective shells and on the interplay and extension of the convective shells [51]. There still are uncertainties in the cross section of elements and in the reaction rate of the afore mentioned $^{22}\text{Ne}(\alpha, n)^{25}\text{Mg}$ process, which among other reactions determines the strength of the s-process [50]. It is important whether processes like mass loss by stellar

wind, neutrino loss, mixing (for example due to Rayleigh–Taylor instabilities) or rotation are included and whether the full evolution from the end of the main sequence to the explosion is covered or only parts of it [29]. And yet the full explosion mechanism is still not understood exactly. An example is the mass cut, which means how much of the mass is ejected into space and how much falls back on the star.

Since the work by Woosley & Heger [50] and Karakas & Lattanzio [23] is the most recent one, we will use them for further calculations. We will see later in this work, that the stars with masses between 8 - 11 $M_{\odot}$ are the most important ones, but unfortunately the associated yields are still missing as it is shown in Fig. 1.4. For this work we fitted them linearly, which is a bit risky, because the nucleosynthesis processes vary also for different types of stars like AGB or the missing S-AGB. For now it is the only possibility to derive the yields, but for future work they should be available as a research group (Lattanzio et al.) is already computing these missing yields.

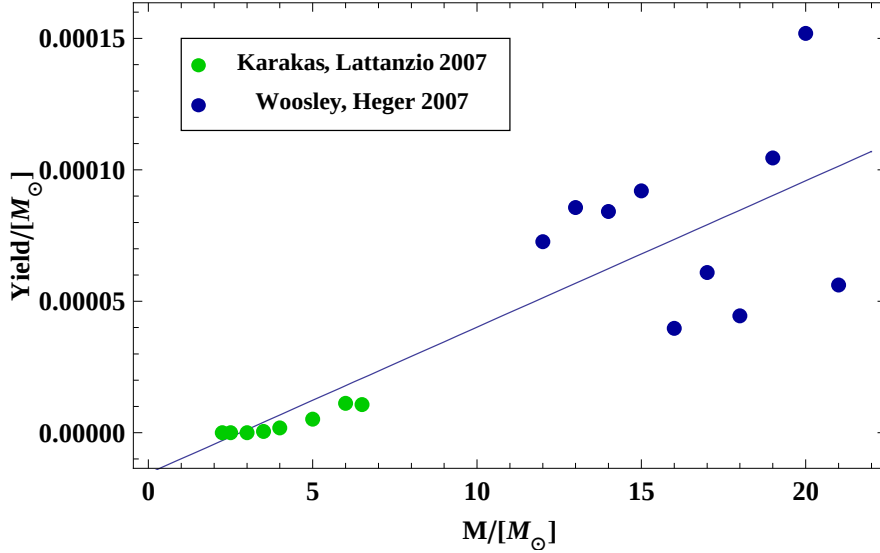


Figure 1.4: The ^{60}Fe abundances from Karakas & Lattanzio and Woosley & Heger in the mass range of 2 - 20 $M_{\odot}$. Because the yields in the range of 7 - 11 $M_{\odot}$ are missing it was fitted linearly.

1.5 The amount of ^{60}Fe arriving on Earth

Because we know the yields of ^{60}Fe , we can calculate the amount that will arrive on Earth, if a supernova explodes in the solar vicinity (from [15])

$$F = \frac{U}{4} \frac{M_{ej}}{4\pi A m_p r^2} e^{-\frac{t}{\tau}}. \quad (1.1)$$

Here, F is the fluence, this means the number of atoms that reach the Earth per cm^2 . What we do is spread the particles over a spherical shell with radius r , which is the distance of the supernova explosion to Earth. The fluence also depends on the ejected mass M_{ej} of ^{60}Fe , its mass number A and the proton mass m_p , but also on the mean life

$$\tau = \frac{t_{1/2}}{\ln 2}, \quad (1.2)$$

where $t_{1/2}$ is the half-life of ^{60}Fe . The time t is the time the supernova remnant takes to reach the Earth, which we will examine in Chapter 5, plus the time the isotope remained on Earth. The factor $1/4$ results from the ratio of the Earth's cross section (πr^2) to its surface area ($4\pi r^2$). The last, but very important factor that is included in the formula is U , the uptake-efficiency. It determines how much of the ^{60}Fe is actually incorporated into the crust. The uptake-factor is obtained from ^{53}Mn measurements. ^{53}Mn is mostly produced by the action of cosmic rays on dust and meteorites. Since it can not be calculated directly, the uncertainty in the factor U is high.

The procedure is described by Knie et al. [24]. First the uptake factor for ^{53}Mn ($t_{1/2}=3.7$ Myr) is calculated, which is the fraction of the flux $\phi_{53,crust}$ in the crust to the total flux $\phi_{53,tot}$

$$U_{Mn} = \frac{\phi_{53,crust}}{\phi_{53,tot}}. \quad (1.3)$$

The flux in the crust was measured by Knie et al., who obtained the value $\phi_{53,crust} = 1.7 \times 10^8 \text{ atoms cm}^{-2} \text{ Myr}^{-1}$. The total flux can be compared with the flux in the Antarctic ice $\phi_{53,ice} = (6.4 \pm 1.4) \times 10^9 \text{ atoms cm}^{-2} \text{ Myr}^{-1}$ determined by Bibron et al. [6] and the flux in deep-sea sediments $\phi_{53,sed} = (2.0 \pm 0.9) \times 10^9 \text{ atoms cm}^{-2} \text{ Myr}^{-1}$ given by Imamura et al. [20]. Then Knie et al. obtain an average value of $\phi_{53,tot} = 4 \times 10^9 \text{ atoms cm}^{-2} \text{ Myr}^{-1}$ for the total flux. Therefore an uptake-efficiency $U_{Mn} \approx 0.4 \%$ is obtained. The uptake factor for ^{60}Fe is then calculated by using

$$U_{Fe} = U_{Mn} \frac{C_{Mn,water}/C_{Fe,water}}{C_{Fe,crust}/C_{Mn,crust}}, \quad (1.4)$$

where $C_{i,water}$ are the concentrations of the dissolved elements in the water around the crust and $C_{i,crust}$ the concentrations in the manganese crust. For our crust, which will be described later in Chapter 2, the concentrations are $C_{Mn,water} \approx 0.15 \text{ nmol kg}^{-1}$, $C_{Fe,water} \approx 0.6 \text{ nmol kg}^{-1}$, $C_{Mn,crust} = 26 \%$ and $C_{Fe,crust} = 15 \%$. Therefore only $U_{Fe} \approx 0.6 \%$ of ^{60}Fe that reaches the Earth is incorporated into the crust.

1.6 Measurement of radionuclides

The concentration of ^{60}Fe was measured with the AMS (accelerator mass spectrometry) at the Munich MP - Tandem Accelerator. As an example for an AMS we will describe VERA (Vienna Environmental Research Accelerator) in Vienna. The ^{10}Be dating of the manganese crust 237KD by Fitoussi et al. [16] was carried out with this instrument.

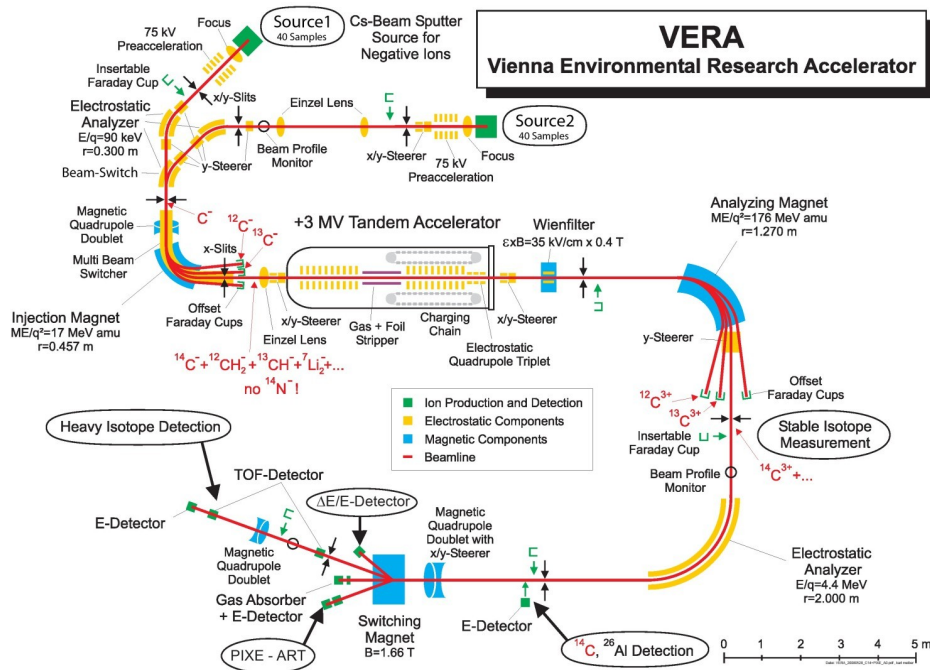


Figure 1.5: The setup of the Vienna Environmental Research Accelerator. Courtesy of the Isotope Research Center, Vienna, <http://isotopenforschung.univie.ac.at>.

Fig. 1.5 shows the setup and operation mode of the accelerator, and Fig. 1.6 shows a photograph. The concentration of a radioisotope is always measured relative to its stable isotope. For ^{60}Fe this is ^{56}Fe , for ^{10}Be it is ^9Be .



Figure 1.6: The Tandem Accelerator in Vienna.

In a simplified demonstration the AMS works as follows [26]: First the sample is scattered in a Sputter Source for Negative Ions. Only negative ions will proceed into the accelerator, that means isobars, that have the same mass number and do not form negative ions, will be already sorted out. In the Electrostatic Analyzer the isotopes are separated by their energies. The Injection Magnet selects the masses, because ions with different masses will be deflected differently. After this procedure there might still exist molecules with similar masses. The negative ions are accelerated in the Tandem Accelerator which is filled with a gas. As they pass through the gas the electrons will be stripped off and the molecules will break. Then they are again separated by velocity, mass and energy. Finally the detector counts the remaining ions.

1.7 The aim of this thesis

An ^{60}Fe anomaly was found in a deep-sea hydrogenetic ferromanganese crust by Knie et al. [24]. This means that more ^{60}Fe isotopes than expected were measured, the signal being significantly above the background. The time-resolved profile will be shown in Chapter 2. The peak position of this anomaly belongs to a time of about 2 - 3 Myr ago. In this work we will analyze a star cluster that is linked to the Local Bubble and connect it with this ^{60}Fe increase in the oceans crust. We test the hypothesis that the ^{60}Fe peak occurs due to multiple SN explosions with an alternative SN model derived by Kahn 1998 [22].

First, in Chapter 2, we will discuss the anomalies that have been measured. Chapter 3 will study the previously mentioned star cluster and calculate the explosion times of already deceased stars. These stars might be the cause for the increase of ^{60}Fe on Earth. We try to reproduce the profile plotted in Fig. 2.4 in the following chapters with analytical calculations. In Chapter 4 and 5, models of supernova remnant expansions will be described. The results will be presented and discussed in Chapter 6. Chapter 7 gives the conclusions and refers to other reservoirs, where ^{60}Fe signals could be found. The future work will be presented in Chapter 8 and other possible candidates of supernova produced radionuclides will be examined. In the last Chapter 9 consequences of near Earth supernova explosions will be discussed.

CHAPTER 2

Supernova relics in the deep ocean manganese crust

Here we focus on important findings of isotope anomalies in hydrogenetic crusts. Later we will link these anomalies to the origin of the Local Bubble.

2.1 A sample in the South Pacific

The first evidence for a nearby supernova explosion in the time span of 2 - 3 Myr ago was found by Knie et al. in 1999 [25]. The idea that it is possible to find isotope anomalies on Earth is not a new one. In 1996, Ellis, Fields and Schramm [13] suggested to search for possible supernova relics and encouraged the examination of deep-ocean's crusts.

To find traces of supernova explosions on Earth, scientists examine different territories. Analysis of the Antarctic Ice, the deep-ocean's crust, the ocean's sediment and even the moon were made. For this work the most interesting case is the discovery of ^{60}Fe isotopes in the deep-sea hydrogenetic ferromanganese crust made by Knie et al. These crusts are particularly suitable for observing the concentration of radionuclides, because due to the low growth rate we can look back millions of years into the past (see Chapter 1.2). And since they grow by obtaining the elements from the surrounding water, they incorporate also particles that come from outside.

The first sample Knie et al. [25] chose originates from Mona Pihoa in the

South Pacific 19° S, 149° W near French Polynesia at a depth of 1300 m. The positions of this hydrogenetic ferromanganese crust as well as those of following samples are marked in Appendix A.

Three layers of the sample were measured with the Munich MP - Tandem Accelerator. Layer 1 covers the depth of 0 - 3 mm, corresponding to a time span of 0 - 2.8 Myr, Layer 2 with a depth of 5 - 10 mm has an age of 3.7 - 5.9 Myr and Layer 3 with 10 - 20 mm relates to 5.9 - 13 Myr. With the ratios of $^{60}\text{Fe}/\text{Fe} = 2.1 \times 10^{-15}$ and $^{60}\text{Fe}/\text{Fe} = 1.4 \times 10^{-15}$ the Layers 1 and 2 clearly showed signals above the background which was estimated to be $^{60}\text{Fe}/\text{Fe} = 0.25 \times 10^{-15}$ by a blank sample.

2.2 The ^{60}Fe anomaly in the crust 237KD

The next hydrogenetic deep ocean ferromanganese crust to be analyzed originates from the top of a submarine mountain (depth 4830 m) in the equatorial Pacific at a position 9°18' N and 146°03' W. This crust with the name 237KD was collected in 1976 with the research ship Valdivia during the cruise VA13/2 (see Fig. 2.1).



Figure 2.1: On the left: the research vessel Valdivia, today called Kommandor Jack. On the right: A piece of the ferromanganese crust 237KD, which is one of the biggest crusts ever recovered. Courtesy of: Left: Dr. Detlef Quadfasel, right: M. Poutivtsev [36].

In 2004 Knie et al. published a paper [24] with the measurements that are the basis for the calculations in this work. They found a definite $^{60}\text{Fe}/\text{Fe}$ peak as shown in Fig. 2.2.

The crust was measured two times. The first measurement, displayed by the blue dots, seems to indicate two anomalies. One at a time of about 3 Myr and another one at about 9 Myr. Those were taken from layers with a thickness of 2 mm whereas the in second measurement the layers had a thickness of 1 mm, here with red dots. The second measurements confirm

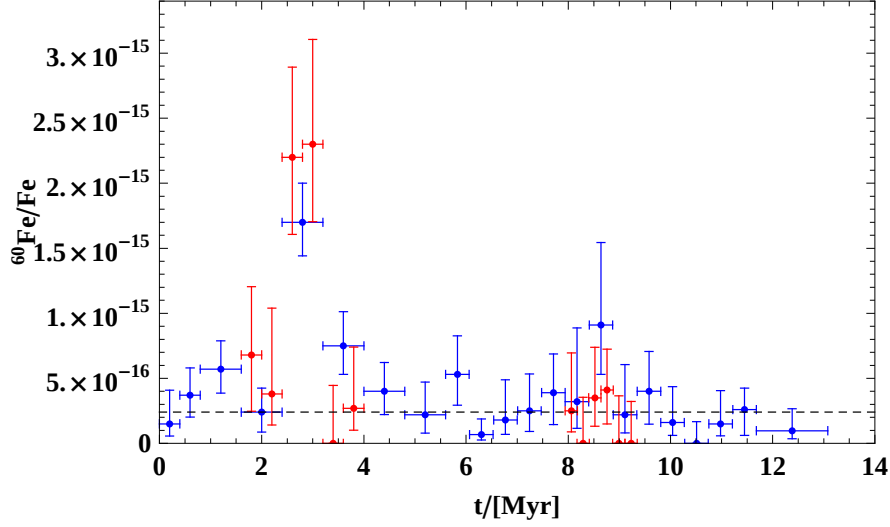


Figure 2.2: The $^{60}\text{Fe}/\text{Fe}$ ratio plotted against the age of the crust. The blue dots originate from core 1, the red dots from core 2. A peak is discernable at a time 3 Myr ago. The dashed line indicates the background.

the peak at 3 Myr, which corresponds to the layers at a depth of 6 - 8 mm, but the peak at 9 Myr turns out to be a runaway-value. The data has not been corrected for radioactive decay, background and the uptake factor. Here the background is considered to have the $^{60}\text{Fe}/\text{Fe}$ ratio of 2.4×10^{-16} . With a ratio of 1.9×10^{-15} the fluence corresponding to the peak was calculated to be $\Phi_{60,crust} = (2.9 \pm 1.0) \times 10^6 \text{ atoms cm}^{-2}$ when corrected for background and radioactive decay.

The age of the crust was derived by Segl et al. 1984 [41] with ^{10}Be measurements, because a constant ^{10}Be flux is assumed. The half-life of this isotope was estimated to be 1.51 Myr. New measurements from Nishiizumi et al. 2007 [35] yield a slightly lower half-life of 1.36 Myr. With this discovery the dating of the crust changes. The peak in the layer of 6 - 8 mm was suggested to arise from a time span of 2.4 - 3.2 Myr. The new time span can be estimated to be 1.69 - 2.53 Myr as announced in Fitoussi et al. from 2008 [16]. As we shall see later in Chapter 6, our analysis fits better with this new data. In [16] Fitoussi et al. describe how to generate the age of the layers. They use the equation

$$T_d = t_{1/2}(^{10}\text{Be}) \frac{\ln(C_0/C_d)}{\ln 2}, \quad (2.1)$$

where T_d is the age of the depth $d \geq 2 \text{ mm}$, $t_{1/2}(\text{Be})$ the half-life of ^{10}Be , C_0 the

^{10}Be concentration measured at a depth $d = 2$ mm and C_d the concentration measured at a depth d . The topmost 2 mm show a lower ^{10}Be concentration. Fitoussi et al. suggest a rapid diffusion or a higher growth rate near the surface. The flux in a depth of 2 - 15 mm is relatively constant giving a growth rate of 2.37 mm/Myr with the new half life of ^{10}Be . In Knie et al. 2004 [24] a growth rate of 2.5 mm/Myr was used. At a depth of about 15 mm the growth rate suddenly slows down as measured by Segl et al. [41], who suggest that a slower bottom water circulation pattern is responsible. In Fig. 2.3, the ^{10}Be concentration is plotted logarithmically against the depths of the crust as it was done in [16]. Here we recognize that the flux between 0 and 10 mm was nearly constant. The slope corresponds to the growth rate of the crust.

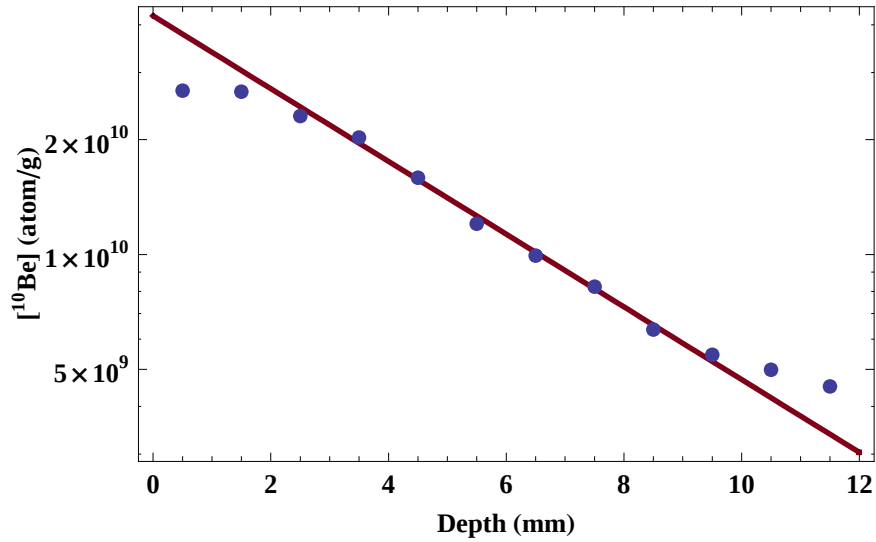


Figure 2.3: The ^{10}Be concentration plotted against the depth of the crust. The concentration was fitted exponentially where the slope indicates a growth rate of 2.37 mm/Myr.

With the new dating the peak shifts about 0.7 Myr. Fig. 2.4 shows the newly dated crust as well as a confirmation of the data made in [16]. The green dots are the new measurements which were taken from the same crust, but from a different drill hole. The black dots represent the old measurements where the mean values of core 1 and core 2 were taken around the original 9 Myr where the runaway-value was located.

We will use this plot in Chapter 6 to fit in our data and try to reproduce the peak.

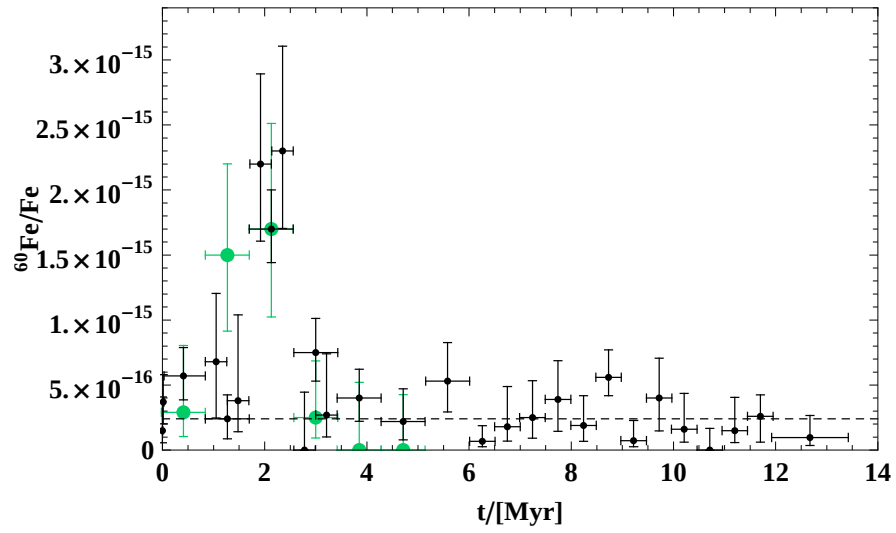


Figure 2.4: The $^{60}\text{Fe}/\text{Fe}$ ratio plotted against the age of the crust based on the new ^{10}Be dating. The black dots correspond to the measurements from 2004 by Knie et al. the green dots to the confirmation of the ^{60}Fe signal made in 2008 by Fitoussi et al.

CHAPTER 3

The formation of the Local Bubble

Because an increase of ^{60}Fe was measured, we need to find stars that might be responsible for its production. This chapter discusses the search of a star cluster in the solar vicinity where some stars have already exploded. Later the explosion times are calculated.

3.1 Searching for OB associations in the solar vicinity

In 2002 T.W. Berghöfer and D. Breitschwerdt (BB02) published a paper [4] linking the ^{60}Fe deposit in the oceans crust to the formation of the Local Bubble. They analyzed a young stellar moving group, the subgroup B1 of the Pleiades moving group, that could have provided enough supernova explosions to blow a cavity into the interstellar medium as big as our Local Bubble. B1 consists of 27 B stars, which are now part of the Scorpius – Centaurus association, with the center of mass lying at a distance of about 135 ± 15 pc away from the Sun. They appear to be the youngest stars found in the solar vicinity. Berghöfer and Breitschwerdt calculated that about 19 SNe must have already occurred in that star cluster. With the age of the Local Bubble estimated to be about 13 Myr the rate of the explosions would be $f_{\text{SN}} \sim 1/(6.5 \times 10^5)$ yr.

Four years later, in 2006, Fuchs et al. [17] went further into detail by examining all stars within a radius of 200 pc centered around the Sun. By searching the Hipparcos catalog for stars bluer than $(B-V) < -0.05$ with

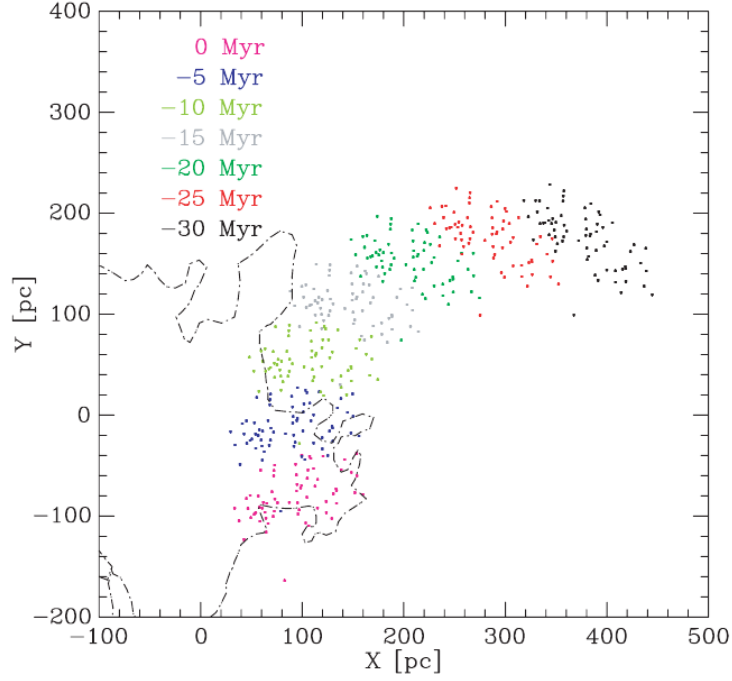


Figure 3.1: The trajectories of UCL and LCC over a time span of 30 Myr into the past [17]. The dashed line implicates the outline of the region occupied by the Local Bubble today. The first stars entered that region around 15 Myr ago (gray).

parallaxes larger than 5 mas they could extract 762 B stars within the given distance. For 610 stars the radial velocities are available in the catalog. The kinematic development was traced backwards for a time span of 30 Myr (see Fig. 3.1) by using the epicyclic equations of motion from Lindblad 1959 [30] and Wielen 1982 [49]

$$\begin{aligned}
 X(t) &= X(0) - \frac{V(0)}{-2B} [1 - \cos(\kappa t)] + \frac{U(0)}{\kappa} \sin(\kappa t), \\
 Y(t) &= Y(0) + 2A \left[X(0) - \frac{V(0)}{-2B} \right] t + \frac{\Omega_0}{-B\kappa} V(0) \sin(\kappa t) \\
 &\quad + \frac{2\Omega_0}{\kappa^2} U(0) [1 - \cos(\kappa t)], \\
 Z(t) &= \frac{W(0)}{\nu} \sin(\nu t) + Z(0) \cos(\nu t),
 \end{aligned} \tag{3.1}$$

$$\begin{aligned}
U(t) &= U(0) \cos(\kappa t) - \frac{\kappa}{-2B} V(0) \sin(\kappa t), \\
V(t) &= \frac{-2B}{\kappa} U(0) \sin(\kappa t) + V(0) \cos(\kappa t), \\
W(t) &= W(0) \cos(\nu t) - Z(0) \nu \sin(\nu t).
\end{aligned} \tag{3.2}$$

Here $X(t)$, $Y(t)$ and $Z(t)$ are the spatial coordinates of the stars and $U(t)$, $V(t)$ and $W(t)$ their velocity components. The epicyclic frequency $\kappa = \sqrt{-4\Omega_0 B}$ consists of the angular frequency $\Omega_0 = V_{LSR}/R_\odot$ of the rotation of the local standard of rest around the Galactic Center and has a value of $220 \text{ km s}^{-1}/8 \text{ kpc}$. The Oort constants A and B are related by $A = -B = \Omega_0/2$, so we obtain $\kappa = 4 \times 10^{-8} \text{ yr}^{-1}$. The vertical oscillation frequency is $\nu = \sqrt{4\pi G \rho_0}$ and with $\rho_0 = 0.1 \text{ M}_\odot \text{ pc}^{-3}$ the local density and G the gravitational constant $\nu = 7.5 \times 10^{-8} \text{ yr}^{-1}$.

Tracing the paths of the stars backwards in time one can recognize that not all stars belong to the group observed today. In Fig. 3.2 we can see that most stars were spread into a wide cloud. There are 302 stars lying in the dashed box. These are stars which stayed together for the last 30 Myr and could be possible members of the Sco-Cen association.

Looking at their distribution today, Fuchs et al. excluded another 66 stars. Finally they compared the $(B-V)_J$ colors with the theoretical isochrones from the literature, after they transformed the colours from the Tycho system in the Hipparcos catalog into the Johnson system. They found out that almost all stars are members of the Sco-Cen subgroups Upper Scorpius (US), Upper Centaurus Lupus (UCL) and Lower Centaurus Crux (LCC), which are the youngest OB associations in the solar vicinity. Altogether there were 79 stars that remained after this procedure.

The ages of UCL and LCC are conjectured to be between 20 - 30 Myr, whereas the age of US was not possible to determine. The trajectories of the remaining 79 stars were calculated with (3.1) and (3.2). Since there are errors in the velocity components U , V and W , the paths of additional pseudo-stars were computed for every star.

In Fig. 3.3 the trajectories of the 79 objects and the pseudo-stars are plotted. The pseudo star that comes closest to the Sun with the minimal distance of 64.77 pc is shown in red color in Fig. 3.3, the star with the largest distance is given in blue color. The original star the minimal extreme trajectory corresponds to has the number HD 113902 or HIP 64053 and is identified as a B8V star in the SIMBAD database.

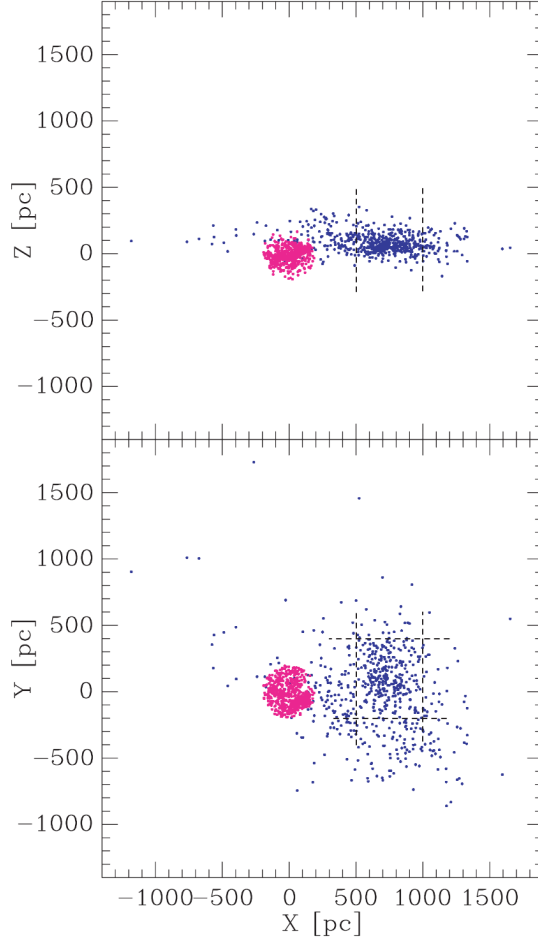


Figure 3.2: Possible stars that caused the origin of the Local Bubble [17]. The pink dots are the positions of the stars today and the blue dots are the locations 30 Myr ago.

With the mass-to-magnitude relation, which Fuchs et al. derived from Schaller et al. [40]

$$M_*/M_\odot = 3.587 - 1.453M_V + 0.183M_V^2 + 0.069M_V^3. \quad (3.3)$$

Fuchs et al. calculated the masses of the 79 stars. In Appendix B the table of all stars is given together with the stars magnitudes, coordinates, velocities and memberships of the subclusters. With a magnitude of $M_V = 0.64$ mag, HIP 64053 then has a mass of $3.02 M_\odot$.

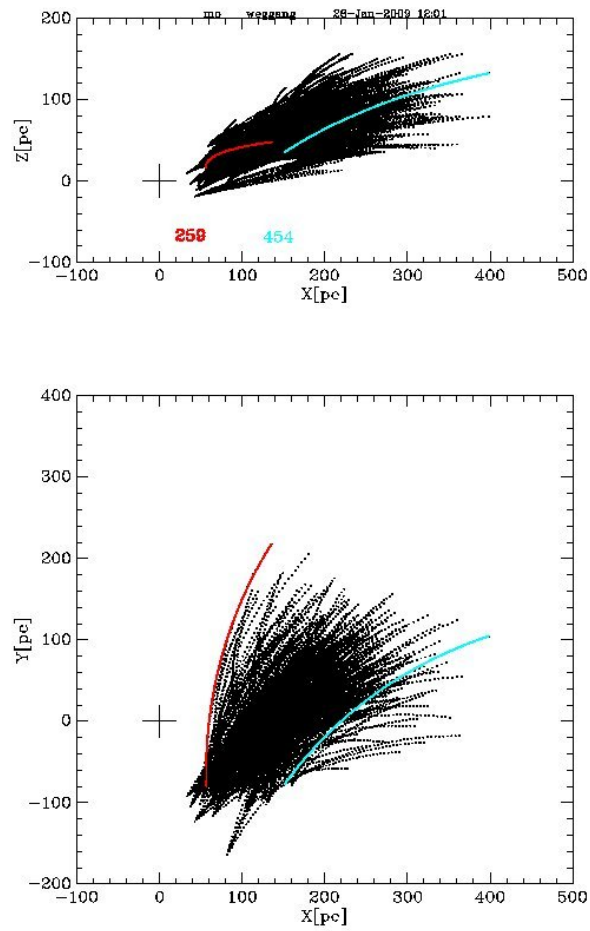


Figure 3.3: The trajectory of 711 objects. The red line marks the path of the pseudo-star that came closest to the Sun, the blue line is the pseudo-star farthest away. Courtesy of Dipl.-Phys. C. Dettbarn.

3.2 Possible supernovae

To estimate how many stars in the UCL and LCC have exploded already to create the Local Bubble we will again follow Fuchs et al. They fitted an initial mass function (IMF) from Massey et al. 1995 [33]

$$\frac{dN}{dM} = \frac{dN}{dM}\bigg|_0 M^{\Gamma-1}. \quad (3.4)$$

The IMF describes the distribution of stellar masses in a star population. The index Γ is -1.1 ± 0.1 for young massive stars. As we can see from table B.1/B.2 in Appendix B there are 42 stars in the UCL and 27 in the LCC. To calculate the number of missing stars, the normalization constants $dN/dM|_0$ were obtained by computing the masses for A0 stars at the lower end of the main sequence and for B0 stars at the upper end. A0 stars have a magnitude $M_V = 1$ mag which corresponds to $M_l = 2.6 M_\odot$ by (3.3) and B0 stars with $M_V = -3.7$ mag have a mass of $M_u = 8.2 M_\odot$. With

$$N = \int_{2.6}^{8.2} \frac{dN}{dM}\bigg|_0 M^{-2.1} = 0.228 \frac{dN}{dM}\bigg|_0 \quad (3.5)$$

Fuchs et al. got $dN/dM|_0 = 42/0.228 = 184$ and $dN/dM|_0 = 27/0.228 = 118$ for the UCL and the LCC respectively. The main sequence lifetime of a massive star was derived to be

$$\tau = \tau_0 M^{-\alpha} \quad (3.6)$$

for masses between 2 and 67 $M_\odot$. Here $\tau_0 = 1.6 \times 10^8$ yr and $\alpha = 0.932$. With $\tau_{M_{\Delta\tau}} = \tau_u - \Delta\tau$ denoting the lifetime of the first and most massive star that exploded in the Local Bubble, τ_u the lifetime of the star with the mass M_u , and $\Delta\tau$ the time the association entered the present LB volume, the masses of the most massive stars that exploded are

$$M_{\Delta\tau} = \left(M_u^{-\alpha} - \frac{\Delta\tau}{\tau_0} \right)^{-1/\alpha}. \quad (3.7)$$

If the Cluster entered the LB at a time $\Delta\tau = 10$ Myr then $M_{\Delta\tau} = 15.4 M_\odot$, if $\Delta\tau = 15$ Myr, then $M_{\Delta\tau} = 26.6 M_\odot$. Finally the number of missing stars were calculated by

$$N_{SN} = \int_{8.2}^{M_{\Delta\tau}} \frac{dN}{dM}\bigg|_0 M^{-2.1} dM. \quad (3.8)$$

This gives about 8 - 12 supernovae in the UCL and 6 - 8 in the LCC for $\Delta\tau$ between 10 and 15 Myr. Fuchs et al. even went further and estimated that about 5 - 12 SNe went off before entering the volume occupied by the LB today.

3.3 Calculating the explosion times

With this information we will now evaluate the masses and times of individual supernovae that might have caused the ^{60}Fe peak in the manganese crust. The aim is to achieve a distribution of supernova explosions during the last 13 Myr, which is the estimated age of the LB. The age of the cluster can be derived by computing the lifetime of an $8.2 M_{\odot}$ star, which gives us a $\tau = 2.25 \times 10^7 \text{ yr}$ with (3.6). This is in good agreement with BB02 and Fuchs et al. One possibility to distribute the supernovae is to let them explode in equidistant time intervals. For about 20 SNe in 13 Myr years this would give one SN in $6.5 \times 10^5 \text{ yr}$ as was mentioned before. Because this scenario is not very probable, we calculated the explosion times based on of the IMF in (3.8). The procedure is to divide the IMF for the UCL and the LCC into intervals so that every bin contains one star. In Fig. 3.4 and Fig. 3.5 it is illustrated how the 12 stars of the UCL and the 8 stars of the LCC can be distributed.

Next we calculate the masses of each star in the diagram. One way to do it is to take the masses at the boundaries, another way is to evaluate the mean masses. In Tab. 3.1 and 3.2 the values of the masses, corresponding lifetimes and explosion times for the UCL and LCC are listed.

In Fig. 3.6 and 3.7 the explosion times with the masses of the stars for UCL and LCC with boundary and mean masses are plotted. We notice that the explosion times are not equidistant. This is in agreement with our expectations. Since the stars with the masses of $8.2 M_{\odot}$ were our reference to calculate the age of the cluster we will leave them out for further calculations. With equation (3.7) we obtain a mass of $20.66 M_{\odot}$ for an entering time of the cluster into the region of the LB of 13 Myr. Therefore we will also omit stars with masses bigger than $21 M_{\odot}$.

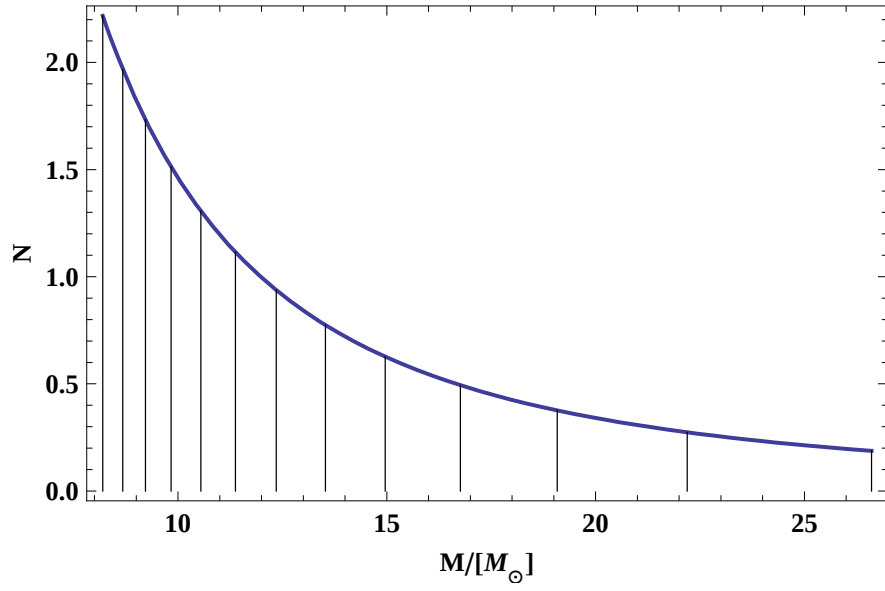


Figure 3.4: The IMF of the UCL divided into intervals that each contain one star.

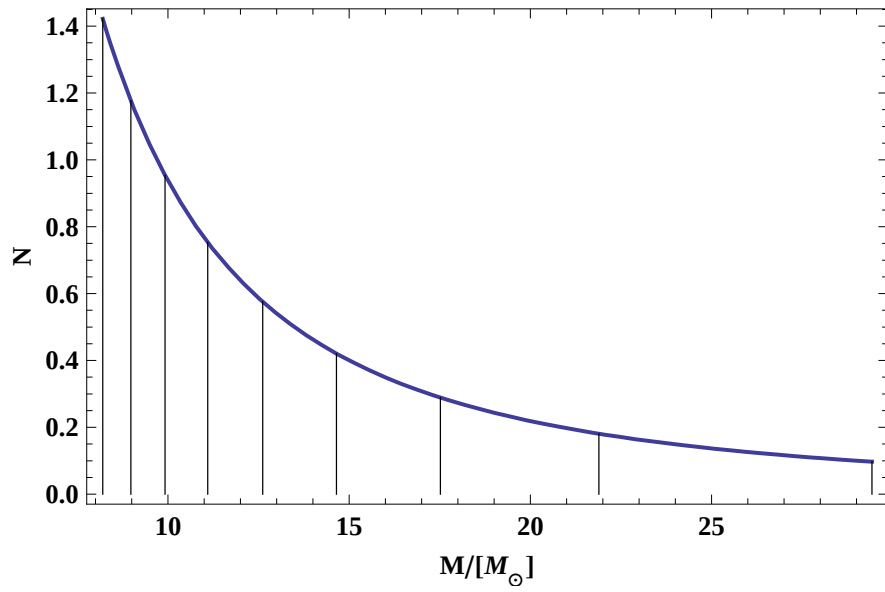


Figure 3.5: The IMF of the LCC divided into intervals that each contain one star.

$M_b/M_\odot$	$\tau_b/[\text{Myr}]$	$t_b/[\text{Myr}]$	$M_m/M_\odot$	$\tau_m/[\text{Myr}]$	$t_m/[\text{Myr}]$
8.20	22.51	0	8.44	21.92	-0.60
8.68	21.35	-1.16	8.95	20.75	-1.76
9.22	20.18	-2.33	9.53	19.57	-2.94
9.84	19.00	-3.51	10.19	18.38	-4.13
10.55	17.80	-4.71	10.96	17.17	-5.34
11.38	16.59	-5.92	11.87	15.95	-6.56
12.36	15.36	-7.15	12.94	14.71	-7.80
13.53	14.12	-8.40	14.25	13.45	-9.06
14.96	12.85	-9.66	15.86	12.17	-10.34
16.76	11.56	-10.95	17.92	10.86	-11.65
19.08	10.24	-12.27	20.64	9.53	-12.99
22.19	8.90	-13.61	24.40	8.15	-14.37

Table 3.1: Masses, lifetimes and explosion times of the stars in the UCL for the boundary values (b) and mean values (m) of the mass.

$M_b/M_\odot$	$\tau_b/[\text{Myr}]$	$t_b/[\text{Myr}]$	$M_m/M_\odot$	$\tau_m/[\text{Myr}]$	$t_m/[\text{Myr}]$
8.20	22.51	0	8.59	21.57	-0.95
8.97	20.70	-1.81	9.44	19.74	-2.78
9.92	18.86	-3.66	10.51	17.87	-4.64
11.10	16.98	-5.53	11.86	15.97	-6.55
12.61	15.07	-7.44	13.63	14.02	-8.49
14.65	13.11	-9.40	16.08	12.02	-10.50
17.52	11.10	-11.42	19.70	9.95	-12.57
21.89	9.42	-13.50	25.66	7.78	-14.74

Table 3.2: Masses, lifetimes and explosion times of the stars in the LCC for the boundary values (b) and mean values (m) of the mass.

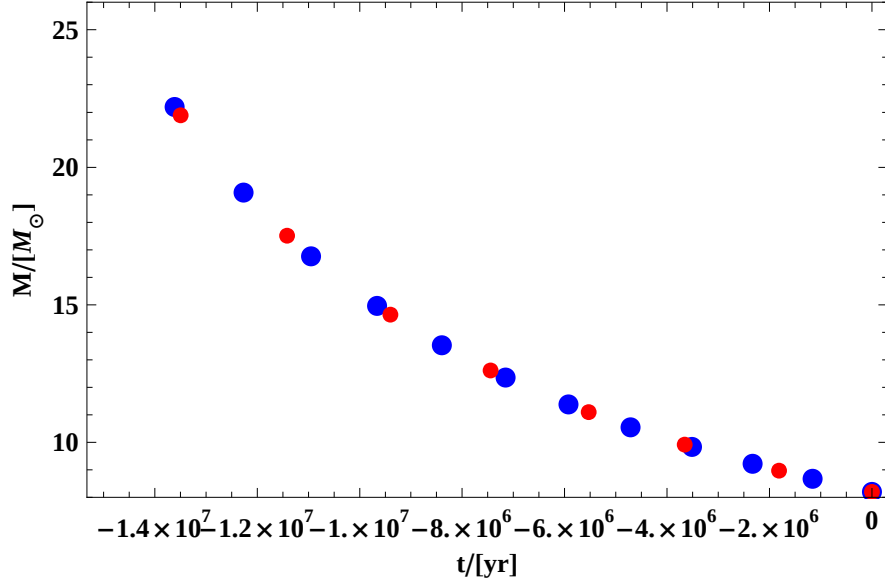


Figure 3.6: The explosion times of the stars of the UCL (blue) and the LCC (red) in correlation with their masses. Here the masses were derived by taking the boundary values of the bins in the IMF above.

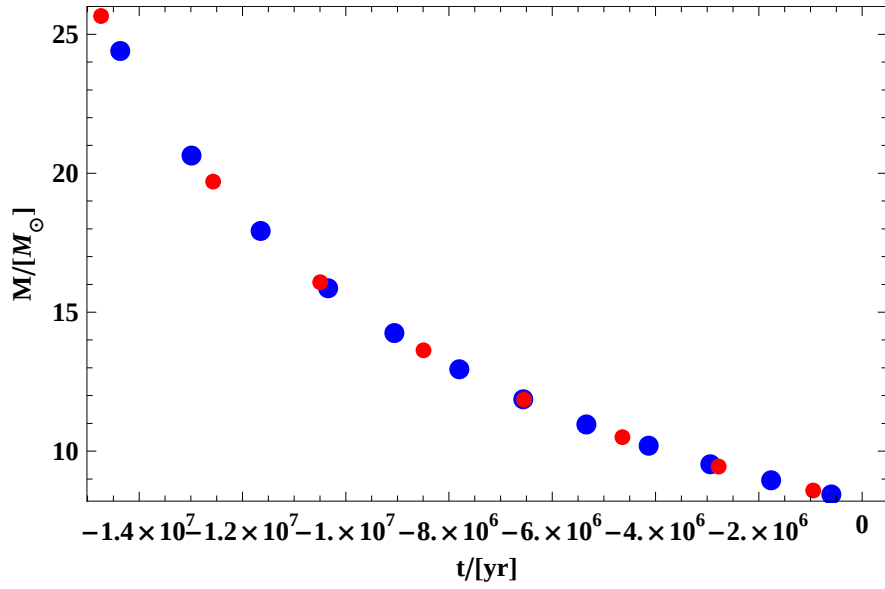


Figure 3.7: The explosion times of the stars of the UCL (blue) and the LCC (red) in correlation with their masses. Here the masses were derived by taking the mean values of the bins in the IMF above.

CHAPTER 4

The theory of supernova remnant expansion

This chapter will cover the basics of hydrodynamics in astrophysics starting with the Rankine-Hugoniot shock jump conditions followed by a physical explanation of the expansion of supernova remnants into a homogeneous ambient medium. The information is based on *Principles of Astrophysical Fluid Dynamics* by Clarke and Carswell [10]. Later we will present a mathematical description of the motion of a supernova remnant as derived by Sedov in 1946.

4.1 The Rankine-Hugoniot Conditions

These are the shock-jump conditions that relate the parameters u_0 , ρ_0 and P_0 before the shock (the upstream) to the parameters u_1 , ρ_1 and P_1 behind the shock (the downstream), where u is the velocity, ρ the density and P the pressure of the medium. These variables change discontinuously across the shock and without a magnetic field the equations, that can be derived from the Euler-equations, take the form

$$\rho_0 u_0 = \rho_1 u_1, \quad (4.1a)$$

$$\rho_0 u_0^2 + P_0 = \rho_1 u_1^2 + P_1, \quad (4.1b)$$

$$\frac{1}{2} u_0^2 + \epsilon_0 + \frac{P_0}{\rho_0} = \frac{1}{2} u_1^2 + \epsilon_1 + \frac{P_1}{\rho_1}. \quad (4.1c)$$

The first equation states that the mass flux is conserved across the shock. When the shock wave compresses the medium the velocity must go down. We will see later how much the density increases and the velocity decreases for a strong shock. The second and the third equation describe the conservation of the momentum flux and the specific total energy, where

$$\epsilon = \frac{1}{\gamma - 1} \frac{P}{\rho} \quad (4.2)$$

is the internal energy per unit mass and γ the adiabatic parameter, which has the value 5/3 for monoatomic and 7/5 for diatomic gas. Inserting this into (4.1c) the equation can be rewritten as

$$\frac{1}{2}u_0^2 + \frac{\gamma}{\gamma - 1} \frac{P_0}{\rho_0} = \frac{1}{2}u_1^2 + \frac{\gamma}{\gamma - 1} \frac{P_1}{\rho_1}. \quad (4.3)$$

Next we introduce the Mach number, which is the ratio of the velocity of the flow to the sound velocity

$$M^2 = \frac{u^2}{c_s^2}. \quad (4.4)$$

If the Mach number is smaller than 1 we have a subsonic flow, if $M = 1$ the flow is sonic and with $M > 1$ it is supersonic. In a fixed reference frame, which will be discussed shortly, $u_0 = -v_s$, where v_s is the shock velocity. This means that M_0 is the ratio of the shock speed over the sound speed. Thus strong shocks can be described by their Mach number, for example a Mach-10 shock has the Mach number 10. The velocity, density and pressure of the gas before and after the shock can now be related by the Mach-number M_0 in the upstream medium

$$\frac{\rho_1}{\rho_0} = \frac{u_0}{u_1} = \frac{(\gamma + 1)M_0^2}{(\gamma + 1) + (\gamma - 1)(M_0^2 - 1)}, \quad (4.5)$$

$$\frac{P_1}{P_0} = \frac{2\gamma}{\gamma + 1}M_0^2 - \frac{\gamma - 1}{\gamma + 1}. \quad (4.6)$$

For a strong shock, $M_0 \rightarrow \infty$, we get

$$\frac{\rho_1}{\rho_0} = \frac{u_0}{u_1} = \frac{\gamma + 1}{\gamma - 1}, \quad (4.7)$$

$$\frac{P_1}{P_0} = \frac{2\gamma}{\gamma + 1}M_0^2. \quad (4.8)$$

The sound speed in an adiabatic flow is

$$c_s^2 = \frac{\gamma P}{\rho}. \quad (4.9)$$

Then the Mach number M_0 becomes

$$M_0^2 = \frac{\rho_0 u_0^2}{\gamma P_0}. \quad (4.10)$$

With these informations we rewrite P_1/P_0 and obtain

$$P_1 = \frac{2}{\gamma + 1} \rho_0 u_0^2. \quad (4.11)$$

For a monoatomic gas the parameters change as following

$$\rho_1 = 4\rho_0, \quad (4.12a)$$

$$u_1 = \frac{1}{4}u_0, \quad (4.12b)$$

$$P_1 = \frac{3}{4}\rho_0 u_0^2. \quad (4.12c)$$

Now we have a short look to what happens with the post-shock temperature. The pressure in an ideal gas can be described as

$$P = \rho \frac{kT}{\mu m_u}, \quad (4.13)$$

where k is the Boltzmann constant, μ the molecular weight and m_u the atomic mass unit. The post-shock temperature therefore is

$$T_1 = \frac{\mu m_u}{k} \frac{P_1}{\rho_1}. \quad (4.14)$$

With $\rho_1 = 4\rho_0$ the pressure in the shocked gas becomes

$$P_1 = \frac{3}{16}\rho_1 u_0^2, \quad (4.15)$$

and by inserting this into the post-shock temperature we can express it in terms of pre-shock variables

$$T_1 = \frac{3}{16} \frac{\mu m_u}{k} u_0^2. \quad (4.16)$$

4.2 Supernova Remnants

As a supernova explodes into the surrounding medium it can be described in three different phases, which characterize how the remnant moves and how its propagation velocity will decrease over the time. For all the approximations that follow we will assume the interstellar medium has a homogeneous density ρ_0 .

4.2.1 Phase I: Free Expansion

A supernova releases about $E = 10^{51}$ erg, part of which, at the time of the explosion, is converted into kinetic energy. The density of the ejected matter is much larger than the density of the surrounding medium, so the medium does not influence the expansion and we get a linear velocity profile

$$r = v_{ej}t. \quad (4.17)$$

Here v_{ej} is the velocity of the ejected mass. With the kinetic supernova energy

$$E = \frac{1}{2}M_{ej}v_{ej}^2, \quad (4.18)$$

the velocity can be computed with

$$v_{ej} = \left(\frac{2E}{M_{ej}} \right)^{1/2}. \quad (4.19)$$

The phase ends when the ejected mass is comparable to the swept-up interstellar medium, i.e. when

$$M_{ej} = M_{sw} = \frac{4\pi}{3}r_{sw}^3\rho_0. \quad (4.20)$$

Say, a supernova ejects about $5 M_\odot$, then the velocity of the mass is $v_{ej} = 4.5 \times 10^8 \text{ cm s}^{-1} = 4484 \text{ km s}^{-1}$. With a particle density of $n = 1 \text{ cm}^{-3}$ and $\rho_0 = nm_H m_u$ with $m_H \sim 1$ and $m_u = 1.66 \times 10^{-24} \text{ g}$ the interstellar medium has a density of $\rho_0 = 1.7 \times 10^{-24} \text{ g cm}^{-3}$. But since the ISM does not only consist of hydrogen, the density will be assumed to have a higher value of $2 \times 10^{-24} \text{ g cm}^{-3}$ during this work. At the end of Phase I the remnant has a radius of $r_{sw} = 10^{19} \text{ cm} = 3.4 \text{ pc}$ after a time $t_{sw} = 2.4 \times 10^{10} \text{ s} = 750 \text{ years}$.

4.2.2 Phase II: The Sedov - Taylor Phase

At the end of Phase I a reverse shock starts to move inwards converting some of the kinetic energy into thermal energy, so the medium gets heated up to very large temperatures. Now the further expansion is due to the pressure from the inside. For further calculations we will make a short remark on the two different reference frames in which the explosion can be observed. One is the lab frame, where the observer is at rest e.g. in the star or in the surrounding ISM. The other is the shock frame, with respect to the shock velocity, which is $v_s = 0$. In the lab frame pre- and post-shock velocities will be denoted by v_0 and v_1 , in the shock frame it is u_0 and u_1 (for visualization see Fig. 4.1). We change between those frames by

$$u_0 = v_0 - v_s, \quad (4.21a)$$

$$u_1 = v_1 - v_s. \quad (4.21b)$$



Figure 4.1: Here the difference between the lab and the shock frame is illustrated. As we approximate the ISM to be at rest compared to the star we can set the velocity $v_0 = 0$ in the lab frame.

Assuming that the velocity of the interstellar medium in the lab frame is $v_0 = 0$, the velocity of the medium in the shock frame becomes $u_0 = -v_s$. Relative to the star the shocked gas then has a velocity

$$v_1 = u_1 + v_s = \left(\frac{\gamma - 1}{\gamma + 1} \right) u_0 + v_s = v_s - \left(\frac{\gamma - 1}{\gamma + 1} \right) v_s = \frac{2}{\gamma + 1} v_s. \quad (4.22)$$

For $\gamma = 5/3$ the shocked gas will move with $3/4$ of the shock speed. We approximate the moving shell to be thin. Let the shell have a thickness d and the whole mass be contained in it, then we can write

$$\frac{4\pi}{3} \rho_0 r^3 = 4\pi \rho_1 r^2 d. \quad (4.23)$$

On substituting $\rho_1 = 4\rho_0$ we calculate the ratio of $d/r = 1/12 \ll 1$, which means that the approximation of a thin shell is justified. The momentum Mv_1 the shell gains is provided by the pressure P on the inside of the shell. With $M = (4\pi/3)r^3\rho_0$ being the swept up mass in the shell we get the equation of motion

$$\frac{d}{dt} \left(\frac{4\pi}{3} r^3 \rho_0 v_1 \right) = 4\pi r^2 P. \quad (4.24)$$

We suppose that the outside pressure is $P_0 = 0$ by setting $T_0 = 0$. For the pressure P acting on the shell we assume that it scales with the pressure in the shell by setting

$$P = \alpha P_1. \quad (4.25)$$

But P_1 is known from the shock jump conditions, and substituting v_1 by the shock velocity which has been derived before we write the equation of motion as

$$\frac{d}{dt} \left(\frac{4\pi}{3} r^3 \rho_0 \frac{2}{\gamma + 1} v_s \right) = 4\pi r^2 \alpha \frac{2}{\gamma + 1} \rho_0 v_s^2. \quad (4.26)$$

Many terms cancel and the shock velocity is $v_s = \dot{r}$

$$\frac{d}{dt} (r^3 \dot{r}) = 3\alpha r^2 \dot{r}^2. \quad (4.27)$$

To solve this differential equation we use the ansatz

$$r = At^\beta \quad (4.28)$$

in equation (4.27), which leads to

$$\frac{d}{dt} (A^3 t^{3\beta} \beta A t^{\beta-1}) = 3\alpha A^2 t^{2\beta} A^2 \beta^2 t^{2\beta-2}. \quad (4.29)$$

Then A cancels out and

$$\beta(4\beta - 1)t^{4\beta-2} = 3\alpha\beta^2 t^{4\beta-2}. \quad (4.30)$$

If $\beta = 0$ then r is constant and the solution is uninteresting. For $\beta \neq 0$ we obtain

$$\beta = \frac{1}{4 - 3\alpha}. \quad (4.31)$$

Now we know that r is a function of this quantity α . To determine α we take a look at the energy. Since some of the explosion energy was converted into thermal energy the total energy is

$$E = E_k + E_t, \quad (4.32)$$

where

$$E_k = \frac{1}{2} M v_1^2 \quad (4.33)$$

is the kinetic energy in the shell and the internal energy per unit volume of the cavity is

$$E_t = \frac{4\pi}{3} r^3 \frac{P}{\gamma - 1} = \frac{4\pi}{3} r^3 \frac{2\alpha \rho_0 v_s^2}{(\gamma - 1)(\gamma + 1)}. \quad (4.34)$$

Substituting v_1 in (4.33) with (4.22) we obtain

$$E = \frac{4\pi}{3} r^3 \rho_0 \left(\frac{1}{2} \left(\frac{2}{\gamma + 1} \right)^2 v_s^2 + \frac{2\alpha}{(\gamma - 1)(\gamma + 1)} v_s^2 \right) \quad (4.35)$$

for the energy. Using the ansatz $r = A t^{1/(4-3\alpha)}$ again we estimate

$$E \propto r^3 \dot{r}^2 \propto t^{(6\alpha-3)/(4-3\alpha)}. \quad (4.36)$$

The energy is conserved, so $6\alpha - 3 = 0$ and

$$\alpha = \frac{1}{2}. \quad (4.37)$$

This means that the radius is proportional to $t^{2/5}$. To obtain the factor A we insert the value for α into the energy equation. Then

$$E = \frac{27}{32} \frac{4\pi}{3} \rho_0 r^3 \dot{r}^2 \quad (4.38)$$

and inserting the ansatz (with $\gamma = 5/3$) results in

$$A = \left(\frac{3}{4\pi} \frac{32}{27} \frac{25}{4} \right)^{1/5} \left(\frac{E}{\rho_0} \right)^{1/5}. \quad (4.39)$$

The factor inside the brackets is about 1.7. With the mathematical method shown in the following section 4.3 we will obtain a factor of 2. Then the radius and the velocity of the shell in the Sedov - Taylor Phase become

$$r = \left(\frac{1.7E}{\rho_0} \right)^{1/5} t^{2/5}, \quad (4.40)$$

$$\dot{r} = \frac{2}{5} \left(\frac{1.7E}{\rho_0} \right)^{1/5} t^{-3/5}. \quad (4.41)$$

This way of deriving the equations illustrate very well the physics behind the problem. But we can not be sure how good the approximations are since there are some rather heuristic assumptions, for example the assumption of the constant α . To be sure that this is the right solution we must derive the equations by using something called the similarity solution, which we will show in the next section 4.3.

When the gas in the cavity cools down to a temperature of about 10^6 K most of the elements start to recombine. Then the cooling becomes more efficient and the pressure on the inside of the shell reduces. We have already derived the equation for the post-shock temperature in (4.16) and with (4.41) we get the equation for the reduction of temperature with time during the adiabatic expansion

$$T = \frac{3}{100} \frac{\mu m_u}{k} \left(\frac{1.7E}{\rho_0} \right)^{2/5} t^{-6/5}. \quad (4.42)$$

Inserting the explosion energy $E = 10^{51}$ erg, the density of the interstellar medium $\rho_0 = 2 \times 10^{-24}$ g cm $^{-3}$, the molecular weight $\mu = 0.61$ for a fully ionized medium with helium abundance of $n_{\text{He}}/n_{\text{H}} = 0.1$, the atomic mass unit $m_u = 1.66 \times 10^{-24}$ g, the Boltzmann-constant $k = 1.38 \times 10^{-16}$ erg K $^{-1}$ the Boltzmann-constant and a temperature of $T = 10^6$ K we obtain a time $t = 9 \times 10^{11}$ s or about 3×10^4 yr. With (4.40) we calculate an increase of the radius of approximately $r = 6 \times 10^{19}$ cm or about 20 pc until that time and then the remnant has a velocity of 2.5×10^7 cm s $^{-1}$.

To determine the cooling time of the system we use the energy equation

$$\frac{d}{dt} \left(\frac{3}{2} n k T \right) = -n^2 \Lambda(T), \quad (4.43)$$

where the left hand side is the change of the total thermal energy for a monoatomic ideal gas with time, and the right hand side describes the radiated energy with $\Lambda(T)$ the cooling function published by F.D. Kahn in 1976 [21]

$$\Lambda(T) = 1.33 \times 10^{-19} T^{-1/2} [\text{erg cm}^3 \text{ s}^{-1}]. \quad (4.44)$$

This cooling law is valid in the temperature range $5 \times 10^4 - 5 \times 10^7$ K. Then the cooling time can be written as

$$t_c \simeq \frac{3kT}{2n\Lambda(T)}. \quad (4.45)$$

The effect of the cooling becomes important when the dynamical time becomes comparable to the cooling time

$$t_c \simeq t_{\text{dyn}} = \frac{r}{\dot{r}} = \frac{5}{2} t, \quad (4.46)$$

with r and $\dot{r}$ being the radius and velocity of the remnant in the Sedov-Taylor Phase. Then we get

$$\frac{3kT}{2n\Lambda(T)} = \frac{5}{2}t. \quad (4.47)$$

Substituting the temperature from (4.42) and the cooling function from (4.44) and taking $n = 4n_0$ for a strong shock, we obtain the following expression

$$t_c = \left(\frac{3}{20}\right)^{5/14} \left(\frac{k}{1.33 \times 10^{-19} n_0}\right)^{5/14} \left(\frac{3}{100} \frac{\mu m_u}{k} \left(\frac{1.7E}{\rho_0}\right)^{2/5}\right)^{15/28}. \quad (4.48)$$

With the values from above we calculate a cooling time of $t_c = 4.6 \times 10^{11} \text{ s} = 1.5 \times 10^4 \text{ yr}$. Still moving according to the Sedov equation the radius will cover a distance of $4.5 \times 10^{19} \text{ cm} = 14.5 \text{ pc}$ during that time and have a velocity of $2 \times 10^7 \text{ cm s}^{-1}$ in the end.

4.2.3 Phase III: The Snowplough Phase

If we assume that the bubble has cooled completely we can neglect the pressure from the inside, and the shell velocity will decrease. The evolution of the shell in this phase is only due to momentum-conservation

$$M\dot{r} = \frac{4\pi}{3}r^3\rho_0\dot{r} = p. \quad (4.49)$$

Here p is the momentum in the shell at the end of Phase II, i.e. also at the beginning of Phase III, where the radius is $r = r_0$, then

$$p = \frac{4\pi}{3}r_0^3\rho_0\dot{r}_0. \quad (4.50)$$

Since the momentum in the shell is conserved, $p = \text{const}$, the time derivative is 0,

$$\frac{d}{dt} \left(\frac{4\pi}{3}r^3\rho_0\dot{r} \right) = 0. \quad (4.51)$$

All the constants cancel out and the equation of motion for Phase III becomes

$$\frac{d}{dt}(r^3\dot{r}) = 0. \quad (4.52)$$

Here we can make the ansatz $r \propto t^\beta$ again and obtain $r \propto t^{1/4}$, but in order to get the constants it is easier to integrate (4.49)

$$\int_{r_0}^r \frac{4\pi}{3}(r')^3 \rho_0 dr' = \int_{t_0}^t p dt', \quad (4.53)$$

where t_0 is the time at the beginning of Phase III. Now we derive for the evolution of the radius

$$r = r_0 \left(1 + \frac{4\dot{r}_0(t - t_0)}{r_0} \right)^{1/4}. \quad (4.54)$$

In the end the remnant will merge with the ISM.

4.3 Self-similar Sedov Solution for a strong point explosion

In this section we will derive the Sedov-solution mathematically. We will also find out how the density distribution of the medium changes after the remnant propagated through the medium. This is important for calculations made in the next Chapter. To obtain the solution we will follow Landau and Lifschitz' *Fluid Mechanics* [28].

A self-similar problem is present, if the quantities that characterize the system do not contain explicit length or time scales. Here we are able to reduce a system of partial differential equations to a system of ordinary differential equations with the help of dimensionless variables. Since we have a point explosion the energy E is released into a homogeneous medium with the density ρ_0 . Together with the independent variables r and t the dimensionless form of the problem becomes

$$r \left(\frac{\rho_0}{Et^2} \right)^{1/5}. \quad (4.55)$$

Then the radius R and the velocity v_s of the remnant can be written as

$$R = \beta \left(\frac{Et^2}{\rho_0} \right)^{1/5}, \quad (4.56)$$

$$v_s = \frac{2R}{5t} = \frac{2}{5} \beta \left(\frac{E}{\rho_0 t^3} \right)^{1/5}, \quad (4.57)$$

with β being a constant which we will determine later. Now we introduce the dimensionless variable ξ so that $\xi = 1$ at the shock radius R

$$\xi = \frac{r}{R} = \frac{r}{\beta} \left(\frac{\rho_0}{Et^2} \right)^{1/5}. \quad (4.58)$$

To describe the profile of the flow behind the shock we write the velocity v , the density ρ and the sound velocity c_s (which substitutes the pressure) in terms of the dimensionless variables G , Z and V

$$\rho = \rho_0 G(\xi), \quad (4.59)$$

$$c_s^2 = \frac{4r^2}{25t^2} Z(\xi), \quad (4.60)$$

$$v = \frac{2r}{5t} V(\xi). \quad (4.61)$$

To obtain the boundary conditions at $\xi = 1$ we need the Rankine-Hugoniot conditions for a strong shock derived in (4.7), (4.11) and (4.22)

$$\rho_1 = \frac{\gamma + 1}{\gamma - 1} \rho_0, \quad (4.62)$$

$$c_s^2 = \frac{\gamma p_1}{\rho_1} = \frac{2\gamma}{\gamma + 1} \frac{\rho_0}{\rho_1} v_s^2 = \frac{2\gamma(\gamma - 1)}{(\gamma + 1)^2} \frac{4R^2}{25t}, \quad (4.63)$$

$$v_1 = \frac{2}{\gamma + 1} v_s = \frac{2}{\gamma + 1} \frac{2R}{5t}. \quad (4.64)$$

Then the boundary conditions become

$$G(1) = \frac{\gamma + 1}{\gamma - 1}, \quad (4.65)$$

$$Z(1) = \frac{2\gamma(\gamma - 1)}{(\gamma + 1)^2}, \quad (4.66)$$

$$V(1) = \frac{2}{\gamma + 1}. \quad (4.67)$$

The Euler equations for a spherically symmetric flow are

$$\frac{\partial \rho}{\partial t} + \frac{\partial(\rho v)}{\partial r} + \frac{2\rho v}{r} = 0, \quad (4.68)$$

$$\frac{\partial v}{\partial t} + v \frac{\partial v}{\partial r} = -\frac{1}{\rho} \frac{\partial p}{\partial r}, \quad (4.69)$$

$$\left(\frac{\partial}{\partial t} + v \frac{\partial}{\partial r} \right) \ln \frac{p}{\rho^\gamma} = 0. \quad (4.70)$$

The Eulerian equations describe the motion of fluid elements where (4.68) represents the conservation of mass and (4.69) is the equation of motion.

The last equation (4.70) follows because the flow is assumed to be adiabatic, which means no loss of energy, so that the entropy is conserved

$$\frac{D}{Dt} \left(\frac{p}{\rho^\gamma} \right) = 0, \quad (4.71)$$

with

$$\frac{D}{Dt} = \frac{\partial}{\partial t} + v \frac{\partial}{\partial r} \quad (4.72)$$

the relation between Lagrangian and Eulerian coordinates. Equations (4.68), (4.69) and (4.70) are a system of partial differential equations. If we employ the dimensionless variables G , Z and V we will get a system of ordinary differential equations as mentioned before. The substitution is done in the following way:

If we differentiate (4.58) with respect to t and r we have

$$\frac{\partial}{\partial t} = -\frac{2}{5} \frac{\xi}{t} \frac{d}{d\xi}, \quad (4.73)$$

$$\frac{\partial}{\partial r} = \frac{\xi}{r} \frac{d}{d\xi}. \quad (4.74)$$

Taking this information into account we are able to write down the partial derivatives of ρ and v from (4.59) and (4.61) with respect to t and r

$$\frac{\partial \rho}{\partial t} = -\frac{2}{5} \frac{\xi}{t} \frac{d}{d\xi} \rho_1 G(\xi), \quad (4.75)$$

$$\frac{\partial \rho}{\partial r} = \frac{\xi}{r} \frac{d}{d\xi} \rho_1 G(\xi), \quad (4.76)$$

$$\frac{\partial v}{\partial t} = -\frac{2r}{5t^2} \left(V(\xi) + \frac{2}{5} \xi \frac{d}{d\xi} V(\xi) \right), \quad (4.77)$$

$$\frac{\partial v}{\partial r} = \frac{2}{5t} \left(V(\xi) + \xi \frac{d}{d\xi} V(\xi) \right). \quad (4.78)$$

Converting the first Euler-equation to the form

$$\frac{\partial \rho}{\partial t} + \rho \frac{\partial v}{\partial r} + v \frac{\partial \rho}{\partial r} + \frac{2\rho v}{r} = 0, \quad (4.79)$$

we can now use the dimensionless variables. This gives

$$\begin{aligned} & -\frac{2\xi}{5t} \rho_1 \frac{d}{d\xi} G + \rho_1 G \left(\frac{2}{5t} \left(V + \xi \frac{d}{d\xi} V \right) \right) \\ & + \frac{2r}{5t} V \frac{\xi}{r} \frac{d}{d\xi} \rho_1 G + \frac{2}{r} \rho_1 G \frac{2r}{5t} V = 0. \end{aligned} \quad (4.80)$$

Then we get an equation where t and r cancel out

$$-(1 - V)\xi \frac{dG}{d\xi} + G\xi \frac{d}{d\xi} V = -3GV. \quad (4.81)$$

Dividing by G and using that

$$\frac{\xi}{G} \frac{dG(\xi)}{d\xi} = \frac{d \ln G(\xi)}{d \ln \xi}, \quad (4.82)$$

we obtain the final form as given in Landau and Lifschitz

$$\frac{dV}{d \ln \xi} - (1 - V) \frac{d \ln G}{d \ln \xi} = -3V. \quad (4.83)$$

Now we will transform the third Euler-equation. We compute p from (4.59) and (4.60) with $c_s^2 = (\gamma p / \rho)$ and get

$$p = \frac{4r^2}{25\gamma t^2} \rho_1 G Z. \quad (4.84)$$

If we insert this into (4.70) we obtain the following expression

$$\left(\frac{\partial}{\partial t} + v \frac{\partial}{\partial r} \right) \left(\ln \left(\frac{2r}{5\sqrt{\gamma}t} \right)^2 Z - (\gamma - 1) \ln(\rho_1 G) \right) = 0. \quad (4.85)$$

Transforming into the dimensionless variables will again result in r and t cancelling. Using relation (4.82) the final result is

$$\frac{d \ln Z}{d \ln \xi} - (\gamma - 1) \frac{d \ln G}{d \ln \xi} = -\frac{5 - 2V}{1 - V}. \quad (4.86)$$

We transformed only the first and third equation. To solve this system of differential equations we can derive an integral which reduces the number of the three independent equations to two and then we can replace e.g. the momentum equation with that integral.

This integral can be derived in the following way. Because we have a point explosion the total energy in the system (a sphere with the radius R) is constant and equals the value that is released by the explosion. This is the case because we also assume that the pressure of the ISM is negligible compared to the pressure inside the remnant, and so is the energy. Because we have a similarity flow the energy within smaller spheres with $\xi = \text{const}$ must be also constant. The radial velocity at these surfaces is $v_n = 2r/(5t)$ (from (4.57)). Then we can set up the equation

$$dt \, 4\pi r^2 \rho v \left(w + \frac{v^2}{2} \right) = dt \, 4\pi r^2 \rho v_n \left(\epsilon + \frac{v^2}{2} \right). \quad (4.87)$$

The term on the left hand side means that the energy flows through the surface of a sphere with radius r in the time dt , but at the same time the volume of the sphere increases by $dt v_n 4\pi r^2$ (term on the right hand side). Equation (4.87) states that the energy stays constant. Here

$$w = \frac{c_s^2}{\gamma - 1} \quad (4.88)$$

is the enthalpy and

$$\epsilon = \frac{1}{\gamma - 1} \frac{P}{\rho} = \frac{c_s^2}{\gamma(\gamma - 1)} \quad (4.89)$$

the internal energy (compare (4.2)). If we transform (4.87) to dimensionless variables and use (4.88) and (4.89) we have

$$\left(\frac{2r}{5t}\right)^3 \frac{VZ}{\gamma - 1} + \left(\frac{2r}{5t}\right)^3 \frac{V^3}{2} = \left(\frac{2r}{5t}\right)^3 \frac{Z}{\gamma(\gamma - 1)} + \left(\frac{2r}{5t}\right)^3 \frac{V^2}{2}. \quad (4.90)$$

Solving this equation for Z gives us the first integral which only depends on the dimensionless quantity V and satisfies the boundary conditions automatically

$$Z = \frac{\gamma(\gamma - 1)(1 - V)V^2}{2(\gamma V - 1)}. \quad (4.91)$$

Equations (4.83), (4.86) and (4.91) form now our system we have to solve. Details are given in Appendix C. The resulting equations are

$$\xi^5 = \left(\frac{\gamma + 1}{2} V\right)^{-2} \left(\frac{\gamma + 1}{7 - \gamma} (5 - (3\gamma - 1)V)\right)^{v_1} \left(\frac{\gamma + 1}{\gamma - 1} (\gamma V - 1)\right)^{v_2}, \quad (4.92)$$

$$G = \frac{\gamma + 1}{\gamma - 1} \left(\frac{\gamma + 1}{\gamma - 1} (\gamma V - 1)\right)^{v_3} \left(\frac{\gamma + 1}{7 - \gamma} (5 - (3\gamma - 1)V)\right)^{v_4} \left(\frac{\gamma + 1}{\gamma - 1} (1 - V)\right)^{v_5}, \quad (4.93)$$

with

$$v_1 = -\frac{13\gamma^2 - 7\gamma + 12}{(3\gamma - 1)(2\gamma + 1)}, \quad (4.94)$$

$$v_2 = \frac{5(\gamma - 1)}{2\gamma + 1}, \quad (4.95)$$

$$v_3 = \frac{3}{2\gamma + 1}, \quad (4.96)$$

$$v_4 = \frac{v_1}{2 - \gamma}, \quad (4.97)$$

$$v_5 = -\frac{2}{2-\gamma}. \quad (4.98)$$

Next we evaluate the constant β . We can calculate it by using that the total energy in the blast waves consist of kinetic and thermal energy and remains constant in the sphere

$$E = \int_0^R \rho \left(\frac{v^2}{2} + \frac{\epsilon}{\rho} \right) 4\pi r^2 dr. \quad (4.99)$$

Inserting (4.89), the dimensionless coordinates (4.59), (4.60), (4.61) and

$$r = \xi \beta \left(\frac{Et^2}{\rho_1} \right)^{1/5}, \quad (4.100)$$

$$dr = \beta \left(\frac{Et^2}{\rho_1} \right)^{1/5} d\xi, \quad (4.101)$$

we write for the energy

$$E = 4\pi \int_0^1 \frac{4r^2}{25t^2} \left(\frac{1}{2} V^2 + \frac{1}{\gamma(\gamma-1)Z} \right) \xi^2 \beta^3 \left(\frac{Et^2}{\rho_1} \right)^{3/5} d\xi. \quad (4.102)$$

Then we will obtain for β

$$\beta^5 \frac{16\pi}{25} \int_0^1 G \left(\frac{1}{2} V^2 + \frac{1}{\gamma(\gamma-1)Z} \right) \xi^4 d\xi = 1. \quad (4.103)$$

To solve this integral we will transform it and integrate over the dimensionless velocity V

$$\beta^5 \frac{16\pi}{25} \int_{\frac{1}{\gamma}}^{\frac{2}{\gamma+1}} G(V) \left(\frac{1}{2} V^2 + \frac{1}{\gamma(\gamma-1)Z(V)} \right) \xi^5 \frac{d \ln \xi}{dV} = 1. \quad (4.104)$$

The lower limit of the integral is due to $V \rightarrow (1/\gamma)$ as $\xi \rightarrow 0$ which can be calculated from (4.92), and the upper limit is again the boundary condition $V(1)$. We can solve the integral numerically (for example with Mathematica), if we first replace $d \ln \xi / dV$ by (C.8), which is derived in Appendix C. We can choose different values for the adiabatic index γ . For example

$$\gamma = \frac{7}{5} \Rightarrow \beta = 1.033 \text{ (for air)}, \quad (4.105)$$

$$\gamma = \frac{5}{3} \Rightarrow \beta = 1.152 \text{ (for monoatomic gas)}. \quad (4.106)$$

This means when a supernova remnant expands into a monoatomic gas, the radius will evolve according to

$$R(t) = \left(\frac{2.026E}{\rho_0} \right)^{1/5} t^{2/5}. \quad (4.107)$$

If $\xi \rightarrow 0$ then $V \rightarrow (1/\gamma)$ and $G \rightarrow 0$ (from 4.93), thus we can write

$$V - \frac{1}{\gamma} \propto \xi^{5/v_2}, \quad (4.108)$$

$$G \propto \xi^{5v_3/v_2}. \quad (4.109)$$

Then we calculate the ratios v/v_2 and ρ/ρ_2 as a function of ξ for $\xi \rightarrow 0$

$$\frac{v}{v_2} \propto \frac{r}{R}, \quad (4.110)$$

$$\frac{\rho}{\rho_1} \propto \left(\frac{r}{R} \right)^{3/(\gamma-1)}. \quad (4.111)$$

For the monoatomic gas the exponent $3/(\gamma - 1)$ has a value of $9/2$. This will be used later in the alternative supernova remnant model developed by Kahn 1998 [22] in Chapter 5. We notice that the density distribution after the explosion is not homogeneous, but behaves like a power law. In the explosion center the medium is very thin, but with $\xi \rightarrow 1$ most matter is concentrated in a shell.

CHAPTER 5

An alternative model

In the following we calculate analytically the expansion of a supernova remnant in an inhomogeneous medium, which results from the explosion of a previous SN. Our analysis follows closely the treatment of Kahn 1998 [22], also correcting some errors in his treatment and deriving a more precise equation for the expansion of a supernova remnant in Phase III.

5.1 The late evolution of supernova remnants

The standard Sedov - Taylor solutions for the shock radius and shock speed after a time τ are

$$r = \left(\frac{2E}{\rho_0} \right)^{1/5} \tau^{2/5}, \quad (5.1)$$

$$\dot{r} = \frac{2}{5} \left(\frac{2E}{\rho_0} \right)^{1/5} \tau^{-3/5}, \quad (5.2)$$

with P the pressure and ρ_0 the density of the ISM. The adiabatic parameter κ is given by $\kappa = \frac{P}{\rho_0^{5/3}}$ (which was defined by Kahn 1976). With $P_1 = \frac{3}{4}\rho_0\dot{r}^2$ the post-shock pressure derived from the Rankine-Hugoniot conditions for a strong shock, $\rho_1 = 4\rho_0$ and (5.2), the adiabatic parameter immediately after the shock becomes

$$\begin{aligned} \kappa_s^{3/2} &= 0.02 \frac{\dot{r}^3}{\rho_0} \\ &= 0.002 \frac{E^{3/5}}{\rho_0^{8/5}} \tau^{-9/5}. \end{aligned} \quad (5.3)$$

This relation remains valid until cooling becomes non-negligible. The mass of the gas in the sphere that has been shocked is, with (5.1),

$$M_s = \frac{4\pi}{3} \rho_0 r^3 = \frac{4\pi}{3} \rho_0 \left(\frac{2E}{\rho_0} \right)^{3/5} \tau^{6/5}. \quad (5.4)$$

From (5.3) and (5.4) we can derive a relation between the value of κ_s and the mass enclosed within the shock

$$\begin{aligned} \kappa_s M_s &= \left(0.002 \frac{E^{3/5}}{\rho_0^{8/5}} \tau^{-9/5} \right)^{2/3} \frac{4\pi}{3} \rho_0 \left(\frac{2E}{\rho_0} \right)^{3/5} \tau^{6/5} \\ &= 0.002^{2/3} \frac{E^{2/5}}{\rho_0^{16/15}} \tau^{-6/5} \frac{4\pi}{3} \rho_0 \frac{2^{3/5} E^{3/5}}{\rho_0^{3/5}} \tau^{6/5} \\ &= 0.10 \frac{E}{\rho_0^{2/3}}. \end{aligned} \quad (5.5)$$

Gas that has been shocked at time τ will have cooled completely at a time

$$\begin{aligned} t &= \tau + t_{cool} = \tau + \frac{\kappa^{2/3}}{q} \\ &= \tau + 0.002 \frac{E^{3/5}}{q \rho_0^{8/5}} \tau^{-9/5}, \end{aligned} \quad (5.6)$$

with $t_{cool} = \kappa^{3/2}/q$, where the parameter q has a value of $4 \times 10^{32} \text{ cm}^6 \text{ g}^{-1} \text{ s}^{-4}$ (Kahn 1976) and depends on the composition of the gas. The given value is valid for solar abundances.

The first element to cool must have passed through the shock at time

$$t - \tau = \tau_* = 0.002 \frac{E^{3/5}}{q \rho_0^{8/5}} \tau_*^{-9/5},$$

and thus

$$\begin{aligned} \tau_*^{14/5} &= 0.002 \frac{E^{3/5}}{q \rho_0^{8/5}}, \\ \tau_* &= \frac{0.108 E^{3/14}}{q^{5/14} \rho_0^{4/7}}. \end{aligned} \quad (5.7)$$

The gas will have cooled completely when

$$\begin{aligned}
t = t_{II} &= \tau_* + 0.002 \frac{E^{3/5}}{q \rho_0^{8/5}} \tau_*^{-9/5} \\
&= \frac{0.108 E^{3/14}}{q^{5/14} \rho_0^{4/7}} + \frac{0.002 E^{3/5}}{q \rho_0^{8/5}} \cdot \left(\frac{0.108 E^{3/14}}{q^{5/14} \rho_0^{4/7}} \right)^{-9/5}, \\
t_{II} &= 0.216 \frac{E^{3/14}}{q^{5/14} \rho_0^{4/7}}. \tag{5.8}
\end{aligned}$$

This marks the end of the Sedov - Taylor Phase, Phase II. Now we compute the radius at that time:

$$\begin{aligned}
r_{II} &= \left(\frac{2E}{\rho_0} \right)^{1/5} t_{II}^{2/5} \\
&= \left(\frac{2E}{\rho_0} \right)^{1/5} \left(\frac{0.216 E^{3/14}}{q^{5/14} \rho_0^{4/7}} \right)^{2/5} \\
&= 0.622 \frac{E^{2/7}}{q^{1/7} \rho_0^{3/7}}. \tag{5.9}
\end{aligned}$$

With the released energy of a supernova explosion being about $E = 10^{51}$ erg, the parameter $q = 4 \times 10^{32} \text{ cm}^6 \text{ g}^{-1} \text{ s}^{-4}$ and the density $\rho_0 = 2 \times 10^{-24} \text{ g cm}^{-3}$, we get a time and a radius with the values of

$$\begin{aligned}
r_{II} &= 7.32 \times 10^{19} \text{ cm}, \\
t_{II} &= 1.45 \times 10^{12} \text{ s}.
\end{aligned}$$

Much later in the Snowplough Phase only the gas with $t < t_{cool} = \kappa_s^{3/2}/q$ will be hot. With

$$\kappa_s^{3/2} > qt \tag{5.10}$$

and (5.5) we derive the mass of the hot gas as

$$M_h = 0.10 \frac{E}{\rho_0^{2/3} \kappa_s} = 0.10 \frac{E}{\rho_0^{2/3} q^{2/3}} t^{-2/3}. \tag{5.11}$$

The adiabatic parameter in the hot gas for partial mass M in a spherical symmetrical volume is

$$\kappa = 0.10 \frac{E}{\rho_0^{2/3} M}. \tag{5.12}$$

With this equation and

$$\kappa = \frac{P}{\rho_0^{5/3}} \Rightarrow \rho_0 = \left(\frac{P}{\kappa} \right)^{3/5},$$

we can now compute the specific volume v of the gas

$$\begin{aligned} v &= \frac{1}{\rho_0} = \frac{1}{\left(\frac{P}{\kappa}\right)^{3/5}} = \kappa^{3/5} P^{-3/5} \\ &= \left(\frac{0.10E}{\rho_0^{2/3}M}\right)^{3/5} P^{-3/5} = \frac{0.25E^{3/5}}{\rho_0^{2/5}M^{3/5}} P^{-3/5}. \end{aligned} \quad (5.13)$$

The whole volume occupied by the hot gas can be derived by integrating over the mass and using the equation above and (5.11)

$$\begin{aligned} V &= \int_0^{M_h} v dM = \int_0^{M_h} \frac{0.25E^{3/5}}{\rho_0^{2/5}M^{3/5}} P^{-3/5} dM \\ &= \frac{0.25E^{3/5}}{\rho_0^{2/5}} P^{-3/5} \int_0^{M_h} M^{-3/5} dM = \frac{0.25E^{3/5}}{\rho_0^{2/5}P^{3/5}} \cdot \frac{5}{2} M_h^{2/5} \\ &= \frac{0.25E^{3/5}}{\rho_0^{2/5}P^{3/5}} \cdot \frac{5}{2} \left(\frac{0.10E}{\rho_0^{2/3}q^{2/3}} t^{-2/3}\right)^{2/5}, \\ V &= \frac{0.25E}{\rho_0^{2/3}P^{3/5}q^{4/15}} t^{-4/15}. \end{aligned} \quad (5.14)$$

The pressure P in a spherical volume $V = (4\pi/3)r^3$ of radius r and time t is therefore

$$\begin{aligned} P^{3/5} &= \frac{0.25E}{\rho_0^{2/3}q^{4/15}V} t^{-4/15} = \frac{0.25E}{\rho_0^{2/3}q^{4/15}\frac{4\pi}{3}r^3} t^{-4/15}, \\ P &= \frac{0.0091E^{5/3}}{\rho_0^{10/9}q^{4/9}} r^{-5} t^{-4/9}. \end{aligned} \quad (5.15)$$

Now Kahn introduces the dimensionless radius R and the dimensionless time θ , normalized by the previous calculated constants r_{II} and t_{II}

$$\begin{aligned} r &= Rr_{II}, \\ t &= \theta t_{II}. \end{aligned} \quad (5.16)$$

With these dimensionless quantities and (5.8) and (5.9) the pressure P becomes

$$\begin{aligned} P &= \frac{0.0091E^{5/3}}{\rho_0^{10/9}q^{4/9}} (Rr_{II})^{-5} (\theta t_{II})^{-4/9} \\ &= \frac{0.0091E^{5/3}}{\rho_0^{10/9}q^{4/9}} \left(R \frac{0.622E^{2/7}}{q^{1/7}\rho_0^{3/7}}\right)^{-5} \left(\theta \frac{0.216E^{3/14}}{q^{5/14}\rho_0^{4/7}}\right)^{-4/9} \\ &= 0.192E^{1/7} q^{3/7} \rho_0^{9/7} R^{-5} \theta^{-4/9}, \\ P &= p_* R^{-5} \theta^{-4/9}, \end{aligned} \quad (5.17)$$

with the value p_* being

$$p_* = 0.192E^{1/7}q^{3/7}\rho_0^{9/7}. \quad (5.18)$$

In Phase III of the supernova explosion, the Snowplough phase, the shell moves by momentum conservation, but it still will be pushed by the hot interior. The pressure by the hot gas on the shell is $F = 4\pi r^2 P$, while the momentum conservation is described by $F = \frac{d}{dt}(M\dot{r})$ with $M = \frac{4\pi}{3}r^3\rho_0$. The equation of motion can now be written as

$$\begin{aligned} \frac{d}{dt}(M\dot{r}) &= 4\pi r^2 P, \\ \frac{d}{dt}\left(\frac{4\pi}{3}r^3\rho_0\dot{r}\right) &= 4\pi r^2 P, \\ \frac{1}{3}\rho_0\frac{d}{dt}(r^3\dot{r}) &= Pr^2. \end{aligned} \quad (5.19)$$

This looks like the equation for the Sedov - Taylor phase as derived in (4.24). But here we have a different pressure acting on the inside of the shell. By differentiating with respect to t , we get

$$\begin{aligned} \frac{1}{3}\rho_0(3r^2\dot{r}^2 + r^3\ddot{r}) &= Pr^2, \\ \frac{1}{3}\rho_0(r\ddot{r} + 3\dot{r}^2) &= P, \\ \frac{1}{3}\rho_0 r\ddot{r} + \rho_0\dot{r}^2 &= P. \end{aligned} \quad (5.20)$$

This equation is valid only long after time t_{II} . The usual equation of motion neglects the push of the shell by the hot gas and does not have the pressure term on the right hand side of the equation (see (4.52)). Now we will express equation (5.19) in terms of the dimensionless variables R and θ from (5.16). The term on the left hand side can be written as

$$\frac{1}{3}\rho_0\frac{d}{dt}(r^3\dot{r}) = \frac{1}{3}\rho_0\frac{d}{dt}\left(r^3\frac{d}{dt}r\right) = \frac{1}{3}\rho_0\frac{d}{dt}\left(\frac{1}{4}\frac{d}{dt}r^4\right) = \frac{1}{12}\rho_0\frac{d^2}{dt^2}r^4.$$

Inserting in the dimensionless variables and using (5.17) the equation of motion (5.19) becomes

$$\begin{aligned}\frac{1}{12}\rho_0\frac{d^2}{dt^2}r^4 &= Pr^2, \\ \frac{1}{12}\rho_0\frac{d^2}{d(\theta t_{II})^2}(Rr_{II})^4 &= P(Rr_{II})^2, \\ \frac{1}{12}\rho_0\frac{r_{II}^4}{t_{II}^2}\frac{d^2}{d\theta^2}R^4 &= p_*R^{-5}\theta^{-4/9}R^2r_{II}^2, \\ R^3\frac{d^2R^4}{d\theta^2} &= \frac{12t_{II}^2p_*}{r_{II}^2\rho_0}\theta^{-4/9}.\end{aligned}$$

Now we can rewrite the right hand side of the last equation by means of (5.8), (5.9) and (5.18)

$$\begin{aligned}12\frac{t_{II}^2p_*}{r_{II}^2\rho_0}\theta^{-4/9} &= 12\frac{0.192E^{1/7}q^{3/7}\rho_0^{9/7}\cdot(\frac{0.216E^{3/14}}{q^{5/14}\rho_0^{4/7}})^2}{\rho_0\cdot(0.622\frac{E^{2/7}}{q^{1/7}\rho_0^{3/7}})^2}\theta^{-4/9} \\ &= 12\frac{0.192E^{1/7}q^{3/7}\rho_0^{9/7}\cdot(0.216)^2E^{3/7}q^{1/7}\rho_0^{6/7}}{\rho_0\cdot(0.622)^2E^{4/7}q^{5/7}\rho_0^{8/7}}\theta^{-4/9} \\ &= 0.278\frac{E^{4/7}q^{5/7}\rho_0^{15/7}}{E^{4/7}q^{5/7}\rho_0^{15/7}}\theta^{-4/9} \\ &= 0.278\theta^{-4/9}.\end{aligned}$$

With the definition

$$S \equiv R^4, \quad (5.21)$$

the equation of motion in dimensionless variables becomes

$$S^{3/4}\frac{d^2S}{d\theta^2} = 12\frac{t_{II}^2p_*}{r_{II}^2\rho_0}\theta^{-4/9} = 0.278\theta^{-4/9}. \quad (5.22)$$

The term on the right hand side of the equation represents again the pressure on the shell by the hot gas. Without this term we get the ordinary Snowplough Phase solution and S is linear in θ (since we have replaced R by $S^{1/4}$).

During Phase II the motion is described by the Sedov - Taylor equation

(5.1). Replacing time and radius by the dimensionless variables gives

$$\begin{aligned}
r &= \left(\frac{2E}{\rho_0} \right)^{1/5} t^{2/5}, \\
Rr_{II} &= \left(\frac{2E}{\rho_0} \right)^{1/5} \theta^{2/5} t_{II}^{2/5}, \\
R &= \left(\frac{2E}{\rho_0} \right)^{1/5} \frac{t_{II}^{2/5}}{r_{II}} \theta^{2/5} = S^{1/4}, \\
S^{1/4} &= \frac{2^{1/5} E^{1/5}}{\rho_0^{1/5}} \cdot \left(\frac{0.216 E^{3/14}}{q^{5/14} \rho_0^{4/7}} \right)^{2/5} \cdot \frac{q^{1/7} \rho_0^{3/7}}{0.622 E^{2/7}} \theta^{2/5} \\
&= \underbrace{\frac{2^{1/5} 0.216^{2/5}}{0.622}}_1 \underbrace{\frac{E^{1/5} E^{3/35} q^{1/7} \rho_0^{3/7}}{E^{2/7} q^{1/7} \rho_0^{1/5} \rho_0^{8/35}}}_1 \theta^{2/5}, \\
S &= \theta^{8/5}. \tag{5.23}
\end{aligned}$$

The solution for (5.22) for Phase III matches smoothly onto that for Phase II above, which means that the initial conditions must fit for both equations. Since the end of Phase II is reached at $t = t_{II}$, $\theta = 1$ at that point of time and the initial conditions are

$$\begin{aligned}
S(1) &= 1, \\
S'(1) &= 8/5.
\end{aligned}$$

The solution for the differential equation Kahn obtained in his paper is somewhat different. Firstly there is a factor 0.278 and secondly he left out some higher order terms in his calculation. In Appendix D the method Kahn used to obtain his result is discussed and a more precise method is presented. Then the dimensionless radius becomes

$$S = 2.50\theta(1 - 0.356\theta^{-7/36} - 0.063x^{-7/18} - 0.030x^{-7/12} - \dots), \tag{5.24}$$

$$R = 1.258\theta^{1/4}(1 - 0.089\theta^{-7/36} - 0.016\theta^{-7/18} - 0.007\theta^{-7/12} - \dots). \tag{5.25}$$

The next question is, whether the radius really expands faster than without the pressure term. The usual equation of motion in our variables reads

$$\frac{d^2 S}{d\theta^2} = 0. \tag{5.26}$$

Integrating twice we obtain the linear solution

$$S = a\theta + b, \tag{5.27}$$

and using again our initial conditions the result is

$$S = \frac{8}{5}\theta - \frac{3}{5}, \quad (5.28)$$

$$R = \left(\frac{8}{5}\theta - \frac{3}{5}\right)^{1/4}. \quad (5.29)$$

To compare the solutions, we have plotted them in a diagram (Fig. 5.1), where we can clearly notice that the remnant moves faster with a pressure from inside, as desired.

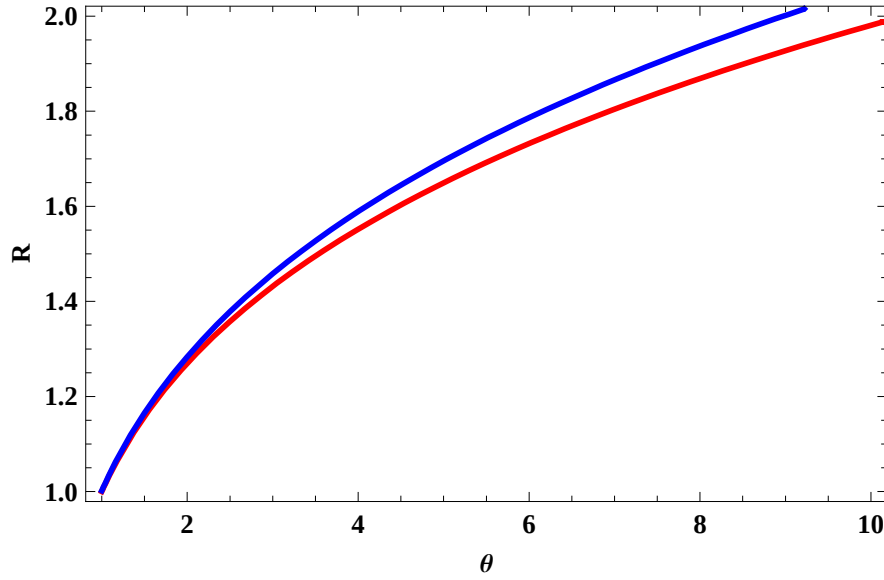


Figure 5.1: This plot demonstrates the difference between the Snowplough Phase without pressure (red line) and the solution from Kahn with pressure (blue line). Here we observe that the radius doubles when $\theta \sim 10$ without pressure and when $\theta \sim 9$ with pressure.

Because we got a slightly different value for the radius it will influence the following calculations and we will very often get other factors than Kahn. Next he introduces the dimensionless volume

$$V = \frac{4\pi}{3}R^3 = \frac{4\pi}{3}(1.258\theta^{1/4})^3 = 8.33\theta^{3/4}. \quad (5.30)$$

It is estimated that there are about three supernova explosions in the galaxy in a century, so the rate at which they occur in dimensionless variables can be expressed as

$$N \equiv \nu r_{II}^3 t_{II}, \quad (5.31)$$

with

$$v = (30 \text{ years})^{-1} (6 \times 10^{10} \text{ pc}^3)^{-1} = 6 \times 10^{-76} \text{ cm}^{-3} \text{ s}^{-1}. \quad (5.32)$$

Then it follows with the values of t_{II} and r_{II}

$$N = 6 \cdot 10^{-76} (7.32 \cdot 10^{19})^3 \cdot 1.45 \cdot 10^{12} = 0.00034.$$

The fractional volume occupied by supernova remnants for a time θ in a homogeneous ISM is

$$f = N \int_0^\theta V d\theta = 8.33N \int_0^\theta \theta^{3/4} d\theta = 4.76N\theta^{7/4}, \quad (5.33)$$

where $f = 1$ means that all the available space is filled with the remnants. In this case we get for $\theta = 0.41N^{-4/7} = 39.2$ which corresponds to a time $t = \theta t_{II} = 1.8 \times 10^6$ years. Then the radius of the remnant will have a value of $R = (\frac{3 \cdot 8.33}{4\pi} \theta^{3/4})^{1/3} = 3.15$, which is equivalent to $r = R r_{II} = 2.3 \times 10^{20} \text{ cm}$ or 74.7 pc. But, as Kahn says, these estimates are inaccurate. In simulations by Avillez, Berry and Kahn (1997) it is shown that the ISM never returns to its initial state and that all SN remnants evolve in an inhomogeneous medium.

The thermal energy is given by

$$\epsilon_{ther} = \frac{1}{\gamma - 1} \frac{P}{\rho_0},$$

with $\gamma = 5/3$ and the specific volume $v = 1/\rho_0 = (4\pi/3)r^3$ the equation becomes

$$\epsilon_{ther} = \frac{3}{2} \frac{4\pi}{3} P r^3 = 2\pi P r^3.$$

Now we insert P from (5.17), r_{II} from (5.9) and use $R^{-2} = (\frac{3 \cdot 8.33}{4\pi} \theta^{3/4})^{-2/3} = 0.63\theta^{-1/2}$ to compute the thermal energy content of a remnant in the Snow-plough Phase

$$\begin{aligned} \epsilon_{ther} &= 2\pi P R^3 r_{II}^3 \\ &= 2\pi \cdot 0.192 E^{1/7} q^{3/7} \rho_0^{9/7} R^{-5} \theta^{-4/9} R^3 \frac{(0.622)^3 E^{6/7}}{q^{3/7} \rho_0^{9/7}} \\ &= 0.291 E R^{-2} \theta^{-4/9} = 0.291 \cdot 0.63 E \theta^{-1/2} \theta^{-4/9} \\ &= 0.184 E \theta^{-17/18}. \end{aligned} \quad (5.34)$$

The thermal energy is always small when compared to E . For the dimensionless times $\theta = 10$ and $\theta = 20$ the values become $\epsilon_{ther} = 0.021E$ and $\epsilon_{ther} = 0.011E$.

5.2 A supernova exploding in a low density region

Since the Local Bubble evolved by many supernova explosions, it is necessary to see what happens with the equations, if a supernova explodes inside an existing SN remnant. The gas is already shocked by a previous supernova and has a density distribution of the form

$$\rho = \Omega r^n, \quad (5.35)$$

because the mass is swept up in the shell and therefore more concentrated as we approach the shell. In this case $n = 9/2$ (see 4.111). How this density distribution looks like, we can see in Fig. 5.2.

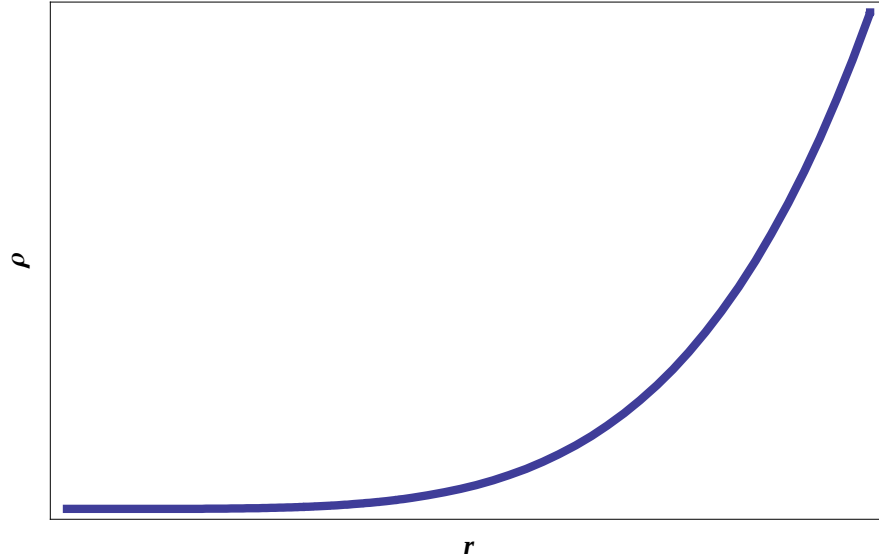


Figure 5.2: This Figure shows the density distribution $\rho \approx r^n$ with $n = 9/2$. We can see that most of the matter is swept up into a shell.

The event of a second supernova exploding is more likely to occur during Phase III than Phase II, because Phase III lasts much longer than Phase II. We will compute Ω by assuming that all the mass of the hot gas M_h from (5.11) is contained within a sphere with the radius $r = 1.258r_{II}\theta^{1/4}$ as derived before. Therefore we can write (5.11) in dimensionless variables

$$M_h = \frac{0.10E}{\rho_0^{2/3} q^{2/3}} \left(\frac{0.216E^{3/14}}{\rho_0^{4/7} q^{5/14}} \theta \right)^{-2/3} = \frac{0.277E^{6/7}}{\rho_0^{2/7} q^{3/7}} \theta^{-2/3}, \quad (5.36)$$

and express the hot mass in terms of the density distribution (5.35)

$$\begin{aligned}
dM_h &= 4\pi r^2 \rho dr = 4\pi r^{n+2} \Omega dr, \\
M_h &= \frac{4\pi}{n+3} r^{n+3} \Omega \\
&= \frac{8\pi}{15} \Omega r^{15/2} = \frac{8\pi}{15} \Omega (1.258 r_{II} \theta^{1/4})^{15/2} \\
&= \frac{8\pi}{15} \cdot 5.577 \Omega r_{II}^{15/2} \theta^{15/18}.
\end{aligned}$$

These masses are equal and by taking

$$r_{II}^3 = \left(\frac{0.622 E^{2/7}}{q^{1/7} \rho_0^{3/7}} \right)^3 = \frac{0.241 E^{6/7}}{q^{3/7} \rho_0^{9/7}},$$

and replacing it on the left hand side we will compute Ω

$$\begin{aligned}
\frac{0.277 E^{6/7}}{\rho_0^{2/7} q^{3/7}} \theta^{-2/3} &= \frac{8\pi}{15} \cdot 5.577 \Omega r_{II}^{15/2} \theta^{15/18}, \\
\frac{1.149}{\rho_0^{-1}} r_{II}^3 \theta^{-2/3} &= \frac{8\pi}{15} \cdot 5.577 \Omega r_{II}^{15/2} \theta^{15/18}, \\
\Omega &= 0.123 \rho_0 r_{II}^{-9/2} \theta^{-61/24}. \tag{5.37}
\end{aligned}$$

In the following Kahn states that we can find the effect of the explosion analytically, but we have to make some drastic approximations. A similarity solution requires the second explosion to occur at the same place as the first. This is not the case in the Local Bubble, because the star cluster that probably caused the bubble moved through the galaxy on the trajectory we described before. But it does not really make a difference if the second explosion is off-center, because the shock takes only little time to move through the hot thin gas compared to the time it takes to move through the thick shell. This means that the period of asymmetrical expansion is only brief and when it moves through the shell the shock is becoming more spherical.

Another approximation is that the energy content of the remaining hot gas is small compared to E , the energy of the explosion, but this assumption is easily met. The similarity solution also only applies, when the ejecta mass of the explosion is negligible. With $M_{II} = 4000 M_\odot$ ($1 M_\odot = 1.989 \times 10^{33} g$) and again the values $E = 10^{51} erg$, $\rho_0 = 2 \times 10^{-24} g cm^{-3}$ and $q = 4 \times 10^{32} cm^6 g^{-1} s^{-4}$ we can show from (5.36) that the mass of hot gas in the remnant is

$$M_h = \frac{9.026 \cdot 10^{35}}{4000 \cdot 1.989 \cdot 10^{33}} \theta^{-2/3} = 0.113 M_{II} \theta^{-2/3}. \tag{5.38}$$

If a supernova ejects, e.g., $5 M_\odot$ and the second explosion takes place about $t = 500000$ years after the first explosion, when $\theta = 10.82$, then the similarity solution holds if

$$5 \ll 453.8f \times 10.86^{-2/3} = 92.54f. \quad (5.39)$$

Here, f is the fraction of the hot diffuse gas that has been overrun by the second shock, by looking at (5.39) it has to exceed about 6 percent. This value is higher than the 2 percent from Kahn's paper, but it is still a small number, which means that in general the approximation holds.

The shock path can be found via the virial theorem

$$\frac{1}{2} \frac{d^2 I}{dt^2} = 2T + U_{\text{pot}},$$

where I is the moment of inertia, the potential energy is $U_{\text{pot}} = 0$ and the total kinetic energy is $T = E_{\text{ther}} + E_{\text{kin}} = E$. Then we get

$$\frac{1}{2} \ddot{I} = 2(E_{\text{ther}} + E_{\text{kin}}) \equiv E. \quad (5.40)$$

The moment of inertia is expressed as

$$I \equiv \int_0^{M(t)} r^2 dM. \quad (5.41)$$

Now we define the dimensionless radius and the dimensionless mass which are normalized by the shock radius r_s and the mass M_s that the shock encloses

$$\xi \equiv \frac{r}{r_s},$$

$$\mu \equiv \frac{M}{M_s}.$$

The motion of the shock leads to a change of the moment of inertia I in the shell, because the material within the shock radius has been displaced outwards. The moment of inertia I_1 becomes, with $r = \xi r_s$, $M = \mu M_s$ and therefore $dM = M_s d\mu$,

$$\begin{aligned} I_1 &= \int_0^{M_s} r^2 dM = \int_0^1 \xi^2 r_s^2 M_s d\mu \\ &= M_s r_s^2 \int_0^1 \xi^2 d\mu. \end{aligned}$$

The upper limit of the integral M_s was replaced by $\mu = M_s/M_s = 1$, because we normalized with the value of M_s . The moment of inertia I_2 , after the shock moved through the shell, can be derived by taking $dM = 4\pi r^2 \rho dr$ with $\rho = \Omega r^n$ in case of a general n and we get

$$\begin{aligned} I_2 &= \int_0^{M_s} r^2 dM = \int_0^{r_s} 4\pi r^4 \Omega r^n dr \\ &= 4\pi \Omega \int_0^{r_s} r^{n+4} dr. \end{aligned}$$

We will calculate this second integral by means of

$$\begin{aligned} M_s &= \frac{4\pi}{n+3} r_s^{n+3} \Omega \\ \Rightarrow \Omega &= \frac{n+3}{4\pi} \frac{M_s}{r_s^{n+3}}, \end{aligned} \tag{5.42}$$

and substitute $r^{n+4} = \xi^{n+4} r_s^{n+4}$ and $dr = r_s d\xi$. This leads to

$$\begin{aligned} I_2 &= 4\pi \frac{n+3}{4\pi} \frac{M_s}{r_s^{n+3}} \int_0^1 \xi^{n+4} r_s^{n+4} r_s d\xi \\ &= (n+3) M_s \frac{r_s^{n+5}}{r_s^{n+3}} \frac{1}{n+5} \xi^{n+5} \Big|_0^1 \\ &= M_s r_s^2 \frac{n+3}{n+5}. \end{aligned}$$

Then the change in the moment of inertia becomes

$$\begin{aligned} \Delta I &= I_1 - I_2 \\ &= M_s r_s^2 \int_0^1 \xi^2 d\mu - 4\pi \Omega \int_0^{r_s} r^{n+4} dr \\ &= M_s r_s^2 \int_0^1 \xi^2 d\mu - M_s r_s^2 \frac{n+3}{n+5} \\ &= M_s r_s^2 \left[\int_0^1 \xi^2 d\mu - \frac{n+3}{n+5} \right]. \end{aligned} \tag{5.43}$$

In a similarity solution, $M_s r_s^2$ varies like t^2 , because

$$\begin{aligned} M_s &\sim r_s^3 \sim t^{6/5}, \\ r_s^2 &\sim t^{4/5}, \\ \Rightarrow M_s r_s^2 &\sim t^{6/5} t^{4/5} = t^2, \end{aligned}$$

but we can also see from the virial theorem that

$$\frac{1}{2} \frac{d^2 I}{dt^2} = 2E \Rightarrow \frac{1}{2} \frac{dI}{dt} \sim 2Et \Rightarrow \frac{1}{2} I \sim Et^2 \Rightarrow 2E \sim \frac{I}{t^2}.$$

Then, with (5.42) and (5.43), it follows that

$$\begin{aligned} 2E &= \frac{1}{2} \frac{d^2 I}{dt^2} = \frac{M_s r_s^2}{t^2} \left[\int_0^1 \xi^2 d\mu - \frac{n+3}{n+5} \right] \\ &= \frac{4\pi\Omega r_s^{n+5}}{(n+3)t^2} \left[\int_0^1 \xi^2 d\mu - \frac{n+3}{n+5} \right]. \end{aligned} \quad (5.44)$$

We know from the Rankine-Hugoniot conditions that the density for a strong shock goes up by a factor 4. Then we write for the mass

$$\frac{M}{M_s} = 4 \frac{4\pi}{n+3} \frac{n+3}{4\pi} \frac{\Omega}{\Omega} \left(\frac{r}{r_s} \right)^{n+3},$$

which can be rewritten as $\mu = 4\xi^{n+3}$, and the derivative is $d\mu = 4(n+3)\xi^{n+2}d\xi$. Immediately behind the shock, where $r = r_s$ it follows that $\xi = 1$ and the boundary condition is therefore

$$\frac{d\mu}{d\xi} = 4(n+3). \quad (5.45)$$

Kahn then approximates the integral in (5.44) by assuming that μ increases with ξ by a power law as ξ approaches unity ($\mu = \xi^{4(n+3)}$ by integrating (5.45)) which can be substituted into

$$\begin{aligned} \int_0^1 \xi^2 d\mu &= \int_0^1 \xi^2 4(n+3) \xi^{4(n+3)-1} d\xi \\ &= 4(n+3) \int_0^1 \xi^{2+(4n+11)} d\xi \\ &= \frac{2(n+3)}{2n+7}. \end{aligned} \quad (5.46)$$

Then (5.44) leads to

$$2E = \frac{4\pi\Omega r_s^{n+5}}{(n+3)t^2} \left[\frac{2(n+3)}{2n+7} - \frac{n+3}{n+5} \right]. \quad (5.47)$$

Rearranging the equation we will compute r_s

$$\begin{aligned}
2E &= \frac{4\pi\Omega r_s^{n+5}}{(n+3)t^2} \frac{2(n+3)(n+5) - (n+3)(2n+7)}{(2n+7)(n+5)} \\
&= \frac{4\pi\Omega r_s^{n+5}}{t^2} \frac{3}{(2n+7)(n+5)} \\
\Rightarrow r_s^{n+5} &= \frac{(n+5)(2n+7)Et^2}{6\pi\Omega}.
\end{aligned} \tag{5.48}$$

The Sedov-solutions (5.1) and (5.2) combined with the density distribution (5.35) give us

$$\begin{aligned}
r_s &= \left(\frac{2E}{\Omega r_s^n} \right)^{1/5} t^{2/5} = \left(\frac{2E}{\Omega} \right)^{1/5} r_s^{-n/5} t^{2/5} \\
&= \left(\frac{2E}{\Omega} \right)^{1/(n+5)} t^{2/(n+5)}, \\
\dot{r}_s &= \frac{2}{n+5} \left(\frac{2E}{\Omega} \right)^{1/(n+5)} t^{2/(n+5)-1} = \frac{2r_s}{(n+5)t}.
\end{aligned} \tag{5.49}$$

Then with

$$\dot{r}_s^2 = \frac{4r_s^2}{(n+5)^2 t^2} \Rightarrow r_s^2 = \frac{\dot{r}_s^2 (n+5)^2 t^2}{4}$$

and

$$r_s^{n+5} = r_s^2 r_s^{n+3},$$

we rewrite (5.48)

$$\begin{aligned}
\frac{\dot{r}_s^2 (n+5)^2 t^2}{4} r_s^{n+3} &= \frac{(n+5)(2n+7)Et^2}{6\pi\Omega}, \\
\dot{r}_s^2 &= \frac{2(2n+7)E}{3\pi(n+5)\Omega} r_s^{-(n+3)},
\end{aligned} \tag{5.50}$$

and with $n = 9/2$ we find for the shock-velocity of the second shock

$$\dot{r}_s^2 = \frac{2}{3\pi} \frac{16}{9.5} \frac{E}{\Omega} r_s^{-15/2} = 0.36 \frac{E}{\Omega} r_s^{-15/2}. \tag{5.51}$$

The second shock reaches the edge of the original remnant when

$$r_s = 1.258 r_{II} \theta^{1/4}, \tag{5.52}$$

which is the solution of the differential equation we got in (5.25). Now we insert this equation and Ω from (5.37) into (5.51)

$$\begin{aligned} \dot{r}_s^2 &= 0.36 \frac{E}{0.123 \rho_0 r_{II}^{-9/2} \theta^{-61/24}} (1.258 r_{II} \theta^{1/4})^{-15/2} \\ &= 0.521 \frac{E}{\rho_0 r_{II}^3} \theta^{2/3}. \end{aligned} \quad (5.53)$$

Continuing the calculation, taking t_{II} and r_{II} from (5.8) and (5.9), we derive for the shock speed

$$\begin{aligned} \dot{r}_s^2 &= 0.521 \frac{E}{\rho_0} \frac{r_{II}^2}{r_{II}^5} \theta^{2/3} = 0.521 \frac{E}{\rho_0} r_{II}^2 \frac{q^{5/7} \rho_0^{15/7}}{(0.622)^5 E^{10/7}} \theta^{2/3} \\ &= \frac{0.521}{(0.622)^5} r_{II}^2 \frac{q^{10/14} \rho_0^{8/7}}{E^{6/14}} \theta^{2/3} = \frac{0.521}{(0.622)^5} r_{II}^2 \left(\frac{0.216}{0.216} \right)^2 \left(\frac{q^{5/14} \rho_0^{4/7}}{E^{3/14}} \right)^2 \theta^{2/3} \\ &= \frac{0.521 (0.216)^2}{(0.622)^5} \left(\frac{r_{II}}{t_{II}} \right)^2 \theta^{2/3} = 0.261 \left(\frac{r_{II}}{t_{II}} \right)^2 \theta^{2/3}, \end{aligned} \quad (5.54)$$

and obtain a value of about $258 \theta^{1/3} \text{ kms}^{-1}$ there. The dynamical time can be derived from (5.49) with $t = t_{dyn}$ and $n = 9/2$

$$t_{dyn} = \frac{4}{19} \frac{r_s}{\dot{r}_s}. \quad (5.55)$$

With (5.52) and (5.54) the dynamical time becomes

$$t_{dyn} = \frac{4}{19} \frac{1.258 r_{II} \theta^{1/4}}{(0.261)^{1/2} \left(\frac{r_{II}}{t_{II}} \right) \theta^{1/3}} = 0.519 t_{II} \theta^{-1/12}. \quad (5.56)$$

The dynamical time has to be compared with the cooling time. If $t_{cool} \gg t_{dyn}$ then the ISM can reach hydrostatic equilibrium. But with the present values of physical parameters the left hand side is smaller and the ISM can never attain equilibrium unless it is heavily depleted in metals. Taking the adiabatic parameter κ_s from (5.3) and $\rho = \Omega r^{9/2}$ we get for t_{cool}

$$t_{cool} = \frac{\kappa_s^{3/2}}{q} = \frac{0.02 \dot{r}_s^3}{q \rho} = \frac{0.02 \dot{r}_s^3}{q \Omega r_s^{9/2}}. \quad (5.57)$$

Now we insert r_s , $\dot{r}_s$ and Ω from (5.52), (5.54) and (5.37)

$$\begin{aligned}
t_{cool} &= \frac{0.02(0.261)^{3/2} \left(\frac{r_{II}}{t_{II}}\right)^3 \theta}{q \cdot 0.123 \rho_0 r_{II}^{-9/2} \theta^{-61/24} \cdot (1.258)^{9/2} r_{II}^{9/2} \theta^{9/8}} \\
&= \frac{0.02(0.261)^{3/2}}{0.123(1.258)^{9/2}} \frac{r_{II}^3}{q \rho_0 t_{II}^3} \theta^{29/12} \\
&= \frac{0.02(0.261)^{3/2}}{0.123(1.258)^{9/2}} \frac{r_{II}^3}{q \rho_0} t_{II}^{-4} \theta^{29/12} \\
&= \frac{0.02(0.261)^{3/2}}{0.123(1.258)^{9/2}} \frac{t_{II}}{q \rho_0} \frac{(0.622)^3 E^{6/7}}{q^{3/7} \rho_0^{9/7}} \frac{q^{20/14} \rho_0^{16/7}}{(0.216)^4 E^{12/14}} \theta^{29/12} \\
&= 0.854 t_{II} \theta^{29/12}.
\end{aligned} \tag{5.58}$$

Then with these values found for t_{dyn} and t_{cool} we will do the comparison

$$\frac{t_{cool}}{t_{dyn}} = \frac{0.854}{0.519} \frac{t_{II}}{t_{II}} \frac{\theta^{29/12}}{\theta^{-1/12}} = 1.646 \theta^{5/2}. \tag{5.59}$$

As we derived in (4.16) the post shock temperature is

$$T_s = \frac{3}{16} \frac{\mu m_u}{k} \dot{r}_s^2.$$

With (5.54) it becomes

$$T_s = \frac{3}{16} \frac{\mu m_u}{k} \cdot 0.261 \left(\frac{r_{II}}{t_{II}}\right)^2 \theta^{2/3}.$$

Inserting $r_{II} = 7.32 \times 10^{19}$ cm, $t_{II} = 1.45 \times 10^{12}$ s, $\mu = 0.61$, $m_u = 1.66 \times 10^{-24}$ g and $k = 1.38 \times 10^{-16}$ erg K⁻¹ the post shock temperature is about

$$T_s = 0.92 \times 10^6 \theta^{2/3} K. \tag{5.60}$$

This value can be quite high if θ is large enough, that is if the second supernova is delayed sufficiently long.

CHAPTER 6

Results and Discussion

In the previous chapters we derived the methods needed to reproduce the measured ^{60}Fe abundance in the ferromanganese crust 237KD as was shown in Fig. 2.4. The only issue not addressed yet is how much of the total mass is ejected by a supernova explosion. For this work we have used the values Woosley and Heger [50] provided, together with the ^{60}Fe yields. The ejected mass depends linearly on the mass of the star as we can see in Fig. 6.1, where the green dots represent the mass that was ejected by the supernova explosion and the red dots the mass that is lost before the explosion due to stellar wind. The amount of ^{60}Fe that is lost by stellar wind in the used mass range of 8 - 20 $M_{\odot}$ is negligible.

Now we are prepared to calculate diagram 2.4 analytically by using the steps listed below:

1. Compute the explosion times. This was done in Chapter 3.
2. Calculate the radius and the duration of the free expansion as it is shown in Chapter 4.2.1.
3. Use formula (5.48) in Chapter 5 to derive the remaining time the remnant takes to hit the Earth.
4. With the information about the total time the remnant needs and the ^{60}Fe yield from Chapter 1.4 the fluence can be calculated with formula (1.1) from Chapter 1.5.

Eight different cases have been considered, with varying densities and the two different mass models explained in Chapter 3. These were obtained by

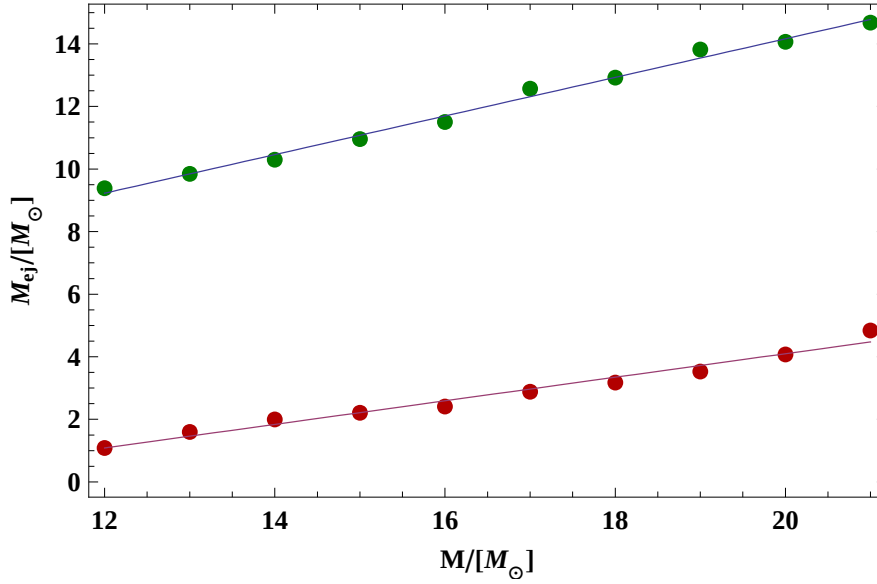


Figure 6.1: The mass loss of stars due to stellar wind (red) and the supernova explosion (green) as a function of the initial masses of the stars.

taking the boundary and the mean values of the bins in the IMF resulting into the mass distributions shown in Fig. 3.6 and 3.7. The values of the density were taken to be 0.01, 0.1, 1 and 10 atoms/cm³, since the ISM density can vary, because in reality the ISM not uniform, but turbulent and inhomogeneous. The resulting graphs are given below starting with the boundary mass model and ascending densities followed by the mean mass model. The computed data for the incoming ⁶⁰Fe are plotted on the measured data for comparison.

Together with every plot, possible trajectories for every explosion are shown in the x-y and x-z plane. For comparison the trajectories of the still existing stars are also plotted. The two (in the case of particle density $n = 1$ atom/cm³ in the mean mass model three (see Fig. 6.8)) explosions that give the highest amount of ⁶⁰Fe are shown in a slightly different color. As before in Chapter 3 the stars are separated by their present-day membership of the subgroups UCL (blue) and LCC (red) of the Scorpius OB2 association. The pseudotrajectories result from the errors of the velocities of the 79 stars that were extracted by Fuchs et al [17] as was also described in Chapter 3.

As we see, the analytical data fits very well to the measurements. It was also possible to find trajectories to the positions of the explosions, so that

enough ^{60}Fe arrives on Earth. Only two times we used a pseudo-trajectory of the extreme trajectory for the minimal distance. This was the case for the boundary models with the densities of 0.1 and 10 atoms/cm³. For every explosion there was more than one alternative for a pseudo-trajectory. We did not use any paths twice to show a variety of possibilities. The positions of the trajectories are well distributed, but show a tendency to reside at the solar-facing side of the cluster. Especially at the peak-positions we had to use closer explosion distances (72 - 80 pc) in order to achieve a sufficient amount of ^{60}Fe . For comparison of the influence of the ISM density on the time and the amount of ^{60}Fe arriving on Earth, we varied the densities with the trajectories. This is shown in Appendix E.

One has to note that there are quite a few uncertainties in the model. The largest uncertainty arises from the unknown ^{60}Fe yields as we had to fit the existing abundances of AGB stars in the mass range of 2 - 6.5 $M_{\odot}$ from Karakas and Lattanzio [23] with the yields of massive stars in the mass range of 12 - 20 $M_{\odot}$ from Woosley and Heger [50]. It was discussed in Chapter 1.4 that the resulting yields are very sensitive to the initial conditions in numerical calculations. They depend also on the different star types. It is possible that the ejected ^{60}Fe abundances in SAGB stars are totally different from the abundances in stars with larger masses and thus the interpolations we used. Further computations in the important mass range of 8 - 11 $M_{\odot}$ will shed light on this subject.

Another large uncertainty occurs from the uptake factor mentioned in Chapter 1.5, because it was derived by ^{53}Mn measurements from 1974 and 1979 in the Antarctic ice [6] and the deep-sea sediment [20]. New, more detailed analyses would either confirm the uptake factor or show that with modern methods a different value is derived.

The IMF that was taken from [17] to obtain the masses of the exploding stars also has uncertainties. But since we have the mass range of 8.2 - 20 $M_{\odot}$ we already used the minimal possible masses, because stars with masses less than 8 $M_{\odot}$ do not explode as supernovae. That means if the IMF would change the masses could only be higher and therefore the yields would increase.

The model from Kahn that was used to compute a supernova remnant exploding in a region where the ISM has already been processed by a previous supernova, was derived only for the case when two supernovae occur. The first supernova runs into a homogeneous ambient density and the second one runs into the medium where the mass has been swept up into a shell with the density profile $\rho \sim r^{9/2}$. This profile was taken for every supernova. To compute the time the remnant takes in Phase II one

should calculate how the density profile changes after each supernova. This could be done with the similarity solution described in Chapter 4.3 and then the new density profile must be inserted into the Kahn solution.

The following Figs. 6.2 - 6.9 show the analytical data plotted with the ^{60}Fe measurements. We calculated the data with the boundary and the mean mass model. Both models were computed with different ISM densities, varying from 0.01 - 10 atoms/cm³.

In Figs. 6.10 - 6.17 the pseudotrajectories (green) for each model and density are plotted. For comparison the trajectories of the still existing stars of the stellar moving group are also shown (grey).

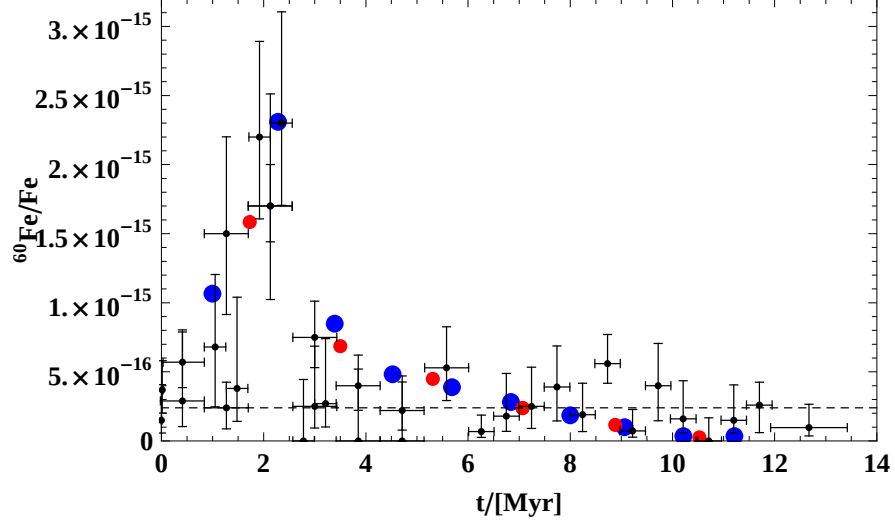


Figure 6.2: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 0.01 \text{ atoms/cm}^3$.

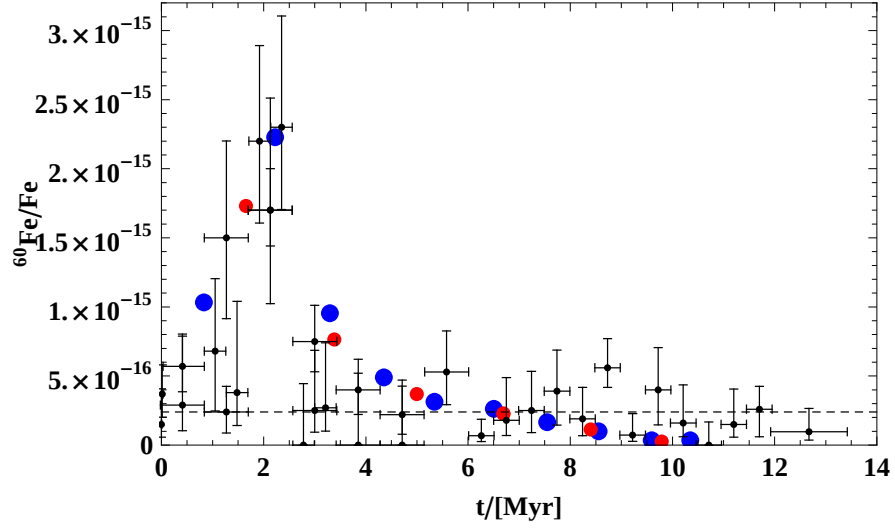


Figure 6.3: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 0.1 \text{ atoms/cm}^3$.

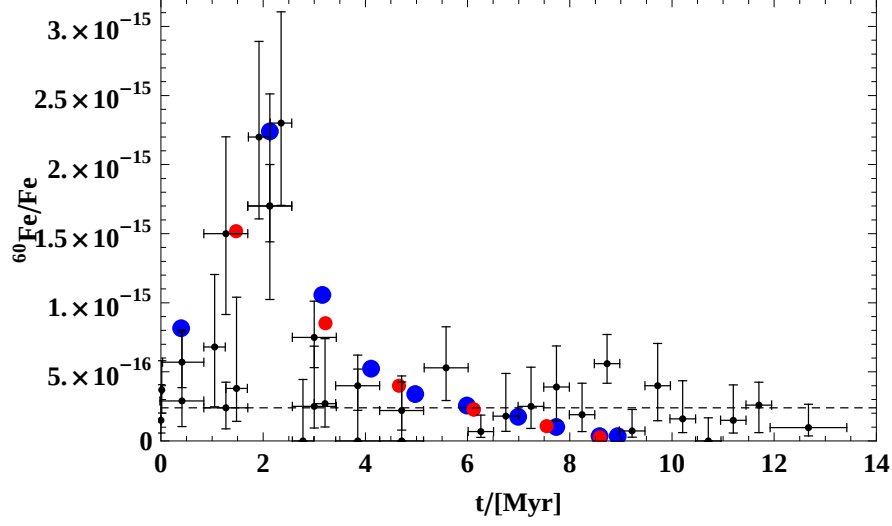


Figure 6.4: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 1 \text{ atom/cm}^3$.

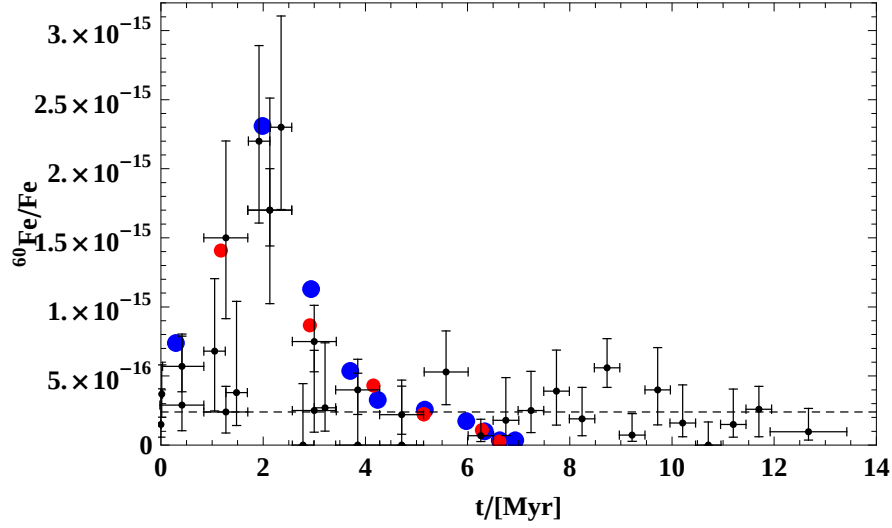


Figure 6.5: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 10 \text{ atom/cm}^3$.

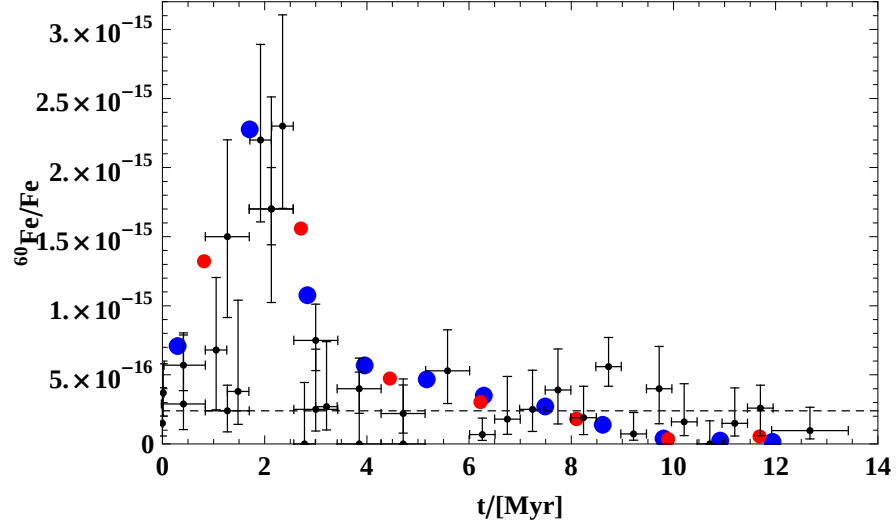


Figure 6.6: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 0.01 \text{ atoms/cm}^3$.

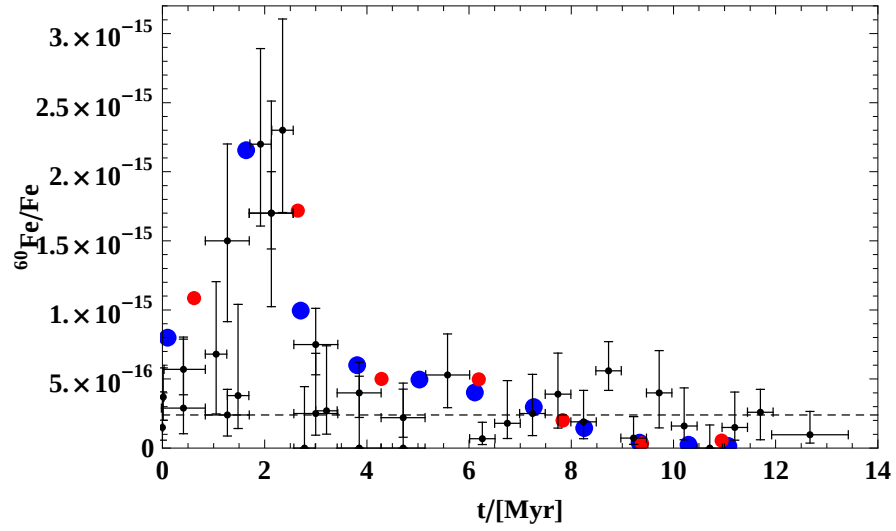


Figure 6.7: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 0.1 \text{ atoms/cm}^3$.

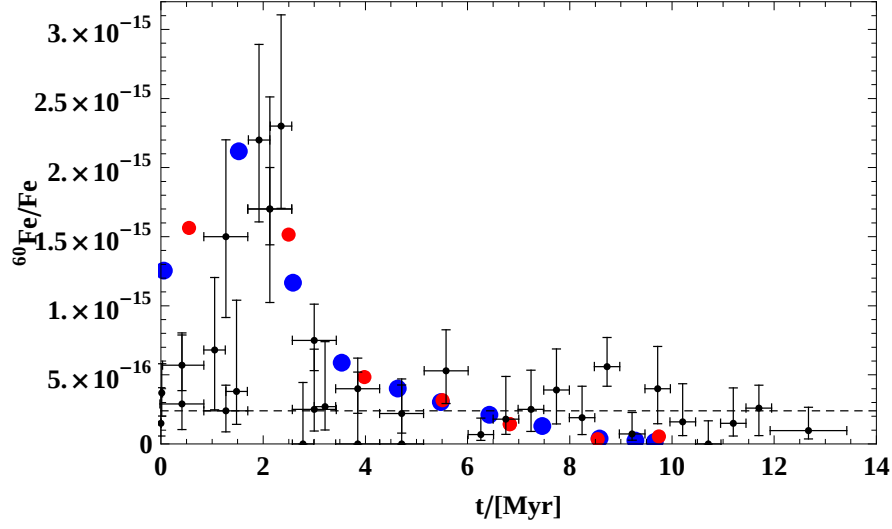


Figure 6.8: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 1 \text{ atom/cm}^3$.

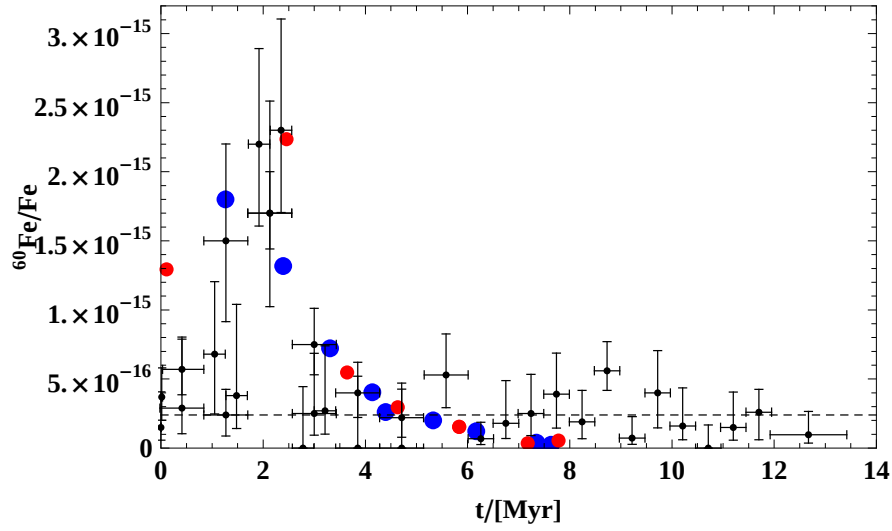


Figure 6.9: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 10 \text{ atoms/cm}^3$.

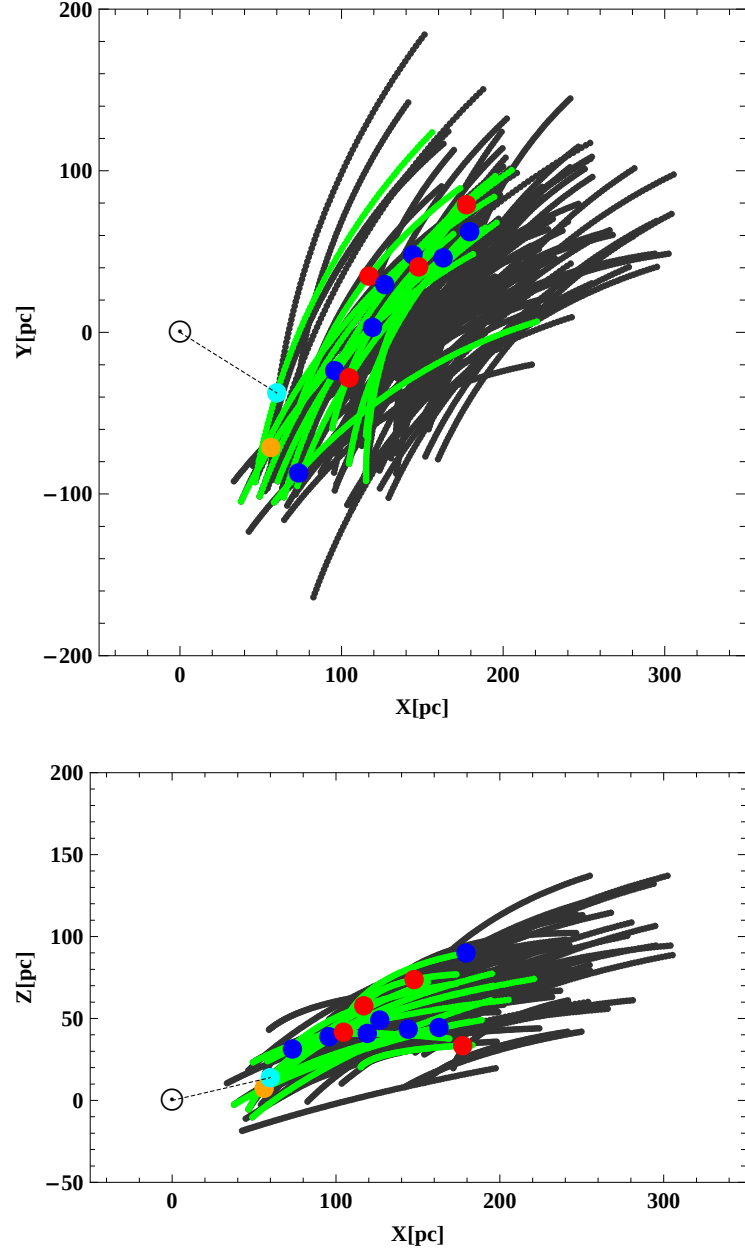


Figure 6.10: The trajectories for the explosions which are responsible for the ^{60}Fe profile of Fig. 6.2 in the x-y and x-z plane of our galaxy. The Sun is located at the coordinates (0,0,0). The two explosions, which contribute most to the ^{60}Fe distribution, are depicted in a slightly different color. The closest explosion, marked with a dashed line, occurred at a distance of 71.8 pc to the sun.

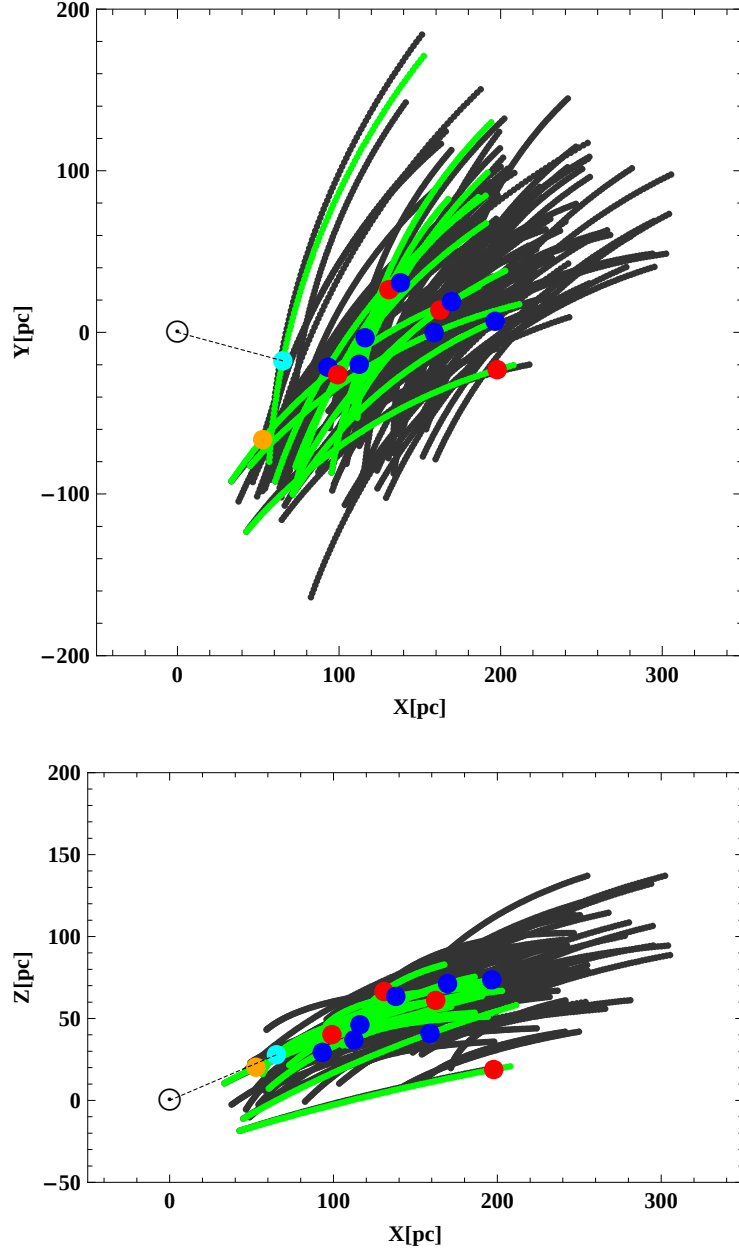


Figure 6.11: The trajectories for the explosions which are responsible for the ^{60}Fe profile of Fig. 6.3 in the x-y and x-z plane of our galaxy. The Sun is located at the coordinates (0,0,0). The two explosions, which contribute most to the ^{60}Fe distribution, are depicted in a slightly different color. The closest explosion, marked with a dashed line, occurred at a distance of 73.1 pc to the sun.

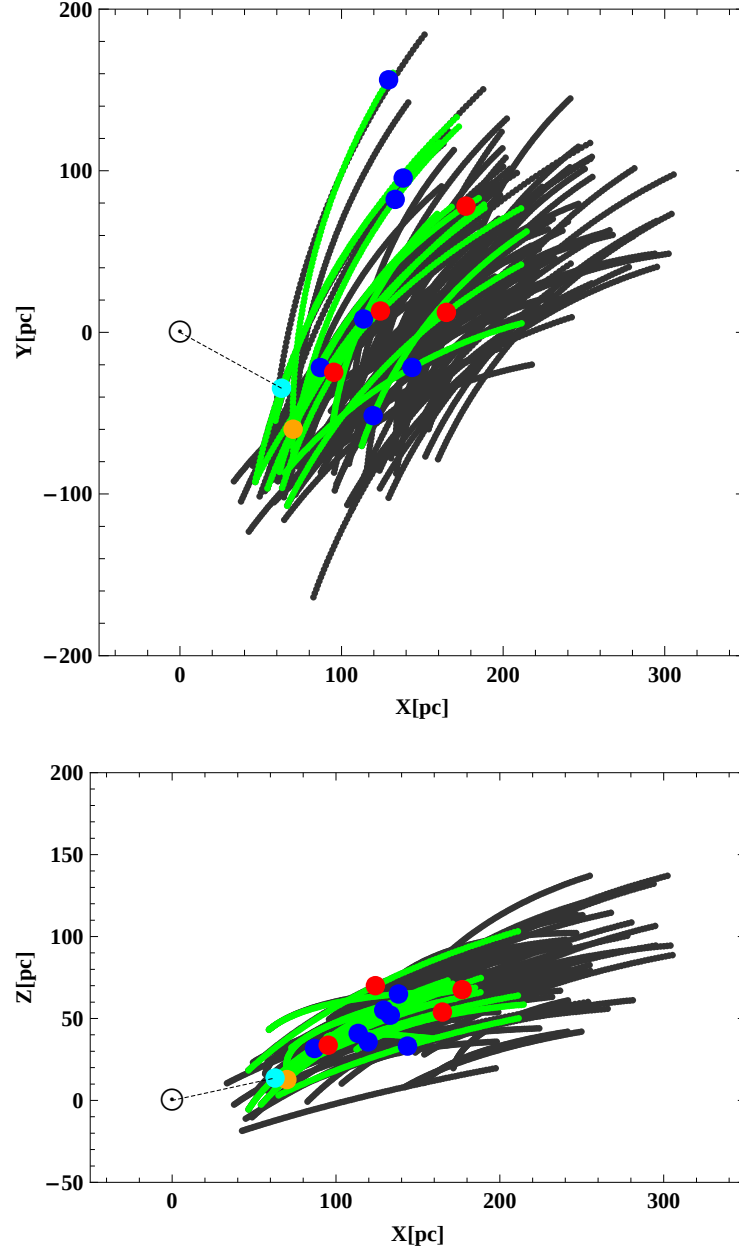


Figure 6.12: The trajectories for the explosions which are responsible for the ^{60}Fe profile of Fig. 6.4 in the x-y and x-z plane of our galaxy. The Sun is located at the coordinates (0,0,0). The two explosions, which contribute most to the ^{60}Fe distribution, are depicted in a slightly different color. The closest explosion, marked with a dashed line, occurred at a distance of 72.9 pc to the sun.

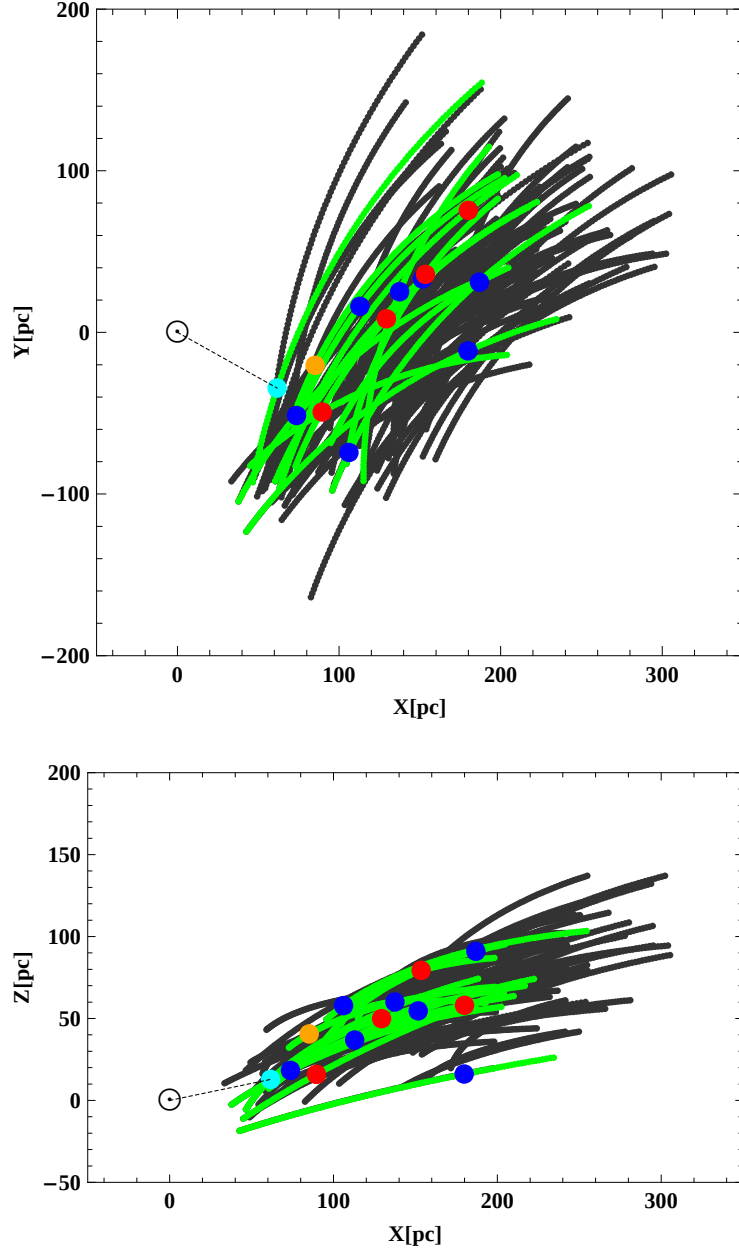


Figure 6.13: The trajectories for the explosions which are responsible for the ^{60}Fe profile of Fig. 6.5 in the x-y and x-z plane of our galaxy. The Sun is located at the coordinates (0,0,0). The two explosions, which contribute most to the ^{60}Fe distribution, are depicted in a slightly different color. The closest explosion, marked with a dashed line, occurred at a distance of 71.8 pc to the sun.

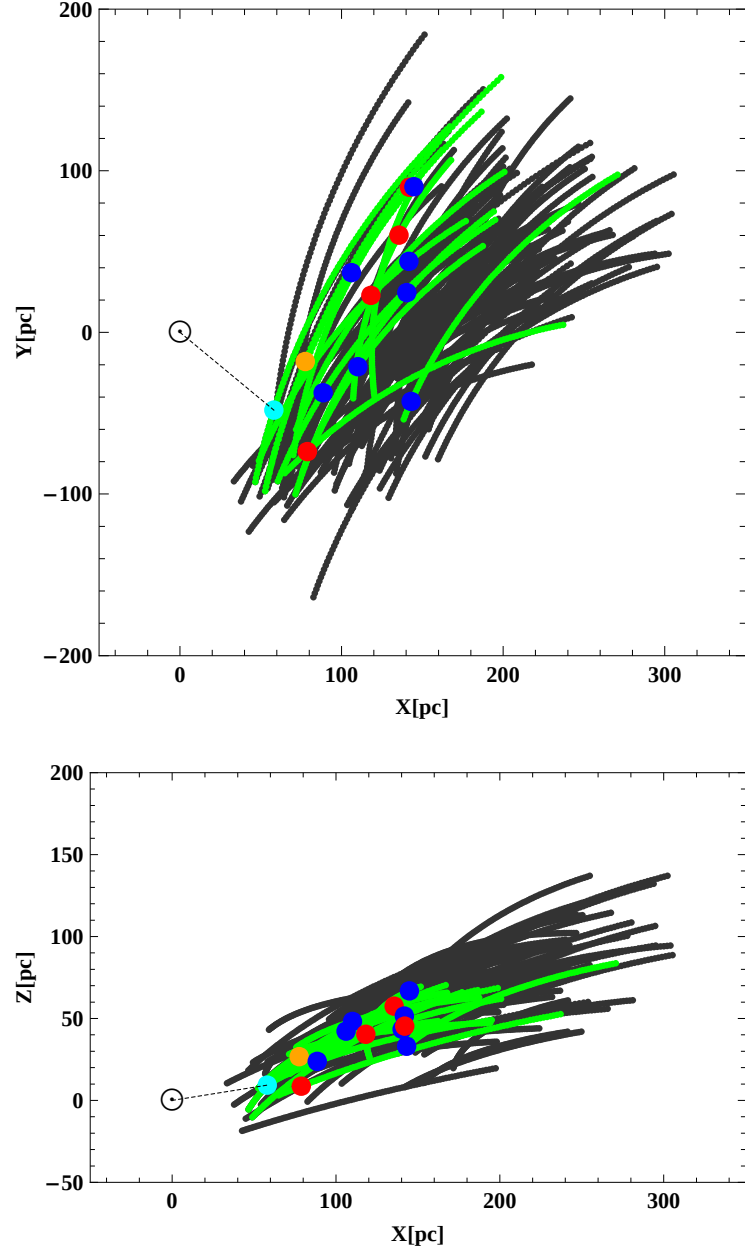


Figure 6.14: The trajectories for the explosions which are responsible for the ^{60}Fe profile of Fig. 6.6 in the x-y and x-z plane of our galaxy. The Sun is located at the coordinates (0,0,0). The two explosions, which contribute most to the ^{60}Fe distribution, are depicted in a slightly different color. The closest explosion, marked with a dashed line, occurred at a distance of 76.3 pc to the sun.

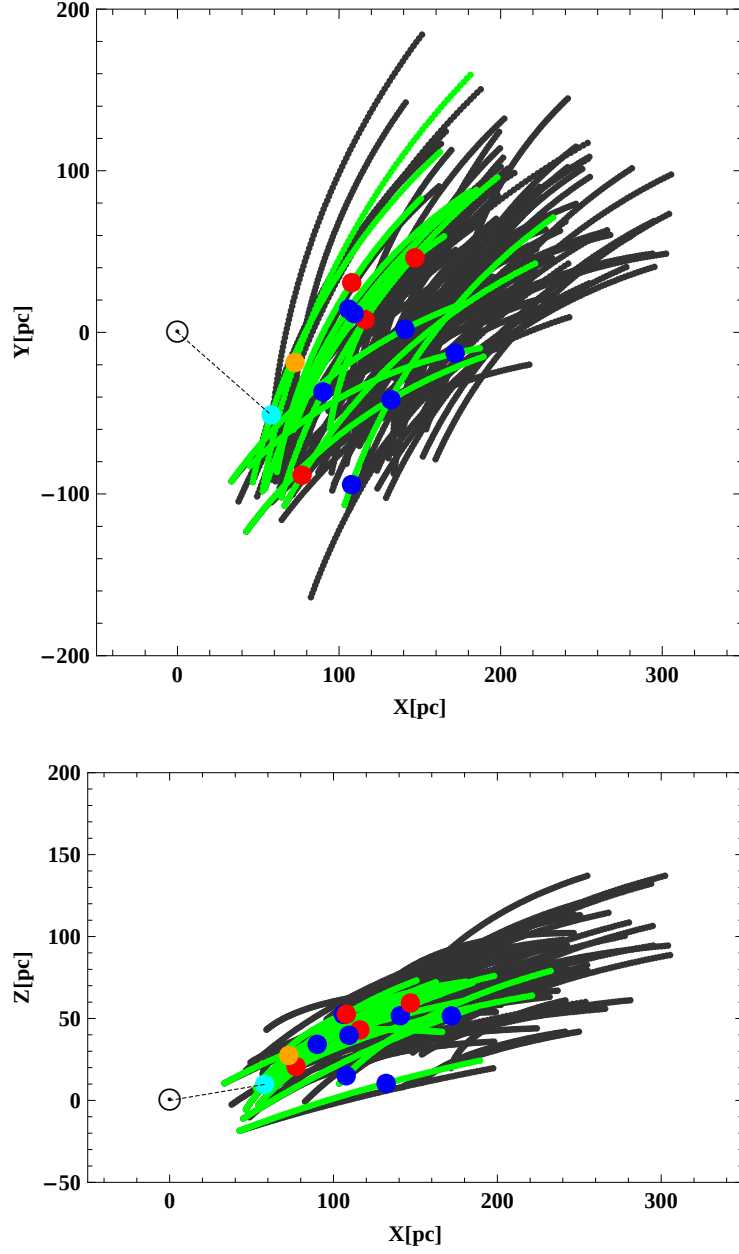


Figure 6.15: The trajectories for the explosions which are responsible for the ^{60}Fe profile of Fig. 6.7 in the x-y and x-z plane of our galaxy. The Sun is located at the coordinates (0,0,0). The two explosions, which contribute most to the ^{60}Fe distribution, are depicted in a slightly different color. The closest explosion, marked with a dashed line, occurred at a distance of 78.4 pc to the sun.

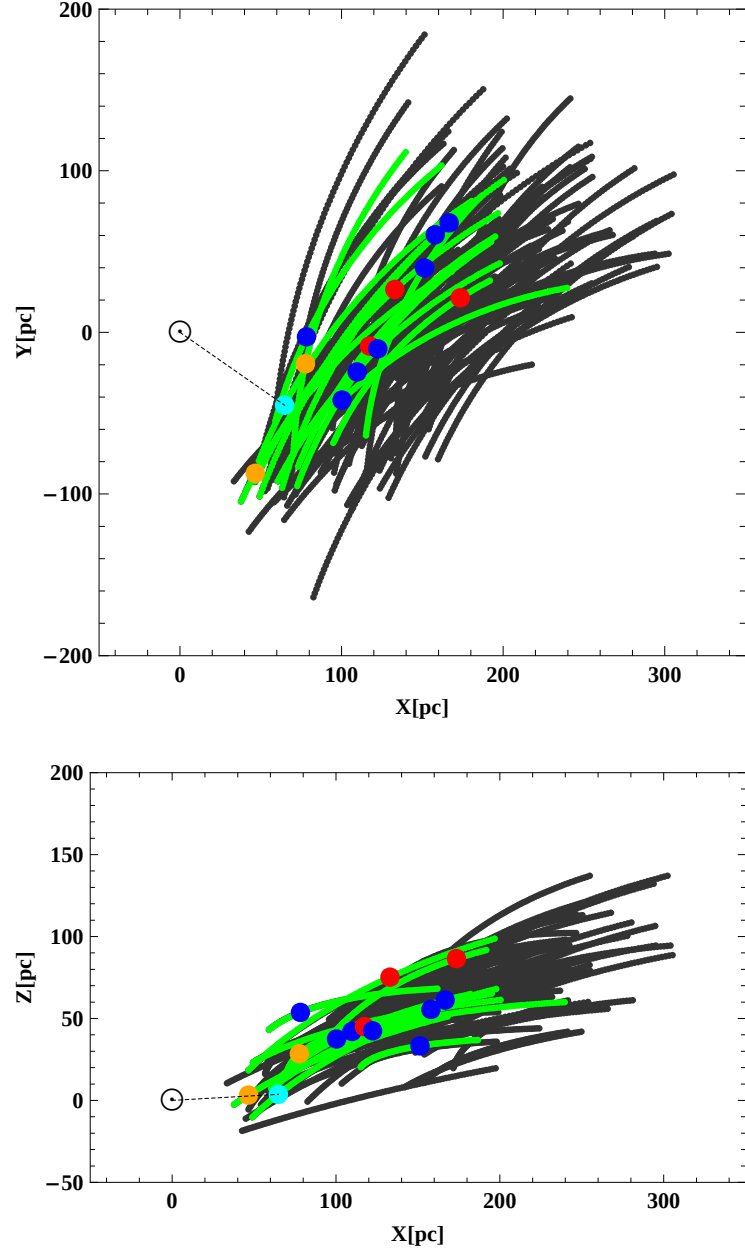


Figure 6.16: The trajectories for the explosions which are responsible for the ^{60}Fe profile of Fig. 6.8 in the x-y and x-z plane of our galaxy. The Sun is located at the coordinates (0,0,0). The three explosions, which contribute most to the ^{60}Fe distribution, are depicted in a slightly different color. The closest explosion, marked with a dashed line, occurred at a distance of 79.1 pc to the sun.

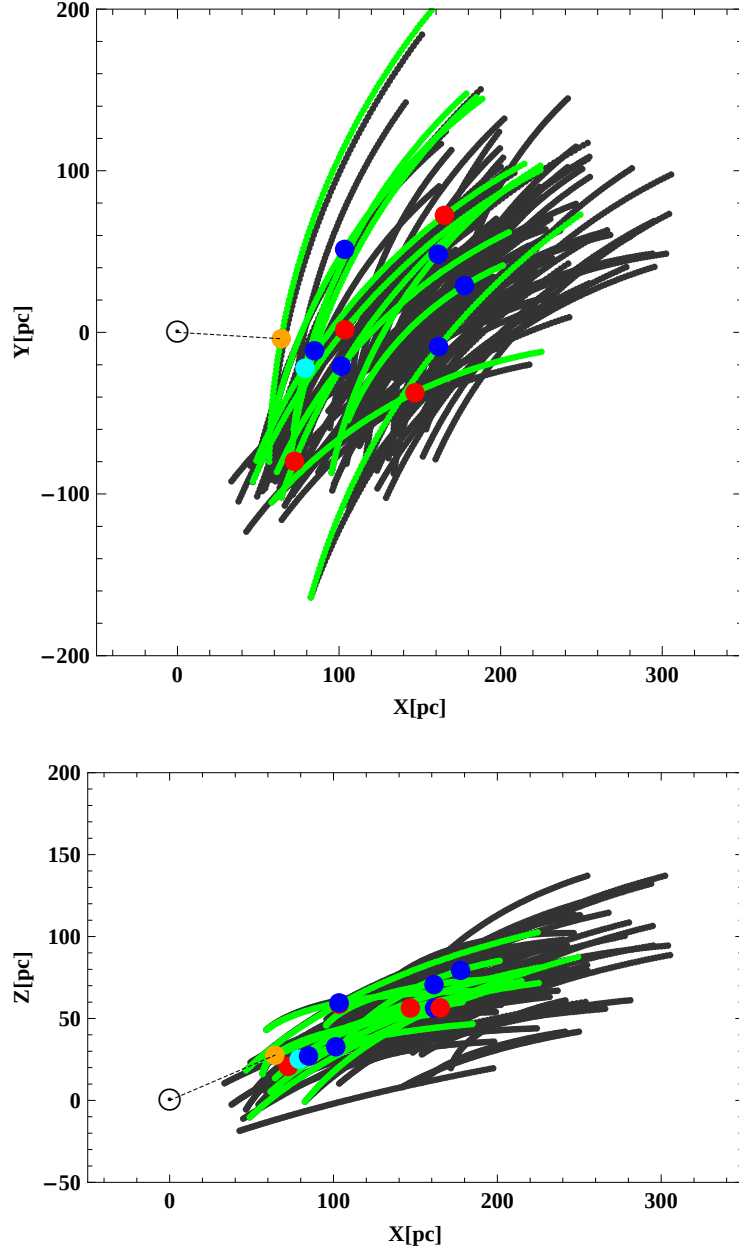


Figure 6.17: The trajectories for the explosions which are responsible for the ^{60}Fe profile of Fig. 6.9 in the x-y and x-z plane of our galaxy. The Sun is located at the coordinates (0,0,0). The two explosions, which contribute most to the ^{60}Fe distribution, are depicted in a slightly different color. The closest explosion, marked with a dashed line, occurred at a distance of 69.9 pc to the sun.

CHAPTER 7

Conclusions

We have used the IMF derived by Fuchs et al. [17] to calculate the explosion times of the Supernovae of the stars that belong to the Scorpius OB2 association today. Then we have computed the time the remnant takes to hit the Earth with the free expansion and the Kahn model. The remnant propagates undamped into the surrounding medium until the ejected mass equals the swept up mass and then slows down as it enters the second phase. It should be emphasized that, instead of the common Sedov-solution, we used the alternative Kahn model (1998) [22] to evaluate the expansion of the SN remnants in Phase II. This model includes a medium that has been shaped by a previous Supernova explosion, whereas the Sedov-model uses a homogeneous ambient medium. This means the Kahn model describes SNe exploding inside an existing supernova remnant, which in this case is the Local Bubble. One has to realize that standard superbubble models (e.g. Mac Low and McCray, 1988 [31]) are not applicable here, because we do not have to describe the time-dependent evolution of the Local Bubble shell, but the shock (containing the ejected ^{60}Fe) of an individual remnant within the bubble, when it hits the surface of the Earth, long before it catches up with the outer shell.

With the ^{60}Fe yields provided by Karakas and Lattanzio [23] and Woosley and Heger [50] we have been able to calculate the amount of this isotope that reaches the Earth.

Figs. 6.2 - 6.9 show the result of these calculations. The analytically computed data mostly fits the measurements fairly well. Not only coincides the peak position of the measurements very well with the closest point of

the stellar moving group to Earth, we are also able to fit the whole ^{60}Fe profile. This means it is very likely that not only one, but multiple supernovae produced this distribution.

With an ISM density of 10 atoms/cm^3 the most recent SN in the mean mass model did not reach the Earth anymore, because the remnant expanded too slowly. Therefore we can conclude, that the medium must be sufficiently thin (less than 1 atom/cm^3) for the remnant to arrive on Earth in time.

Despite all the uncertainties that were discussed in Chapter 6 we obtained the given ^{60}Fe profile, especially the peak position. It was possible to find pseudo-trajectories associated with the distances of the explosions to account for sufficient ^{60}Fe reaching the Earth. These are plotted in Figs. 6.10 - 6.17.

In Appendix E the trajectories we found for the special cases are plotted with varying densities to compare the influence on the surrounding density on the expansion of the remnant. We find, that the results still fit very good to the measurements, which means that there is a great variety of locations for the explosions in the stellar moving group.

As the pseudotrajectories were computed from the trajectories of existing stars it is likely they may have existed, because they all are situated within the paths of the real stars. The explosions that caused the peak positions are all located on the edge of the star cluster. This is not a problem, because it is not unlikely that one or more Supernovae occurred in the solar vicinity at the distances of 72 - 80 pc.

7.1 Further Measurements

To confirm the ^{60}Fe peak in the hydrogenetic ferromanganese crust 237KD, signatures have been searched in other reservoirs.

7.1.1 ^{60}Fe in a marine sediment

A marine sediment has the advantage that the time resolution is much higher due to a larger growth rate. The disadvantage is that the concentration of external isotopes is diluted and so lower in every layer, because the sediment grows faster.

Fitoussi et al. 2008 [16] analyzed a sediment that originated from the North Atlantic at the coordinates 66°55.5' N and 6°27' W. The position is marked in the figure in Appendix A. The authors measured the sediment in the time interval of 1.68 - 3.2 Myr, the same time the anomaly was found in the crust. But the signal they discovered was much lower than expected. Around the peak position the sediment was measured longer, but it did not show a signal that differed from the other time ranges. Fitoussi et al. gave the following possible explanations:

- The fluence in the crust was overestimated, because the uptake efficiency was derived wrong.
- The deposition time is very long.
- The sediment did not incorporate any ^{60}Fe isotopes.
- The ^{60}Fe signal in the crust did not increase due to supernova explosions, but there is another source.

The last explanation was considered by Basu et al. 2007 [2] who suggest micrometeorites as a possible source. But this was disproved by Fitoussi et al., because the peak in the crust ^{237}Kd was found at different drill holes and the micrometeorites should be a local phenomenon.

At least the eventual conclusion by Fitoussi et al. is reassuring: if the lowest signal is taken as the background in the sediment there appears to be a marginal increase of the $^{60}\text{Fe}/\text{Fe}$ ratio at 2.4 Myr.

7.1.2 ^{60}Fe on the Moon

The Moon has the advantage that there is no atmosphere, so the sedimentation rate is negligible. The concentration should be high, because the signal is not diluted.

Two lunar cores were analyzed for possible ^{60}Fe anomalies by Cook et al. 2009 [11]. Altogether 12 samples were measured, but the predicted ^{60}Fe fluences were 9 - 15 times larger than the observed ones. Only one of the 12 samples showed a signal above the background, which is produced by the reaction of galactic cosmic rays with iron meteorites. As an explanation Cook et al. suggest that the predicted fluence was overestimated by Fitoussi et al. [16] or that the supernova explosions occurred at a certain angle that was inclined to the plane of the ecliptic.

7.1.3 ^{60}Fe in another crust

An increase of the ^{60}Fe concentration was also found in the crust 29DR-32 originating from the position $28^{\circ}13'\text{N}$ and $177^{\circ}22'\text{W}$, 4000 km away from 237KD in the North Pacific. The measurement time was short compared to 237KD resulting in only few data points with high errors. The data has not been published yet (Korschinek et al. , private communication). A clear peak is visible in the required time span (Fig. 7.1). A confirmation of this anomaly would corroborate the theory that it was caused by SN explosions.

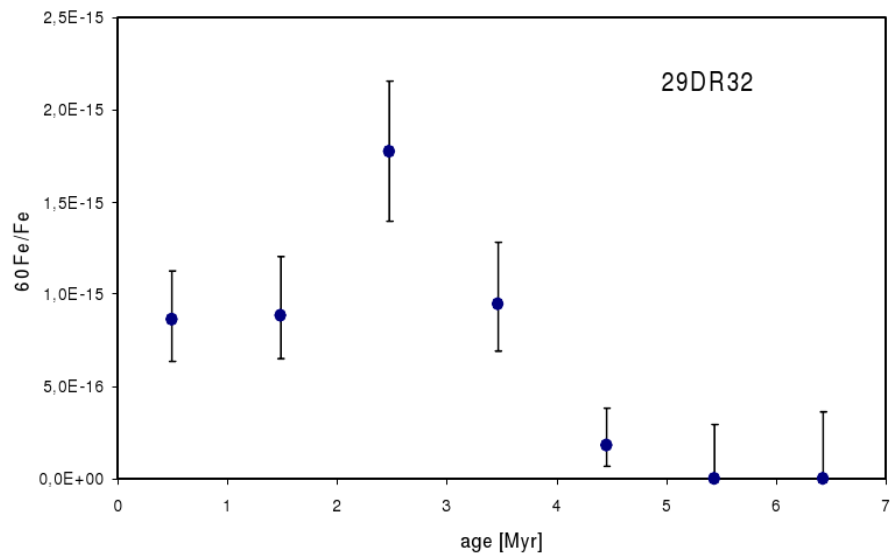


Figure 7.1: Unpublished measurements of ^{60}Fe in the ferromanganese crust 29DR-32, where a clear signal is visible. Courtesy of Korschinek et al.

CHAPTER 8

Future Work

If we want to improve the calculations we must use modeled yields for the mass range of 8 - 20 $M_{\odot}$, which will hopefully be available soon. It would also be interesting to analyze how probable the pseudotrajectories are, that means how high the possibility is that an explosion goes off at a certain trajectory. Computations concerning this problem are already in progress. More details could be added to the calculations, for example diffusion, the magnetic field or the impact of the solar wind. The Kahn model could be modified, as it was explained in Chapter 6. This work studied the expansion of Supernova remnants analytically. Numerical calculations will be carried out soon.

But what is most important is that the peak should be confirmed in another crust. More measurements should be carried not only with ^{60}Fe , but also with other radionuclides that will be mentioned in Chapter 9. Without proof that these anomalies exist elsewhere we cannot say if our model is the most probable one, although it undoubtedly explains the measured data very well.

8.1 Other isotopes of interest

The search of supernova produced radionuclides involves not only ^{60}Fe . Other promising radionuclides have been measured already and more will be measured in the future.

8.1.1 ^{10}Be

This cosmogenic radionuclide has a half-life of 1.36 Myr [35]. Although it is produced in massive stars (around $10^{-9} M_{\odot}$ for 11 - 20 $M_{\odot}$ [51]) it is mainly created by the spallation of oxygen and nitrogen by cosmic rays in the high Earth's atmosphere [36]. Enhanced cosmic ray flux coming from supernova explosions could provide a higher ^{10}Be fluence on Earth [14]. An increase of ^{10}Be was measured for example in Antarctic ice cores from Dome C and Vostok in 1987 by Raisbeck et al. [37]. The concentration was found to be two times larger than expected. More detailed analysis of the Vostok core gave two peaks at ~ 35 kyr and ~ 60 kyr ago. The peak at ~ 35 kyr was later confirmed at Byrd Station in Antarctica, in ice cores from Greenland (Camp Century) (Beer et al. 1992 [3]), in deep-sea sediments from the Gulf of California (McHargue et al. 1995 [34]) and in the Mediterranean sea (Cini Castagnoli et al. 1995 [9]). The other peak at ~ 60 kyr could not be confirmed again. Ellis et al. [13] argue that this ^{10}Be anomaly could stem from a supernova explosion in the solar vicinity at a distance of 20 - 40 pc and suggest Geminga as a possible source.

8.1.2 ^{26}Al

According to Woosley and Weaver [51] the production of ^{26}Al ($t_{1/2} = 0.72$ Myr [46]) in massive stars is quite high with 10^{-3} - $10^{-2} M_{\odot}$ for stars with masses between 11 and 20 $M_{\odot}$. Like ^{10}Be it is a cosmogenic radionuclide and is produced by spallation processes of cosmic rays with argon, but the ^{26}Al concentration in the troposphere is about 500 times lower than the concentration of ^{10}Be . Detailed measurements of ^{26}Al in a manganese crust or a sediment will be carried out soon.

8.1.3 ^{53}Mn

With a half-life of 3.7 Myr [19] ^{53}Mn is mainly synthesized in massive stars (around $10^{-5} M_{\odot}$ in stars with the mass range of 11-20 $M_{\odot}$ [51]). Another significant extraterrestrial component is created by nuclear reactions of cosmic rays on iron in meteorites and dust particles. The in-situ production of ^{53}Mn in the Earth's lithosphere is very low and occurs due to the spallation of high-energy neutrons, that originate from reactions of cosmic rays with the atmosphere [36]. Several manganese crusts have been measured by Poutivtsev 2007 [36] (the positions are also marked in Appendix A), but they do not show a clear signal. The background is very high and a possible input from SN explosions lies inside the error limit.

8.1.4 ^{244}Pu and other heavy isotopes

Since heavy r-process isotopes are expected to form only in core collapse supernovae, they could be an indicator for explosions in the solar vicinity. Ellis, Fields and Schramm 1996 [13] suggested to search for radionuclides such as ^{129}I , ^{182}Hf , ^{146}Sm and ^{244}Pu in the deep-ocean's crust. The hydro-genetic ferromanganese crust 237KD was analyzed by Wallner et al. 2004 [47] to detect ^{244}Pu ($t_{1/2} = 81.2$ Myr [1]). One event was observed in this crust. However, they had to remove the upper 2mm of the crust to avoid detection of anthropogenic plutonium isotopes, because this could result from atomic bomb tests on Earth.

CHAPTER 9

Epilogue

This last short chapter will consider the consequences for the Earth's climate and biosphere if a supernova occurs at a small distance.

9.1 Consequences of near-Earth supernova explosions

The increase of cosmic ray flux could have an effect on the Earth's climate. A report on the effect is given by Fields and Ellis 1999 [14]. Svensmark et al. [44, 45] examined the correlation of the solar activity with the total cloud cover of the Earth. If the Sun-spot cycle reaches its maximum the solar wind, hence the cosmic ray flux increases, cloud formation might be triggered. An increased cloud cover could result in a cooling of the Earth's temperature [18]; a "cosmic-ray winter" could last for several thousand years.

Such a climate change was observed in Africa around 2.8 Myr ago. The climate became more arid which resulted from a remote forcing by cold sea-surface temperatures from the North Atlantic [12]. It is even suspected that the evolution of mankind was triggered by this climate shift. But the time period corresponds more to the ^{60}Fe anomaly before the peak was shifted to 2.2 Myr due to the newly derived half-life of ^{10}Be .

However, an increase of cosmic rays might have caused some minor extinctions in the past [14], one about 13 Myr ago and another at about 3 Myr ago [42]. Phytoplankton can be reduced due to UV radiation. The impact would propagate to marine animal families which are above in the food

chain. Species of zooplankton, bivalves, gastropods and echinoids were affected [14].

Whether such impact happened also 2.2 Myrs ago is an interesting question. It would give additional evidence to the theory that supernova explosions caused the ^{60}Fe peak in the hydrogenetic ferromanganese crust.

APPENDIX A

Positions of selected crusts and sediments that have been analyzed



Figure A.1: Locations of crusts and sediments that have been analyzed for isotope anomalies. The radionuclide ^{60}Fe was measured in the hydrogenetic ferromanganese crusts from Mona Pihoa [25] and in 237KD and 29DR-32 from Pacific ocean [24], but also in a Marine sediment in the North Atlantic [16]. The crusts 29DR-32, 29DR-45, 4DR in the North Pacific as well as 237KD have been analyzed for ^{53}Mn by [36]. Courtesy of Google Maps.

APPENDIX B

Table of selected stars

Table A1. Identified members of US (ass = 1), UCL (ass = 2) and LCC (ass = 3). Positional and velocity errors are given by the ϵ_i .

HIP- no.	ass	M_V (mag)	$(B - V)_0$ (mag)	X (pc)	Y (pc)	Z (pc)	U (km s ⁻¹)	V (km s ⁻¹)	W (km s ⁻¹)	ϵ_X	ϵ_Y	ϵ_Z	ϵ_U	ϵ_V	ϵ_W
50847 ¹³	3	-0.63	-0.132	42.7	-123.3	-18.6	-11.06	-16.11	-3.62	2.8	8.1	1.2	1.5	3.5	0.6
53701 ¹⁵	3	0.82	-0.098	37.8	-104.6	-2.6	-8.03	-19.46	-5.77	2.5	6.9	0.2	1.6	3.5	0.6
55425 ¹⁵	3	-1.07	-0.157	33.4	-92.1	10.4	-10.80	-14.61	-5.61	1.8	4.8	0.6	1.5	3.5	0.6
57851	3	-0.31	-0.156	46.6	-92.5	-5.5	-4.95	-25.11	-8.09	2.7	5.4	0.3	1.2	1.6	0.5
58326	3	-0.73	-0.157	82.6	-163.8	-0.7	-8.65	-26.16	-8.81	7.9	15.6	0.1	2.4	3.4	1.0
58720	3	1.01	-0.080	45.0	-82.5	-11.1	-11.33	-13.92	-7.20	2.3	4.2	0.6	1.9	3.2	0.6
59173	3	-0.94	-0.192	49.4	-101.5	23.2	-7.84	-23.69	-4.98	4.1	8.4	1.9	1.5	1.8	0.9
59449	3	-1.17	-0.171	46.6	-92.2	18.3	-9.60	-23.79	-9.51	3.6	7.1	1.4	2.1	3.4	1.2
59747	3	-2.45	-0.237	52.7	-98.1	7.4	-5.49	-28.71	-6.86	3.5	6.6	0.5	1.4	1.7	0.6
60009	3	-1.20	-0.187	54.2	-96.5	-2.6	-6.81	-21.74	-8.00	3.3	5.9	0.2	1.0	0.8	0.6
60710	3	-0.68	-0.162	58.2	-105.2	23.9	-12.13	-13.98	-6.34	4.9	8.9	2.0	2.2	3.3	1.0
60823	3	-1.76	-0.198	64.5	-115.9	29.3	-12.72	-18.47	-7.94	5.9	10.6	2.7	2.3	3.3	1.3
61585	3	-2.17	-0.212	48.9	-79.3	-10.3	-8.39	-19.51	-7.89	2.2	3.6	0.5	2.0	3.2	0.5
62058	3	0.47	-0.079	66.4	-107.2	14.8	-9.90	-19.18	-6.21	5.4	8.7	1.2	1.5	1.7	0.7
62327	3	-0.86	-0.183	64.3	-102.0	13.5	-6.65	-24.85	-6.89	4.7	7.4	1.0	1.5	1.7	0.7
62434 ¹⁶	3	-3.92	-0.240	57.9	-91.1	6.0	-12.21	-26.68	-6.05	3.8	6.0	0.4	1.5	1.1	0.6
63003	3	-1.29	-0.180	63.3	-96.2	11.5	-6.03	-21.39	-5.81	4.2	6.5	0.8	1.0	0.9	0.6
63005	3	-0.29	-0.148	60.6	-92.0	11.0	-6.83	-20.64	-4.21	4.1	6.2	0.7	2.2	3.2	0.6
63007	3	-0.63	-0.166	60.5	-92.0	7.2	-7.81	-20.05	-6.70	4.0	6.1	0.5	2.2	3.2	0.6
63945	3	-0.83	-0.147	71.3	-100.1	31.4	-10.48	-16.96	-6.03	6.5	9.1	2.9	2.4	3.1	1.2
64004 ¹³	3	-1.35	-0.227	71.5	-100.2	28.2	-4.21	-21.89	-3.11	7.3	10.2	2.9	2.5	3.1	1.1
64053	3	0.64	-0.093	56.8	-80.0	16.1	-1.27	-29.22	-4.81	3.7	5.2	1.0	2.3	3.1	0.9
64425 ¹⁸	3	-0.56	-0.081	61.6	-86.3	5.3	-5.04	-18.79	-6.99	9.9	13.9	0.9	2.9	3.4	1.4
65112	3	0.05	-0.132	72.6	-95.0	20.8	-8.05	-20.15	-5.56	5.9	7.7	1.7	2.5	3.1	0.9
65271	3	-0.78	-0.173	64.9	-87.1	3.1	-11.63	-16.36	-5.38	4.1	5.5	0.2	2.4	3.1	0.4
66454	3	0.52	-0.112	75.7	-86.6	32.4	-7.46	-17.85	-3.78	6.9	7.9	2.9	2.3	2.6	1.1
67036 ⁰³	3	1.18	0.098	73.0	-83.0	21.5	-6.08	-19.46	-5.23	6.2	7.0	1.8	1.5	1.7	0.8
67464 ¹³	2	-2.40	-0.229	95.8	-97.8	49.5	-7.33	-22.19	-6.12	10.7	11.0	5.6	1.5	2.0	1.1
67472 ⁰⁵	2	-2.58	-0.180	106.5	-109.4	52.9	-6.86	-23.38	-6.42	12.2	12.5	6.1	2.4	2.7	1.4
67669	2	-0.54	-0.149	59.1	-54.6	43.1	-4.63	-20.69	-3.96	4.7	4.4	3.5	1.3	1.6	1.1
67973	2	0.51	-0.090	70.2	-75.9	17.3	-0.22	-24.56	-3.88	4.8	5.2	1.2	3.4	3.7	1.0
68245	2	-1.94	-0.219	96.9	-93.7	46.6	-7.62	-19.12	-6.37	9.3	9.0	4.5	1.5	1.8	1.0
68282	2	-1.68	-0.209	87.3	-86.4	36.3	-9.09	-18.99	-6.43	7.8	7.7	3.2	1.6	1.8	0.9
68862	2	-1.33	-0.199	95.4	-86.7	45.8	-4.18	-22.32	-4.88	9.8	8.9	4.7	2.3	2.5	1.4
69618 ⁰⁶	2	-1.07	-0.138	103.5	-106.7	10.3	-9.63	-19.90	-8.16	9.1	9.4	0.9	5.3	5.5	1.0
70300	2	-1.16	-0.199	94.5	-75.0	44.0	-7.24	-18.06	-5.30	9.2	7.3	4.3	1.5	1.8	0.8

Figure B.1

Table A1 – continued

HIP- no.	ass	M_V (mag)	$(B - V)_0$ (mag)	X (pc)	Y (pc)	Z (pc)	U (km s ⁻¹)	V (km s ⁻¹)	W (km s ⁻¹)	ϵ_X	ϵ_Y	ϵ_Z	ϵ_U	ϵ_V	ϵ_W
70455	2	0.92	-0.089	115.1	-91.9	50.0	-1.80	-24.44	-3.46	15.4	12.3	6.7	4.7	4.2	2.2
70626	2	0.62	-0.087	104.6	-81.3	46.7	-3.76	-19.72	-6.36	11.6	9.0	5.2	2.1	2.4	1.4
71352 ⁰⁷	2	-2.73	-0.270	72.2	-54.8	27.1	-10.77	-17.42	-6.55	5.7	4.3	2.1	1.5	1.7	0.8
71453	2	0.19	-0.117	98.5	-72.0	40.5	-7.53	-16.98	-5.52	9.8	7.1	4.0	3.0	2.6	1.4
71536 ⁰³	2	-0.84	-0.150	72.0	-60.1	16.3	-3.79	-18.67	-5.44	4.9	4.1	1.1	5.6	4.8	1.4
71724 ¹⁵	2	1.16	-0.086	94.9	-68.5	36.9	-5.83	-15.03	-6.25	9.7	7.0	3.8	1.1	1.6	1.0
71727 ¹⁵	2	0.57	-0.124	123.7	-96.5	36.2	-9.74	-20.50	-8.01	16.9	13.2	5.0	5.9	5.1	2.2
71860 ⁰⁹	2	-3.87	-0.222	129.1	-102.3	33.3	-8.77	-22.92	-9.13	16.5	13.1	4.3	1.7	2.5	1.6
71865	2	-0.88	-0.187	72.7	-49.2	32.1	-6.09	-17.31	-5.17	4.9	3.3	2.2	1.2	1.3	0.7
72683 ¹⁰	2	-1.21	-0.167	99.2	-69.7	30.5	-4.44	-21.75	-5.35	9.7	6.8	3.0	1.1	1.7	0.8
72800	2	-0.34	-0.167	94.3	-59.1	38.6	-2.09	-16.26	-1.16	8.4	5.3	3.5	1.6	1.5	0.7
73334	2	-2.99	-0.216	133.9	-87.4	42.1	-2.94	-22.31	-5.53	16.1	10.5	5.1	1.3	2.1	1.0
73807 ¹⁵	2	-2.01	-0.186	123.5	-85.4	26.3	-8.33	-21.29	-3.97	18.1	12.5	3.8	3.5	3.4	1.2
74066 ¹⁵	2	0.24	-0.150	104.9	-62.6	33.2	-8.53	-23.00	-5.73	11.3	6.7	3.6	6.2	4.3	2.1
74100	2	0.18	-0.118	112.5	-70.2	31.1	-7.15	-17.28	-4.37	12.6	7.8	3.5	3.4	2.7	1.3
74479 ¹⁵	2	0.83	-0.091	93.5	-48.6	35.3	-8.96	-14.13	-4.21	8.6	4.5	3.2	3.2	2.2	1.3
74950 ¹¹	2	-0.40	-0.103	133.7	-74.6	38.0	0.26	-24.24	-2.81	15.2	8.5	4.3	6.3	4.1	2.0
75141 ¹³	2	-2.79	-0.238	133.3	-72.9	37.4	-11.02	-18.73	-7.28	17.9	9.8	5.0	2.7	3.0	1.2
75151 ¹⁵	2	1.11	-0.113	105.6	-54.4	33.6	-8.85	-13.54	-5.13	14.6	7.5	4.7	5.5	3.4	1.9
75264	2	-2.61	-0.192	130.6	-77.8	27.7	-3.20	-21.98	-2.56	12.3	7.3	2.6	3.3	2.6	0.9
75304	2	-1.83	-0.158	159.8	-78.5	53.6	-8.17	-22.61	-5.98	22.9	11.2	7.7	2.7	3.5	1.3
75647	2	-0.08	-0.148	111.3	-52.8	36.1	-7.74	-17.37	-4.39	12.0	5.7	3.9	8.7	4.6	2.9
76297	2	-3.41	-0.217	151.9	-76.7	35.8	-8.26	-21.78	-8.70	32.8	16.5	7.7	8.9	6.4	2.9
76371	2	-1.07	-0.184	115.1	-63.8	20.3	-2.86	-19.47	-2.27	11.2	6.2	2.0	1.7	1.8	0.6
76395	2	0.95	-0.103	99.4	-39.3	35.1	-8.24	-14.60	-4.54	9.3	3.7	3.3	3.3	2.0	1.3
76945	2	-0.67	-0.145	108.5	-42.8	33.6	-3.21	-22.09	-3.65	11.1	4.4	3.4	1.7	2.2	1.0
77286	2	0.23	-0.118	107.2	-40.8	32.0	-0.87	-19.16	-2.89	11.6	4.4	3.5	6.7	3.2	2.1
77635	1	-1.77	-0.170	144.3	-35.7	59.2	-5.12	-19.74	-7.55	21.0	5.2	8.6	4.3	3.0	2.0
77840 ⁰¹	1	-1.31	-0.160	120.4	-28.1	49.0	-10.45	-15.61	-8.52	18.9	4.4	7.7	1.5	2.6	1.1
77900	1	-0.01	-0.096	148.2	-38.1	55.7	-3.92	-21.17	-7.88	19.5	5.0	7.3	2.4	3.0	1.3
77909	1	0.04	-0.097	126.7	-29.0	51.2	-9.40	-15.78	-9.41	21.2	4.8	8.6	3.4	2.8	1.6
78207 ⁰²	1	-1.03	-0.089	137.7	-8.7	75.3	-6.74	-15.01	-5.07	17.1	1.1	9.4	1.6	2.0	1.0
78246 ¹³	1	-0.60	-0.140	136.6	-28.7	54.1	-12.34	-15.49	-10.56	16.8	3.5	6.6	3.1	2.2	1.5
78265 ¹³	1	-2.85	-0.249	128.9	-29.2	48.7	-12.31	-15.32	-10.47	15.2	3.5	5.8	6.8	2.6	2.7
78384	2	-2.48	-0.226	138.4	-53.8	28.9	-5.80	-21.12	-6.28	16.3	6.3	3.4	3.6	2.8	1.2
78655	2	-1.14	-0.141	148.4	-56.6	29.2	-9.36	-22.89	-7.30	18.2	6.9	3.6	2.8	3.0	1.1
78756 ¹⁵	2	0.83	-0.056	153.9	-59.9	28.0	-6.36	-22.90	-4.40	24.5	9.5	4.5	6.9	4.5	1.7
78877 ⁰⁴	1	-0.02	-0.096	137.5	-23.4	53.1	-6.35	-19.40	-10.64	17.6	3.0	6.8	3.4	2.7	1.7
79044	2	1.15	-0.077	120.5	-40.4	25.9	-2.31	-20.04	-5.47	12.5	4.2	2.7	6.9	3.0	1.7
79404	1	-1.32	-0.197	134.4	-28.3	41.6	-5.18	-15.43	-6.86	15.2	3.2	4.7	6.9	2.4	2.3
81914	2	0.15	-0.119	141.9	-44.0	8.1	-6.73	-14.66	-3.91	18.0	5.6	1.0	1.8	2.0	0.6
82545	2	-2.47	-0.223	153.6	-37.7	10.7	-3.26	-19.78	-3.70	20.9	5.1	1.5	0.9	2.7	0.8
84970	1	-2.95	-0.223	171.6	1.4	19.7	-1.37	-20.10	-5.05	20.5	0.2	2.3	3.6	2.3	1.0

¹multiple, ²emml./variable, ³variable, ⁴rotnl. variable, ⁵Be, ⁶emml./binary, ⁷Be/neb.emm., ⁸variable/ β Cep, ⁹binary, ¹⁰ecl.binary, ¹¹ecl.binary, ¹²variable/ β Cep, ¹³spec.binary, ¹⁴double, ¹⁵variable/ β Cep/double?, ¹⁶ellips.variable/double?

Figure B.2: Hipparcos numbers of all 79 stars as listed in [17] with positions, velocities and errors

APPENDIX C

The Sedov Solution

The equations

$$-3V = \frac{dV}{d \ln \xi} - (1 - V) \frac{d \ln G}{d \ln \xi}, \quad (\text{C.1a})$$

$$-\frac{5 - 2V}{1 - V} = \frac{d \ln Z}{d \ln \xi} - (\gamma - 1) \frac{d \ln G}{d \ln \xi}, \quad (\text{C.1b})$$

$$Z = \frac{\gamma(\gamma - 1)(1 - V)V^2}{2(\gamma V - 1)} \quad (\text{C.1c})$$

from our system of ordinary differential equations we want to solve. To do so, we have to transform the equations in such a way that the resulting derivatives $d \ln \xi / dV$ and $d \ln G / dV$ only depend on V . We start by eliminating $d \ln G / d \ln \xi$ from equations (C.1a) and (C.1b). We obtain

$$\frac{d \ln G}{d \ln \xi} = \left(3V + \frac{dV}{d \ln \xi} \right) \left(\frac{1}{1 - V} \right), \quad (\text{C.2})$$

$$\frac{d \ln G}{d \ln \xi} = \left(\frac{d \ln Z}{d \ln \xi} + \frac{5 - 2V}{1 - V} \right) \frac{1}{\gamma - 1}, \quad (\text{C.3})$$

and by equating (C.2) and (C.3) we get

$$\frac{dV}{d \ln \xi} - \frac{1 - V}{\gamma - 1} \frac{d \ln Z}{d \ln \xi} - \frac{5 - 3(\gamma - 1)V}{\gamma - 1} = 0. \quad (\text{C.4})$$

Multiplying by $d \ln \xi / dV$ results in

$$\frac{d \ln \xi}{dV} = \left(1 - \frac{1 - V}{\gamma - 1} \frac{d \ln Z}{dV} \right) \left(\frac{\gamma - 1}{5 - 3(\gamma - 1)V} \right). \quad (\text{C.5})$$

To integrate we have to achieve that $d \ln \xi / dV$ only depends on V . This is done with the help of the first integral (C.1c). To replace $d \ln Z / dV$ we must take the logarithm of Z

$$\ln Z = \ln \left(\frac{\gamma(\gamma - 1)}{2} \right) - \ln(1 - V) + 2 \ln V - \ln(\gamma V - 1), \quad (\text{C.6})$$

and then differentiate with respect to V

$$\frac{d \ln Z}{dV} = -\frac{1}{1 - V} + \frac{2}{V} - \frac{\gamma}{\gamma V - 1}. \quad (\text{C.7})$$

This can be inserted into (C.5) and gives

$$\begin{aligned} \frac{d \ln \xi}{dV} &= \frac{\gamma}{5 - (3\gamma - 1)V} - \frac{2(1 - V)}{V(5 - (3\gamma - 1)V)} \\ &\quad + \frac{\gamma(1 - V)}{(\gamma V - 1)(5 - (3\gamma - 1)V)}. \end{aligned} \quad (\text{C.8})$$

The first term can easily be integrated and becomes

$$\int \frac{\gamma}{5 - (3\gamma - 1)V} dV = -\frac{\gamma}{3\gamma - 1} \ln(5 - (3\gamma - 1)V) + \text{const}_1. \quad (\text{C.9})$$

For the second and third term we have to evaluate partial fractions and the integrals become

$$\begin{aligned} \int -\frac{2(1 - V)}{V(5 - (3\gamma - 1)V)} dV &= \int \left(-\frac{2}{5V} - \frac{6}{5} \frac{\gamma - 2}{5 - (3\gamma - 1)V} \right) dV \\ &= -\frac{2}{5} \ln V + \frac{6}{5} \frac{\gamma - 2}{3\gamma - 1} \ln(5 - (3\gamma - 1)V) + \text{const}_2, \end{aligned} \quad (\text{C.10})$$

$$\begin{aligned} &\int \frac{\gamma(1 - V)}{(\gamma V - 1)(5 - (3\gamma - 1)V)} dV \\ &= \int \left(-\frac{\gamma - \gamma^2}{(2\gamma + 1)(\gamma V - 1)} + \frac{3(-2\gamma + \gamma^2)}{(2\gamma + 1)(5 - (3\gamma - 1)V)} \right) dV \\ &= \frac{\gamma - 1}{2\gamma + 1} \ln(\gamma V - 1) - \frac{3(-2\gamma + \gamma^2)}{(2\gamma + 1)(3\gamma - 1)} \ln(5 - (3\gamma - 1)V) \\ &\quad + \text{const}_3. \end{aligned} \quad (\text{C.11})$$

Putting all the terms together results in

$$\ln \xi = -\frac{2}{5} \ln V + \frac{\gamma - 1}{2\gamma + 1} \ln(\gamma V - 1) - \frac{13\gamma^2 - 7\gamma + 12}{5(2\gamma + 1)(3\gamma - 1)} + c_1. \quad (\text{C.12})$$

We can compute the constant by inserting the boundary condition

$$V(1) = \frac{2}{\gamma + 1}, \quad (\text{C.13})$$

and get the final solution for ξ

$$\xi^5 = \left(\frac{\gamma + 1}{2} V \right)^{-2} \left(\frac{\gamma + 1}{7 - \gamma} (5 - (3\gamma - 1)V) \right)^{v_1} \left(\frac{\gamma + 1}{\gamma - 1} (\gamma V - 1) \right)^{v_2}, \quad (\text{C.14})$$

with

$$v_1 = -\frac{13\gamma^2 - 7\gamma + 12}{(3\gamma - 1)(2\gamma + 1)}, \quad (\text{C.15})$$

$$v_2 = \frac{5(\gamma - 1)}{2\gamma + 1}. \quad (\text{C.16})$$

To compute G we take the first Euler-equation (C.1a) and express the fraction $d \ln G / d \ln \xi$ as

$$\frac{d \ln G}{d \ln \xi} = \frac{1}{1 - V} \frac{dV}{d \ln \xi} + \frac{3V}{1 - V}. \quad (\text{C.17})$$

Then we multiply with $d \ln \xi / dV$ and the equation becomes

$$\frac{d \ln G}{dV} = \frac{1}{1 - V} + \frac{3V}{1 - V} \frac{d \ln \xi}{dV}. \quad (\text{C.18})$$

But the expression for $d \ln \xi / dV$ we know already from (C.8). Inserting this gives

$$\begin{aligned} \frac{d \ln G}{dV} &= \frac{1}{1 - V} + \frac{3\gamma V}{(1 - V)(5 - (3\gamma - 1)V)} - \frac{6}{5 - (3\gamma - 1)V} \\ &\quad + \frac{3\gamma V}{(\gamma V - 1)(5 - (3\gamma - 1)V)}. \end{aligned} \quad (\text{C.19})$$

We integrate this the same way we did before and get the solution

$$\begin{aligned} \ln G &= \frac{3}{2\gamma + 1} \ln(\gamma V - 1) + \frac{13\gamma^2 - 7\gamma + 12}{(2\gamma + 1)(3\gamma - 1)(2 - \gamma)} \ln(5 - (3\gamma - 1)V) \\ &\quad + \frac{-2}{2 - \gamma} \ln(1 - V) + c_2. \end{aligned} \quad (\text{C.20})$$

Here again we calculate the constant using the boundary conditions (C.13) and

$$G(1) = \frac{\gamma + 1}{\gamma - 1}. \quad (\text{C.21})$$

Thus the final solution for G reads

$$G = \frac{\gamma + 1}{\gamma - 1} \left(\frac{\gamma + 1}{\gamma - 1} (\gamma V - 1) \right)^{v_3} \left(\frac{\gamma + 1}{7 - \gamma} (5 - (3\gamma - 1)V) \right)^{v_4} \left(\frac{\gamma + 1}{\gamma - 1} (1 - V) \right)^{v_5}, \quad (\text{C.22})$$

with

$$v_3 = \frac{3}{2\gamma + 1}, \quad (\text{C.23})$$

$$v_4 = \frac{v_1}{2 - \gamma'}, \quad (\text{C.24})$$

$$v_5 = -\frac{2}{2 - \gamma}. \quad (\text{C.25})$$

APPENDIX D

Solving the differential equation for the expansion of a supernova remnant in Phase III with non-negligible inside pressure

In the following we will derive a solution of Kahn's (1998) [22] differential equation

$$S^{3/4} \frac{d^2 S}{d\theta^2} = 0.278 \theta^{-4/9}. \quad (\text{D.1})$$

In addition we will present a more precise solution.

To make the equation look more familiar we substitute θ with x and S with y

$$y'' = \alpha x^{-4/9} y^{-3/4},$$

where α is a positive constant. We use the ansatz

$$y = c_0 x^{a_0} + \text{l.o.t.},$$

where l.o.t. stands for lower order terms. We do not know the highest power a_0 , but we can compute it by inserting y into the differential equation and comparing the power. To do this we have to differentiate y twice and assume that $a_0 \neq 0$ and $a_0 \neq 1$, because otherwise the first term of the derivatives would become zero. In this step we leave out the lower order

terms, because we only need to compare the first term

$$\begin{aligned}y' &= a_0 c_0 x^{a_0-1}, \\y'' &= a_0(a_0-1)c_0 x^{a_0-2}, \\y^{-3/4} &= c_0^{-3/4} x^{(-3/4)a_0}.\end{aligned}$$

Inserting this into the equations gives

$$\begin{aligned}a_0(a_0-1)c_0 x^{a_0-2} &= \alpha x^{-4/9} c_0^{-3/4} x^{(-3/4)a_0} \\&= \alpha c_0^{-3/4} x^{-4/9-(3/4)a_0}.\end{aligned}$$

Now we can compare the power of x to compute a_0

$$\begin{aligned}a_0 - 2 &= -\frac{4}{9} - \frac{3}{4}a_0 \\ \Rightarrow a_0 &= \frac{8}{9}.\end{aligned}$$

Now we have to check if this is the right solution by doing the test

$$\begin{aligned}y &= c_0 x^{8/9}, \\y'' &= c_0 \frac{8}{9} \left(-\frac{1}{9}\right) x^{-10/9} = -\frac{8}{81} c_0 x^{-10/9}, \\y^{3/4} &= c_0^{3/4} x^{2/3},\end{aligned}$$

$$\begin{aligned}c_0^{3/4} x^{2/3} \left(-\frac{8}{81}\right) c_0 x^{-10/9} &= \alpha x^{-4/9} \\ \underbrace{c_0^{3/4} c_0}_{<0} \left(-\frac{8}{81}\right) x^{-4/9} &= \alpha x^{-4/9}.\end{aligned}$$

The left hand side is negative, since c_0 must be positive, but the right hand side is positive. This leads to a contradiction, so the power of $8/9$ can not be right. Because we assumed before that a_0 should not be 0 or 1 it has to be one of these numbers.

Which one it actually is we can find out by comparing the next power of the second derivative, because the first term will be zero. This power will be $a_0 - a_1 - 2$ with $a_1 < a_0$. If $a_0 = 0$ then $x^{-a_1-2} \sim x^{-4/9}$, but this is not of the same order. That means that $a_0 = 1$, which is clear, because the homogeneous solution of the differential equation should give us a linear term.

If we insert the ansatz $y = c_0x + \text{l.o.t.}$ into the equation and integrate twice we get the solution (which checks with Kahn (1998) [22] for $\alpha = 0.32$)

$$\begin{aligned}
y'' &= 0.32x^{-4/9}c_0^{-3/4}x^{-3/4} + \text{l.o.t.} = 0.32c_0^{-3/4}x^{-43/36} + \text{l.o.t.}, \\
y' &= 0.32c_0^{-3/4}\left(-\frac{36}{7}\right)x^{-7/36} + c_1 + \dots, \\
y &= 0.32c_0^{-3/4}\left(-\frac{36}{7}\right)\frac{36}{29}x^{29/36} + c_1x + c_2 + \dots \\
&= c_0x - 0.32c_0^{-3/4}\left(-\frac{36^2}{7 \cdot 29}\right)x^{29/36} + c_2 + \dots \\
&= c_0x(1 - 0.32c_0^{-7/4}\left(-\frac{36^2}{7 \cdot 29}\right)x^{-7/36} + \frac{c_2}{c_0}x^{-1} + \dots).
\end{aligned}$$

After the second integration the term c_1x appeared, which could be replaced by the term c_0x from the ansatz. With the initial conditions $y(1) = 1$ and $y'(1) = 8/5$ we are now able to compute the coefficients c_0 and c_2 which have the following values

$$\begin{aligned}
c_0 &= 2.44, \\
c_2 &= -0.39.
\end{aligned}$$

Then we obtain exactly the same solution Kahn got, which is

$$S = 2.44\theta(1 - 0.43\theta^{-7/36} - 0.16\theta^{-1} \dots),$$

and because $R = S^{1/4}$, the dimensionless radius becomes

$$R = 1.25\theta^{1/4}(1 - 0.11\theta^{-7/36} - 0.041\theta^{-1}).$$

But Kahn (1998) left out many terms that are larger than the term with x^{-1} which comes from the integration, because he lost the lower order terms during the integration. These terms we will now derive as follows. The ansatz is

$$y = c_0x(1 + c_1x^{-a_1}).$$

Taking the second derivative and computing $y^{3/4}$ by using a Taylor expansion $(1 + \epsilon)^{-3/4} = 1 - \frac{3}{4}\epsilon + \frac{21}{32}\epsilon^2 - \frac{77}{128}\epsilon^3 + \dots$ around $\epsilon = 0$ the exponents of the underlined terms can be compared

$$\begin{aligned}
y'' &= c_0c_1a_1(a_1 - 1)\underline{x^{-a_1-1}}, \\
y^{-3/4} &= c_0^{-3/4}x^{-3/4}(1 - \frac{3}{4}c_1x^{-a_1}), \\
x^{-4/9}y^{-3/4} &= c_0^{-3/4}\underline{x^{-4/9-3/4}}(1 - \frac{3}{4}c_1x^{-a_1}).
\end{aligned}$$

For the exponent a_1 we obtain the value

$$\begin{aligned} -a_1 - 1 &= -\frac{4}{9} - \frac{3}{4} \\ \Rightarrow a_1 &= 7/36, \end{aligned}$$

which agrees with Kahn's solution. Now we will develop the ansatz a bit further to evaluate the next exponent

$$\begin{aligned} y &= c_0 x(1 + c_1 x^{-7/36} + c_2 x^{-a_2}), \\ y'' &= \frac{29}{36}(-\frac{7}{36})c_0 c_1 x^{-43/36} + c_0 c_2 a_2(a_2 - 1)x^{-a_2-1}, \\ x^{-4/9} y^{-3/4} &= c_0^{-3/4} x^{-43/36} (1 - \frac{3}{4} c_1 x^{-7/36} - \frac{3}{4} c_2 x^{-a_2} + \frac{21}{32} c_1^2 x^{-14/36}), \end{aligned}$$

$$\begin{aligned} -a_2 - 1 &= -\frac{43}{36} - \frac{7}{36} \\ \Rightarrow a_2 &= \frac{14}{36}. \end{aligned}$$

This exponent $-14/36$ is bigger than -1 . Computing the other exponents the solution looks like

$$\begin{aligned} y &= c_0 x(1 + c_1 x^{-7/36} + c_2 x^{-14/36} + c_3 x^{-21/36} + c_4 x^{-28/36} \\ &\quad + c_5 x^{-35/36} + d_0 x^{-1} + c_6 x^{-42/36} + d_1 x^{-43/36} + \dots). \end{aligned}$$

As we can see, four additional terms appear until the part from the integration occurs. The coefficients of the terms that come from the integration are marked with d 's. This ansatz we can use to evaluate the coefficients by taking again the second derivative, expanding $y^{3/4}$ into a power series and inserting this into our differential equation. In Mathematica it is easier to substitute $x = z^{36}$, divide by $c_0 x$ and develop the Taylor-series around infinity, because we are dealing with large x . Then, by comparing the coefficients we get

$$\begin{aligned} c_1 &= -1.773c_0^{-7/4}, \quad c_2 = -1.554c_0^{-7/2}, \quad c_3 = -3.687c_0^{-21/4}, \\ c_4 &= -15.635c_0^{-7}, \quad c_5 = -372.949c_0^{-35/4}, \\ c_6 &= -544.316c_0^{-21/2}, \quad d_1 = -0.897d_0c_0^{-7/4}, \dots \end{aligned}$$

As we can see all the c -coefficients depend on c_0 , whereas the first coefficient that depends on d_0 is d_1 . Developing the ansatz further one will notice that

all the d 's depend on d_0 and c_0 . To compute these two numbers we use the initial conditions and the factor 0.278 (instead of 0.32) derived earlier and get

$$\begin{aligned}c_0 &= 2.50082, \\d_0 &= -0.06094.\end{aligned}$$

Now we can use c_0 and d_0 to compute the other coefficients and eventually get the solution

$$S = 2.50\theta(1 - 0.356\theta^{-7/36} - 0.063x^{-7/18} - 0.030x^{-7/12} - \dots). \quad (\text{D.2})$$

This solution differs from Kahn's (1998), already in the leading term. Now the dimensionless radius R becomes

$$R = 1.258\theta^{1/4}(1 - 0.089\theta^{-7/36} - 0.016\theta^{-7/18} - 0.007\theta^{-7/12} - \dots). \quad (\text{D.3})$$

APPENDIX E

Comparison of the trajectories for different densities

To compare the influence of the ISM density on the time and the amount of ^{60}Fe arriving on Earth, we varied the densities with the trajectories we found in Chapter 6. For example the trajectories used for a density with 1 atom/cm^3 are now plotted with densities of 10, 0.1 and 0.01 atoms/cm^3 . We find, that the results still fit very well to the measurements, which means that there is a great variety of locations for the explosions in the stellar moving group.

This is important, since it increases the probability that one of the trajectories corresponds to the stars which exploded inside the Local Bubble.

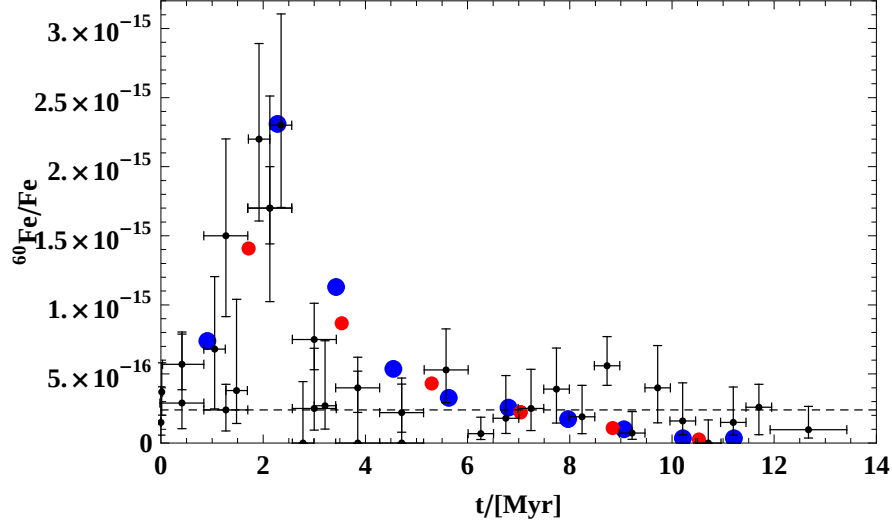


Figure E.1: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 0.01 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 10 atoms/cm^3 .

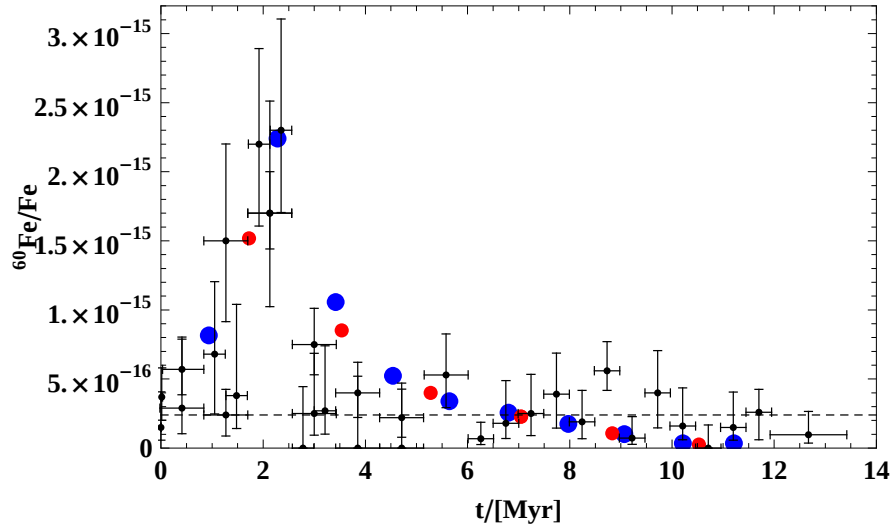


Figure E.2: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 0.01 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 1 atom/cm^3 .

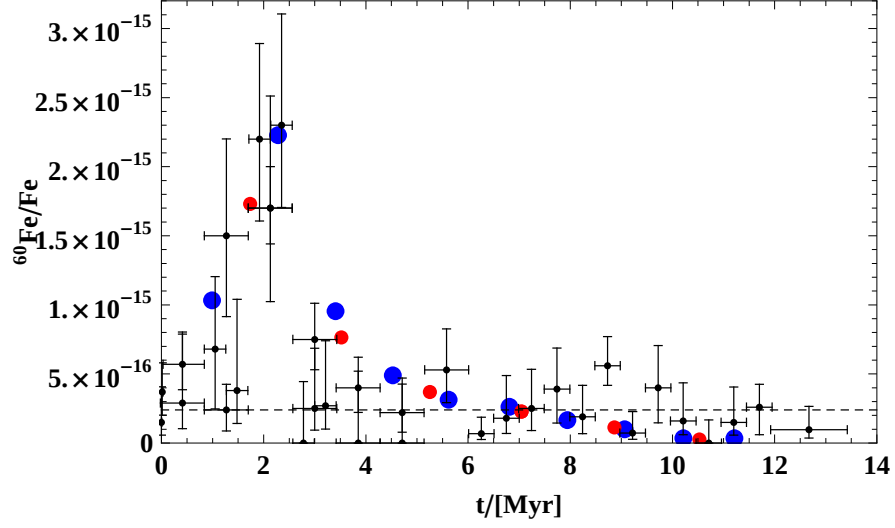


Figure E.3: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 0.01 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 0.1 atoms/cm^3 .

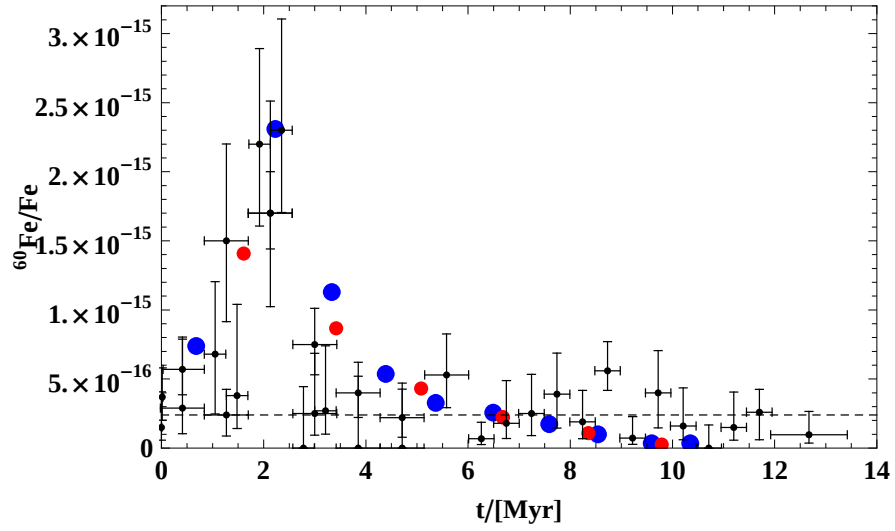


Figure E.4: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 0.1 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 10 atoms/cm^3 .

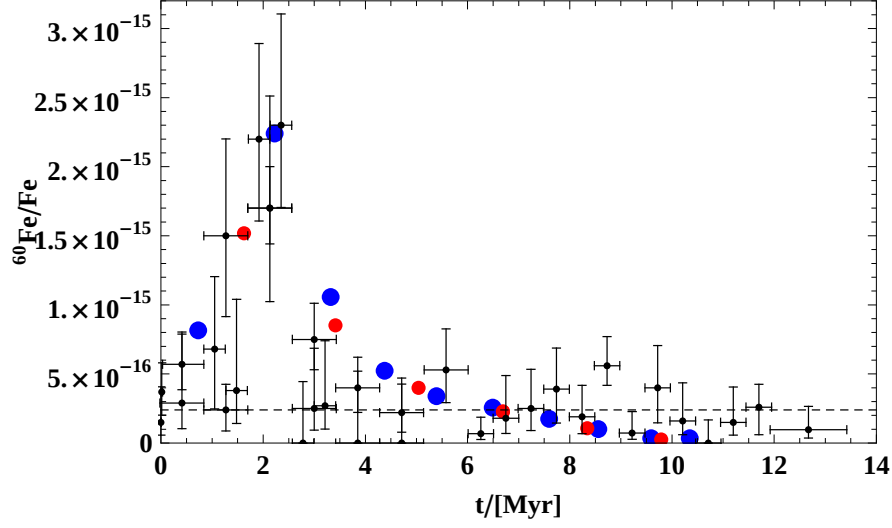


Figure E.5: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 0.1 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 1 atom/cm^3 .

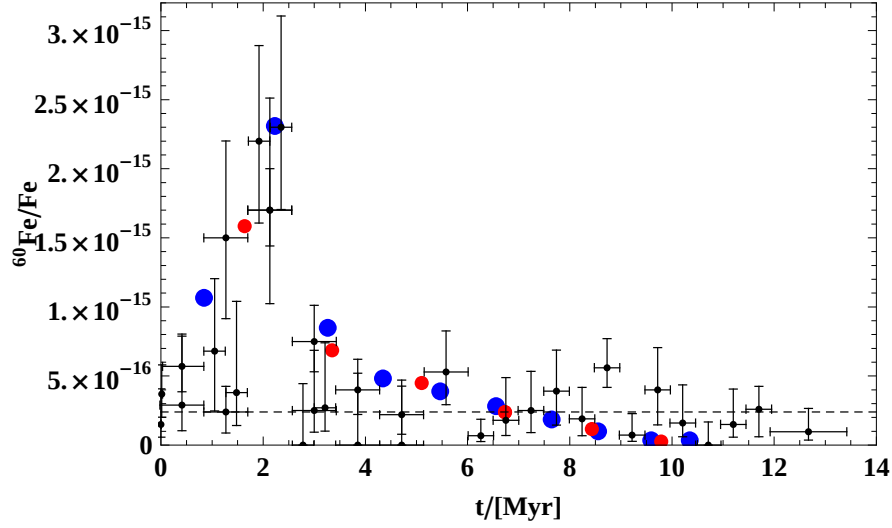


Figure E.6: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 0.1 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 0.01 atoms/cm^3 .

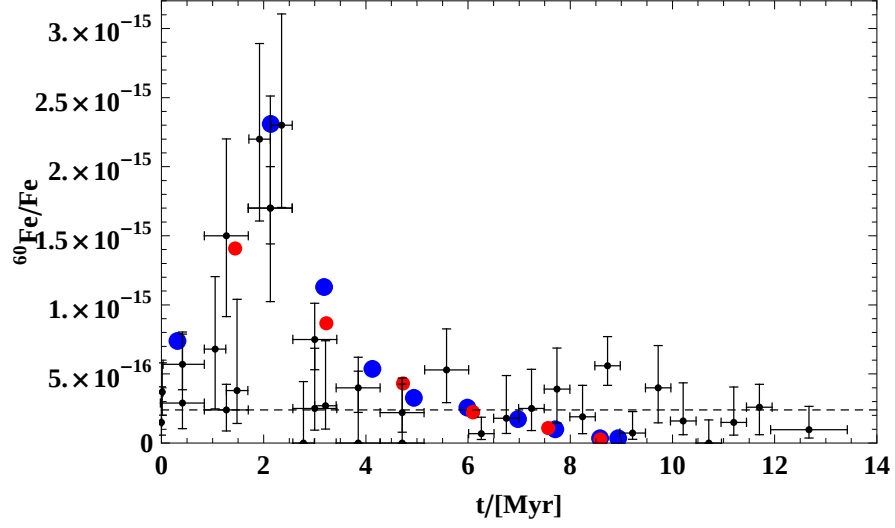


Figure E.7: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 1 \text{ atom/cm}^3$. Here we plot the trajectories we found for a density of 10 atoms/cm 3 .

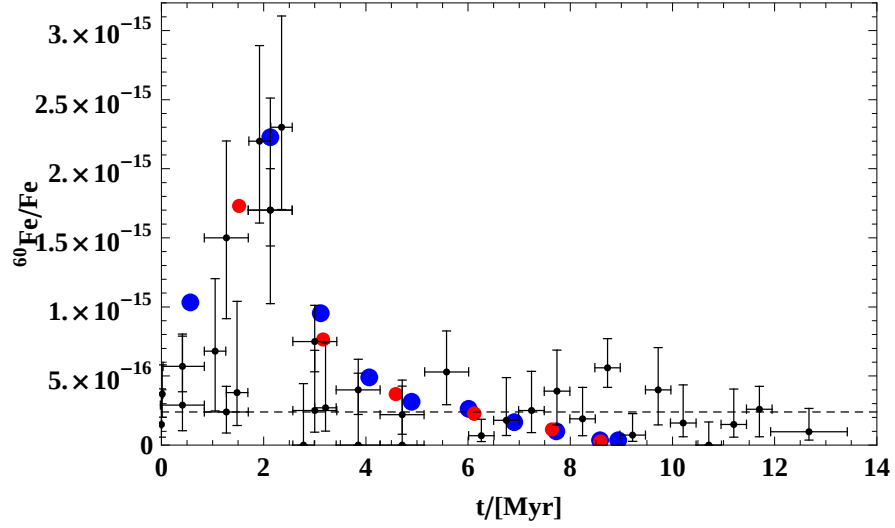


Figure E.8: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 1 \text{ atom/cm}^3$. Here we plot the trajectories we found for a density of 0.1 atoms/cm 3 .

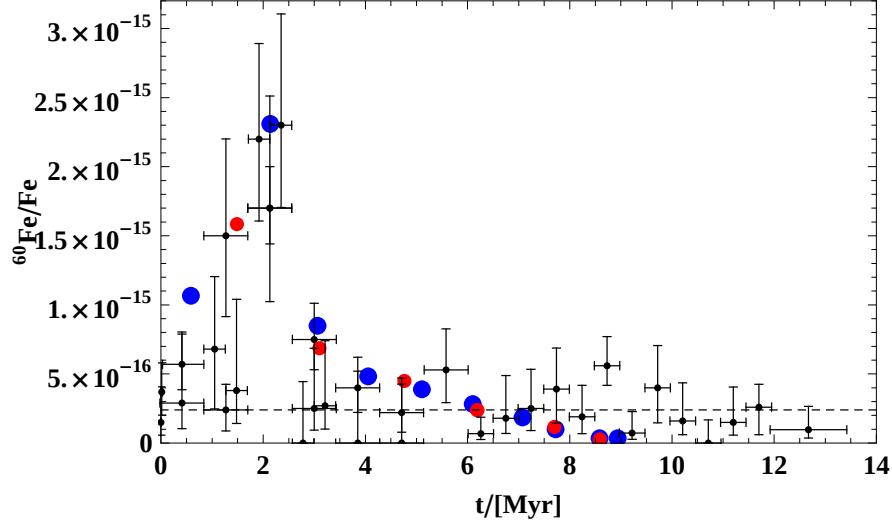


Figure E.9: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 1 \text{ atom/cm}^3$. Here we plot the trajectories we found for a density of 0.01 atoms/cm^3 .

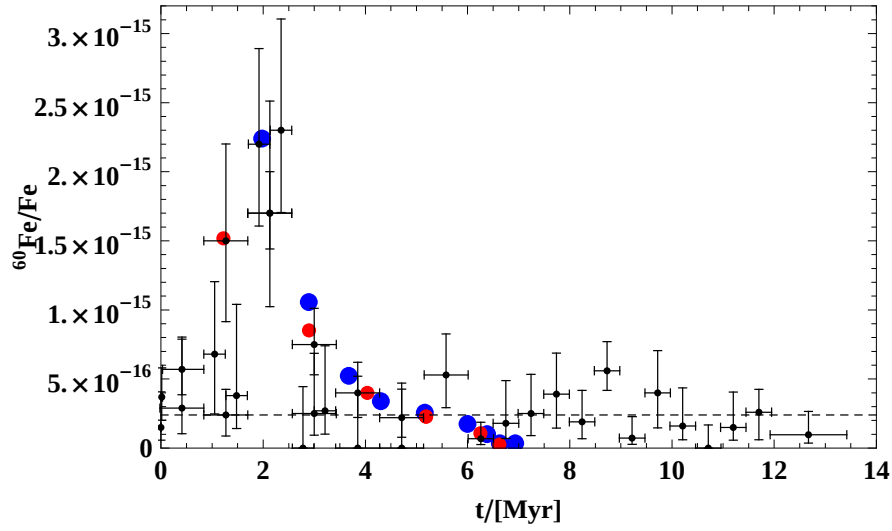


Figure E.10: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 10 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 1 atom/cm^3 .

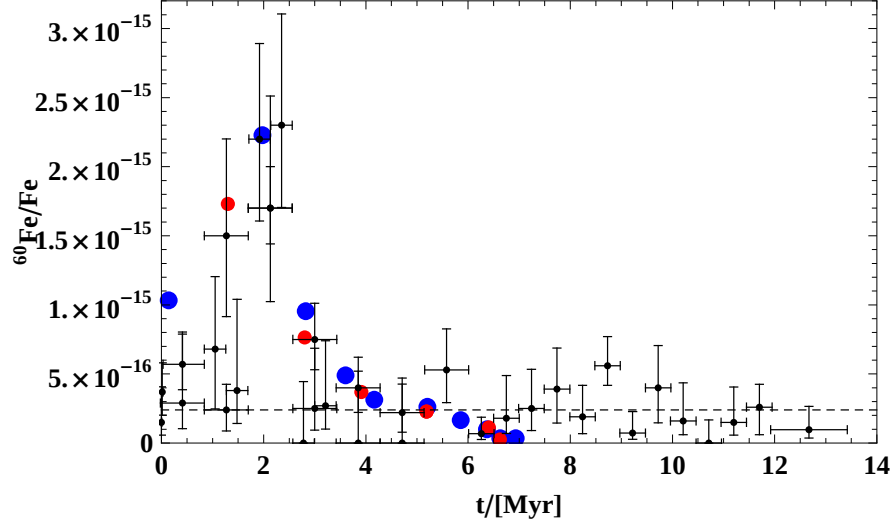


Figure E.11: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 10 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 0.1 atoms/cm^3 .

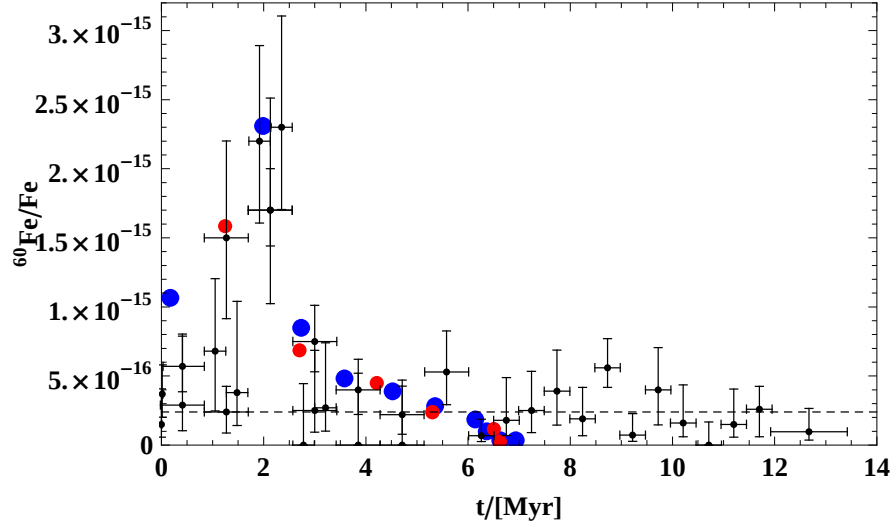


Figure E.12: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the boundary mass model with an ISM density of $n = 10 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 0.01 atoms/cm^3 .

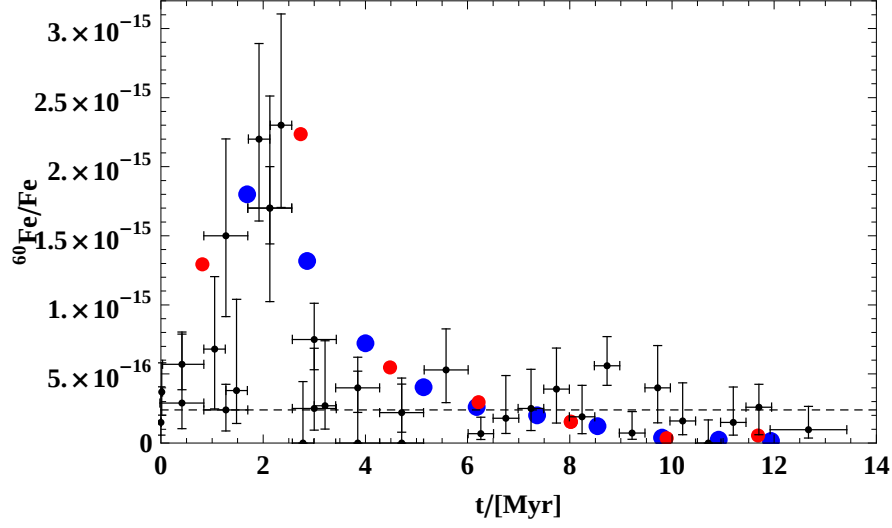


Figure E.13: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 0.01 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 10 atoms/cm^3 .

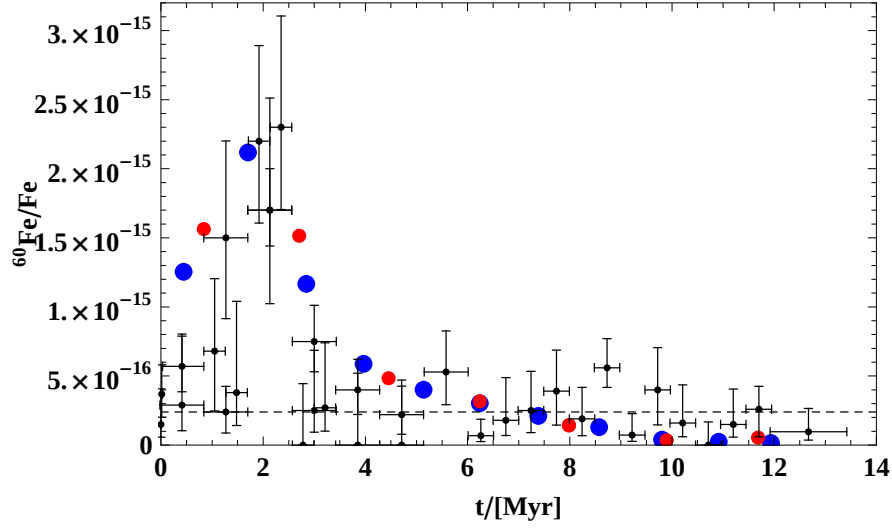


Figure E.14: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 0.01 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 1 atom/cm^3 .

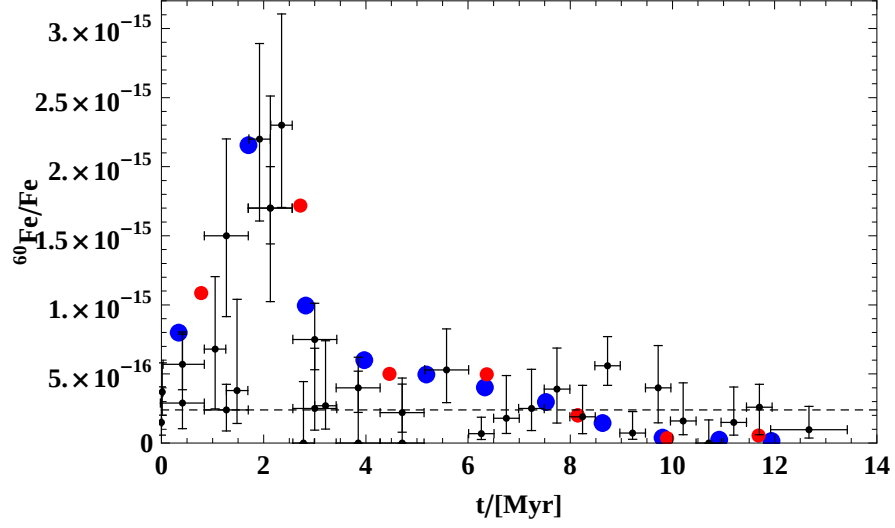


Figure E.15: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 0.01 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 0.1 atoms/cm^3 .

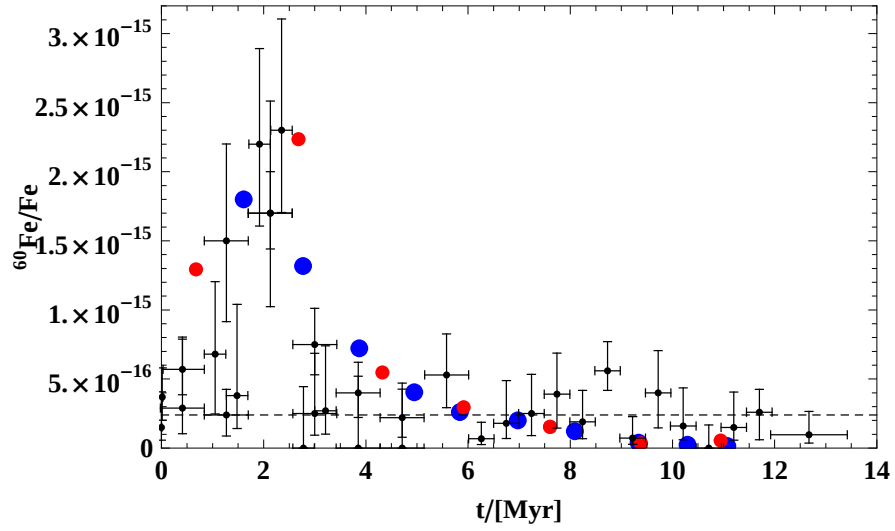


Figure E.16: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 0.1 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 10 atoms/cm^3 .

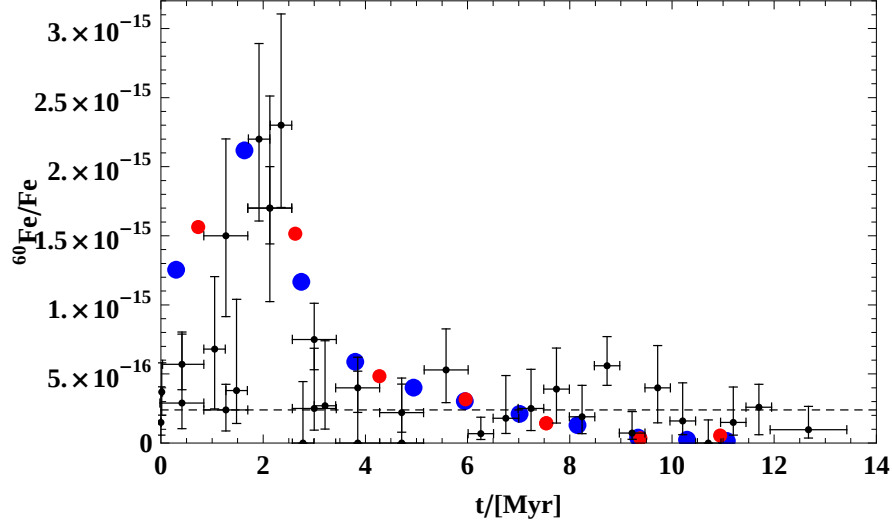


Figure E.17: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 0.1 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 1 atom/cm^3 .

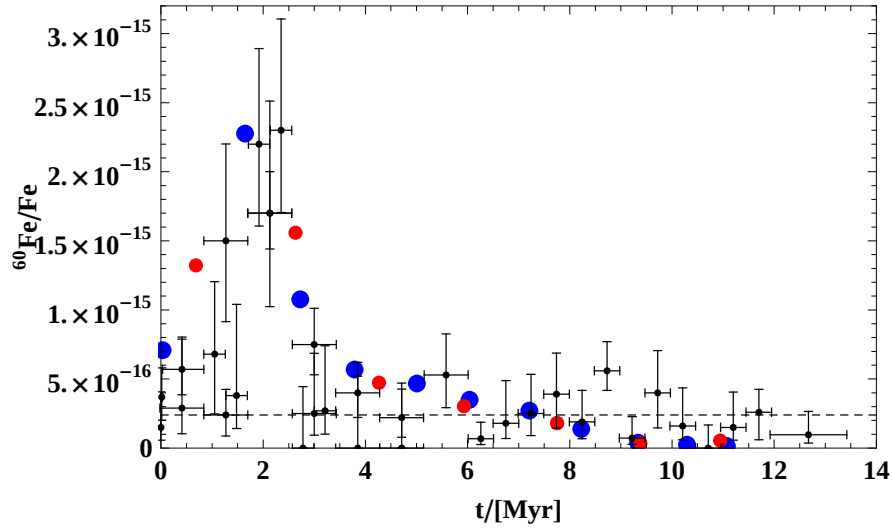


Figure E.18: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 0.1 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 0.01 atoms/cm^3 .

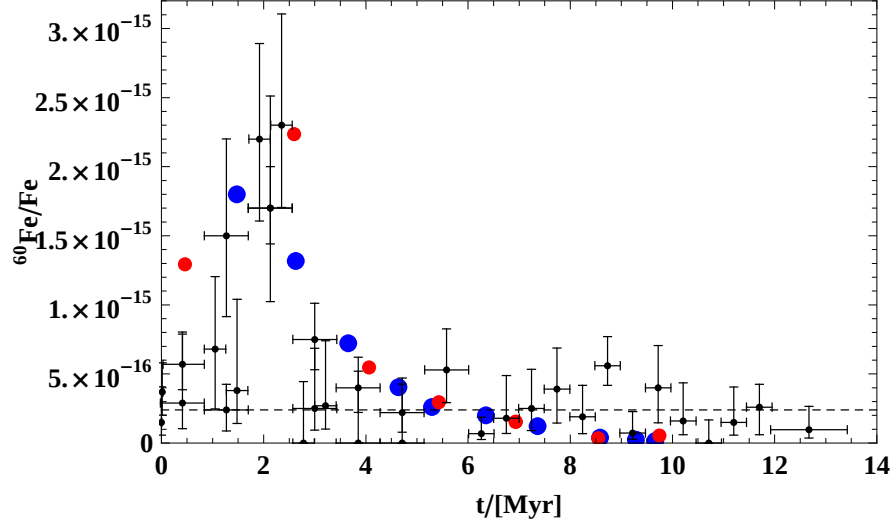


Figure E.19: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 1 \text{ atom/cm}^3$. Here we plotted the trajectories we found for a density of 10 atoms/cm^3 .

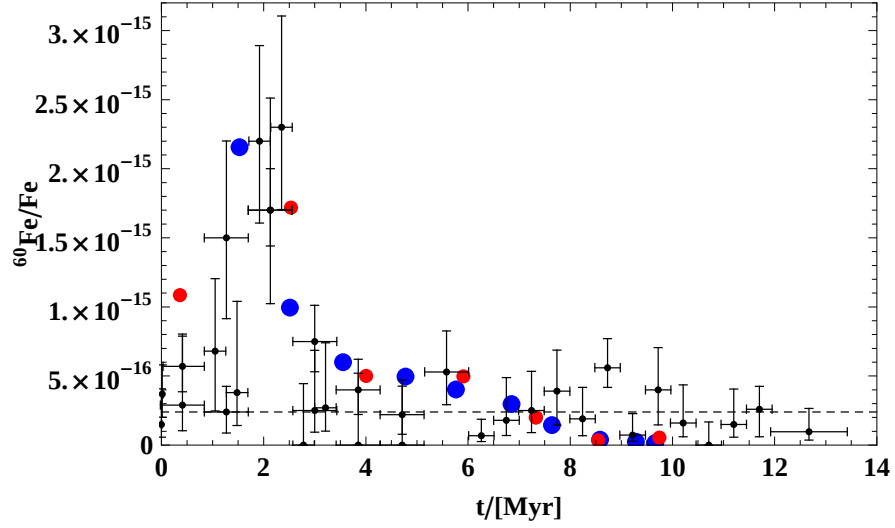


Figure E.20: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 1 \text{ atom/cm}^3$. Here we plotted the trajectories we found for a density of 0.1 atoms/cm^3 .

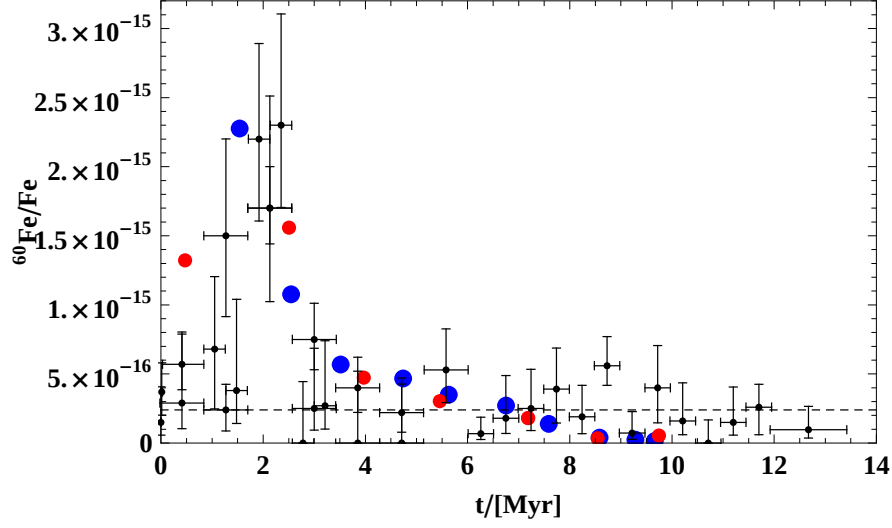


Figure E.21: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 1 \text{ atom/cm}^3$. Here we plot the trajectories we found for a density of 0.01 atoms/cm^3 .

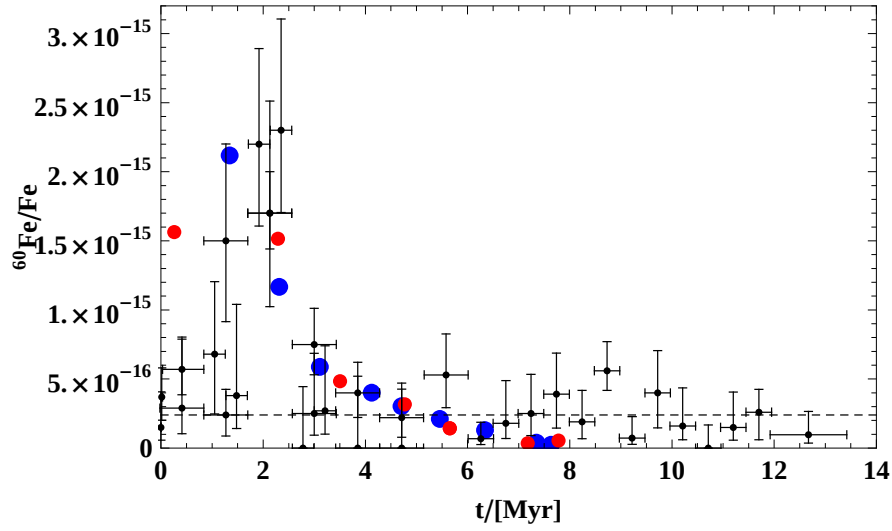


Figure E.22: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 10 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 1 atom/cm^3 .

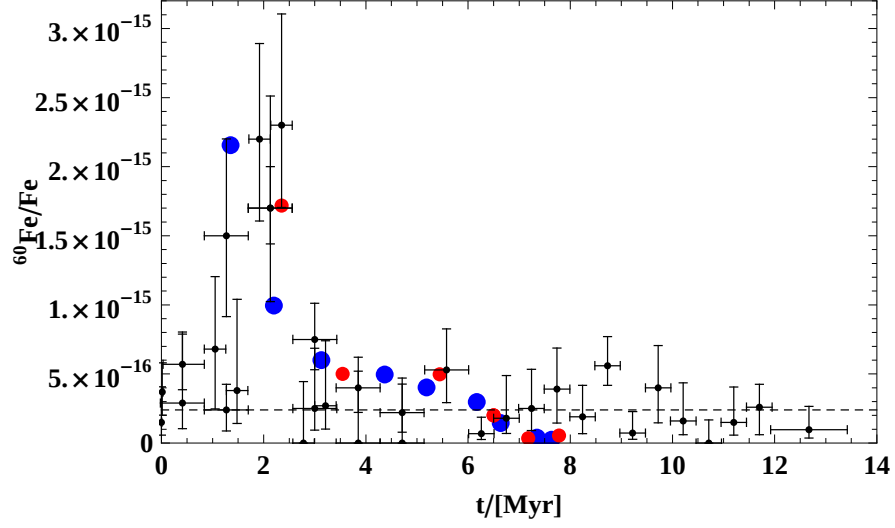


Figure E.23: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 10 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 0.1 atoms/cm^3 .

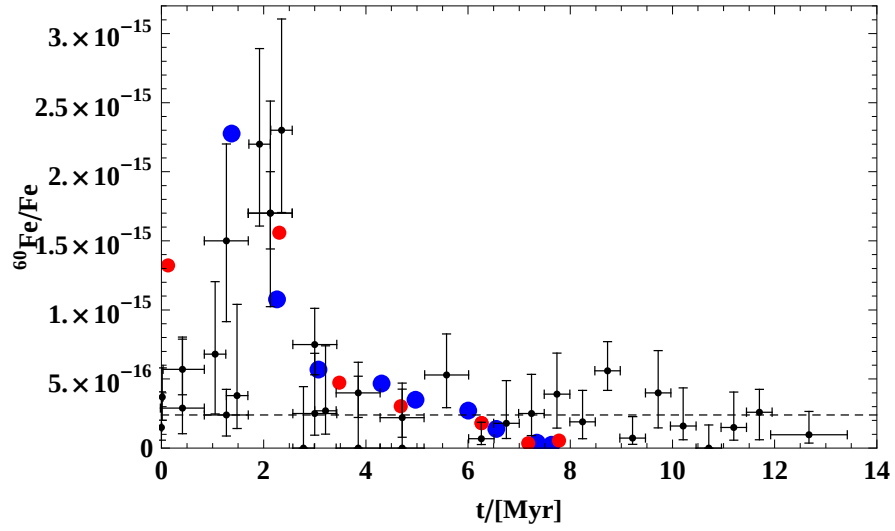


Figure E.24: The computed data (UCL: blue, LCC: red) plotted over the ^{60}Fe measurements (black points) for the mean mass model with an ISM density of $n = 10 \text{ atoms/cm}^3$. Here we plot the trajectories we found for a density of 0.01 atoms/cm^3 .

BIBLIOGRAPHY

- [1] S. K. Aggarwal. Precise and accurate determination of alpha decay half-life of ^{244}Pu by relative activity method using thermal ionization mass spectrometry and alpha spectrometry. *Radiochimica Acta, Oldenbourg Wissenschaftsverlag*, 94:397 – 401, 2006.
- [2] S. Basu, F.M. Stuart, C. Schnabel, and V. Klemm. Galactic-Cosmic-Ray-Produced ^3He in a Ferromanganese Crust: Any Supernova ^{60}Fe Excess on Earth? *Physical Review Letters*, 98(141103), 2007.
- [3] J. Beer, S.J. Johnsen, G. Bonani, R.C. Finkel, C.C. Langway, H. Oeschger, B. Stauffer, M.J.F. Suter, and W. Wölfli. ^{10}Be peaks as time markers in polar ice cores. *The Last Deglaciation: Absolute and Radiocarbon Chronologies*, pages 141–153, 1992.
- [4] T.W. Berghöfer and D. Breitschwerdt. The origin of the young stellar population in the solar neighborhood – a link to the formation of the Local Bubble? *Astronomy and Astrophysics*, 390:299–306, 2002.
- [5] K. Bethge. *Kernphysik*. Springer, 1996.
- [6] R. Bibron, R. Chesselet, G. Crozaz, G. Leger, J.P. Mennessier, and E. Picciotto. Extra-terrestrial ^{53}Mn in Antarctic ice. *Earth and Planetary Science Letters*, 21:109–116, 1974.
- [7] D. Breitschwerdt. Die Lokale Blase und ihr Ursprung. *Sterne und Weltraum*, (8):30–39, 2005.

- [8] D. Breitschwerdt and M.A. de Avillez. The history and future of the Local and Loop I bubbles. *Astronomy and Astrophysics*, 452:L1–L5, 2006.
- [9] G. Cini Castagnoli, G. Bonino, B. Lehmann, and C. Taricco. Cosmogenic Isotopes, Geomagnetic and Geochemical Signals in a Mediterranean Sea Sediment at 35000 yr BP. *24th International Cosmic Ray Conference*, 4:1204, 1995.
- [10] C.J. Clarke and R.F. Carswell. *Principles of Astrophysical Fluid Dynamics*. 2007.
- [11] D.L. Cook, E. Berger, T. Faestermann, G. Herzog, K. Knie, G. Korschinek, M. Poutivtsev, G. Rugel, and F. Serefidin. ^{60}Fe , ^{10}Be , and ^{26}Al in Lunar Cores 12025/8 and 60006/7: Search for a Nearby Supernova. *40th Lunar and Planetary Science Conference*, (1129), 2009.
- [12] P.B. deMenocal. Plio-Pleistocene African Climate. *Science*, 270:53 – 59, 1995.
- [13] J. Ellis, B.D. Fields, and D.N. Schramm. Geological Isotope Anomalies as Signatures of nearby Supernovae. *The Astrophysical Journal*, 470:1227–1236, 1996.
- [14] B.D. Fields and J. Ellis. On deep-ocean ^{60}Fe as a fossil of a near-earth supernova. *New Astronomy*, 4:419–430, 1999.
- [15] B.D. Fields, K.A. Hochmuth, and J. Ellis. Deep-Ocean Crusts as Telescopes: Using Live Radioisotopes to Probe Supernova Nucleosynthesis. *The Astrophysical Journal*, 621:902–907, 2005.
- [16] C. Fitoussi, G. Raisbeck, K. Knie, G. Korschinek, T. Faestermann, S. Goriely, D. Lunney, M. Poutivtsev, G. Rugel, C. Waelbroeck, and A. Wallner. Search for Supernova-Produced ^{60}Fe in a Marine Sediment. *Physical Review Letters*, 101(12), 2008.
- [17] B. Fuchs, D. Breitschwerdt, M.A. de Avillez, C. Dettbarn, and C. Flynn. The search of the origin of the Local Bubble redivivus. *Monthly Notices of the Royal Astronomical Society*, 373:993–1003, 2002.
- [18] D.L. Hartmann. Radiative Effects of Clouds on Earth’s Climate. *Aerosol-Cloud-Climate Interactions*, page 151, 1993.
- [19] M. Honda and M. Imamura. Half-life of ^{53}Mn . *Physical Review C*, 4:1182–1188, 1971.

- [20] M. Imamura, T. Inoue, Nishiizumi, and S. Tanaka. ^{53}Mn in Deep-Sea Sediment Cores - An Indicator of Past Solar Activity. *16th International Cosmic Ray Conference*, 2:304–307, 1979.
- [21] F.D. Kahn. The Temperature in very old Supernova Remnants. *Astronomy and Astrophysics*, 50:145–148, 1976.
- [22] F.D. Kahn. The Galactic Fountain. *Lecture Notes in Physics*, 506:483–494, 1998.
- [23] A. Karakas and J.C. Lattanzio. Stellar Models and Yields of Asymptotic Giant Branch Stars. *Publications of the Astronomical Society of Australia*, 24:103–117, 2007.
- [24] K. Knie, G. Korschinek, T. Faestermann, E.A. Dorfi, G. Rugel, and A. Wallner. ^{60}Fe Anomaly in a Deep-Sea Manganese Crust and Implications for a Nearby Supernova Source. *Physical Review Letters*, 93(17), 2004.
- [25] K. Knie, G. Korschinek, T. Faestermann, C. Wallner, J. Scholten, and W. Hillebrandt. Indication for Supernova Produced ^{60}Fe Activity on Earth. *Physical Review Letters*, 83:18–21, 1999.
- [26] W. Kutschera. Das Sortieren von Atomen "One by One". *Physik in unsere Zeit*, 5:203 – 208, 2000.
- [27] R. Lallement, B.Y. Welsh, J.L. Vergely, F. Crifo, and D.M. Sfeir. 3D mapping of the dense interstellar gas around the Local Bubble. *Astronomy and Astrophysics*, 411:447–464, 2003.
- [28] L.D. Landau and E.M. Lifschitz. *Lehrbuch der theoretischen Physik*, Band VI, Hydrodynamik.
- [29] M. Limongi and A. Chieffi. The Nucleosynthesis of ^{26}Al and ^{60}Fe in Solar Metallicity Stars Extending in Mass from 11 to 120 $M_{\odot}$: The Hydrostatic and Explosive Contributions. *The Astrophysical Journal*, 647:483–500, 2006.
- [30] B. Lindblad. *Handbuch der Physik*, 53:21, 1959.
- [31] M. Mac Low and R. McCray. Superbubbles in disk galaxies. *The Astrophysical Journal*, 324:776 – 785, 1988.

- [32] A. Mangini, M. Segl, G. Bonani, H.J. Hofmann, E. Morenzoni, M. Nessi, M. Suter, W. Wölfli, and K.K. Turekian. Mass-spectrometric ^{10}Be dating of deep-sea sediments applying the Zürich Tandem Accelerator. *Nuclear Instruments and Methodes in Physics Research B5*, pages 353–358, 1984.
- [33] P. Massey, K.E. Johnson, and K. Degioia-Eastwood. The Initial Mass Function and Massive Star Evolution in the OB Associations of the Northern Milky Way. *Astrophysical Journal*, 454:151–171, 1995.
- [34] L.R. McHargue, P.E. Damon, and D.J. Donahue. Enhanced cosmic-ray production of ^{10}Be coincident with the Mono Lake and Laschamp geomagnetic excursions. *Geophysical Research Letters*, 22:659–662, 1995.
- [35] K. Nishiizumi, M. Imamura, M.W. Caffee, J.R. Southon, R.C. Finkel, and J. McAninch. Absolute calibration of ^{10}Be AMS standards. *Nuclear Instruments and Methods in Physics Research Section B*, 258:403–413, 2007.
- [36] M. Poutivtsev. Extraterrestrisches ^{53}Mn in hydrogenetischen Mangankrusten. 2007.
- [37] G.M. Raisbeck, F. Yiou, D. Bourles, C. Lorius, and J. Jouzel. Evidence for two intervals of enhanced ^{10}Be deposition in Antarctic ice during the last glacial period. *Nature*, 326:273–277, 1987.
- [38] T. Rauscher, A. Heger, R.D. Hoffman, and S.E. Woosley. Nucleosynthesis in Massive Stars with Improved Nuclear and Stellar Physics. *The Astrophysical Journal*, 576:323–348, 2002.
- [39] G. Rugel, T. Faestermann, K. Knie, G. Korschinek, M. Poutivtsev, D. Schumann, N. Kivel, I. Günter Leopold, R. Weinreich, and M. Wohlmuther. New Measurement of the ^{60}Fe Half-Life. *Physical Review Letters*, 103(7), 2009.
- [40] S. Schaller, D. Schaerer, G. Meynet, and A. Maeder. New grids of stellar models from 0.8 to 120 solar masses at $Z = 0.020$ and $Z = 0.001$. *Astronomy and Astrophysics Supplement Series*, 96:269–331, 1992.
- [41] M. Segl, A. Mangini, G. Bonani, H.J. Hofmann, M. Nessi, M. Suter, W. Wölfli, G. Friedrich, W.L. Plger, A. Wiechowski, and J. Beer. ^{10}Be - dating of a manganese crust from Central North Pacific and implications for ocean paleocirculation. *Nature*, 309:540–543, 1984.

- [42] J.J. Sepkoski. Phanerozoic overview of mass extinction. *Patterns and Processes in the History of Life*, pages 277 – 295, 1986.
- [43] D.M. Sfeir, R. Lallement, F. Crifo, and B.Y. Welsh. Mapping the contours of the Local Bubble: preliminary results. *Astronomy and Astrophysics*, 346:785–797, 1999.
- [44] H. Svensmark. Influence of Cosmic Rays on Earth’s Climate. *Physical Review Letters*, 81:5027 – 5030, 1998.
- [45] H. Svensmark and E. Friis-Christensen. Variation of cosmic ray flux and global cloud coverage - a missing link in solar-climate relationships. *J. Atmos. Sol.-Terr. Phys*, 59:1225 – 1232, 1997.
- [46] J.H. Thomas, R.L. Rau, R.T. Skelton, and R.W. Kavanagh. Half-life of ^{26}Al . *Physical Review C (Nuclear Physics)*, 30:385–387, 1984.
- [47] C. Wallner, T. Faestermann, U. Gerstmann, K. Knie, G. Korschinek, C. Lierse, and G. Rugel. Supernova produced and anthropogenic ^{244}Pu in deep sea manganese encrustations. *New Astronomy Reviews*, 48:145–150, 2004.
- [48] B.Y. Welsh, D.M. Sfeir, M.M. Sirk, and R. Lallement. EUV mapping of the local interstellar medium: the Local Chimney revealed? *Astronomy and Astrophysics*, 352:308–316, 1999.
- [49] R. Wielen. Landolt - Börnstein: Numerical Data and Functional Relationships in Science and Technology - New Series., in K. Schaifers, H.H. Voigt, eds. *Astronomy and Astrophysics*, 2:208, 1982.
- [50] S.E. Woosley and A. Heger. Nucleosynthesis and remnants in massive stars of solar metallicity. *Physics Reports*, 442:269–283, 2007.
- [51] S.E. Woosley and T.A. Weaver. The Evolution and Explosion of Massive Stars. II. Explosive Hydrodynamics and Nucleosynthesis. *The Astrophysical Journal Supplement Series*, 101:181–235, 1995.

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