Reaction parameters for rp-process calculations above $Z = 32$

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We discuss nuclear structure and nuclear reaction input parameters for rp-process nucleosynthesis network calculations beyond $Z = 32$. Due to the lack of experimental data for neutron-deficient nuclei in this region, we investigate currently available model predictions for nuclear masses and deformations as well as methods of calculating reaction rates (Hauser-Feshbach) and β-decay rates (QRPA and shell model). In addition, we apply a valence nucleon ($N_{\rm p}N_{\rm n}$) correlation scheme for the prediction of masses and deformations. We also describe the calculations of 2p-capture reactions, which had not been considered before in this mass region. Using this information as input, we then present nuclear reaction network calculations to investigate the influence of nuclear structure on the rp-process between Ge and Sn in various scenarios. We find that in X-ray bursts 2p-capture reactions accelerate the reaction flow into the $Z \geq 36$ region considerably. Therefore, the rp-process in most X-ray bursts does not end in the $Z = 32$–36 region as previously assumed and overproduction factors of 10^7–10^8 are reached for some light p-nuclei in the $A = 80$–100 region. This might be of interest in respect of the yet unexplained large observed solar system abundances of these nuclei. Nuclei in this region can also be produced via the rp-proces in accretion disks around low mass black holes. Our results indicate that the rp-process energy production in the $Z > 32$ region cannot be neglected in these scenarios. We discuss in detail the influence of the various nuclear structure input parameters and their current uncertainties on these results. It turns out that rp-process nucleosynthesis is mainly determined by nuclear masses and β-decay rates of nuclei along the proton drip line. We present a detailed list of nuclei for which mass or β-decay rate measurements would be crucial to further constrain the models.

I. INTRODUCTION

Energy generation and nucleosynthesis in explosive hydrogen burning are characterized at lower temperatures, $T \leq 3 \cdot 10^8$ K by the hot CNO cycles [1] and at higher temperatures, $T \geq 3 \cdot 10^8$ K by the rp-process [2] and the αp-process [3]. A recent systematic study of the reaction flow and the time scales associated with the rp- and the αp-process [4] showed that at sufficiently high temperature and density conditions, $T \geq 8 \cdot 10^8$ K and $\rho \geq 10^4 \mathrm{g/cm}^3$, the reaction path may well proceed beyond mass $A = 64$ and $Z = 32$ depending on the time scale of the explosive event. Extreme hydrogen burning conditions of this kind have been proposed for various scenarios.

Accretion onto black holes is a very promising scenario for the explanation of observed quasars, active galactic nuclei, and radio jets. For a recent review on the various models describing the accretion process and jet formation see Chakrabarti [5]. Thin or thick accretion disks around black holes with masses $M \leq 10^4$ $M_\odot$ are also a possible site for nucleosynthesis via the rp-process, if the viscosity parameter α [6] is low, $\alpha \leq 10^{-4}$ [7–9]. This is important, since information on the disk composition can be obtained by analysis of X-ray spectra (see for example [7] and references therein). Furthermore, radio jets, powered either by threading material out along magnetic field lines or by the pressure of the anisotropic radiation field, can expel significant fractions of the processed material into the interstellar medium, thus making these scenarios possible contributors to galactic nucleosynthesis [8]. Originally, the simulation of nucleosynthesis in accretion disks has concentrated on massive black holes with $M \geq 10^6$ $M_\odot$ where only fairly low temperature and density conditions are anticipated [9]. However, the accretion onto low mass black holes leads to more compact systems. For these conditions, high temperatures between $1{\cdot}10^9$ and $4{\cdot}10^9$ K can be reached in the accretion disk with densities between 10^3 and 10^5 g/cm^3 and inflow timescales between $4{\cdot}10^4$ and $4{\cdot}10^5$ s. This is quite sufficient to process light mass isotopes $A \leq 20$ towards heavier masses with $A \geq 70$ via the rp-process [4]. Recently, nucleosynthesis in a thick accretion disk around a 10 $M_\odot$ black hole has been modeled with a reaction network up to Ge [10]. For the cases with low density, nucleosynthesis via the rp-process proceeds up to the medium mass range between $A = 40$ and $A = 60$ within the given inflow time. The reaction flow is mainly hindered by inverse photodisintegration processes. For the high-density models, the hydrogen content in the accretion disk is exhausted and nucleosynthesis continues via α-capture processes. The reaction networks used for the previous calculations were limited to $Z \leq 32$ (Ge) [9,10]. This mass range is quite sufficient for the relatively low temperatures and densities in the accretion disk of a massive black hole, but it might be too limited for nucleosynthesis simulations in the accretion disk of a black hole with $M \leq 10^2$ $M_\odot$.

The possibility of hydrogen burning in the accretion disk of a neutron star has been addressed as well [11]. The models discussed are characterized by a rather small viscosity, and a low accretion rate. At a mass transfer rate of approximately 10^{-14} $M_\odot \mathrm{yr}^{-1}$, temperatures greater than $T \approx 10^8$ K at a density of $\rho \approx 10^4$ g/cm^4 are reached in the center of the accretion disk. For these conditions, the nuclear energy generation is entirely based on the β-limited CNO cycle and becomes temperature independent. It should however be noted that for these calculations, the reaction network was limited to the pp-chains, the CNO cycles, and the triple-α-reaction only. Possible contributions to the energy generation and nucleosynthesis via the rp-process have not been considered. Because this process is extremely temperature sensitive, its possible influence on hydrogen burning in accretion disks around neutron stars should be investigated.

Type I X-ray bursts have been pointed out as possible sites for high temperature hydrogen burning via the rp- and αp-processes [2,3,12–16]. The standard models are based on accretion in a close binary system onto the surface of a neutron star with an approximate accretion rate of 10^{-8} to 10^{-10} $M_\odot \mathrm{yr}^{-1}$. Nuclear burning is ignited at the base of the accreted envelope at high densities, $\rho \geq 10^5$ g/cm^3, via the pp-chains, the CNO-cycles, and the triple-α-process. The released energy triggers a thermal runaway near the surface of the neutron star at highly degenerate conditions, and peak temperatures up to $3{\cdot}10^9$ K can be reached before the degeneracy is lifted when the Fermi temperature is reached. These temperatures are sufficiently high to trigger the rp- and the αp-process which cause rapid nucleosynthesis towards heavier nuclei and produce the energy ultimately observed as an X-ray burst. The time scale for the thermal runaway varies between 10 s up to 100 s [3]. Within this time scale, the rp-process can proceed well up to ^{56}Ni [2] or even further [3,4,17]. Previous attempts to extend the network calculations beyond $Z = 32$ suffered considerably from the lack of reliable nuclear physics information on masses, lifetimes, and reaction rates [17]. The simulation of rp-process nucleosynthesis in these scenarios which is crucial for the understanding of energy generation, fuel consumption, and eventually contributions to galactic nucleosynthesis, therefore requires a substantially improved input data set.

The rp-process also plays a role in novae. These events occur in close binary systems when matter from a main sequence star is accreted onto the surface of a white dwarf companion and ignites under degenerate conditions (e.g. [18–21]). Although the main energy sources are the CNO cycles, the rp-process is important in the determination of the isotopic composition of the ejecta, especially in terms of galactic γ-ray emitters like ^{22}Na and ^{26}Al. However, at densities of 10^3–10^4 g/cm^3, peak temperatures of maximal 0.4 GK (for O-Ne-Mg white dwarfs), and burning timescales of the order of 100–1000 s, the rp-process does not proceed beyond mass 40 [4,21]. Scenarios of this type are therefore not considered in this study.

The rp-process has also been proposed as the dominant energy source in the convective envelope of Thorne-Zytkow objects [22–24] which are predicted to appear like red giants, but they contain a degenerate neutron core [25]. It was suggested that most of the luminosity of the object originates from rp-process burning in the convective envelope rather than from the accretion onto the core. The convection zone extends continually from the surface down to the bottom of the envelope. The rp-process is dominant at non-degenerate conditions just above the neutron core at temperatures in the order of $T \approx$ (1–2) GK and densities between $\approx 10^2$–10^3 g/cm^3. Large scale convection mixes fresh material down from the outer layers of the envelope into the hot temperature zone and transports the processed material out towards the cooler surface. The macroscopic time scale for the turnover process through the hot bottom regions ranges from 0.01 s to 0.1 s [22]. Within this time scale, the initial material can be processed via the rp-process towards higher masses causing enrichment of the long-lived β^+-emitters. During the long periods outside the burning zone $\approx 10^8$ s, the material decays back to stability before being processed again with the next turnover. It has been suggested that within the lifetime of the star, heavy nuclei up to to $A \approx 150$ could be produced by continuous convective exposure of matter passing through the hot bottom zone [22–24,26]. It has been shown however that (γ,α)-reactions which had not been considered by the previous authors strongly reduce the reaction flow to nuclei beyond ^{68}Se [4]. Furthermore, recent calculations that include previously neglected neutrino emission indicate that

Thorne-Zytkow objects are unlikely to form and are also gravitationally unstable [27,28]. Therefore Thorne-Zytkow objects are not a probable astrophysical scenario. However, the mentioned exploratory studies were important and useful for all types of rp-process environments.

A reliable investigation of the nucleosynthesis and the energy generation in the discussed scenarios clearly requires the extension of the reaction networks towards heavier nuclei with $Z \geq 32$ as well as a detailed analysis of the reaction path and the reaction flow in this mass range. An additional motivation of the present work was to investigate the possibility that light p-nuclei in the $92 \leq A \leq 98$ mass range are produced via the rp-process. The origin of the large abundances of these nuclei, especially ^{92}Mo, ^{94}Mo, ^{96}Ru and ^{98}Ru, as observed in the solar system is a longstanding problem and the rp-process might offer a solution. p-nuclei are nuclei in the mass range $74 \leq A \leq 196$ on the more proton rich side of stability that cannot be synthesized by neutron capture in the s- or r-process (see for example [29]). The commonly accepted scenario for the production of most of these nuclei is the γ-process (photodisintegration of s- and r-nuclei) during explosive oxygen and neon burning in type II supernovae [30–33]. This scenario however cannot account for the large abundances of the p-nuclei in the $92 \leq A \leq 98$ mass range. Several scenarios for the production of additional amounts of these nuclei had been proposed, like the γ-process in type I supernovae [34] or the rp-process in Thorne Zytkow objects [35]. However, later studies showed that both of these scenarios are unlikely to produce p-nuclei (for type I supernovae see [36], for Thorne-Zytkow objects the previous paragraph).

Particularly important for rp-process nuclear reaction network calculations above $Z \geq 32$ are the nuclear masses of the relevant neutron-deficient isotopes, the proton-capture reaction rates, as well as their inverse photodisintegration rates, and finally the β-decay and electron-capture rates. First attempts to study the nucleosynthesis in X-ray bursts [17] and in Thorne-Zytkow objects [22,23,35] have been made. The calculations presented by Wallace and Woosley [17] were based on a 16 element approximation network (see also [2]), mainly designed to investigate the relevance of the rp-process beyond ^{56}Ni and the resulting energy generation in an X-ray burst. The nuclear masses were based on early macroscopic-microscopic model predictions [37], the proton-capture rates were calculated by an analytical approximation to Coulomb barrier penetration assuming one dominant exit channel [38], and the β^+-decay rates were estimated from the Gross Theory [39].

The initial calculations of the rp-process in Thorne-Zytkow objects [22] also relied on a simplified network based on similar nuclear input parameters. The unknown masses were predicted using a mass model based on the liquid drop model with an empirical shell correction [40] while proton-capture rates and β^+-decay rates in the upper mass regions were the same as in [17]. The nucleosynthesis of heavier isotopes was not calculated dynamically but was treated in the waiting point approximation, assuming a (p,γ)-(γ,p) equilibrium within the turnover timescale.

In a recent study of the rp-process in Thorne-Zytkow objects, nucleosynthesis has been calculated dynamically within the mass range between Fe and In [35]. The calculations were performed for a temperature grid between $7{\cdot}10^8$ K and $1.4{\cdot}10^9$ K. The nuclear masses were obtained from published data [41] or were based on predictions using a semi-empirical mass formula [42]. The proton-capture rates were again approximated using the formalism outlined in Ref. [38]. The β-decay rates were based on recent experimental results [43] and on calculations of the Gamow-Teller β^+-strength function using the proton-neutron quasi particle random-phase approximation (QRPA) [44–46].

The quality and predictive power of all these astrophysical models for energy generation as well as for nucleosynthesis depend clearly on the quality of the input parameters. Unfortunately, only a limited amount of experimental information is available for neutron-deficient nuclei with masses above $A = 64$. All calculations are therefore based on various model predictions of limited quality and accuracy, partly due to the rather complex nuclear structure in this particular mass range. It therefore seems neccessary to perform a systematic study to investigate the influence of various input parameters on the reaction path and the reaction flow of the rp-process at high temperatures and densities in the range of nuclei between $Z = 32$, $A = 64$ and $Z = 50$, $A = 100$.

In the following section (section II), we give an overview over several global mass models and compare the uncertainties for the predictions of masses, proton binding energies, α-binding energies, and Q_β-values. In this context, the influence of nuclear deformation is studied in detail. In section III, we describe the calculations of the rates for particle-induced nuclear reactions, photodisintegration processes, and β-decay half-lives of neutron-deficient isotopes. This includes a detailed discussion of the calculation of 2p-capture rates. In section IV, we present the results of extensive network calculations using reaction and decay rates based on different mass model predictions. The purpose of these network calculations is the following. First, we want to investigate the influence of the input parameters like reaction rates and nuclear masses on nucleosynthesis results. This leads to an outline of future experiments that will be most important to improve the accuracy of these calculations. Second, we want to demonstrate that in some of the discussed scenarios, rp-process nucleosynthesis above $Z = 32$ plays an important role in terms of galactic nucleosynthesis and energy generation.

II. MASS MODELS

The quality and reliability of the rp-process reaction network calculations depend critically on the accuracy of the nuclear masses which determine the position of the proton drip line, $M(Z+1,N) - M(Z,N) - m_\mathrm{p} = 0$, and the Q-values for proton-capture processes, β-decays, and other reactions. While most of the masses for light nuclei and for nuclei close to stability are experimentally well known [47], experimental information on masses for neutron deficient nuclei far from stability above $Z = 32$ is rather limited. Fig. 1 gives an overview over the available mass data in this region. The extension of rp-process network calculations beyond $Z = 32$ discussed in this paper requires reliable predictions for the experimentally unknown masses of isotopes close to the proton drip line. The presently available mass models [48–50] differ considerably in their predictions. Also, the agreement with the experimental data [47] is not quite satisfying. This is demonstrated in Fig. 2 which shows for several mass models the error in predicting experimentally known masses for the nuclei of the neutron deficient side of β-stability in each isotopic chain. The theoretical error σ_th is calculated using the recursive formula [50]:

$$\sigma_\mathrm{th}{}^2 = \frac{\sum_{i=1}^{n} w_i \left[(M^i_\mathrm{exp} - M^i_\mathrm{th})^2 - \sigma^i_\mathrm{exp}{}^2 \right]}{\sum_{i=1}^{n} w_i} \qquad \text{with} \qquad w_i = (\sigma^i_\mathrm{exp}{}^2 + \sigma_\mathrm{th}{}^2)^{-2}, \tag{1}$$

where M^i_exp are the experimental masses with experimental errors σ^i_exp and M^i_th the theoretical masses. This procedure decouples the theoretical errors from the experimental errors. Fig. 2 indicates typical deviations in σ_th between 0.1 and 1 MeV for all models. These errors represent a considerable improvement in comparison to earlier models (see [40,48]). However, mass model uncertainties of this order of magnitude are still dominating the overall uncertainties in astrophysical nucleosynthesis calculations (see section IV). The influence of nuclear masses has been studied in great detail for questions of r-process nucleosynthesis [51–54] and it had been demonstrated that the use of different mass models can modify the predicted r-process composition considerably.

Mass models are also important for the prediction of nuclear shapes, which affect shell structure and level density, and are therefore crucial for an accurate theoretical treatment of reaction rates. However, only some of the available mass models include shape degrees of freedom and can therefore be used to calculate nuclear deformations.

A. Nuclear masses

We have concentrated in this study on four mass models that are either frequently used in nucleosynthesis calculations or have been suggested to be reliable for predictions in the mass region discussed here. Two of the models are based on the macroscopic-microscopic approach where the total ground-state energy of the nucleus is calculated as the sum of a macroscopic term E_mac, and a microscopic term E_mic

$$M_\mathrm{ex}(Z,N)c^2 = E_\mathrm{tot}(Z,N,\mathrm{shape}) = E_\mathrm{mac}(Z,N,\mathrm{shape}) + E_\mathrm{mic}(Z,N,\mathrm{shape}). \tag{2}$$

The two models in this category are the ones by Hilf et al. from 1976 [55], and the FRDM (1992) by Möller et al. [50].

The classical mass formula by Hilf et al. is still frequently used for nucleosynthesis calculations [4,56] especially in r-process models, where it is more successful than some of the recent mass models [51]. The Hilf et al. mass formula is based on the work of von Groote et al. [57] where the macroscopic term is derived from the droplet model [40]. Hilf et al. improved the model considerably to predict the observed steep slope of the mass parabola far from stability. The microscopic term is based on semi-empirical shell corrections which were constructed from the gross single particle level densities [57].

The global FRDM (1992) mass model by Möller et al. [37,50,58,59] is also of the macroscopic-microscopic type. This model was successfully applied in r-process calculations but leads to some problems near shell closures [51]. For the present calculations, the latest mass tabulations [50] have been used. The macroscopic term is based on the finite range droplet model (FRDM) and the microscopic term is determined from calculated single-particle levels. In this particular model a folded-Yukawa potential is used. The details of the model terms are discussed extensively in ref. [50]. The calculations have improved over earlier versions [58,59] in the treatment of the pairing terms, and by including additional shape degrees of freedom.

In addition we used a valence nucleon correlation scheme to predict unknown masses. This scheme is based on the correlation of the microscopic mass correction, that is determined by the nuclear structure, with the average number of p-n interactions given by the parameter P [60,61]. P can be simply expressed with the number of valence protons N_p and neutrons N_n:

$$P = \frac{N_{\rm p} N_{\rm n}}{N_{\rm p} + N_{\rm n}} \quad (3)$$

Valence nucleon correlation schemes have been very successful in parametrizing nuclear shape changes and predicting the onset of collectivity in various regions of nuclei. They may therefore serve as an empirical method for the prediction of nuclear structure effects in the microscopic mass correction.

This scheme is applied by characterizing a particular region of nuclei using experimentally known masses. Semi-empirical microscopic mass corrections are obtained from experimental mass data by subtracting a macroscopic component that describes general nuclear properties. The macroscopic mass component we used was from the FRDM (1992) mass model. The semi-empirical microscopic mass corrections obtained with this method were plotted as a function of P. This is shown in Fig. 3 for nuclei with proton and neutron numbers between the shell closures at $N, Z = 28$ and $N, Z = 50$. We restrict this study to this range of nuclei, but the method can in principle be applied to any mass region. Fig. 3 demonstrates the strong correlation that had already been shown for the actinides and the rare-earth nuclei [61]. A cubic fit through the data can then be used to characterize the mass region and in turn be used for interpolation of microscopic mass corrections of unknown nuclear masses by simply using the P-values of the respective nuclei. Note that no extrapolation is necessary with this method, since enough experimental data are available to span the whole possible range of P (which is 0.5–5.5 for $28 < N, Z < 50$). A more detailed description of this method along with other nuclear structure results will be published separately [62].

The mass formula by Jänecke and Masson [42] is based on the Garvey-Kelson relation which allows the prediction of masses by extrapolation from known masses in the same region or from the masses of the mirror nuclei.

The four different methods presented above are applied to derive nuclear masses for all stable or neutron-deficient isotopes in the $Z = 32$ to $Z = 50$ range. The results are compared with experimentally known masses. Fig. 4 shows the difference between the theoretical and the experimental masses for the isotopic chains $Z = 34$, $Z = 36$, $Z = 38$, $Z = 39$, $Z = 40$, $Z = 42$, $Z = 44$, and $Z = 48$ as a function of the mass number A. Fig. 5 gives the corresponding theoretical errors (see Eq. 1) for the mass predictions for all nuclei considered here. All plots clearly indicate that the predictions of the Jänecke et al. formula are in good agreement with the experimental masses (within 100 keV) while the other mass models show considerably higher deviations (400–600 keV). Two features are noteworthy; first, all of the here considered mass models overpredict the ^{80}Y mass by 2 MeV. For the Jänecke et al. approach this is the only deviation larger than 0.7 MeV in this mass region. The mass of ^{80}Y was measured with the β-γ coincidence method in two experiments [63,64] and the resulting uncertainty is only 129 keV [47]. However, a recently performed measurement based on cyclotron mass spectroscopy yields a 2.18(17) MeV lower ^{80}Y mass [65] in better agreement with the mass model predictions. Second, for Zr ($Z = 40$), Nb ($Z = 41$) and Mo ($Z = 42$) the FRDM (1992) mass model shows especially large deviations (up to 2 MeV) for nuclei within 3 mass units from the $N = Z$ line. This is a region of highly deformed nuclei and the deviations correlate directly with deviations in the prediction of nuclear deformation. This point is discussed in paragraph II B. In addition, the FRDM (1992) mass model and to some extent also the Hilf et al. model exhibit large deviations around neutron number 55 for the same group of elements.

B. Nuclear deformations

The rp-process reaction path includes nuclei with significantly deformed shapes, especially in the $Z = 36$–40 region. Compared to spherical nuclei, permanently deformed nuclei have different single particle levels and additional collective degrees of freedom, that increase level densities. Deformation is also an important parameter in mass models, since the potential energy of a nucleus depends strongly on its shape. Single particle levels, level densities and nuclear masses are however the most important ingredients to calculate nuclear reaction rates and β-decay half-lives. For a reliable calculation of these quantities it is therefore important, to take the shape of the nucleus in the gound state into account.

There are several shape parametrizations commonly used in the literature that are all based on the assumption of a constant matter and charge density within the nucleus and a reasonably sharp decrease at the edges. We will use the ϵ-parametrization throughout this paper, which is also used in the FRDM (1992) mass model and in our QRPA calculations for β-decay rates. For a definition of this parametrization and how it is related to the other shape descriptions (for example the β-parametrization) see for example the recent review by Nazarewicz and Ragnarsson [66].

We will concentrate in the following discussion on quadrupole deformations, which are the most important ones. They are described by a nonzero deformation parameter ϵ_2, which is negative for oblate and positive for prolate shapes. Nuclear deformation is not directly accessible by experiments, but there are measurable quantities that can be related to the nuclear shape under certain conditions. One of these conditions is the axial symmetry of the deformed nucleus.

In the following we will assume axial symmetry although there has been a claim that some of the nuclei we consider here might in fact be asymmetric rotors [67].

One possibility to derive information on the nuclear shape is to measure the spectroscopic ground-state quadrupole moment Q_s. The Q_s is related to the intrinsic quadrupole moment Q_i by

$$Q_s = \frac{3K^2 - J(J+1)}{(J+1)(2J+3)} Q_i, \tag{4}$$

where J is the angular momentum of the nucleus and K is its projection on the symmetry axis. For $J = 0$ the spectroscopic quadrupole moment vanishes. Ground state spectroscopic quadrupole moments can therefore only be measured for odd-even, even-odd and odd-odd nuclei ($K = J$ in the ground-state). The intrinsic quadrupole moment is directly related to the nuclear charge distribution and therefore to the nuclear shape. Considering terms up to second order we can calculate ϵ_2 using

$$\epsilon_2 = -1 + \sqrt{1 + \frac{5}{2ZR_0^2} Q_i}. \tag{5}$$

where Z is the charge number and R_0 is the mean nuclear radius, which can be approximated by $R_0^2 = 0.0144 \cdot A^{2/3} b$.

Another possibility to derive the absolute value of the intrinsic ground-state quadrupole deformation is to measure the reduced E2 γ-transition strength $B(E2:(JK)_{initial} \rightarrow (JK)_{final})$ within a rotational band. This is mostly done in even-even nuclei, using the transition from the first excited 2^+ state to the 0^+ ground-state ($K = 0$). Assuming that the 2_1^+ state is predominantely rotational, we can derive the permanent intrinsic quadrupole moment using

$$|Q_i| = \sqrt{\frac{16\pi}{5} \frac{B(E2:2^+ \rightarrow 0^+)}{e^2}}. \tag{6}$$

Eq. 5 allows then to derive the quadrupole deformation parameter. By using $B(E2:2_1^+ \rightarrow 0_1^+)$ data it is not possible to distinguish between prolate and oblate shapes, because the sign of Q_i in Eq. 6 remains undetermined.

For some even-even nuclei, especially the more unstable ones, no experimental $B(E2:2_1^+ \rightarrow 0_1^+)$ values are available, but the energy of the first excited 2^+ state $E(2_1^+)$ is known. In these cases the quadrupole deformation can be estimated from $E(2_1^+)$, using the global systematics established by Grodzins et al. [68,69]:

$$\epsilon_2 \approx 0.95 \cdot \sqrt{\frac{1228}{A^{7/3} E(2_1^+)}} \tag{7}$$

A is here the mass number and $E(2_1^+)$ the excitation energy of the 2_1^+ state in MeV. We used here the relation $\epsilon_2 \approx 0.95 \cdot \beta_2$, which neglects higher order β-defomations. Eq. 7 seems to work rather well in this mass region, see for example [70].

Fig. 6 shows the experimental quadrupole deformation parameters for the isotopic chains Z = 34, 36, 38, 40, 42, 44 plotted as functions of mass number. The experimental quadrupole moments were taken from the compilation of Raghavan [71], updated with recent results from laser spectroscopy experiments (for Z = 36 [72], and for Z = 38 [73,74]). The experimental $B(E2:2_1^+ \rightarrow 0_1^+)$ values are from the compilation of Raman et. al. [75]. The $E(2_1^+)$ data are from ref. [76] for ^{72}Kr, [67] for ^{76}Sr, [70] for ^{80}Zr and Nuclear Data Sheets for the other nuclei. The experimental quadrupole deformations in Fig. 6 show that neutron-deficient nuclei with $36 \leq Z \leq 42$ are strongly deformed with extreme deformations of up to $\epsilon_2 = 0.4$ occuring for some nuclei close to the $N = Z$ line. However, the experimental data do not extend yet to the rp-process nuclei (see section IV) near the proton drip line. In the case of ^{72}Kr, ^{76}Sr and ^{80}Zr estimates from $E(2_1^+)$ data are possible, but for all other nuclei along the proton drip line deformation parameters have to be determined with theoretical models. While many theoretical shape calculations are available for selected isotopes in this region, mostly Rb and Sr nuclei, we want to concentrate in this study on global models that allow the consistent determination of mass and deformation for the whole nuclear reaction network.

Fig. 6 also shows the predicted quadrupole deformation of the two macroscopic-microscopic models by Hilf et al. [55] and Möller et al. [50] (FRDM 1992) described in section II A. We have also applied a valence nucleon correlation scheme ($N_\mathrm{p}N_\mathrm{n}$ scheme) similar to the one used to calculate masses. The Jänecke et al. mass model does not yield any information about the deformation which is treated implicitly by extrapolating the masses from neighboring or mirror nuclei.

The Hilf et al. model predicts only positive quadrupole deformation terms and therefore considers only prolate deformation for the ground-state of the nucleus. The deformation is calculated by minimizing the total energy of the nucleus as a function of the quadrupole deformation term ϵ_2, using for the macroscopic term a shape dependent droplet model. The shape dependence of the microscopic term is approximated by multiplying the shell correction with a shape dependent function. The predictions shown in Fig. 6 indicate deformation only for isotopes in the mass range $A = 75$–80 for the $Z = 32$ to $Z = 38$ chains.

In the FRDM (1992) approach deformation is treated up to higher multipole orders, including in addition to the quadrupole deformation, the octupole ϵ_3, the hexadecapole ϵ_4, and hexacontatetrapole ϵ_6 deformations. The total ground-state energy is determined by minimizing the potential energy surface using the ϵ-shape parametrization with the folded Yukawa potential. The procedure is described in detail in ref. [50]. Fig. 6 shows the resulting quadrupole deformation parameters ϵ_2, indicating rapid changes between prolate ($\epsilon_2 > 0$) and oblate deformation ($\epsilon_2 < 0$) in the ground-state for $Z \leq 37$. While the predicted deformation for isotopes with $Z \geq 43$ is rather small but slightly prolate along the isotopic chain, it is particularly noticable that for proton numbers in the range $34 \leq Z \leq 42$ strong deformations are expected for the neutron deficient nuclei $A \leq 84$. Oblate deformation dominates for $Z < 36$, while for $Z > 36$ mostly prolate shapes are predicted.

We also obtained the deformation parameters ϵ_2 from an $N_\mathrm{p}N_\mathrm{n}$ correlation scheme. This scheme is based on expressing the quadrupole deformation as a function of the product of valence protons and neutrons $N_\mathrm{p}N_\mathrm{n}$ (see section II A). The fit through the experimental data is used to predict quadrupole deformation parameters for unknown nuclei. This method was also used by Raman et al. [75], who showed that the best correlation is not obtained for the quadrupole deformation itself but for $\epsilon_2/\epsilon_{2(sp)}$ with

$$\epsilon_{2(sp)} = 0.95 \cdot \beta_{2(sp)} = 0.95 \cdot \frac{1.59}{Z}, \tag{8}$$

which introduces an additional charge number dependence. β_2 in Eq. 8 is the deformation parameter used by Raman et al. Fig. 7 shows the experimental $\epsilon_2/\epsilon_{2(sp)}$ values as a function of $N_\mathrm{p}N_\mathrm{n}$. The data from experimental $B(E2 : 2_1^+ \rightarrow 0_1^+)$ values and most of the data from measurements of ground-state quadrupole moments are well correlated. There are a few nuclei, for example ^{79}Rb ($N_\mathrm{p}N_\mathrm{n} = 72$) or ^{72}As ($N_\mathrm{p}N_\mathrm{n} = 55$) where the spectroscopic ground-state quadrupole moments give much smaller deformations than expected from their $N_\mathrm{p}N_\mathrm{n}$ value. This might indicate that the assumption of axial symmetry is not valid in these cases. We therefore use only the $B(E2 : 2_1^+ \rightarrow 0_1^+)$ measurements to obtain a linear fit, which allows us to predict unknown quadrupole deformations from their $N_\mathrm{p}N_\mathrm{n}$ value. Note that these unknown nuclei lie within the $N_\mathrm{p}N_\mathrm{n}$ range of the experimental data points, so that the predictions are obtained by interpolation, not by extrapolation. As shown in Fig. 6, the deformation increases smoothly towards the neutron-deficient side of the line of stability and reaches typically maximum deformation close to the $N = Z$ line.

The onset of deformation as predicted by the Hilf et al. mass formula is for $Z = 34$ and 38 in rough agreement with experimental data, though close to the proton drip line spherical shapes are predicted. This is in strong disagreement with the very high deformations indicated by experiments and other models. Furthermore, for $Z \geq 40$ the Hilf et al. model predicts spherical nuclei though it is now known that in the $Z = 38$–42 region the highest known ground-state deformations occur.

Overall good agreement can be observed between the experimental data and the absolute values of the deformation parameter $|\epsilon_2|$ predicted by the FRDM (1992) mass model, though it predicts a steeper increase in deformation towards lower mass numbers than indicated by the experimental data. This might lead to large deviations in the close vicinity of the onset of deformation. The decrease of deformation with increasing charge number (towards the closed shell at $Z = 50$) also seems to be too steep. For $Z = 44$ the FRDM (1992) mass model predicts hardly any deformation while $E(2_1^+)$ data and the $N_\mathrm{p}N_\mathrm{n}$ scheme seem to indicate still substantial deviations from spherical shape for some nuclei. Because of its ability to predict the sign of deformation, the FRDM (1992) mass model has in general an advantage. However, for a prolate minimum at ϵ_2 in the potential surface there is typically an oblate minimum close to $-\epsilon_2$. In some cases the energy difference between the oblate and the prolate minimum is smaller than the model uncertainties, which obviously means that the predicted sign of ϵ_2 is uncertain. This is for example the case for ^{77}Kr. The FRDM (1992) mass model predicts a strong oblate deformation ($\epsilon_2 = -0.22$), while experiments indicate a strong prolate deformation of about the same size. There is however another minimum in the potential energy surface at $\epsilon_2 = +0.35$, which is only 130 keV higher. The sign can therefore not be determined and within the model uncertainties the predicted deformation agrees with the experimental value (though Fig. 6 might suggest a large discrepancy).

The $N_\mathrm{p}N_\mathrm{n}$ scheme is limited to the prediction of the absolute value of quadrupole deformation. The comparison between the model results and the experimental values shows good agreement. Both the proposed increase of deformation and its slow decrease with neutron deficiency is reflected in the data. For $Z \geq 42$ only very limited data on ground-state deformations are available. Clearly, measurements of quadrupole moments or $B(E2 : 2_1^+ \rightarrow 0_1^+)$ values of neutron-deficient nuclei with $Z \geq 42$ are necessary to determine the onset of deformation towards the proton drip line and to test the various models in that range.

C. Reaction Q-values

The input parameters for the rp-process calculations are not the absolute nuclear masses, but mass differences, i.e. the proton separation energies S_{p}, the α-separation energies S_α and the the β^+-decay Q-values $Q_{\beta+}$. The proton separation energies are determined from the nuclear masses $M(Z,N)$, $S_{\mathrm{p}} = (M(Z,N) - M(Z-1,N) - M_{\mathrm{p}})c^2$. The proton separation energies determine the proton-capture processes and the inverse photodisintegration processes, as described in section III B. The α-separation energies, $S_\alpha = (M(Z,N) - M(Z-2,N-2) - M_\alpha)c^2$ are important for proper treating of backprocessing via α emission or (γ,α) photodisintegration processes as discussed in section III B. The reaction Q-value for a nuclear reaction $i(j,k)l$ that determines the energy production (or consumption) in the reaction network can be calculated from the separation energies in the entrance and exit channel: $Q_{i(j,k)l} = S_l - S_i$ ($S_\gamma = 0$ of course). The $Q_{\beta+}$-values, $Q_{\beta+} = (M(Z,N) - M(Z-1,N+1) - m_{\mathrm{e}})c^2$ are important for the treatment of the β-decay processes along the rp-process reaction path, in particular for calculating the Fermi integrals and the β-decay lifetimes (see section III D). The $Q_{\beta+}$-values are often expressed as electron capture Q-values $Q_{EC} = Q_{\beta+} + 2m_{\mathrm{e}}c^2$.

Fig. 8 shows the proton separation energies for the neutron-deficient odd Z isotopes between $Z = 33$ and $Z = 49$ as predicted by the four different mass models. Also indicated are the experimentally known proton separation energies [47]. The solid line indicates the proton drip line S_{p}=0. As can be seen in Fig. 8, reasonable agreement exists between the different predictions within the range of experimental data. Strong deviations between the experimental values and the predictions by the FRDM (1992) mass model [50] occur for the deformed Rb ($Z = 37$), Sr ($Z = 38$), Y ($Z = 39$), Zr ($Z = 40$), and Nb($Z = 41$) isotopes near $N = 42$, as already pointed out in section II A. Extending the calculations towards the proton drip line, the proton separation energies predicted by the various models diverge considerably. This divergency is typically much larger than the average deviations to experimental values indicating that the uncertainties of at least some mass models increase considerably towards the proton drip line. Particularly noticeable is also that the Hilf et al. prediction [55] is generally softer and that the drip line is therefore shifted towards the neutron deficient side compared to the other, more recent predictions. This is typical for almost the entire range of nuclei considered here. Only for $Z \geq 47$ the various predicted drip lines are in better agreement (see also Fig. 9). The deviations between the various mass predictions near the proton drip line are typically ≤ 2 MeV. A better agreement of ≤ 1 MeV is observed, however, between the latest macroscopic-microscopic prediction, which takes deformation effects into account [50], and the predictions based on the $N_{\mathrm{p}}N_{\mathrm{n}}$-scheme and the Jänecke et al. Garvey Kelson relation.

Fig. 10 shows the α-separation energies S_α for the even Z isotopic chains from $Z = 34$ to $Z = 48$ as a function of the mass number A. The predicted values show a reasonable agreement within ≤ 1 MeV with the experimental values. The plots also indicate clearly the increase in α-binding energy around the closed neutron shell $N = 50$. Substantial deviations between the various predictions for the isotopic chains $Z = 38$ to $Z = 44$ do, however, occur near $N = 42$. The FRDM (1992) mass model [50] predicts substantially smaller α-binding energies than the other models. In the case of $Z = 42$, 43, and 44, the ^{84}Mo, ^{85}Tc, ^{84}Ru, and ^{85}Ru isotopes are even predicted to be α unbound. Though due to the Coulomb barrier β-decay will still be the dominant decay channel, compound nuclei with low α-binding energies lead to a strongly enhanced (p,α)-cross section. This is important for the ^{83}Nb(p,α)^{80}Zr reaction, which is part of the rp-process reaction flow (see section IV). The strong enhancement of this reaction branch derived on the basis of the FRDM (1992) mass mass model might lead under some conditions to considerable backprocessing of the rp-process reaction flow and to the formation of a Zr-Nb cycle. (see section IV F). This predicted drop in α-binding energy is however not observed in the experimental data available for the $Z = 38$ to $Z = 41$ chains. Furthermore, there is a direct correlation between the drip in α-binding energy and the ϵ_2 ground-state deformation of the respective isotopes. While at $N = 42$ the FRDM (1992) mass model predicts only a small deviation from the spherical shape with $\epsilon_2 = 0.05$ and a rapid change towards decreasing neutron number, the experimental data indicate strong ground-state deformations with $\epsilon_2 \approx 0.2$–0.4. The drop in α-binding energy predicted by the FRDM (1992) mass model may therefore be an artifact due to problems with the prediction of nuclear shapes.

Fig. 11 shows the $\mathrm{Q_{EC}}$-values for the even Z isotopic chains from $Z = 34$ to $Z = 48$ plotted as functions of mass number A. The predicted values agree typically within ≤ 0.5 MeV with the experimental results. As discussed already in section II A and above in connection with proton and α-separation energies, substantially larger deviations up to 1 MeV occur for the deformed $Z = 39$, 40, 41, 42 isotopes close to the $N = Z$ line. The values predicted by the Hilf et al. model [55] are typically lower by approximately 2 MeV near the $N = Z$ line than the predictions of the other models. This may have a considerable influence on the prediction of β-decay lifetimes as discussed in section III D.

D. Summary of mass model properties

Fig. 5 shows the theoretical errors (see Eq. 1) of the various mass models for the nuclei on the neutron-deficient side of β-stability between As and Sn. Here, the Jänecke et al. model shows uncertainties of $\approx$ 100 keV, while the other models show larger errors (about 500–700 keV).

As discussed in section II C the relevant parameters for the rp-process calculations are not the masses itself but mass differences. We therefore show in Fig. 5 also the theoretical errors for proton separation energies, α-separation energies, and electron capture Q-values. For the FRDM (1992), Hilf et al. and $N_{\rm p}N_{\rm n}$-models the theoretical errors for mass differences are significantly smaller, especially for the proton separation energies, while for the Jänecke et al. model the uncertainties get slightly larger. For proton separation energies the errors range from 110 keV for the Jänecke et al. mass model to 400 keV for the FRDM (1992) mass model. It is remarkable that the much simpler $N_{\rm p}N_{\rm n}$ scheme predicts proton separation energies with the same accuracy than the FRDM (1992) mass model although deviations for single masses are much larger.

An important consideration, however, is the predictive power of the various mass models for nuclei close to the proton drip line, where no experimental data are available. The increasing differences between the various mass model predictions towards the proton drip line in Fig. 8 indicate that uncertainties become larger for extrapolations towards more neutron-deficient nuclei. This systematic behaviour is also visible in Fig. 12, which shows the difference between predicted and measured masses as a function of the z-component of the isospin $T_z = (N-Z)/2$. There is a systematic increase in deviations towards the proton drip line (lower T_z), even if the larger experimental errors are taken into account. This increase is rather moderate for the Hilf et al. mass model, but quite pronounced for the other models. The FRDM (1992) mass model shows for most nuclei with $T_z \leq 2$ only a weak increase in uncertainty, while on the other hand very large deviations of more than 1 MeV occur for a group of 13 nuclei. With the exception of ^{64}Ge, ^{66}As, and ^{101}In, all these nuclei lay at the transition from spherical shape to strong deformation in the Sr-Mo region. The large deviations in the masses predicted by the FRDM (1992) are therefore correlated with the large deviations in the predicted shapes at the onset of deformation as discussed in sections II B and II C. To get an estimate of the predictive power we can also compare the theoretical errors of the various mass models to the errors we obtain by considering only masses that have been measured since 1988 (23 nuclei). These mass data could not be used to fit parameters in the Hilf et al. and Jänecke et al. models and are also the most neutron-deficient masses that are known experimentally. Fig. 5 shows the results. For all models there is a significant increase in the theoretical errors for the most recently measured nuclei. This shows that all these models have problems extrapolating towards neutron-deficient nuclei in this mass region. However, the increase in error varies strongly among the mass models. It is relatively small for the Hilf et al. mass formula (30%), moderate for the FRDM (1992) mass model (70%) and the $N_{\rm p}N_{\rm n}$ model (100%), but quite strong for the Jänecke et al. mass model (450%). Though the Jänecke et al. mass model has still a theoretical error comparable to the other models, this steep increase in uncertainty indicates, that this model is less useful for extrapolating towards neutron deficient nuclei in this mass region. These trends are confirmed by recent mass measurements of ^{100}Sn and ^{100}In [77] that had been performed after the compilation of this study.

An important consideration for nuclear reaction network calculations is the consistency of the different input parameters. It was shown in r-process calculations (see for example [56]) that when input data are taken from different theoretical sources, occasionally nuclear structure effects vanish or artificial effects show up. The FRDM (1992) mass model offers here the advantage of a unified approach, allowing the prediction of masses, deformations, and together with the QRPA code described in section III D also β-decay rates with a consistent description of nuclear structure. Most network calclations in section IV will therefore be performed with an input parameter set based on the FRDM (1992) mass model. For comparison we will also use the Hilf et al. mass formula, which predicts the proton drip line to be further away from stability (see Fig. 8 and 9), to investigate the effects on the reaction path, processing timescale, and nucleosynthesis.

III. INPUT FOR NETWORK CALCULATIONS

Nuclear reaction network calculations are necessary to follow the time evolution of the isotopic abundances $Y_i = X_i/A_i$ (mass fraction divided by mass number, which has here the units g/mole, see for example [78]), to determine the amount of energy released by nuclear reactions and to find the reaction flux which defines the actual reaction path for the nucleosynthesis process. This depends strongly on the temperatures and densities in the explosive event, which may change rapidly.

The reaction network is defined as a set of differential equations for the various isotopic abundances. The time derivative of the abundance for each isotope is expressed in terms of the reaction rates of the different production and depletion reactions,

$$\frac{dY_i}{dt} = \sum_j N^i_j \lambda_j Y_j + \sum_{j,k} N^i_{j,k} \rho N_A < \sigma v >_{(j+k)} Y_j Y_k + \sum_{j,k,l} N^i_{j,k,l} \rho^2 N_A^2 < \sigma v >_{(j+k+l)} Y_j Y_k Y_l. \tag{9}$$

The rates of the reactions represented by the three terms on the right side of the equation are categorized in decay and photodisintegration rates λ_j, two particle capture rates $\rho N_A < \sigma v >_{(j+k)}$, and three particle interaction rates $\rho^2 N_A^2 < \sigma v >_{(j+k+l)}$. The particle induced reaction rates depend on the density ρ as defined in [78], which can be set equal to the mass density (with the mass unit gram, when the definiton of Y from above is used) in all practical calculations. N_A represents Avogadro's number. The individual N^i's are given by $N^i_j = N_i$, $N^i_{j,k} = N_i/(N_j!N_k!)$, and $N^i_{j,k,l} = N_i/(N_j!N_k!N_l!)$. The N_i's represent positive or negative numbers which specify how many particles of species i are created or destroyed in the reaction. The denominators, including factorials, avoid double counting of the number of reactions when identical particles react with each other. For a detailed discussion see [79].

The time integrated net reaction flow between two isotopes i and j is defined by,

$$F_{i,j} = \int \left[\frac{dY_i}{dt}_{(i\rightarrow j)} - \frac{dY_j}{dt}_{(j\rightarrow i)} \right] dt. \tag{10}$$

The time evolution of the isotopes, and the time integrated reaction flux depend critically on the reaction rates. For the extension of the rp-process network beyond $Z = 32$, β-decays, photodisintegrations, proton, and α-captures have to be considered and the respective reaction rates have to be calculated. The calculations presented here are based on the masses and deformations which were derived in the framework of the models presented in the previous section. Therefore, the reaction rates are subject to the uncertainties inherent in these input parameters.

A. Particle induced reactions

The rate for a nuclear reaction $i(j,k)l$, $N_A\!<\!\sigma v\!>_{(j,k)}$, can be expressed in terms of the particle energy E, the energy dependent reaction cross section $\sigma(E)$, and the stellar temperature T by,

$$N_A < \sigma v >_{(j,k)} = \left[\frac{8}{\mu \pi} \right]^{1/2} (kT)^{-3/2} \int_0^\infty E\sigma(E) \cdot \exp(-E/kT) dE. \tag{11}$$

Here μ denotes the reduced mass of the target projectile system. The reaction cross section depends critically on the Coulomb barrier between the two interaction particles (for charged particle reactions) and on the nuclear structure of the compound nucleus.

In the cases of a low level density in the compound system, the reaction cross section is determined by single resonant Breit Wigner terms and by nonresonant contributions. The resonant terms correspond to unbound states in the compound nucleus, the nonresonant terms to transitions to bound states in the final nucleus

$$N_A < \sigma v > = N_A < \sigma v >_{\mathrm{r}} + N_A < \sigma v >_{\mathrm{nr}} . \tag{12}$$

Low level density conditions are predominatly observed in light mass nuclei or in nuclei near closed shells. For the reaction network calculations in this study most of the reaction rates for proton-capture on nuclei with $A \leq 40$ are based either on experimental data [80,81] or on recent calculations of the nuclear level structure of the compound system [4,82]. The implications of these rates on low temperature nucleosynthesis have been discussed extensively before [4,81–83] and will not be addressed in the present study.

For rp-process reactions on nuclei with $A > 40$, the level density in the compound system is typically high and the cross section is then dominated by a multitude of overlapping resonances and appears nonresonant [84]. The cross sections can then be approximated using the Hauser-Feshbach approach, a statistical model of compound nuclear reactions. This approach is well established and widely used for the calculation of thermonuclear reaction rates [56,85,86]. The reaction cross section for a reaction $i^\delta(j,k)l^\nu$ from the target i in state state δ to the final nucleus l in state ν is expressed in terms of the transmission coefficients T

$$\begin{aligned} \sigma^{\delta\nu}_{jk} = & \frac{\pi \hbar^2}{2\mu_{ij} E_{ij}} \cdot \frac{1}{(2J^\delta_i + 1)(2J_j + 1)} \\ & \cdot \sum_{J,\pi} (2J+1) \frac{T^\delta_j(E, J, \pi, E^\delta_j, J^\delta_j, \pi^\delta_j) \cdot T^\nu_k(E, J, \pi, E^\nu_k, J^\nu_k, \pi^\nu_k)}{\sum_m T_m(E, J, \pi)}, \end{aligned} \tag{13}$$

where E_{ij} is the center of mass energy, and μ_{ij} the reduced mass for the target projectile system. $T_k^\nu(E,J,\pi,E_k^\nu,J_k^\nu,\pi_k^\nu)$ is the transmission coefficient for the transition from a state (E,J,π) in the compound nucleus into the state ν of the final nucleus under emission of the particle k. $T_j^\delta(E,J,\pi,E_j^\delta,J_j^\delta,\pi_j^\delta)$ describes the entrance channel accordingly. $T_m(E,J,\pi)$ is the total transmission coefficient for the emission of a particle m via the transition from the excited state (E,J,π) in the compound nucleus to all possible states in the respective final nuclei. It can be expressed in terms of the individual transmission coefficients and depends strongly on the level density $\rho(E_m,J_m,\pi_m)$ in the final nucleus

$$T_m(E,J,\pi) = \sum_{\lambda=0}^{\omega} T_m^\lambda(E,J,\pi,E_m^\lambda,J_m^\lambda,\pi_m^\lambda) + \int_{E_m^\omega}^{E-S_m} \sum_{J_m,\pi_m} T_m(E,J,\pi,E_m,J_m,\pi_m)\rho(E_m,J_m,\pi_m)dE_m. \tag{14}$$

The first part represents the sum over all experimental known states in the final nucleus (labelled by the emitted particle m) and the second part represents the integration over the level density above the highest experimentally known level energy E_m^ω up to the maximum energy $E-S_m$ with S_m representing the respective channel separation energy. It should be pointed out that in most of the cases discussed here no information is available about the excited states; the transmission coefficients are therefore only determined by transitions to the ground-state $\omega=0$ and the second term in Eq. 14.

To calculate the total cross section into all final states ν

$$\sigma_{jk}^\delta = \sum_\nu \sigma_{jk}^{\delta\nu}, \tag{15}$$

the transmission coefficient for the exit channel T_k^ν in Eq. 13 has to be replaced by the total transmission coefficient T_k as defined in Eq. 14.

To calculate the cross section σ_{jk} in an astrophysical plasma, reactions on thermally populated states δ in the initial nucleus have also to be taken into account

$$\sigma_{jk} = \frac{\sum_\delta (2J_i^\delta+1)\exp(-E_i^\delta/kT)\sigma_{j,k}^\delta}{\sum_\delta (2J_i^\delta+1)\exp(-E_i^\delta/kT)}. \tag{16}$$

The critical parameters for calculating reliably the Hauser-Feshbach cross sections for proton-capture processes are the nuclear masses, which determine the proton separation energies, the transmission coefficients for the particle and γ reaction channels and the level densities of the initial and final nuclei. The cross sections were calculated using the code SMOKER [56,87,88]. The individual transmission coefficients for the particle channels, $T_m^\lambda(E,J,\pi,E^\lambda,J^\lambda,\pi^\lambda)$ are determined by solving the Schrödinger equation with an optical potential for the particle nucleus interaction. In the case of capture reactions, the γ transition coefficients are dominated by E1 and M1 dipole transitions,

$$T_\gamma^\lambda(E,J,\pi,E^\lambda,J^\lambda,\pi^\lambda) = T_{\mathrm{E1}}^\lambda(E,J,\pi,E^\lambda,J^\lambda,\pi^\lambda) + T_{\mathrm{M1}}^\lambda(E,J,\pi,E^\lambda,J^\lambda,\pi^\lambda), \tag{17}$$

where the final states λ are also states of the compound nucleus. The E1 transmission coefficients are calculated in the frame work of the Lorentzian representation of the Giant Dipole Resonance (GDR). The energy of the GDR is obtained from droplet model predictions taking into account the quadrupole deformation [50,89]. The width of the GDR can be expressed in a phenomenological approximation including the coupling to quadrupole surface vibrations. This approach and the comparison with the experimental data is discussed in more detail in ref. [56,87]

The γ-transition coefficients for M1 magnetic dipole transition are derived from a simple single particle model [56,87].

The level density in the involved nuclei, $\rho(E,J,\pi)$ is derived in terms of the backshifted excitation energy $U=E-\delta$ using the backshifted Fermi gas model [90] which describes the level density in terms of the level density parameter a and the spin cut-off parameter σ,

$$\begin{aligned} \rho(U,J,\pi) &= \frac{1}{2}\rho(U)\cdot f(U,J,\pi), \\ \rho(U) &= \frac{1}{\sqrt{2}\sigma}\frac{1}{12\cdot a^{1/4}}\frac{e^{2\sqrt{aU}}}{U^{5/4}}, \\ f(U,J) &= \frac{2J+1}{2\sigma^2}e^{(-J(J+1)/2\sigma^2)}. \end{aligned} \tag{18}$$

The backshift δ representing the energy to break up the first nucleon pair in order to form an excited state can be best reproduced by the pairing corrections $\Delta(Z,N)$ derived from the droplet mass model. Comparison with experimental data yields, $\delta = \Delta(Z,N) - 10/A$.

For a spherical nucleus the spin cut-off parameter σ can be expressed in terms of the moment of inertia for a rigid body $\Theta = (2/5)m_u A R^2$,

$$\sigma^2 = (\Theta/\hbar^2)\sqrt{U/a}. \tag{19}$$

However, in the region of nuclei discussed here significant deformations occur. Axial deformation was therefore taken into account by modifying the moment of inertia. The spin cut-off parameter then depends on the quadrupole deformation paramter ϵ_2

$$\sigma^2 = (\Theta/\hbar^2)(1+\epsilon_2/3)^{2/3}(1-2\epsilon_2/3)^{1/3}\sqrt{U/a}. \tag{20}$$

The level density parameter a can be obtained from the shell correction term $E_{\rm mic}$ [56,90,91],

$$a = A \cdot (c_0 + c_1 E_{\rm mic}). \tag{21}$$

The parameters c_0 and c_1 are determined from a best fit to experimentally observed level densities. This approach was used in connection with the Hilf et al. mass model. Another improved method for the calculation of a was described by Ignatyuk [91,92]. There, the level density parameter a is directly related to the excitation energy U and the shell correction energy $E_{\rm mic}$,

$$a(U) = a^* \left[1 + \frac{E_{\rm mic}}{U}(1 - e^{-\gamma U})\right] \tag{22}$$

with $a^* = \alpha A(1+\beta A^{-1/3})$ as asymptotic level density parameter at high excitation energies, which is expressed by a least square fit to average level spacing data as a function of mass number A. The fit parameters α and β range between 0.052–0.068 and 3.9–7.9, respectively [91]. The damping parameter, γ is expressed by the empirical relation, $\gamma = 0.4/A^{1/3}$. This method was used together with the FRDM (1992) mass model. A detailed analysis of the procedure is given in a forthcoming paper [88]. It should be noticed that an expansion of $e^{-\gamma U}$ for small excitation energies U gives an expression identical to Eq. 21 for $\beta = 0$, while for large U the asymptotic a^* is obtained. The energy dependence of a is a strong improvement over the prior usage of Eq. 21.

The predicted level densities of the compound nuclei after proton-capture at the here relevant temperatures above 1 GK are generally high enough to justify the statistical model approach. The critical quantity is the number of states in the compound nucleus within the excitation energy range defined by the Gamow window. Fig. 1 shows the level densities based on the FRDM (1992) mass model together with Eq. 22 and on the Hilf et al. mass model with Eq. 21. It displays the number of states available in the Gamow window for proton-capture at $T = 1$ GK for the compound nuclei in the mass range considered here. Fig. 1 shows that at the temperatures above 1 GK considered here, the number of states in the proton capture Gamow window for all displayed nuclei is at least 5–10, which is considered to be the lower limit for the application of the Hauser-Feshbach method. The situation for α-capture is similar. As an extreme example we show in Fig. 13 the ^{64}Ge(p,γ)^{65}As reaction rate that plays an important role in the rp-process. Level density calculations indicate only 7 levels in the Gamow window at 1 GK. However, for temperatures below 2 GK the Hauser-Feshbach reaction rate agrees very well with the rate calculated from the 3 known resonances in ^{65}As [4]. At higher temperatures the experimental information on states in ^{65}As might be too limited. The only nucleus with an even lower level density is ^{82}Mo (5 levels), which is discussed in section III C.

The Hauser-Feshbach calculations for proton and α induced reactions in the mass range above $Z = 32$ have been performed using both, the Hilf et al. model [55] and the FRDM (1992) mass model [50] for the calculation of nuclear masses, deformations and level densities. For the calculations with the Hilf et al. mass model the level density parameter a was derived using Eq. 21. With the FRDM (1992) mass model the level density parameter was calculated with Eq. 21 as well as with the improved Eq. 22 to determine the influence on the reaction rate predictions. As an example, Fig. 13 shows the calculated rates for the ^{81}Zr(p,γ)^{82}Nb, the ^{89}Ru(p,γ)^{90}Rh and the ^{83}Nb(p,α)^{40}Zr reactions. The improved level density treatment leads to substantial differences in the reaction rates. The improved rate is in the case of ^{81}Zr(p,γ)^{82}Nb about a factor of 5 higher, in the case of ^{89}Ru(p,γ)^{90}Rh a factor of 2 smaller at temperatures above 2–3 GK. The ^{83}Nb(p,α)^{80}Zr reaction rate shows discrepancies of up to an order of magnitude at temperatures of several GK. This clearly demonstrates the importance of an adequate level density treatment in Hauser-Feshbach calculations. Fig. 13 shows also the differences between the reaction rates calculated with the Hilf et al. and with the FRDM (1992) mass model. The (p,γ) rates agree well for temperatures below 1–3 GK while for higher temperatures discrepancies reach up to a factor of 4–20. Differences in the case of the ^{83}Nb(p,α)^{80}Zr reaction are even larger. Here, the rate based on the FRDM (1992) mass model is 1–2 orders of magnitude larger over the whole temperature range. The reason for this is the low α-binding energy predicted by the FRDM (1992) mass model in this region of high deformation (as discussed in section II C). This shows that the uncertainties introduced by the choice of mass model can well be larger than typical uncertainties of the statistical model of about a factor of 2–3 [56]. A mass accuracy of at least 500 keV is necessary to keep mass induced uncertainties below a factor of 2, as can be seen in Fig. 14, which shows the ^{81}Zr(p,γ)^{82}Nb reaction rate as a function of reaction Q-value.

Fig. 15 shows the ^{81}Zr(p,γ)^{82}Nb reaction rate for a temperature of 1.5 GK as a function the quadrupole deformation of target and residual nucleus. The dependance on the residual nucleus deformation is very small and symmetric with respect to prolate and oblate shapes. On the other hand there is a change in the reaction rate of an order of magnitude between high prolate and high oblate deformation of the target nucleus. This illustrates how important it is to have accurate predictions of nuclear shapes in regions of strongly deformed nuclei.

B. Inverse reaction rates

Inverse reactions are reactions with negative Q-value. The reactions of this type which are most important for the rp-process are photodisintegration processes, mainly of the type (γ,p) and (γ,α). In the following we describe the calculations of the photodisintegration of a nucleus i into nucleus j with emission of a proton $i(\gamma,\mathrm{p})j$. Other photodisintegration processes as well as particle induced inverse reaction rates were calculated the same way.

The reaction rates for the photodisintegration process, $\lambda_{i(\gamma,\mathrm{p})j}$ can be derived from the inverse proton-capture rate $<\sigma v>_{j(\mathrm{p},\gamma)i}$ using the detailed balance principle. The rate depends also on the Q-value $Q_{j(\mathrm{p},\gamma)}$ for the $j(\mathrm{p},\gamma)i$ reaction and can be expressed as

$$\lambda_{i(\gamma,\mathrm{p})j} = \frac{(2J_\mathrm{p}+1)G_j}{G_i}\left(\frac{\mu kT}{2\pi\hbar^2}\right)^{3/2} \cdot \exp\left(-\frac{Q_{j(\mathrm{p},\gamma)}}{kT}\right) \cdot <\sigma v>_{j(\mathrm{p},\gamma)i} . \tag{23}$$

The normalized partition functions G_i and G_j account for the thermal population of the excited states in initial i and final nucleus j. They are defined as usual, for example for nucleus i with states δ and corresponding excitation energies E_x^δ

$$\begin{aligned} G_i &= \sum_\delta^\omega g_i^\delta \exp\left(-\frac{E_x^\delta}{kT}\right) \\ &+ \int_{E_\omega} \int_{J_i,\pi_i} \exp\left(-\frac{E}{kT}\right) \rho_i(E,J,\pi)\, dE\, dJ\, d\pi \end{aligned} \tag{24}$$

with g^δ=2J_δ+1 as the statistical weight of the state δ with spin J_δ. The first part of the equation accounts for the contribution of known discrete states up to ω, the second part accounts for the thermal population of the excitation range above ω with high level density $\rho(E,J,\pi)$.

Due to the level density dependence of the partition functions the photodisintegration rate is also level density dependent. The dominant parameter, however, is the reaction Q-value, on which the reaction rate depends exponentially. Small deviations in the separation energies will therefore lead to considerable changes in the photodisintegration rate. It is here were nuclear masses play the most important role.

For the rp-process network calculation (γ,p)-reactions are the most important photodisintegration processes. Fig. 16 shows the ratio of the photodisintegration rate and the capture rate

$$\frac{\lambda_{i(\gamma,\mathrm{p})j}}{\lambda_{j(\mathrm{p},\gamma)i}} = \frac{\lambda_{i(\gamma,\mathrm{p})j}}{<\sigma v>_{j(\mathrm{p},\gamma)i} N_A \rho Y_\mathrm{p}} \tag{25}$$

as a function of temperature for various proton-capture Q-values. The shown curves were calculated for typical rp-process conditions with a density $\rho = 10^6$ g/cm^3 and a proton abundance Y_p=0.771 and assuming $G_j = G_i = 1$. They can easily be scaled to different densities. It can be seen that at temperatures of 1–2 GK photodisintegration is dominating for proton-capture Q-values below 1 MeV.

The other photodisintegration process that might play a role in the rp-process is the photoinduced α-emission. This takes place if the α-threshold in proton bound nuclei is low. A relatively low α-threshold, $S_\alpha \leq 1.5$ MeV, is predicted by the FRDM (1992) mass model for ^{80}Sr, ^{82}Zr, and ^{84}Nb. The model predicts also that ^{84}Mo is even slightly α-unbound, $S_\alpha \approx -0.5$ MeV (see section II C). Fig. 17 shows for comparison the ratio between the α-photodisintegration rate and the proton-capture rate as a function of temperature for the above listed nuclei. The proton-capture rate was calculated for the same typical proton density conditions as for Fig. 16. The ratio is shown for the rates based on the mass predictions by Hilf et al. [55] as well as by the FRDM (1992) mass model [50]. The lower α-separation energies of the FRDM (1992) mass model result in a higher rate for photoinduced α-emission compared to the Hilf et al. model, but the (p,γ) branch is still dominant for the temperature and density conditions ($\rho \geq 10^4$ g/cm^3, 1 GK $\leq T \leq$ 2 GK, see section IV) considered here.

C. 2p-capture reactions and particle decay

A limiting parameter for the reaction path of the rp-process is the proton drip line (S_p = 0). Fig. 9 shows the location of the proton drip line as predicted by the various mass models. In all earlier calculations the consideration of proton-capture processes had been limited to particle stable isotopes only. For proton unstable isotopes ($S_\mathrm{p} \leq 0$) immediate proton-decay had been assumed and no further proton-capture had been calculated. Further processing therefore depended on the β-decay of the last proton stable isotone, the so called waiting point nucleus (see ref. [4] and section IV B for a detailed discussion). In this picture the rp-process path cannot cross the proton drip line. This represents certainly an oversimplified view. The lifetime of a proton unstable nucleus $\tau = 1/\lambda_\mathrm{p} = \hbar/\Gamma_{\mathrm{p}_0}$, can be appreciably long due to the Coulomb barrier, especially close to the proton drip line, where the proton is only weakly unbound. This was pointed out before in the case of ^{65}As [93]. At high temperatures however, the proton break-up increases due to photo induced proton emission via the excited states in the proton unbound nucleus, $\tau = \lambda_\mathrm{p}^{-1} + \lambda_{<\gamma,\mathrm{p}>}^{-1}$.

Because of the finite lifetime of the proton unbound nuclei, 2p-capture reactions on the last proton bound isotone with even Z are possible. They allow to bridge the single odd Z proton unstable nucleus finally producing an even Z proton bound nucleus. These processes are possible due to the pronounced difference in the proton drip line between odd and even Z isotopes, that can be seen for example in Fig. 1. 2p-capture processes have been discussed before in the light mass region $A \leq 40$ [94]. It had been shown that at high temperature and density conditions the 2p-capture processes can be considerably faster than the competing β-decay of the target nucleus. In this case the reaction path would run across the proton drip line.

Two proton-capture on the last proton bound isotone (Z,N) can be described as proton-capture on the intermediate proton unbound nucleus $(Z+1,N)$, which is produced by resonant scattering of protons in the stellar plasma with an energy above the threshold. The time derivative of the abundance of the $(Z+1,N)$ nucleus in a state ν, $Y_\nu(Z+1,N)$ can be expressed in terms of the possible production and destruction processes

$$\begin{aligned}\frac{dY_\nu(Z+1,N)}{dt} &= Y(Z,N)Y_\mathrm{p}\rho N_A <\mathrm{p},\mathrm{p}>_\nu - Y_\nu(Z+1,N)\lambda_\mathrm{p}^\nu - Y_\nu(Z+1,N)\lambda_{(\gamma,\mathrm{p})}^\nu \\ &- Y_\nu(Z+1,N)\lambda_\beta^\nu - Y_\nu(Z+1,N)Y_\mathrm{p}\rho N_A <\mathrm{p},\gamma>_\nu . \end{aligned} \tag{26}$$

The reaction rates on the right side of the equation are the spontaneous proton-decay rate λ_p^ν, the photo-induced proton-decay rate $\lambda_{(\gamma,\mathrm{p})}^\nu$, and the β-decay rate λ_β^ν of the state ν , as well as the production rate via resonant proton-capture on the particle stable nucleus (Z,N), $\rho N_A <\mathrm{p},\mathrm{p}>_\nu$, and the depletion rate via proton-capture on the particle unstable nucleus $(Z+1,N)$, $\rho N_A <\mathrm{p},\gamma>_\nu$. Y_p is the proton abundance while $<$p,p$>$ and $<$p,$\gamma>$ denote the stellar reaction rates $<\sigma v>$ for proton scattering and proton-capture (p,γ), respectively. Direct α-decay processes have not been included. While the FRDM (1992) mass model predicts some α unbound isotopes in this mass range, the high Coulomb barrier will prevent α-decay ($S_\alpha \geq -1.0$ MeV).

The 2p-capture rate can be approximated assuming that the proton-decay processes of the particle unstable nucleus are faster than its proton-capture and β-decay rate [94,95]. Within the time scale of the proton-scattering process, an equilibrium abundance for the proton unbound nucleus in state ν, $Y_\nu(Z+1,N)$ is reached:

$$Y_\nu(Z+1,N) = \rho Y(Z,N)Y_\mathrm{p}\frac{N_A <\mathrm{p},\mathrm{p}>_\nu}{(\lambda_\mathrm{p}^\nu + \lambda_{(\gamma,\mathrm{p})}^\nu)}. \tag{27}$$

The total reaction rate r (in numbers of reactions per time) for the 2p-capture is then determined by the proton-capture rate on this equilibrium abundance summed over all states ν in nucleus $(Z+1,N)$

$$r = \sum_\nu \rho^2 N_A^2 Y_\nu(Z+1,N)Y_\mathrm{p} <\mathrm{p},\gamma>_\nu . \tag{28}$$

If we define the stellar 2p-capture reaction rate $N_A^2 <2\mathrm{p},\gamma>$ by writing r as

$$r = \frac{1}{2}\rho^3 N_A^3 Y(Z,N) Y_{\rm p}^2 < 2{\rm p}, \gamma >, \tag{29}$$

we get for the 2p-capture rate

$$N_A^2 < 2{\rm p}, \gamma >= \sum_\nu 2N_A^2 \frac{< {\rm p}, {\rm p} >_\nu < {\rm p}, \gamma >_\nu}{\lambda_{\rm p}^\nu + \lambda_{(\gamma,{\rm p})}^\nu}. \tag{30}$$

$\lambda_{(\gamma,{\rm p})}^\nu$, the photo induced proton-decay rate of the proton unbound nucleus in the state ν via a state λ (in the proton unbound nucleus as well) into the ground-state of nucleus (Z,N) can be calculated using the detailed balance principle:

$$\lambda_{(\gamma,{\rm p})}^\nu = \sum_\lambda \frac{2J_\nu + 1}{2J_\lambda + 1} \cdot \frac{\Gamma_\gamma^{\lambda\nu}\Gamma_{\rm p}^{\lambda 0}}{(\Gamma_\gamma^\lambda + \Gamma_{\rm p}^\lambda)} \hbar^{-1} \cdot \exp(-E_{\nu\lambda}/kT). \tag{31}$$

$E_{\nu\lambda}$ is the energy difference between the states ν and λ ($E_{\nu\lambda} = E_{x\lambda} - E_{x\nu} > 0$) in the proton unbound nucleus $(Z+1,N)$. J_λ, Γ_γ^λ and $\Gamma_{\rm p}^\lambda$ are the spin, the γ width and the proton width of the intermediate states λ. The second index for the decay widths indicates the partial widths for the decay into the corresponding state. Typically this term will be smaller than the rate for the proton-decay, in the case of large excitation energies due to the exponential term and in the case of low excitation energies due to the small proton-decay width of the excited state $\Gamma_{\rm p} \ll \Gamma_\gamma$. More accurate predictions are not possible due to the lack of detailed information especially on the γ-widths of excited states in the nuclei of interest here. We therefore neglect photoinduced proton-decay for the calculations of the 2p-capture reaction rates. For similar reasons we neglect γ-decay of a level ν in nucleus $(Z+1,N)$, before it either proton-decays or captures another proton.

The proton unstable nucleus is produced via resonant proton-capture $< {\rm p}, {\rm p} >_\nu$ on nucleus (Z,N). To calculate the reaction rate in a stellar plasma we have to take into account that low lying states in the target nucleus are thermally populated. The production rate of nucleus $(Z+1,N)$ in the state ν via resonant proton-capture on the target in the state δ can be calculated [94] using

$$< {\rm p}, {\rm p} >_{\delta\nu} = \left(\frac{2\pi}{\mu kT}\right)^{\frac{3}{2}} \hbar^2 \frac{2J_\nu + 1}{(2J_{\rm p}+1)(2J_\delta+1)} \cdot \Gamma_{\rm p}^{\nu\delta} \exp\left(-\frac{E_{\delta\nu}}{kT}\right), \tag{32}$$

where $J_{\rm p}$ is the proton spin, μ the reduced mass in the entrance channel, $\Gamma_{\rm p}^{\nu\delta}$ the width of state ν for the proton-decay into the daughter state δ, and $E_{\delta\nu}$ the energy difference between nucleus $(Z+1,N)$ in state ν and nucleus (Z,N) in state δ. $E_{\delta\nu}$ can be expressed in terms of the (negative) proton-capture Q-value on nucleus (Z,N) and the excitation energies of states δ and ν, E_x^δ and E_x^ν:

$$E_{\delta\nu} = -Q_{({\rm p},\gamma)(Z,N)} + E_x^\nu - E_x^\delta. \tag{33}$$

The proton-decay rate of nucleus $(Z+1,N)$ in the state ν, $\lambda_{{\rm p}\nu}$ in Eq. 30, can be expressed in terms of the proton widths for the decay into the accessable states δ of nucleus (Z,N)

$$\lambda_{\rm p}^\nu = \sum_\delta \frac{\Gamma_{\rm p}^{\nu\delta}}{\hbar}. \tag{34}$$

Inserting expressions 32 and 34 into Eq. 30, and neglecting photo induced proton-decay we obtain for the 2p-capture rate on a target in state δ

$$N_A^2 < 2{\rm p}, \gamma >_\delta = 2N_A^2 \left(\frac{2\pi}{\mu kT}\right)^{\frac{3}{2}} \hbar^3 \sum_\nu \frac{2J_\nu + 1}{(2J_{\rm p}+1)(2J_\delta+1)} \frac{\Gamma_{\rm p}^{\nu\delta}}{\sum_{\delta'} \Gamma_{\rm p}^{\nu\delta'}} \exp\left(-\frac{E_{\nu\delta}}{kT}\right) < {\rm p}, \gamma >_\nu . \tag{35}$$

In a stellar plasma the excited states in the target nucleus are thermally populated and the stellar 2p-capture reaction rate is therefore

$$N_A^2 < 2{\rm p}, \gamma >= N_A^2 \frac{\sum_\delta (2J_\delta + 1) < 2{\rm p}, \gamma >_\delta \exp\left(-\frac{E_x^\delta}{kT}\right)}{\sum_\delta (2J_\delta + 1) \exp\left(-\frac{E_x^\delta}{kT}\right)}, \tag{36}$$

where E_x^δ is the excitation energy of state δ. If we insert Eqs. 35 and 33 we finally obtain for the stellar 2p-capture reaction rate

$$N_A^2 < 2\mathrm{p},\gamma >= 2N_A^2 \left(\frac{2\pi\hbar^2}{\mu kT}\right)^{\frac{3}{2}} \frac{G_{(Z+1,N)}}{(2J_\mathrm{p}+1)G_{(Z,N)}} \exp\left(\frac{Q_{(\mathrm{p},\gamma)(Z,N)}}{kT}\right) < \mathrm{p},\gamma >, \tag{37}$$

where $< \mathrm{p},\gamma >$ is the stellar proton-capture reaction rate on the thermally excited nucleus $(Z+1,N)$. In Eq. 37 we used the nuclear partition functions G as defined in Eq. 24.

Eq. 37 does not depend on the proton widths of the states involved and demonstrates the validity of the Saha equation in this case. Note also that Eq. 37 is identical to the equation one obtains assuming a positive proton-capture Q-value on nucleus (Z,N) and an equilibrium between nuclei (Z,N) and $(Z+1,N)$ via (p,γ)- and (γ,p)-reactions calculated using the detailed balance principle. It is therefore valid for negative and positive Q-values as long as equilibrium between target and intermediate nucleus is established. No other assumptions on γ- or proton widths are made, except for neglecting $\lambda_{(\gamma,\mathrm{p})}$ in Eq. 30, when the proton-capture Q-value is negative.

The proton-capture rate on the proton unstable nucleus $N_A < \mathrm{p},\gamma >$ was calculated using the Hauser-Feshbach model (see section III A). In the case of the 2p-capture rate on ^{80}Zr, the low level density in the final nucleus ^{82}Mo (see Fig. 1) might lead to an overestimated 2p-capture rate, when the Hauser-Feshbach model is applied. However, the uncertainties in the 2p-capture rates are at this point dominated by the experimentally unknown Q-value. The proton separation energies were taken from the FRDM (1992) mass model in all cases. Tab. I lists the calculated 2p-capture rates on the even-even $N=Z$ nuclei between $A=68$ and $A=96$.

The inverse reaction rate for 2p-capture reactions is given by the (γ,p)-rate on the final nucleus. The rate $\lambda_{i(\gamma,2\mathrm{p})}$ of the inverse 2p-capture reaction $i(\gamma,2\mathrm{p})k$ via the intermediate proton unbound nucleus j can then be written in terms of the forward reaction rate $< 2\mathrm{p},\gamma >$ using Eqs. 37 and 23:

$$\lambda_{(\gamma,2\mathrm{p})} = \left[\left(\frac{kT}{2\pi\hbar^2}\right)^{\frac{3}{2}}\right]^2 (\mu_k\mu_j)^{\frac{3}{2}} \frac{(2J_\mathrm{p}+1)^2 G_k}{G_i} \exp\left(-\frac{Q_{k(\mathrm{p},\gamma)}+Q_{j(\mathrm{p},\gamma)}}{kT}\right) \frac{1}{2} < 2\mathrm{p},\gamma >, \tag{38}$$

where μ_k and μ_i are the reduced masses for a proton and nucleus k and i respectively. This is similar to Eq. 23 for single proton-capture and expresses the detailed balance principle for 2p-capture rates.

D. β-decay rates

Another important input parameter for the rp-process reaction network studies are the β^+-decay half-lives of the neutron-deficient nuclei. For the lighter isotopes ($Z \leq 34$), a wide range of experimental ground-state lifetime data are available. For the β-decays of isotopes heavier than $A=20$ at higher stellar temperature conditions, the decay from thermally populated excited levels has to be taken into account. The decay rate becomes then a function of temperature. For the neutron-deficient isotopes up to $Z=30$ and mass $A=60$ those temperature dependent decay rates were taken from Fuller, Fowler, and Newman [96–98]. The enhancement of the decay rate due to continuum electron capture at high stellar densities included in the Fuller, Fowler, and Newman data, is negligible for the density range up to about 10^6 g/cm^3 discussed in this study.

For the extension of the reaction network beyond $Z=32$, $A=64$ the β-decay half-lives of nuclei close to the proton drip line are especially important (see paragraph IV). Fig. 1 shows that for the majority of these nuclei no experimental data are available. Therefore, the ground-state β-decay rates as well as the decay rates from thermally populated excited states have to be calculated. Due to the high Q_β values for nuclei far off the line of stability, the β^+-decays are dominated by Gamow-Teller transitions. For very neutron-deficient nuclei with $Z>N$ the superallowed Fermi transition to the isobaric analog state (IAS) has to be taken into account in principle. However, IAS energies are typically so high, that the lower Q-value does not allow Fermi transitions to compete with the Gamow-Teller decays to low lying states. We therefore limited our discussion to β^+-decay via Gamow-Teller transitions.

The decay rate from the ground-state in the parent nucleus to the ground- and excited states in the daughter is expressed in terms of the β-strength S_β as a function of excitation energy E by

$$\lambda_i = \frac{G_v^2 m_e^5 c^4}{2\pi^3\hbar^7} g_A^2 \int_0^{Q_\beta} S_\beta(E) f_0(Z,Q_\beta - E)dE. \tag{39}$$

In this equation G_v represents the vector coupling constant, g_A the ratio of the axial vector coupling constant to G_v, and $f_0(Z,Q_\beta - E)$ the Fermi function. The β-strength function describes the energy dependence of the transition probability and is expressed in terms of the reduced transition probability $B(E,J,\pi)$ to a final state at excitation energy E with spin and parity J,π, weighted with the corresponding level density

$$S_\beta(E) = \sum_{J,\pi} B(E,J,\pi)\rho(E,J,\pi). \tag{40}$$

With these definitions $B(E,J,\pi)$ represents here the reduced transition matrix element. We have calculated the Gamow-Teller strength functions within the quasiparticle random phase approximation (QRPA). The theoretical basis and the model parameters are the same as in the FRDM (1992) mass model: A folded-Yukawa potential is used to calculate the wave functions for the initial and final states in the mother and daughter nuclei. Pairing and shell correction terms are also identical to the ones used in the FRDM (1992) mass model. A detailed discussion of the model parameters can be found in refs. [46,99,100]. As most other models, the present approach makes the assumption of identical shapes for the ground-state of the mother nucleus and for all populated states in the daughter nucleus. Hence, β-decays between isobars of different nuclear shape, e.g. an oblate mother decaying to a prolate daughter as for ^{82}Mo$(\beta^+\nu)^{82}$Nb, as well as cases of shape coexistence in the final nucleus as discussed by Ryde [101] can only be simulated in specific cases but not treated in a selfconsistent way. It has been shown in the previous chapter that the mass region discussed here is characterized by strong deformation and drastic changes in shape along an isotopic chain. Nevertheless, the QRPA model represents a valuable microscopic approximation for calculating the β-strength function. Moreover, it is based on the same single particle potential, the same pairing terms and the same shell correction terms as the FRDM (1992) mass model. In contrast to other models it represents in conjunction with the FRDM (1992) mass model a unified approach for the description of all nuclear physics parameters.

For some nuclei near the $N = Z$ line along the main reaction path (see section IV), the Gamow-Teller strengths have been calculated using the shell model code OXBASH [102]. For the nuclei with $56 \leq A \leq 72$ we used a model space consisting of the subshells $1f_{5/2}$, $2p_{3/2}$ and $1p_{1/2}$ with an inert ^{56}Ni core. We employ the MSDI interaction developed by Koops and Glaudemans [103] for this model space. In the mass range $84 \leq A \leq 100$ we use a model space consisting of an inert ^{88}Sr core with valence proton and neutron holes limited to the $2p_{1/2}$ and the $1g_{9/2}$ subshells. In principle this model space stretches from ^{76}Sr up to ^{100}Sn. However, in the lower half of this mass range the influence of the $1f_{5/2}$ and the $2p_{3/2}$ subshells is considerable. Therefore we only calculate nuclei with $A \geq 84$ with this model space. The deformations are small (except for ^{84}Mo) so no significant contributions from the higher lying $\pi 1g_{7/2}$ and $\nu 2d_{5/2}$ subshells are expected. However, the higher subshells can renormalize the Hamiltonian and the GT operator. For the calculations of the Gamow-Teller strengths we used the "$T = 0$ fit" interaction described in [104] and the free Gamow-Teller operator. The results might be improved in the future using effective Gamow-Teller operators. The interactions are fitted to the spectra of experimentally known nuclei. Therefore the application of the shell model calculations for very neutron deficient nuclei represents an extrapolation. The reliability of this extrapolation as well as other details of the procedure will be discussed in a separate paper [105].

As example, Fig. 18 shows Gamow-Teller strengths resulting from the QRPA calculations as a function of the excitation energy in the final nucleus for the β-decay of the even-even nuclei ^{76}Sr, ^{80}Zr, ^{84}Mo, ^{88}Ru, ^{92}Pd, and ^{96}Cd. For the latter three cases reliable shell model calculations are available and the corresponding Gamow-Teller strength functions are shown in Fig. 19 for comparison. The shell model strength functions are very similar to the QRPA results for the low excitation energies that play a role for the half-life calculation. In the case of the ^{96}Cd decay, the low lying QRPA Gamow-Teller strength is slightly more spread out. The giant Gamow-Teller resonance that can be seen in the strength functions calculated with the QRPA is missing in the shell model results, since it involves the 1g9/2 and 1g7/2 excitations that are not included in the shell model space. To explore the sensitivity of the β-strength function to the nuclear deformation, QRPA calculations were performed for various ground-state deformations. Fig. 20 shows the β-strength function for the decay of ^{80}Zr, calculated for a spherical nucleus as well as for a quadrupole deformation parameter ϵ_2 of -0.383, which corresponds to an oblate instead of a prolate shape as predicted (in this case ϵ_4 and ϵ_6 were kept at the values predicted by the FRDM (1992) mass model). There is a considerable change in the strength distribution, comparing the spherical case with the predictions for a deformed nucleus. In the case of deformation the dominant part of the strength clearly resides in the transitions to higher excited states, and the strength is distributed over a much wider range of excitation energies. This qualitative behavior is relatively independent on whether the deformation is prolate or oblate. This is in qualitative agreement with a previous study using the simpler Tamm-Dancoff approximation for the same nucleus [106].

Since the Fermi function $f_0(Z, Q_\beta - E)$ in Eq. 39 strongly depends on the excitation energy,

$$f_0(Z, Q_\beta - E) \propto (Q_\beta - E)^5. \tag{41}$$

the β-decay rate is dominated by transitions to the ground-state or to levels with low excitation energy (typically below 2–3 MeV). Eqs. 39 and 41 also demonstrate the strong dependence of β-decay rate calculations on the adopted Q-value. However, in all the cases where the β-decay half-life has to be calculated no experimental information on Q_β or nuclear shapes are available. To explore the uncertainties in the half-life predictions related to the deformation of the initial nucleus and to Q_β, the half-lives of the even-even nuclei ^{80}Zr and ^{84}Mo have been calculated for the entire range of deformation between oblate, $\epsilon_2 = -0.5$ and prolate shape, $\epsilon_2 = +0.5$; in addition the half-lives are calculated for Q_β values ranging from 4.0 to 9.0 MeV. Fig. 21 shows that due to the uncertainty in deformation the calculated half-life varies typically within a factor of 3. This is in good agreement with the results of an earlier study of β-decay half-lives of neutron deficient isotopes in the mass range $A = 68$–82 [107]. Note that for ^{80}Zr the dependency of the half-life on ϵ_2 is quite steep at the predicted ϵ_2 value, which makes in this case the β-decay half-life calculations quite sensitive to deformation. The choice of the Q_β value influences the half-life prediction severely due to the energy dependence of the Fermi function. This is demonstrated in Fig. 22 for the β-decays of ^{80}Zr and ^{92}Pd. Over the range of the various mass predictions the half-life varies by a factor of 2 for ^{92}Pd and a factor of 5 for ^{80}Zr. Therefore, the half-lives predicted here depend strongly on the choice of the mass model. The shell model results depend on Q_β in the same way.

Fig. 23 shows the half-lives as calculated for the even Z isotopic chains from $Z = 38$ to $Z = 48$ plotted as functions of mass number A. All the predictions, the QRPA calculations as well as the shell model calculations, are based on the Q_β values resulting from the FRDM (1992) mass model discussed in section II. Also shown are previous model predictions in this mass range [107,108]. which are also based on the QRPA approach but use the simpler Nilsson potential and treat pairing in the BCS approximation, which might be problematic near closed shells [46]. These calculations show an overall agreement with the results of the present work.

For comparison Fig. 23 shows also the available experimental half-lives. Since we calculated only allowed Gamow-Teller transitions, the comparison should be limited to half-lives smaller than 100 s. The QRPA model predicts those nuclei within a factor of 5.5. The shell model leads here to a considerable improvement with an average deviation of a factor of 2.7. For the subset of nuclei for which shell model calculations are available the QRPA predictions agree on average within a factor of 5.9, which is slightly worse than the overall deviation. The fact that these nuclei are only weakly deformed indicates that deformation is not a significant source of uncertainty in the QRPA predictions. It clearly would be desirable to develop a model space allowing the calculation of all half-lives within the shell model approach.

Calculating β-decay processes in high temperature scenarios requires to take into account the β-decay of thermally excited states. It has been shown before [46] that the β-decay lifetime of excited states can be significantly different from the lifetime of the ground-state. The total temperature dependent β-decay rate $\lambda(T)$ is then expressed in terms of the decay rates of the excited states λ_i and the corresponding excitation energy E_i

$$\lambda(T) = \frac{\sum_i^{E_i<S_\mathrm{p}} \lambda_i \exp\left(-\frac{E_i}{kT}\right)}{\sum_i^{E_i<S_\mathrm{p}} \exp\left(-\frac{E_i}{kT}\right)}. \tag{42}$$

The sum extends over all states up to the proton separation energy S_p. Since the temperature range discussed in this paper (up to $kT = 300$ keV) is relatively low, only a few low lying excited states will be thermally populated. To explore the temperature dependence of the β-decay rates we concentrated on the calculations for the $N = Z$ even-even nuclei listed in Tab. II, which are of particular importance for the rp-process (see section IV). For these nuclei only the population of the first excited 2^+ state plays a role. The experimentally known energies of the lowest 2^+ states ($E(2^+)$) of the nuclei considered here are given in Tab. II. Where no experimental information is available, the energies can be calculated using the fact that they are strongly correlated with nuclear deformation and thus with the product of valence protons and neutrons $N_\mathrm{p}N_\mathrm{n}$ (see also section II B). Fig. 24 shows the excitation energies of the experimentally known first excited 2^+ states in even-even nuclei in the range $28 \leq Z, N \leq 50$ as a function of the corresponding $N_\mathrm{p}N_\mathrm{n}$ value. The scheme works very well for nuclei with $Z > 35$ but shows strong deviations for Zn, Ge and Se. We therefore extrapolated $E(2^+)$ from a fit through the data with $Z > 35$ to calculate $E(2^+)$ for ^{88}Ru, ^{92}Pd and ^{96}Cd (see results in Tab. II). The good agreement between the values obtained for the other nuclei and the experimental data reflects the strong $N_\mathrm{p}N_\mathrm{n}$ correlation (93%). For comparison we show in Tab. II the results from the shell model calculation that assume small deformations. ^{84}Mo is the only strongly deformed nucleus for which a shell model calculation was possible. As expected the discrepancy between the shell model results and the experimental level energy is very large for this nucleus. Also listed in Tab. II, are the experimental and predicted β-decay rates for the ground-states as well as the theoretically predicted rates for the β-decay of the 2^+ excited states as derived from the shell model. For the highly deformed nuclei ^{76}Sr and ^{80}Zr no shell model calculations of β-decay rates from excited states are possible because of the above discussed limitations in the configuration space, but these nuclei have the lowest excitation energies for the first 2^+ state and should therefore show the strongest effect. The only cases where the β-decay half-life from the first excited 2^+ state is predicted to be significantly different to the one from the ground-state are ^{64}Ge, ^{68}Se and ^{72}Kr. The total β-decay rates $\lambda_i(T)$ for these nuclei are shown in Fig. 25 as functions of temperature. For temperatures below 2 GK the rates stay constant and around temperatures of 3 GK the effect is less than approximately 15% and thus negligible. Since 3 GK is about the maximum temperature occuring in the scenarios of this study, we did not include this temperature dependency in our network calculations. However, for temperatures above 3 GK temperature dependence of β-decay half-lives has to be considered.

IV. NETWORK CALCULATIONS

The main purpose of this paper is to determine the influence of the various nuclear structure input parameters (as discussed in the previous sections) on the rp-process nucleosynthesis for $Z \geq 32$. The rp-process for $Z < 32$ has been discussed extensively in the work of Van Wormer et al. [4]. To analyze this influence in a clear fashion, we performed our calculations with constant temperatures and densities. In X-ray bursts, the peak temperature at around 3 GK is reached within 1–10 s. Subsequent cooling due to expansion occurs on timescales of 10–100 s. As we will show in the following, there is only a rather narrow temperature window of $T = 1$–1.8 GK (at densities around $\rho = 10^6$ g/cm^3) where rp-process nucleosynthesis in the $Z > 32$ range is faster than the rather short event timescale. Processing at lower temperatures is hampered by the Coulomb barrier, at higher temperatures by photodisintegration (see section IV C). A calculation with constant temperature of 1.5 GK and a density of 10^6 g/cm^3 should therefore give results similar to calculations using temperature and density profiles. In the following we will refer to this choice of temperature and density describing X-ray burst nucleosynthesis as XRB-conditions. Processing times are of the order of the observed X-ray burst timescales, which vary between 10 and 100 s among different burst sources as well as among different bursts from the same source.

In accretion disks around black holes, density and temperature are continuosly increasing with much longer timescales of 10^4–10^5 s while matter accretes onto the black hole (for a temperature profile see for example [10]). Peak densities of 10^4 g/cm^3 and peak temperatures of several GK are possible, depending on the viscosity parameter assumed for the disk. Due to the long timescale, significant processing will occur at lower temperatures and densities. However, the final nucleosynthesis results for nuclei above $Z = 32$ will again be dominated by the reactions taking place during the temperature window for most effective processing at 1–1.8 GK. We therefore will also show results for network calculations with a constant temperature of $T = 1.5$ GK and a constant density of 10^4 g/cm^3 to approximate nucleosynthesis in accreting low mass black holes. In the following we will refer to this temperature and density as ABH-conditions.

In previous investigations of Thorne-Zytkow objects, rp-process nucleosynthesis took place in repeating very short pulses (0.01–0.1 s), separated by very long time intervals of the order of 10^8 s. Constant temperature and density calculations are not appropriate in this case. Since recent theoretical calculations make the existence of Thorne-Zytkow objects very questionable, we refrain in this paper from presenting nucleosynthesis results approximating these scenarios.

All calculations presented here were performed with an initial solar abundance distribution. As we will show later, the choice of seed abundances is not a critical parameter in this study, except for the H/He ratio.

Our nuclear reaction network includes 637 nuclei from H to Sn. Above Sn the rp-process path enters a region of α-unstable nuclei, which will probably terminate the rp-process. The reaction rates for nuclei with $Z \leq 32$ have been discussed extensively by Van Wormer et al. [4] and have been updated with the shell model based (p,γ)-rates of Herndl et al. [82]. Throughout the network, experimentally known rates were used if available. However, the only experimental rates for the reaction path above $Z = 32$ far away from stability are some β-decay rates (see Fig. 1). The experimentally unknown particle and γ-ray induced rates have been calculated with the Hauser-Feshbach code SMOKER as described in section III A. The β-decay rates were calculated in a few cases using the shell model code OXBASH (see Fig. 23), in all other cases using the QRPA as described in section III D.

We used in this study three different sets of reaction and β-decay rates called FRDM1, FRDM2 and HIL. They differ in the theoretically predicted rates above $Z = 32$. FRDM1 and FRDM2 are based on the FRDM (1992) mass model by Möller et al. [50], while HIL is based on the Hilf et al. mass model [55]. The choice of these mass models was justified in section II D. The Hilf et al. mass formula differs considerably from the FRDM (1992) mass model in its predictions for neutron-deficient isotopes. Therefore, the comparison between the network calculations based on FRDM1 or FRDM2 to the ones based on HIL illustrates the influcence of mass model uncertainties on the results. The only difference between FRDM1 and FRDM2 is that FRDM2 does not include the 2p-capture reactions described in section III C. Therefore, comparison between the FRDM1 and FRDM2 based network calculations allows the investigation of the effect of 2p-capture reactions on rp-process nucleosynthesis. The three reaction and β-decay rate sets are described in more detail in the following paragraphs. In addition, the FRDM1 reaction rates are listed in Appendix A.

In all three parameter sets, we used experimental masses where available (see Fig. 1), however, we did not calculate a particle separation energy or a β-decay Q-value from an experimental and a theoretical mass to avoid discontinuities in the mass surface. For FRDM1 and FRDM2, we used the FRDM (1992) mass model to calculate the Hauser-Feshbach rates, the β-decay rates and the 2p-capture rates (FRDM1 only). In two cases that are important for rp-process nucleosynthesis, exceptions were made on the basis of available experimental information.

The first case is the ^{68}Se(p,γ)^{69}Br reaction. The FRDM (1992) mass model predicts a Q-value of 89 keV while Blank et al. [109] estimate from a recent experiment that ^{69}Br might be proton unbound by at least 450 keV. This is in agreement with previous experiments searching for the ground-state proton decay of ^{69}Br that give upper limits on the proton binding energy of about -350 keV [110] and about -400 keV [111]. However, all these experiments rely on assumptions on the production cross section for ^{69}Br in fragmentation reactions. The MSU group had reported the observation of ^{69}Br in Mohar et al. [112], but a second experiment with a similar setup could not confirm that result [113]. We tentatively adopted for ^{68}Se a proton capture Q-value of -450 keV and calculated the proton capture rate assuming 2p-capture (FRDM1 only).

The other case is the ^{81}Zr(p,γ)^{82}Nb reaction, where the FRDM (1992) mass model predicts a ^{81}Zr mass 846 keV above the experimental value (which has an uncertainty of 297 keV) leading to a reaction Q-value of 299 keV. This Q-value is significantly smaller than the predicted values for the other even-odd $N = Z + 1$ nuclei (see Tab. III) in this mass range. Furthermore, comparison of the experimental $Q_{(\mathrm{p},\gamma)}$ values for ^{65}Ge and ^{73}Kr with the FRDM (1992) model predictions (see Tab. III) indicates, that the FRDM (1992) mass model tends to strongly underestimate $Q_{(\mathrm{p},\gamma)}$ values for even-odd $N = Z + 1$ nuclei in this mass range (by about 1.4 MeV in both cases). We therefore adopted for FRDM1 a Q-value of 1.407 MeV obtained from the experimental mass of ^{81}Zr and the extrapolated value by Audi et al. [47] for ^{82}Nb.

For the HIL reaction rates, the Hilf et al. mass model [55] was used to calculate the rates for particle and γ-ray induced reactions. The β-decay rates were the same as in FRDM1 and FRDM2, since we did not adapt our QRPA model to the nuclear structure parameters used in the Hilf et al. mass formula. For the HIL set the older version of SMOKER based on the level density of Eq. 21 was used. 2p-capture reactions can only compete with slow β-decay rates (below $\approx$0.1 s^{-1}). Therefore, they do not have to be included in the HIL reaction data base, since all long lived waiting points can be bridged with single proton capture reactions (see below).

In the following sections, we will discuss the results of four different network runs. We performed three runs at XRB-conditions ($T = 1.5$ GK and $\rho = 10^6$ g/cm^3) using the three different reaction rate sets FRDM1, FRDM2 and HIL as discussed above. In addition we did one calculation at ABH-conditions ($T = 1.5$ GK, $\rho = 10^4$ g/cm^3) based on the reaction rate set FRDM1. We start the discussion in section IV A with a comparison of the obtained rp-process reaction paths in the Ge-Sn region. In section IV B we give a general description of the parameters that determine branching ratios at waiting points. The results are then used for the interpretation of the obtained rp-process time structure (section IV C), the isotopic abundances (section IV D), and the energy production (section IV E).

A. The reaction flow

Fig. 26 shows the reaction flow for all four network runs integrated over a processing time of 1000 s. First, we discuss the reaction paths obtained under XRB-conditions for all three sets of reaction rates, FRDM1, FRDM2 and HIL. In set FRDM1, the reaction path follows mostly the proton drip line. Among the 2p-capture reactions described in section III C only ^{68}Se(2p,γ)^{70}Kr and ^{72}Kr(2p,γ)^{74}Sr play a significant role. They shift a fraction of the reaction path beyond the proton drip line. For $Z \geq 44$ the last odd Z p-bound nuclei become weaker bound and photodisintegration forces a fraction of the reaction path 1 mass unit away from the drip line. The flow in FRDM2 is identical except for the missing 2p-capture reactions.

The Hilf et al. mass model (set HIL) predicts generally higher proton binding energies. Therefore, the proton drip line is shifted towards more neutron-deficient nuclei compared to the predictions of the FRDM (1992) mass model, in the range $34 \leq Z < 42$ by two mass units (for odd Z), and in the range $42 \leq Z < 46$ by one mass unit. For $Z \geq 46$ the drip line is identical to the one predicted by the FRDM (1992). Despite of this drip line shift, the rp-process reaction path is quite similar to the FRDM1 case. This is due to the relatively low proton binding energies predicted by the Hilf et al. mass model for the last proton bound odd Z nuclei in this mass region. This results in significant photodisintegration which shifts part of the reaction flow one mass unit away from the proton drip line. This compensates for the shifted drip line below at $42 \leq Z < 46$. In the $34 \leq Z < 40$ range the difference between the HIL and the FRDM1 proton drip lines is two mass units, so the HIL reaction path is still one mass unit beyond the FRDM1 drip line. This makes it similar to the FRDM1 path via 2p-capture reactions, which become efficient just in this mass range. Though the reaction paths are quite similar, an important difference is that in the FRDM1 case part of the flow is carried by 2p-capture reactions, which are generally less effective than single proton capture sequences. Nevertheless, the similarity of the FRDM1 and HIL rp-process paths is remarkable in view of the fact that the two underlying mass models are very different in their predictions.

The reaction path for ABH-conditions based on the FRDM1 reaction rates is also shown in Fig. 26. The lower density results in stronger photodisintegration (see Eq. 25), especially for higher Z were proton capture rates become smaller due to the higher Coulomb barrier. This shifts the reaction path above $Z = 40$ to less neutron-deficient nuclei, for $40 \leq Z \leq 44$ by about one mass unit, for $46 \leq Z \leq 48$ by about two mass units.

B. Waiting points

The rp-process is characterized by proton capture reaction rates that are orders of magnitude faster than β-decay rates. The reaction path therefore follows a series of fast (p,γ)-reactions until further proton capture is inhibited, either by proton decay (negative proton capture Q-value) or photodisintegration (small positive proton capture Q-value). Then the reaction flow has to wait for the relatively slow β-decay and the respective nucleus is called a "waiting point". The total lifetimes of the waiting points along the reaction path (including all destructive processes like β-decay and net proton capture) entirely determine the speed of nucleosynthesis towards heavier nuclei (section IV C) and the produced isotopic abundances (section IV D), since at any given time essentially all the material is stored in the waiting points.

Waiting points are characterized by low or negative proton capture Q-values that hamper further proton capture. In a sequence of proton captures, this will occur at an even Z nucleus (Z, N) due to the lower proton binding energy of odd Z nuclei. However, the total lifetime of a waiting point nucleus can be significantly reduced by 2p-capture reactions bridging the nucleus $(Z+1, N)$. This type of reaction has been described for proton unbound intermediate nuclei $(Z+1, N)$ in section III C. In the case of a weakly bound $(Z+1, N)$ nucleus strong photodisintegration will bring the nuclei (Z, N) and $(Z+1, N)$ into equilibrium. The proton capture process will then be a 2p-capture process as well that can be described as in section III C, when the proton decay rate of the intermediate nucleus $(Z+1, N)$ is replaced by the rate for photodisintegration (γ,p). This picture holds as long as the (γ,p) rate on the intermediate nucleus $(Z+1, N)$ is faster than its proton capture rate (which defines "weakly proton bound").

The total lifetime of a waiting point nucleus (Z, N) can therefore be calculated by taking into account 2p-capture and β-decays. Two cases have to be distinguished. At high temperatures (typically above 1.5–2 GK) photodisintegration on nucleus $(Z+2, N)$ causes a strong inverse 2p-capture rate and nucleus (Z, N) and $(Z+2, N)$ are in equilibrium. In this case, the total lifetime $1/\lambda_{\text{total}}$ of nucleus (Z, N) is determined by the β-decays of nucleus (Z, N) and $(Z+2, N)$ and can be calculated using Eq. 23 for bound nuclei $(Z+1, N)$ and and Eqs. 37 and 38 for unbound nuclei $(Z+1, N)$:

$$\lambda_{\text{total}} = \lambda_{\beta(Z,N)} + \left[Y_{\text{p}} \rho N_A \left(\frac{2\pi\hbar^2}{kT} \right)^{\frac{3}{2}} \right]^2 \left(\frac{1}{\mu_{(Z,N)} \mu_{(Z+1,N)}} \right)^{\frac{3}{2}} \frac{G_{(Z+2,N)}(T)}{4G_{(Z,N)}(T)} \exp\left(\frac{Q_{(Z,N)(\text{p},\gamma)} + Q_{(Z+1,N)(\text{p},\gamma)}}{kT} \right) \lambda_{\beta(Z+2,N)} \tag{43}$$

with $\lambda_{\beta(Z,N)}$ being the β-decay rate, $G_{(Z,N)}$ the partition function, $\mu_{(Z,N)}$ the reduced mass and $Q_{(Z,N)(\text{p},\gamma)}$ the proton capture Q-value for nucleus (Z, N). T is the temperature, Y_{p} the hydrogen abundance and ρ the mass density. β-decay of the intermediate nucleus $(Z+1, N)$ is negligible due to its low equilibrium abundance.

Note that apart from temperature and density, the lifetime depends only on the nuclear masses (exponentially) and β-decay rates (linear) but not on the proton capture rates.

If the temperature is low, typically below 1.5–2 GK for the cases discussed here, photodisintegration on nucleus $(Z+2, N)$ becomes negligible and the total lifetime $1/\lambda_{\text{total}}$ of nucleus (Z, N) is determined by its β-decay and the 2p-capture rate (Eq. 37):

$$\lambda_{\text{total}} = \lambda_{\beta(Z,N)} + Y_{\text{p}}^2 \rho^2 N_A^2 \left(\frac{2\pi\hbar^2}{kT} \right)^{\frac{3}{2}} \frac{G_{(Z+1,N)}(T)}{2G_{(Z,N)}(T)} \exp\left(\frac{Q_{(Z,N)(\text{p},\gamma)}}{kT} \right) < \text{p}\gamma >_{(Z+1,N)} \tag{44}$$

It is important to note that it is not the proton capture rate on the waiting point isotope (Z, N), but the proton capture rate on the following isotone $(Z+1, N)$ that determines the net proton capture flow in this case.

We obtain Eqs. 43 and 44 independently of the intermediate nucleus $(Z+1, N)$ being proton bound or unbound either by describing regular proton capture and photodisintegration (section III B) or 2p-capture (section III A). There is no discontinuity in the effective stellar lifetime going from a positive to a negative proton capture Q-value. To identify nucleus (Z, N) as a waiting-point it is therefore not sufficient to determine whether nucleus $(Z+1, N)$ is proton bound or unbound. The exact value of the binding energy of nucleus $(Z+1, N)$ has to be known to determine the effective stellar lifetime of nucleus (Z, N) and thus its impedance on the rp-process. It is here where mass model uncertainties are a limiting factor. This is demonstrated for ^{68}Se in Fig. 27, which shows the effective half-life (including β-decay and proton capture) as a function of the proton capture Q-value, calculated from Eqs. 43 and 44 for XRB-conditions. A change in the Q-value of 200 keV, well within mass model uncertainties, changes the half-life by a factor of 5. Note that 2p-capture leads to a significant lifetime reduction for not too negative Q-values.

For the densities and temperatures discussed in this work there were no waiting points due to low proton capture rates - even when the uncertainties of the Hasuer-Feshbach calculations are taken into account. An exception are the endpoint nuclei of the rp-process under ABH-conditions, where the rp-process ends due to the increasing Coulomb barrier.

C. Time structure

An important question for rp-process nucleosynthesis is the processing speed towards heavier elements. This determines the energy generation, the produced isotopic abundances and how much fuel is left in an astrophysical scenario within a given timescale for rp-process conditions. Especially in X-ray bursts with rather short explosion timescales (10–100 s) it has to be investigated whether processing above Kr occurs at all.

The time required to process material beyond Kr is determined by the total lifetimes of the rather long lived isotopes ^{56}Ni, ^{64}Ge, ^{68}Se and ^{72}Kr (see Tab. IV). Whether these nuclei are waiting points or not depends apart from temperature and density strongly on the proton capture Q-value (see. Eqs. 43 and 44), which is known experimentally only in the case of ^{56}Ni. For ^{64}Ge, ^{68}Se, and ^{72}Kr our data set FRDM1 (see Tab. IV for the proton capture Q-values) is in agreement with recent experiments indicating that ^{65}As is mainly β-decaying [43,93], that ^{69}Br is p-unbound by at least 450 keV [109,113] and that ^{73}Rb is p-unbound [112,114]. In the case of ^{69}Br and ^{73}Rb however, the interpretation of the experimental results depends strongly on assumptions on the production cross sections of weakly p-bound nuclei. The temperature and density conditions for which ^{56}Ni, ^{64}Ge, ^{68}Se, or ^{72}Kr are waiting points can be deduced from Fig. 28, which shows the density required for equal net proton capture and β-decay rate as a function of temperature. For densities above the curve, more than half of the reaction flow proceeds via fast proton capture and the lifetime of the nucleus is considerably smaller than its β-decay lifetime. The density required for effective proton capture is increasing at lower temperatures due to the Coulomb barrier, and at higher temperatures due to strong photodisintegration. Since ^{56}Ni has a relatively long half-life, we also show the density that is required to get a total stellar lifetime of 100 s, which is of the order of the β-decay lifetimes for the other isotopes.

At X-ray burst densities of 10^6–10^7 g/cm^3 ^{56}Ni and ^{64}Ge can be bridged very effectively for a wide range of temperatures, since ^{57}Cu and ^{65}As are predicted to be sufficiently proton bound. On the other hand ^{69}Br and ^{73}Rb are p-unbound and ^{68}Se and ^{72}Kr have to be bridged by 2p-capture, which requires higher densities. However, Fig. 28 shows that there exists a temperature window between 1 GK and 1.8 GK, where typical X-ray burst densities are sufficient to bridge these isotopes. In this temperature range 2p-capture will greatly reduce the timescale for the production of elements heavier than Kr. Fig. 28 shows also the densities and temperatures required to bridge ^{64}Ge, ^{68}Se and ^{72}Kr, when the Hilf et al. mass model is applied (reaction rate set HIL). Since the Hilf et al. mass model predicts significantly higher proton binding energies in this region (see Tab. IV for values), which might be in contradiction with experiments for ^{69}Br and ^{73}Rb (as discussed above), all the long lived isotopes can be bridged for essentially all conditions considered here (see also Fig. 26). Therefore, no significant delay occurs for the rp-process in the $A = 64$–72 region, when the HIL reaction rates are used.

For XRB-conditions, the only possible significant delays above Kr are the β-decays of ^{76}Sr, ^{80}Zr and ^{81}Zr (see Tab. IV and III). The necessary conditions to bridge those isotopes when the reaction rate set FRDM1 is used are also shown in Fig. 28. While ^{80}Zr will be always a waiting point, ^{76}Sr can be bridged for densities above $3 \cdot 10^6$ g/cm^3 and ^{81}Zr can be bridged under most conditions. Therefore, the total delay above Kr will be just 9–16 s, depending on the net proton capture on ^{76}Sr. Using the FRDM (1992) mass model predictions for the ^{81}Zr and ^{82}Nb masses instead of the experimental mass for ^{81}Zr and the extrapolated value for ^{82}Nb from Audi et al. [47] (as adopted in FRDM1) would lead to significantly enhanced photodisintegration on ^{82}Nb under most conditions discussed here. This would make ^{81}Zr a waiting point, adding an additional 15 s delay to the rp-process. A measurement of the ^{82}Nb mass as well as an improved experimental value for the ^{81}Zr mass (the quoted uncertainty is 297 keV [64]) would be very important. Also, the experimental half-life of ^{81}Zr has a large uncertainty (15$\pm$5 s) and a more precise measurement would be desirable. Fig. 28 shows also that with the Hilf et al. mass model (reaction rate set HIL) ^{76}Sr can be bridged already for densities above 10^5 g/cm^3. Similar densities as in the FRDM1/FRDM2 calculations (well above 10^6 g/cm^3) are required to bridge the heavier waiting points.

Fig. 29 shows as an example the abundances of the waiting point isotopes as a function of processing time for XRB-conditions, calculated for all three sets of reaction rates, FRDM1, FRDM2, and HIL. The small initial increase during the first 10^{-2} s is due to rapid proton capture of solar seed abundances of nuclei in the corresponding isotone chain (N = const). Much higher abundances are reached on a longer timescale of seconds by processing He with the triple α-reaction and subsequent rapid proton and α-capture starting on ^{12}C. As can be seen, He is used up faster than hydrogen. This causes the decrease of the abundance at later times because processing continues without supplying new seed nuclei via the triple α-reaction. The speed of the reaction flow is illustrated in Fig. 30, which shows the time required to reach the peak abundance at a certain mass number for all three sets of reaction rates FRDM1, FRDM2, and HIL for $T = 1.5$ GK and $\rho = 10^6$ g/cm^3 as well as for a calculation with FRDM1 at a slightly higher density of

$5{\cdot}10^6$ g/cm^3. For XRB-conditions processing up to $A = 64$ takes only 2–2.5 s. For the FRDM calculations the main delay occurs then at ^{68}Se and ^{72}Kr as discussed above. Without 2p-capture reactions (FRDM2) and at a density of 10^6 g/cm^3 the $A = 92$ peak abundance is reached after 72 s. If 2p-capture reactions are included in the network, this happens already after 55 s. At higher densities 2p-capture is more effective resulting in an even faster flow. This is shown for a density of $5 \cdot 10^6$ g/cm^3, which is still not unrealistic for X-ray bursts, where the $A = 92$ maximum is reached already after 22 s. A slightly less p-unbound ^{69}Br or ^{73}Rb would also make the 2p-capture reactions more effective and would therefore have a similar effect. Since the Hilf et al. mass model predicts ^{69}Br and ^{73}Rb to be proton bound, the HIL results represent the extreme, where all long lived waiting points can be bridged. The $A = 92$ maximum is then reached after 7 s. These results clearly demonstrate the importance of 2p-capture reactions as well as the sensitivity of the processing timescale to the nuclear masses.

For the lower densities of $\rho = 10^4$ g/cm^3 in black hole accretion disks, the long lived isotopes ^{56}Ni, ^{68}Se and ^{72}Kr cannot be bridged efficiently anymore as can be seen from Fig. 28. Especially 2p-capture reactions are negligible at these densities. The reaction flow towards heavier elements is therefore strongly delayed as is shown in Fig. 30. Furthermore, above Kr the reaction path is shifted by 1–2 mass units towards more stable nuclei, which causes additional delays due to the the rather slow β-decays of ^{77}Sr ($T_{1/2} = 9$ s), ^{86}Mo ($T_{1/2}$=19.6 s), ^{90}Ru ($T_{1/2} = 13$ s), ^{94}Pd ($T_{1/2} = 9$ s), and ^{98}Cd ($T_{1/2} = 9.2$ s). As can be seen in Fig. 31 and 30 the processing timescale is much longer compared to the calculations at higher densities and the $A = 92$ peak abundance is reached after 370 s. This is however short compared to typical processing times of 10^4–10^5 s in accretion disks around low mass black holes. Therefore, temperatures of 1–1.8 GK and densities around 10^4 g/cm^3 might well exist for several 100 s.

Fig. 32 shows the contours of the processing time required to reach the $A = 92$ peak abundance in the temperature-density plane for the FRDM1, FRDM2 and HIL calculations. The sharp temperature window for effective processing into the $A = 92$ region can be seen clearly. It ranges from 1 to 1.8 GK for the FRDM1 and FRDM2 calculations and from 1 to 2.2–2.5 GK for the HIL calculations, relatively independent of density. Processing times in the window are always less then 500 s for the density range assumed in this study and drop below 100 s above a density of 10^4 g/cm^3 for the HIL reaction rates and above a density of $2{\cdot}10^5$ g/cm^3 for the FRDM1 and FRDM2 reaction rates. Comparing the plots for FRDM1 and FRDM2 demonstrates the importance of 2p-capture reactions. They strongly enhance the processing speed for densities above about 10^6 g/cm^3.

We can conclude that large amounts of $A = 80$–100 nuclei will be produced in scenarios with densities above roughly 10^6 g/cm^3, with solar H and He abundances and with temperatures between 1 and 2 GK for at least 10 s (and no longer than about 200 s). These conditions are well fulfilled for most X-ray bursts. For lower densities and lower or higher temperatures much longer processing times are required, especially when ^{64}Ge or ^{56}Ni become waiting points (see Fig. 28). To constrain the time structure, mass measurements for ^{65}As, ^{68}Se, ^{69}Br and ^{73}Rb would be desirable (see section IV B). All β-decay half-lives relevant for the processing time are known experimentally and with the exception of ^{81}Zr the accuracy is reasonable. In the lower density scenarios ^{56}Ni becomes the dominating waiting point. The time structure depends therefore critical on the net proton capture rate on ^{56}Ni. In this case the masses of all relevant nuclei, ^{56}Ni, ^{57}Cu, and ^{58}Zn are known experimentally with a precision of 50 keV or better. However, for temperatures below 1.5 GK ^{58}Zn drops out of equilibrium and the ^{57}Cu(p,γ)^{58}Zn reaction rate introduces then a major uncertainty in the effective stellar lifetime of ^{56}Ni and thus in the overall time structure of the rp-process above ^{56}Ni (see Eq. 44).

D. Abundance pattern

The processed material along the rp-process path is stored at all times in the waiting point isotopes because of their slow depletion rates. The total lifetimes of the waiting points determine therefore the produced isotopic abundance pattern. The important waiting points can be identified from the calculated rp-process path in Fig. 26. We begin the discussion with the calculations for XRB-conditions. For the FRDM1 reaction rates the longest lived nuclei with low proton capture Q-values along the rp-process path are the even-even $N = Z$ and the even-odd $N = Z + 1$ nuclei which are listed in Tab. III and IV. The temperature and density conditions at which these nuclei become significant waiting points can be determined from Fig. 28, which shows the density required for equal net proton capture rate and β-decay rate as a function of temperature. For densities below the line β-decay dominates and the isotope is a waiting point. It can be seen that for typical densities of 10^6–10^7 g/cm^3 and temperatures between 1–2 GK 2p-capture on the even-even $N = Z$ nuclei above Kr is not very effective, except for ^{76}Sr, which can be bridged for densities above $3 \cdot 10^6$ g/cm^3. These nuclei will therefore be waiting points under most rp-process conditions (see the β-decay branchings in Tab. IV). In contrast, the even-odd $N = Z + 1$ nuclei can be bridged well for most conditions (see the β-decay branchings in Tab. III). As example we show in Fig. 33 the abundance pattern for XRB-conditions at the time of maximum ^{92}Pd abundance (see Fig. 29). The abundance peaks at the masses of the important waiting points ($A = 72$, 76, 80, 84, 88, 92, 96) can be clearly identified. For comparison we also show the abundance pattern obtained without 2p-capture reactions, which shows more pronounced peaks at $A = 68$ (from ^{68}Se) and $A = 72$ (from ^{72}Kr), since the effective lifetimes of these isotopes are increased.

As discussed above, the reaction path calculated with the Hilf et al. mass model (HIL) is very similar to the FRDM1 path, but 2p-capture reactions are replaced by the more effective single proton capture reactions. The nuclei with long β-decay half-lives and small proton capture Q-values along the rp-process path are therefore also the even-even $N = Z$ nuclei listed in Tab. IV. Fig. 28 shows the density required for equal net proton capture and β-decay rate as a function of temperature when the Hilf et al. mass model is used. Here, branching into β-decay for the very long lived isotopes ^{64}Ge, ^{68}Se, ^{72}Kr and also for ^{76}Sr is negligible under most rp-process conditions. This results in a very fast reaction flow and a shifting of the waiting points into the even-even $N = Z - 2$ isotopes ^{66}Se, ^{70}Kr, ^{74}Sr and ^{78}Zr. Those isotopes are rather short lived but become then the slowest β-decay in the isotopic ($Z =$ const) chain due to the enhancement of β-decay half-lives for even-even isotopes compared to even-odd nuclei (see section III D). On the other hand, for even-even $N = Z$ nuclei above Zr, branching into β-decay is dominant. This results in an abundance pattern similar to the one obtained with the FRDM1 calculations for this region. The resulting abundance pattern for XRB-conditions at the time of the $A = 92$ peak abundance is shown in Fig. 33. Again the peaks at the masses of the waiting points ($A = 74$, 78, 80, 84, 88, 92, 96) can be clearly seen.

For the lower densitiy in ABH-conditions and with the reaction rate set FRDM1 the reaction path is slightly shifted to more stable nuclei. This shift however is small so the even-even $N = Z$ nuclei are still on the reaction path and represent again waiting points, except for ^{96}Cd (see branchings into β-decay in Tab. IV). Since 2p-capture does not occur at these low densities ^{68}Se and ^{72}Kr are waiting-points with their full β-decay lifetime. In contrast to the calculations for higher density, all the even-odd $N = Z + 1$ nuclei listed in Tab. III represent now waiting points as well, since photodisintegration cannot be overcome by proton capture rates anymore (see branchings into β-decay in Tab. V). Furthermore, for $Z \geq 42$ the even-even $N = Z + 2$ nuclei ^{86}Mo, ^{90}Ru, ^{94}Mo, and ^{98}Cd become additional rather long lived waiting points. This allows to interpret the resulting abundance pattern at the time of the $A = 92$ peak abundance shown in Fig. 33. For $A < 84$ it is similar to the pattern obtained for the high density calculations with FRDM2, except for the strong peak at $A = 56$ from ^{56}Ni. The additional peaks at $A = 77$ and $A = 81$ are caused by the even-odd waiting points ^{77}Sr and ^{81}Zr. For $A \geq 84$ however, the abundance pattern is very different from the FRDM2 high density calculations, since the longer lived waiting points closer to stability dominante. Instead of the peaks at $A = 84$, 92, and 96 the abundance maxima now occur at $A = 94$ and 98.

In order to determine possible contributions of rp-process nucleosynthesis to the the solar system abundances, we have to look at the overproduction factors relative to the solar system abundance $T = Y/Y_\odot$. The fraction of the observed solar system abundance of an isotope i, that is made in the rp-process $f^i_{\rm rp}$ is then given by

$$f^i_{\rm rp} = T_i f_{\rm proc} \tag{45}$$

where $f_{\rm proc}$ is the fraction of solar material that had been processed in the respective rp-process scenario. Calculation of $f^i_{\rm rp}$ requires therefore to determine f_{proc} from models for the galactic chemical evolution. However, upper limits on $f^i_{\rm rp}$ can be derived already from the pattern of overproduction factors: Since even for the isotope with the largest overproduction factor $T_{\rm max}$ per definition $f^i_{\rm rp} \leq 1$, it follows that

$$f^i_{\rm rp} \leq \frac{T_i}{T_{\rm max}}. \tag{46}$$

Therefore, nuclei that have not the highest overproduction factor can only partially be produced by the rp-process. The production of s- or r-nuclei can also constrain $f^i_{\rm rp}$, since these processes are rather well understood. If the s- or r- contribution to a nucleus k is known to be $f^k_{\rm sr}$, then we must require $f^k_{\rm rp} \leq (1 - f^k_{\rm sr})$ or using Eq. 45

$$f^i_{\rm rp} \leq (1 - f^k_{\rm sr}) \frac{T_i}{T_k}. \tag{47}$$

Dependent on the overproduction factor obtained for the s- or r-nucleus and the value and accuracy of $f^k_{\rm sr}$ this might put severe constraints on the possible rp-process contributions to all other isotopes.

We calculated the overproduction factor T for the most neutron-deficient stable isotope in each mass chain from the abundance patterns shown in Fig. 33, assuming that at the time the $A = 92$ peak abundance is reached further processing stops and all nuclei β-decay. Thereby we neglect β-delayed proton decay and the non-zero timescale of the density and temperature decrease in an astrophysical event. Both effects will smooth the abundance pattern without changing the general structure (maxima and minima).

The resulting overproduction factors are shown in Fig. 34. We first discuss the results for XRB-conditions. Fig. 34 shows that with all three data bases the p-nuclei ^{84}Sr, ^{92}Mo and ^{96}Ru can be produced with high overproduction factors of comparable order of magnitude. In FRDM1 the non-p nuclei ^{72}Ge, ^{76}Se and ^{80}Kr are produced with similar overproduction factors, since the corresponding waiting points ^{72}Kr, ^{76}Sr and ^{80}Zr are not bridged effectively. However, for higher densities (or less p-unbound ^{73}Rb and ^{77}Y) 2p-capture on ^{72}Kr and ^{76}Sr becomes more effective (see Fig. 28) and the overabundance pattern will look more like in the HIL case: Due to the bridging of ^{72}Kr and ^{76}Sr the production of the non-p nuclei ^{72}Ge and ^{76}Se is strongly reduced, while instead relatively large amounts of the p-nuclei ^{74}Se and ^{78}Kr are produced. Apart from that the structure of overproduction factors is remarkably similar for the FRDM1 and HIL calculations, though they are based on different mass models. The only non-p nucleus that will be produced in significant amounts for a variety of mass models and temperature and density conditions is the s-only nucleus ^{80}Kr. Its progenitor ^{80}Zr can not be bridged very effectively as can be seen in Fig. 28, since both mass models predict relatively low proton capture Q-values (see Tab. IV).

Also shown are the overproduction factors for ABH-conditions. As expected from the discussion of the abundance pattern, the overproduction factors for $A \leq 84$ are similar to the FRDM2, XRB calculations without 2p-capture reactions, except for the additional peaks at $A = 77$ and 81. The $A = 56$ abundance peak is irrelevant in terms of overproduction due to the large solar abundance of ^{56}Fe. Above $A = 84$, however, the overpoduction factors are very different from the ones obtained in XRB calculations. The highest overproduction factors in black hole accretion disks are reached for the p-nuclei ^{94}Mo and ^{98}Ru, which are not produced in X-ray bursts. To summarize, the main product in black hole accretion disks under the discussed conditions will be ^{94}Mo and ^{98}Ru. Within the uncertainties due to the β-decay half-life calculations, the production of significant amounts of the p-nucleus ^{84}Sr and the non p-nucleus ^{80}Kr is also possible.

The dependence of the overproduction factors on density and temperature is illustrated in Fig. 32, which shows the contours of the $A = 92$ peak overproduction factor in the temperature-density plane for the various sets of reaction rates. It also shows how the processing time to reach these overproduction factors varies with temperature and density. Within the temperature window of fastest processing as described in section IV C the obtained overproduction factors are increasing with density, relatively steep below 10^5 g/cm^3 and relatively slowly at higher densities. They also depend only weakly on temperature, which justifies the interpretation of the overproduction patterns shown in Fig. 34 as typical for nucleosyntheis at the respective density. At temperatures below or above the temperature window of effective processing the picture is reversed, with the overproduction factors depending strongly on temperature and hardly at all on density.

The obtained overproduction patterns suggest rp-process nucleosynthesis as a possible alternative scenario for the production of the solar abundances of the light p-nuclei ^{84}Sr, ^{92}Mo, ^{94}Mo, ^{96}Ru, ^{98}Ru eventually also for ^{74}Se and ^{78}Kr. This is interesting, since the large solar abundances observed for ^{92}Mo, ^{94}Mo, ^{96}Ru and ^{98}Ru cannot be produced sufficiently in other p-process scenarios [29,36,33]. For ^{92}Mo, contributions from neutrino driven winds in type II supernovae are possible [115], which might be in agreement with the somewhat lower overproduction factor obtained here (Fig. 34). Interestingly the two scenarios that we discussed in this study seem to be complementary in the production of these 4 p-nuclei. ^{92}Mo and ^{96}Ru can only be produced in X-ray bursts, since the higher density is required to shift the reaction path towards the proton drip line. On the other hand, ^{94}Mo and ^{98}Ru can only be produced at the lower densities of black hole accretion disks. These lower densities do also occur during the decrease of an X-ray burst but then the temperatures will be too low to process material in the $Z = 46$ range within the event timescale. The large overproduction of ^{80}Kr in all the scenarios discussed might cause a conflict with the s-process, but an experimental determination of the β-decay half-life of the progenitor ^{80}Zr would be necessary to determine to what degree that is the case (see also section IV F).

The obtained overproduction patterns depend strongly on the adopted nuclear masses, which determine the waiting points. For the rp-process path under XRB-conditions as well as under ABH-conditions, the important masses are not known experimentally. For X-ray burst scenarios mass measurements for ^{80}Zr, ^{81}Nb, ^{82}Nb, ^{84}Mo, ^{85}Tc, ^{88}Ru, ^{89}Rh, ^{92}Pd, ^{93}Ag, ^{96}Cd, ^{97}In and an improved measurement for ^{81}Zr would greatly reduce uncertainties. This is especially true for the $A = 80$–84 nuclei, where large nuclear deformations influence the mass model predictions. For the rp-process calculations in black hole accretion disks additional mass measurements for ^{85}Mo, ^{86}Tc, ^{87}Tc, ^{89}Ru, ^{90}Rh, ^{90}Ru, ^{91}Rh, ^{93}Pd, ^{94}Ag, ^{94}Pd, ^{95}Ag, ^{97}Cd, ^{98}In, ^{98}Cd and ^{99}In would improve the accuracy of the calculations considerably.

The overproduction factor depends also on the progenitors β-decay half-life. This is demonstrated in Fig. 35, which shows that the overproduction factor for ^{92}Mo is roughly proportional to the β-decay half-life of ^{92}Pd. In the calculations under XRB-conditions these half-lives are not known experimentally for $Z \geq 40$. The typical uncertainties in the shell model β-decay half-lives are about a factor of 2.7, which introduces an uncertainty of the same order of magnitude into the overproduction factors. For ^{80}Zr and ^{84}Mo, where no or no reliable shell model calculations are available, the uncertainties are about a factor of 5.5 (see section III D). For the reaction path in black hole accretion disks most of the important β-decay rates are known experimentally except for ^{80}Zr, ^{84}Mo and ^{85}Mo. Therefore half-life measurements for ^{84}Mo, ^{85}Mo, ^{92}Pd, ^{96}Cd, and especially ^{80}Zr would be very important.

Note that the large overproduction factors are obtained by processing ^{4}He via the 3α-reaction and subsequent

^{4}He- and proton capture on the resulting ^{12}C. The results are therefore independent of the initial seed abundance distribution, except for the H/He ratio. This can be seen in Figs. 29 and 31, which show isotopic abundances as a function of time. Most abundances reach a preliminary plateau after less than 0.01 s, which is the result of processing along N =const chains via proton capture on initial solar seed nuclei. However, the abundance levels produced are typically (for p-nuclei) 6 orders of magnitudes smaller compared to the peak abundances reached after 10–100 s, which are the result of processing H and He only.

E. Energy production

The energy released by the nuclear reactions of the rp-process influences the structure of the accretion disk around low mass black holes and powers type I X-ray bursts. It is therefore the crucial quantity in the astrophysical model descriptions for these scenarios. The calculation of the energy production requires rp-process network calculations using detailed temperature and density profiles, which have to be determined self consistently using the energy release of the reaction network. We can however use our constant temperature and density rp-process calculations to investigate, how the produced energy depends on the reaction network used. Fig. 36 shows the total, time integrated energy produced by nuclear reactions (corrected for neutrino losses) as a function of processing time for XRB-conditions. The total energy released when processing proceeds into the mass 80–100 region is 3.6 10^{18} erg/g. As expected this value is reached at different times for the different mass models and reaction rate sets HIL, FRDM1, and FRDM2 reflecting the different processing speeds as discussed in section IV C. Critical is the size of the reaction network. A "network" limited to the 3α-reaction, as it was used in some previous hydrodynamic studies, produces only about 0.3% of the total energy. Fig. 36 also clearly shows that nucleosynthesis above Kr is important in terms of energy production since it accounts for about 50% of the total energy. This implies that the energy production in X-ray bursts and accreting low mass black holes was underestimated in previous studies, which were all based on smaller reaction networks.

F. The Zr-Nb cycle

The FRDM (1992) mass model predicts very low α-separation energies around $N = 42$ for $40 \leq Z \leq 44$, as discussed in section II C. In the case of ^{84}Mo, which is even predicted to be α-unbound by 0.575 MeV, this might have interesting consequences for the rp-process. Though the α-decay half-life ($\approx 10^{45}$ yr) is still longer than the β-decay half-life, the low α-separation energy leads to strong enhancement of the ^{83}Nb(p,α)^{80}Zr reaction rate. This will lead to the formation of an Zr-Nb cycle via the reaction sequence

$$^{83}\mathrm{Nb(p,}\alpha)^{80}\mathrm{Zr}(\beta^+)^{80}\mathrm{Y(p,}\gamma)^{81}\mathrm{Zr}\left\{\begin{array}{c}(\beta^+)^{81}\mathrm{Y(p,}\gamma)\\(\mathrm{p,}\gamma)^{82}\mathrm{Nb}(\beta^+)\end{array}\right\}{}^{82}\mathrm{Zr(p,}\gamma)^{83}\mathrm{Nb}$$

if net proton capture on ^{83}Nb is sufficiently supressed by photodisintegration on ^{84}Mo (see Fig. 37). Under the typical conditions for X-ray bursts and accretion disks around low mass black holes discussed so far, the net ^{83}Nb(p,γ) rate is still dominating the reaction flow (see Fig. 26). This is also illustrated in Fig. 38, which shows the fraction of the reaction flow branching into the ^{83}Nb(p,α)^{80}Zr reaction as a function of temperature. For XRB- and ABH-conditions (T = 1.5 GK) only a few percent of the reaction flow is processed via the ^{83}Nb(p,α)^{80}Zr branch and the Zr-Nb cycle does not play a significant role. However, at temperatures of 2 GK or higher, essentially all material is processed back to ^{80}Zr and stored in the Zr-Nb cycle, which is then the endpoint of the rp-process (see Fig. 37). At these high temperatures equilibrium between ^{83}Nb and ^{84}Mo is established via proton capture and photodisintegration. Therefore, the leakage out of the cycle is entirely determined by the flow via β-decay of ^{84}Mo, which depends on its β-decay rate and the proton capture Q-value of ^{83}Nb (which determines the ^{84}Mo equilibrium abundance). Eventually all material will be converted into ^{80}Zr and ^{81}Zr, since these are the isotopes with the slowest β-decays in the cycle. Fig. 39 shows the abundances of some important waiting points as well as for ^{80}Zr and ^{81}Zr as a function of processing time for a temperature of 2 GK, which corresponds approximately to the threshold temperature above which the Zr-Nb cycle becomes important. ^{56}Ni is a strong waiting point under these conditions (as discussed in section IV C). It takes therefore about 100 s to process material into the Zr region. As can be seen in Fig. 39 no further processing takes place (see ^{84}Mo abundance) and a large fraction of material (about 40–50%) is stored in ^{80}Zr and ^{81}Zr after about 200 s.

In typical X-ray bursts temperatures above 2 GK are reached, but the temperature increases with a timescale of the order of seconds. As a consequence all material will be locked in ^{56}Ni until the temperature drops again into the 1–2 GK window of most effective processing discussed in section IV and IV C. In this case the Nb-Zr cycle will not be important and the resulting reaction path will be the one shown in Fig. 26. There are however some X-ray bursts observed that show rise times of up to about 10 s. This might be enough time to get material past ^{56}Ni before photodisintegration on ^{57}Cu sets in, and this material will then be stored in the Zr-Nb cycle. In these kind of scenarios the existence of the Zr-Nb cycle might influence rp-process nucleosynthesis considerably. In most scenarios describing the nucleosynthesis in the accretion disks around low mass black holes the rp-process reaction path will be close to the one shown in Fig. 26, as discussed in section IV. However, peak temperatures depend very sensitively on model parameters (especially on the adopted viscosity of the disk). It might well be possible, that the Zr-Nb cycle plays a role under some conditions. This has to be investigated using network calculations based on more realistic temperature and density profiles.

The existence of the Zr-Nb cycle depends critically on the low α-separation energy for ^{84}Mo as predicted by the FRDM (1992) mass model. As discussed in section II C there is some indication, that this drop in α-binding energy might be an effect produced by the mass model. To clarify this, mass measurements for ^{84}Mo and ^{80}Zr would be extremely important. Very important would also be the measurement of the β-decay half-life of ^{84}Mo, which determines the leakage rate out of the cycle, and of ^{80}Zr, which determines together with ^{81}Zr the cycle timescale and the produced abundance pattern.

V. CONCLUSIONS

In this paper, we described the necessary ingredients for rp-process network calculations and performed for the first time calculations above $Z = 32$ using a complete reaction network including (γ,α)-reactions and 2p-capture. We discussed the uncertainties in the various nuclear structure input parameters and how they affect rp-process nucleosynthesis calculations in X-ray bursts and accreting black holes. The follwing conclusions can be drawn from this study:

1. The most important parameter for rp-process nucleosynthesis calculations are nuclear masses, since the reaction path for given temperature, density and composition is almost entirely determined by proton capture Q-values. The difficulty for calculations above ^{64}Ge is that for most of the nuclei along the proposed rp-process path no experimental mass data are available. Furthermore, the reaction path enters around $A = 70$–80 a region of highly deformed nuclei, which makes the theoretical prediction of masses more complicated. The comparison of calculations with the very different Hilf et al. and FRDM (1992) mass predictions indicated that mass model accuracies are sufficient for a rough determination of the reaction path. For the calculation of the time structure, the finally produced isotopic abundances, and the energy generation the uncertainties of the mass models are however too large. Mass measurements to a precision of about 100 keV or better for some selected isotopes are therefore crucial to improve the accuracy of the rp-process calculations. The most important nuclear masses were identified in this work and are listed in sections IV C, IV D, and IV F.

Also crucial for rp-process nucleosynthesis calculations are the β-decay half-lives of the nuclei along a given reaction path since they determine time structure and abundance patterns. Generally, for β-decay half-lives more experimental data are available than for nuclear masses. Especially the half-lives of the longer-lived waiting point isotopes in the Ge-Kr region that determine the processing timescale are all known. However, the actual endpoint of the rp-process and the finally produced abundance patterns are determined by the β-decay rates of proton drip line nuclei in the Sr-Sn region that are not yet known experimentally. Half-life measurements of these nuclei that are listed in section IV D would therefore strongly reduce uncertainties.

Generally, the network calculations did not suffer from the uncertainties in the proton capture rates. In most cases it is sufficient to know that these rates are several orders of magnitude larger than β-decay rates. Only at waiting points, where the reaction flow is hampered by a strong reverse reaction, either photodisintegration or proton decay, the exact value of the remaining net proton capture rate influences the branching ratio. Even in these cases, proton capture rates determine the net proton capture rate only for temperatures below typically 1.5 GK when some isotopes drop out of (p,γ)-(γ,p) equilibrium. Under these conditions, the proton capture rate on the isotone *following* the waiting point nucleus might influence the reaction path. However, the sensitivity to proton binding energies is much larger. Only if these are experimentally known, the typical statistical model uncertainties of the proton capture rates (a factor of 2) become significant. The only important case of this kind in this study is the ^{57}Cu(p,γ)^{58}Zn reaction for bridging the waiting point nucleus ^{56}Ni. An experimental determination of this rate would reduce uncertainties in the rp-process time scale.

2. In this work we included for the first time 2p-capture reactions bridging the proton drip line. We demonstrated that 2p-capture reactions are important for high density environments ($\rho \geq 10^6$ g/cm^3) like X-ray bursts, since they enhance the speed of the reaction flow towards nuclei heavier than Kr considerably. As a consequence, the endpoint

of nucleosynthesis in most type I X-ray bursts is not one of the long lived isotopes in the Ge–Kr range as previously assumed, but rather a range of nuclei in the $A = 80$–100 region. This leads to a factor of 2 increase in the predicted energy production. Also, a single X-ray burst consumes more hydrogen fuel than previously assumed. This will have to be taken into account in models of X-ray burst sequences.

3. Large amounts of p-nuclei are produced in the rp-process, especially ^{92}Mo and ^{96}Ru in X-ray bursts and ^{94}Mo and ^{98}Ru in accreting low mass black holes. This is interesting, since the origin of these isotopes in the solar system is not understood as they are systematically underproduced in all standard p-process scenarios.

Whether rp-process nucleosynthesis can produce the necessary amounts of nuclei to account for the observed solar system abundances depends on how often the corresponding event is occuring in our galaxy, how much mass is processed per event, and how much of the processed material is ejected into the interstellar medium (escape factor). In the case of type I X-ray bursts, the biggest uncertainty is this escape factor since the explosion energy is in general smaller than the gravitational binding energy of the accreted layer. We can estimate the necessary escape factor assuming overproduction factors of 10^7–10^8 from above, a birth rate for an X-ray bursting system of 10^{-5}/yr [116], a lifetime of an X-ray bursting system of 10^9 yr [116], a mass transfer rate of 10^{-8}–10^{-9} $M_{\odot}$/yr [117], a total time for galactic nucleosynthesis of 10^{10} yr, and a galaxy disk mass of $6{\cdot}10^{10}$ $M_{\odot}$. This leads to an escape factor of about 0.3 % required for type I X-ray bursts to account for solar system light p-nuclei abundances. It has been argued before that essentially no burned material escapes the neutron star and that therefore type I X-ray bursts are not nucleosynthesis sites. We showed however that the required escape factor is very small. We also showed that using a complete nuclear reaction network results in a substantial increase in energy production, compared with simpler networks. This was not taken into account in previous studies of mass loss from neutron stars via radiation driven winds [118]. As we will show in a forthcoming paper the the energy production rates in X-ray bursts might even be close to the local Eddington limit in the burning zone. Therefore, it cannot be excluded that radiation driven winds lead to the escape of a small fraction of the burned material. However, accurate X-ray burst model calculations including a full nuclear reaction network would be necessary to decide whether X-ray bursts are potential production sites of light p-nuclei.

In the case of accretion disks around black holes uncertainties are too large to make similar estimates. Nevertheless the production of ^{94}Mo and ^{96}Ru is a striking feature of rp-process nucleosynthesis under the conditions in these scenarios and certainly warrants further investigation.

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APPENDIX A: REACTION RATE TABLES

This appendix lists the astrophysical reaction rates of the reaction rate set FRDM1 as described in section IV. Only reactions necessary to calculate rp-process nucleosynthesis in the $32 \leq Z \leq 50$ range are included. These are (p,γ) rates in Tab. VI, (p,α) rates in Tab. VII, (γ,p) rates in Tab. VIII, (γ,α) rates in Tab. IX, and β^+-decay half-lives in Tab. X. The 2p-capture rates have been listed in Tab. I in section III C.

The charged particle reaction rates have been calculated using the Hauser-Feshbach code SMOKER as described in section III A. For the channel separation energies we used either experimental masses or the mass predictions from the FRDM (1992) mass model, with the two exceptions for the proton capture Q-value of ^{68}Se and the mass of ^{82}Nb as discussed in section IV. These rates are a considerable improvement over previously available Hauser-Feshbach calculations. They represent however only a preliminary stage in the recent efforts to improve SMOKER. A complete set of reaction rates (including all nuclei and all types of reactions) based on a version of the code that includes additional changes in the level density description [88] and the α-potential will be published in a forthcoming paper. Note that the calculation of all these rates is based on the statistical model that requires a sufficient level density in the compound nucleus as well as a negleglible contribution from a direct reaction. We showed in section III A that for all the rates presented here the compound nucleus level density for temperatures above 1 GK is high enough to justify the statistical approach. At significantly lower temperatures, which are not important for the rp-process above $Z = 32$, the level density for some nuclei near the proton drip line (the low level density nuclei indicated in Fig. 1) might become too small.

The β^+-decay rates are experimental or have been calculated either with the shell model code OXBASH or the

QRPA (see section III D). The theoretical rates were all based on the nuclear mass predictions of the FRDM (1992) model. For our reaction rate set FRDM1 we used the shell model rates when available, but we list here also the respective QRPA rates for comparison. A complete set of the QRPA β^+- and β^--decay rates for all nuclei has been published recently [99]. The details of the OXBASH shell model calculations and results for more nuclei will also be presented in a forthcoming paper [105]. Note that all the experimental β-decay rates were measured under laboratory conditions. Also, the calculated β-decay rates are limited to Gamow-Teller transitions from the ground-state of the parent nucleus. At the high temperatures and densities in a stellar plasma decay from excited states and continuum electron capture might play a role. Therefore, for the purpose of rp-process network calculations all the here listed β-decay rates, including the experimental data, represent only estimates for astrophysical β-decay rates (this was discussed in section III D).

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[1] J. Audouze, J. W. Truran, and B. A. Zimmerman, Ap. J. **184**, 493 (1973).
[2] R. K. Wallace and S. E. Woosley, Ap. J. Suppl. **45**, 389 (1981).
[3] S. E. Woosley and T. A. Weaver, in *High Energy Transients in Astrophysics*, Vol. 115 of *AIP Conference Proceedings*, edited by S. E. Woosley (American Institute of Physics, New York, 1984), p. 273.
[4] L. Van Wormer *et al.*, Ap. J. **432**, 326 (1994).
[5] S. K. Chakrabarti, Phys. Rep. **266**, 229 (1996).
[6] N. I. Shakura and R. A. Sunyaev, Astron. Astrophys. **24**, 337 (1973).
[7] S. K. Chakrabarti, L. Jin, and W. D. Arnett, Ap. J. **313**, 674 (1987).
[8] L. Jin, Ap. J. **337**, 603 (1989).
[9] L. Jin, W. D. Arnett, and S. K. Chakrabarti, Ap. J. **336**, 572 (1989).
[10] K. Arai and M. Hashimoto, Astron. Astrophys. **254**, 191 (1992).
[11] R. E. Taam and B. A. Fryxell, Ap. J. **294**, 303 (1985).
[12] A. Ayasli and P. C. Joos, Ap. J. **256**, 637 (1982).
[13] R. E. Taam, Ann. Rev. Nucl. Sci. **35**, 1 (1985).
[14] M. Y. Fujimoto, M. Sztajno, W. H. G. Lewin, and J. van Paradijs, Ap. J. **319**, 902 (1987).
[15] R. E. Taam, S. E. Woosley, T. A. Weaver, and D. Q. Lamb, Ap. J. **413**, 324 (1993).
[16] W. H. G. Lewin, J. van Paradijs, and R. E. Taam, Space Sci. Rev. **62**, 233 (1993).
[17] R. K. Wallace and S. E. Woosley, in *High Energy Transients in Astrophysics*, Vol. 115 of *AIP Conference Proceedings*, edited by S. E. Woosley (American Institute of Physics, New York, 1984), p. 319.
[18] J. W. Truran, in *Essays in Nuclear Astrophysics* (Cambridge University Press, Cambridge, 1982), p. 467.
[19] A. Shankar, W. D. Arnett, and B. A. Fryxell, Ap. J. **394**, L13 (1992).
[20] S. Starrfield *et al.*, Phys. Rep. **227**, 223 (1993).
[21] M. Politano *et al.*, Ap. J. **448**, 807 (1995).
[22] G. T. Biehle, Ap. J. **380**, 167 (1991).
[23] R. Cannon, Mon. Not. R. Astron. Soc. **263**, 817 (1993).
[24] G. T. Biehle, Ap. J. **420**, 364 (1994).
[25] K. S. Thorne and A. N. Zytkow, Ap. J. **212**, 832 (1977).
[26] R. Cannon, in *Nuclei in the Cosmos 1992*, edited by F. Käppeler and K. Wisshak (Institute of Physics Publishing, Bristol, 1993), p. 619.
[27] C. L. Fryer, W. Benz, and M. Herant, Ap. J. **460**, 801 (1996).
[28] R. A. Chevalier, Ap. J. **459**, 322 (1996).
[29] D. L. Lambert, Astron. Astrophys. Rev. **3**, 201 (1992).
[30] S. E. Woosley and W. M. Howard, Ap. J. Suppl. **36**, 285 (1978).
[31] S. E. Woosley and W. M. Howard, Ap. J. **354**, (1990).
[32] M. Rayet, N. Prantzos, and M. Arnould, Astron. Astrophys. **227**, 271 (1990).
[33] M. Rayet *et al.*, Astron. Astrophys. **298**, 517 (1995).
[34] W. M. Howard, B. S. Meyer, and S. E. Woosley, Ap. J. **373**, L5 (1991).
[35] M. Hencheck *et al.*, in *Nuclei in the Cosmos III, Third International Symposium on Nuclear Astrophysics, Assergi, Italy*, Vol. 327 of *AIP Conference Proceedings*, edited by M. Busso, R. Gallino, and C. M. Raiteri (American Institute of Physics, New York, 1994), p. 331.
[36] W. M. Howard and B. S. Meyer, in *Nuclei in the Cosmos*, edited by F. Käppeler and K. Wisshak (Institute of Physics, Bristol, 1993), p. 575.
[37] P. Möller and J. R. Nix, At. Data Nucl. Data Tab. **26**, 165 (1981).

[38] S. E. Woosley, W. A. Fowler, J. A. Holmes, and B. A. Zimmerman, Technical report, OAP-422, Caltech (unpublished).
[39] K. Takahashi, M. Yamada, and T. Kondoh, At. Data Nucl. Data Tab. **12**, 101 (1973).
[40] W. D. Myers and W. J. Swiatecki, Nucl. Phys. **81**, 1 (1966).
[41] J. Tuli, in *Nuclear Wallet Cards* (Brookhaven National Laboratory, Brookhaven, 1990).
[42] J. Jänecke and P. Masson, At. Data Nucl. Data. Tab. **39**, 265 (1988).
[43] J. Winger *et al.*, Phys. Lett. B **299**, 214 (1993).
[44] J. A. Halbleib and R. A. Sorensen, Nucl. Phys. A **98**, 542 (1967).
[45] J. Engel, P. Vogel, and M. R. Zirnbauer, Phys. Rev. C **37**, 731 (1988).
[46] P. Möller and J. Randrup, Nucl. Phys. A **514**, 1 (1990).
[47] G. Audi and A. H. Wapstra, Nucl. Phys. A **595**, 409 (1995).
[48] P. Haustein, At. Data Nucl. Data Tab. **39**, 185 (1988).
[49] Y. Aboussir, J. M. Pearson, A. K. Dutta, and F. Tondeur, Nucl. Phys. A **549**, 155 (1992).
[50] P. Möller, J. R. Nix, W. D. Myers, and W. J. Swiatecki, At. Data Nucl. Data Tab. **59**, 185 (1995).
[51] K. L. Kratz *et al.*, Ap. J. **403**, 216 (1993).
[52] F. K. Thielemann *et al.*, Nucl. Phys. A. **570**, 329c (1994).
[53] J. Dobaczewski, I. Hamamoto, W. Nazarewicz, and J. A. Sheikh, Phys. Rev. Lett. **72**, 981 (1994).
[54] S. Goriely and M. Arnould, Astron. Astrophys. **312**, 327 (1996).
[55] E. R. Hilf, H. von Groote, and K. Takahashi, in *Proc. 3rd Inter. Conf. Nuclei far from Stability, Cargese* (CERN, Geneva, 1976), Vol. 76-13, p. 142.
[56] J. J. Cowan, F.-K. Thielemann, and J. W. Truran, Phys. Rep. **208**, 267 (1991).
[57] H. von Groote, E. R. Hilf, and K. Takahashi, At. Data Nucl. Data Tab. **17**, 418 (1976).
[58] P. Möller and J. R. Nix, At. Data Nucl. Data Tab. **39**, 213 (1988).
[59] P. Möller, W. D. Myers, W. J. Swiatecki, and J. Treiner, At. Data Nucl. Data Tab. **39**, 225 (1988).
[60] R. Casten, Nucl. Phys. A **443**, 1 (1985).
[61] P. E. Haustein, D. S. Brenner, and R. F. Casten, Phys. Rev. C **38**, 467 (1988).
[62] A. Aprahamian, A. Gadala-Maria, and N. Cuka, (1997), to be published.
[63] C. J. Lister, P. E. Haustein, D. E. Alburger, and J. W. Olness, Phys. Rev. C **24**, 260 (1981).
[64] S. Della Negra, H. Gauvin, D. Jacquet, and Y. Le Beyec, Z. Phys. A **307**, 305 (1982).
[65] S. Issmer *et al.*, Z. Phys. A (1997), to be published.
[66] W. Nazarewicz and I. Ragnarsson, in *Handbook of Nuclear Properties*, edited by D. Poenaru and W. Greiner (Clarendon Press, Oxford, 1996), p. 80.
[67] C. J. Lister *et al.*, Phys. Rev. C **42**, R1191 (1990).
[68] L. Grodzins, Phys. Lett. **2**, 88 (1962).
[69] F. S. Stephens *et al.*, Phys. Rev. Lett. **29**, 438 (1972).
[70] C. J. Lister *et al.*, Phys. Rev. Lett. **59**, 1270 (1987).
[71] P. Raghavan, At. Data and Nuc. Dat. Tab. **42**, 189 (1989).
[72] M. Keim *et al.*, Nucl. Phys. A **586**, 219 (1995).
[73] F. Buchinger *et al.*, Phys. Rev. C **41**, 2883 (1990).
[74] P. Lievens *et al.*, Phys. Rev. C **46**, 797 (1992).
[75] S. Raman, C. W. Nestor, S. Kahane, and K. H. Bhatt, At. Data Nucl. Data Tab. **42**, 1 (1989).
[76] B. J. Varley *et al.*, Phys. Lett. B **194**, 463 (1987).
[77] M. Chartier *et al.*, Phys. Rev. Lett. **77**, 2400 (1996).
[78] W. D. Arnett, in *Supernovae and Nucleosynthesis* (Princeton University Press, Princeton, 1996).
[79] W. A. Fowler, G. E. Caughlin, and B. A. Zimmerman, Ann. Rev. Astro. Astrophys. **13**, 69 (1975).
[80] G. R. Caughlan and W. A. Fowler, At. Data Nucl. Data Tab. **40**, 291 (1988).
[81] A. E. Champagne and M. Wiescher, Ann. Rev. Nucl. Part. Sci. **42**, 39 (1992).
[82] H. Herndl *et al.*, Phys. Rev. C **52**, 1078 (1995).
[83] M. Wiescher, J. Görres, F.-K. Thielemann, and H. Ritter, Astron. Astrophys. **160**, 56 (1986).
[84] D. G. Sargood, Phys. Rep. **93**, 61 (1982).
[85] S. E. Woosley, W. A. Fowler, J. A. Holmes, and B. A. Zimmerman, At. Data Nucl. Data Tab. **22**, 371 (1978).
[86] J. A. Holmes, S. E. Woosley, W. A. Fowler, and B. A. Zimmerman, At. Data Nucl. Data Tab. **18**, 305 (1976).
[87] F. K. Thielemann, M. Arnould, and J. W. Truran, in *Capture Gamma Ray Spectroscopy 1987*, edited by K. Abrahams and P. van Assche (Institute of Physics Publishing, Bristol, 1988), p. 730.
[88] T. Rauscher, K.-L. Kratz, and F.-K. Thielemann, Phys. Rev. C (1997), submitted.
[89] W. D. Myers *et al.*, Phys. Rev. C **15**, 2031 (1977).
[90] A. Gilbert and A. G. W. Cameron, Can. J. Phys. **43**, 1446 (1965).
[91] A. Mengoni and Y. Nakajima, J. Nucl. Sci. Technol. **31**, 151 (1994).
[92] A. V. Ignatyuk, K. K. Istekov, and G. N. Smirenkin, Sov. J. Nucl. Phys. **29**, 450 (1979).
[93] J. Winger *et al.*, Phys. Rev. C **48**, 3097 (1993).
[94] J. Görres, M. Wiescher, and F.-K. Thielemann, Phys. Rev. C. **51**, 392 (1995).

[95] K. Nomoto, F.-K. Thielemann, and S. Miyaji, Astron. Astrophys. **149**, 239 (1985).
[96] G. M. Fuller, W. A. Fowler, and M. J. Newman, Ap. J. Suppl. **42**, 447 (1980).
[97] G. M. Fuller, W. A. Fowler, and M. J. Newman, Ap. J. **252**, 715 (1982).
[98] G. M. Fuller, W. A. Fowler, and M. J. Newman, Ap. J. Suppl. **48**, 279 (1982).
[99] P. Möller, J. R. Nix, and K.-L. Kratz, At. Data Nucl. Data Tab. (1997), submitted.
[100] J. Krumlinde and P. Möller, Nucl. Phys. A **417**, 419 (1984).
[101] F. Ryde, master Thesis, University Lund, Lund-MPh-93/16 (unpublished).
[102] B. A. Brown, A. Etchegoyen, W. D. M. Rae, and N. S. Godwin (unpublished).
[103] J. E. Koops and P. W. M. Glaudemans, Z. Phys. A **280**, 181 (1977).
[104] F. J. D. Serduke, R. D. Lawson, and D. H. Gloeckner, Nucl. Phys. A **256**, 45 (1976).
[105] H. Herndl and B. A. Brown, (1997), to be published.
[106] I. Hamamoto and X. Z. Zhang, Z. Phys. A **353**, 145 (1995).
[107] M. Hirsch, A. Staudt, K. Muto, and H. V. Klapdor-Kleingrothaus, At. Data Nucl. Data Tab. **53**, 165 (1993).
[108] G. T. Biehle and P. Vogel, Phys. Rev. C **46**, 1555 (1992).
[109] B. Blank *et al.*, Phys. Rev. Lett. **74**, 4611 (1995).
[110] E. Hourani *et al.*, Z. Phys. A **334**, 277 (1989).
[111] J. D. Robertson *et al.*, Phys. Rev. C **42**, 1922 (1990).
[112] M. F. Mohar *et al.*, Phys. Rev. Lett. **66**, 1571 (1991).
[113] R. Pfaff *et al.*, Phys. Rev. C **53**, 1753 (1996).
[114] A. Jokinen *et al.*, Z. Phys. A **355**, 227 (1996).
[115] S. E. Woosley, R. D. Hoffman, F.-K. Thielemann, and T. A. Weaver, in *Nuclei in the Cosmos '96*, edited by J. Görres, G. Mathews, S. Shore, and M. Wiescher (Nucl. Phys. A, Amsterdam, 1997), to be published.
[116] I. Iben Jr., A. V. Tutukov, and L. R. Yungelson, Ap. J. Suppl. **100**, 233 (1995).
[117] I. Lapidus, L. Nobili, and R. Turolla, Ap. J. **433**, 287 (1994).
[118] L. Nobili, R. Turolla, and I. Lapidus, Ap. J. **433**, 276 (1994).
[119] S. S. L. Ooi *et al.*, Phys. Rev. C **34**, 1153 (1986).
[120] J. Görres, T. Chapuran, D. P. Balamuth, and J. W. Arrison, Phys. Rev. Lett. **58**, 662 (1987).
[121] W. Gelletly *et al.*, Phys. Lett. **253**, 287 (1991).

TABLE I. 2p-capture reaction rates calculated with the FRDM (1992) mass model except for ^{68}Se, where in agreement with recent experiments a proton capture Q-value of -450 keV was assumed (see text). To obtain the reaction rate per second and target nucleus multiply with $(1/2)\rho^2 Y_p^2$.

T (GK)	$N_A^2 < 2\mathrm{p}, \gamma >$ (cm^6 s^{-1} mole^{-2})							
	^{68}Se	^{72}Kr	^{76}Sr	^{80}Zr	^{84}Mo	^{88}Ru	^{92}Pd	^{96}Cd
0.10	4.828E-56	5.268E-50	1.482E-48	1.046E-96	7.320E-72	1.973E-65	1.717E-65	4.944E-59
0.15	1.716E-43	1.497E-39	9.774E-39	3.066E-71	6.244E-55	7.982E-51	5.024E-51	6.642E-47
0.20	1.273E-36	1.104E-33	3.698E-33	8.126E-58	9.484E-46	8.920E-43	5.080E-43	4.748E-40
0.30	5.174E-29	4.962E-27	9.476E-27	1.551E-43	1.155E-35	8.676E-34	4.806E-34	3.208E-32
0.40	9.504E-25	3.064E-23	5.280E-23	7.478E-36	4.784E-30	1.071E-28	6.068E-29	1.085E-27
0.50	5.780E-22	1.324E-20	1.769E-20	5.462E-31	2.324E-26	2.580E-25	1.494E-25	1.193E-24
0.60	5.878E-20	1.096E-18	1.214E-18	1.299E-27	1.023E-23	7.172E-23	4.174E-23	1.953E-22
0.70	2.032E-18	3.084E-17	3.030E-17	3.822E-25	1.035E-21	5.204E-21	3.052E-21	9.666E-21
0.80	3.318E-17	4.274E-16	3.804E-16	2.916E-23	3.910E-20	1.520E-19	8.954E-20	2.118E-19
0.90	3.138E-16	3.612E-15	2.902E-15	8.766E-22	7.406E-19	2.338E-18	1.382E-18	2.592E-18
1.00	1.981E-15	2.092E-14	1.523E-14	1.361E-20	8.370E-18	2.222E-17	1.309E-17	2.076E-17
1.50	6.490E-13	5.738E-12	2.670E-12	5.484E-17	1.834E-14	2.742E-14	1.684E-14	1.622E-14
2.00	1.366E-11	1.151E-10	3.716E-11	3.294E-15	1.131E-12	1.200E-12	7.920E-13	6.190E-13
2.50	8.748E-11	7.438E-10	1.714E-10	3.586E-14	1.460E-11	1.249E-11	8.994E-12	6.410E-12
3.00	2.966E-10	2.646E-09	4.478E-10	1.651E-13	8.048E-11	5.942E-11	4.880E-11	3.332E-11
3.50	6.820E-10	6.656E-09	8.428E-10	4.656E-13	2.580E-10	1.696E-10	1.678E-10	1.139E-10
4.00	1.223E-09	1.330E-08	1.293E-09	9.696E-13	5.694E-10	3.410E-10	4.308E-10	2.934E-10
4.50	1.837E-09	2.296E-08	1.736E-09	1.643E-12	9.548E-10	5.336E-10	8.938E-10	6.122E-10
5.00	2.458E-09	3.484E-08	2.116E-09	2.406E-12	1.309E-09	6.876E-10	1.551E-09	1.109E-09
6.00	3.480E-09	6.070E-08	2.552E-09	3.788E-12	1.652E-09	7.870E-10	3.048E-09	2.550E-09
7.00	4.034E-09	6.990E-08	2.512E-09	4.412E-12	1.514E-09	6.706E-10	3.610E-09	4.004E-09
8.00	4.054E-09	5.560E-08	2.096E-09	4.098E-12	1.164E-09	4.846E-10	3.004E-09	4.578E-09
9.00	3.610E-09	3.518E-08	1.532E-09	3.202E-12	8.064E-10	3.188E-10	2.056E-09	4.160E-09
10.00	2.924E-09	2.022E-08	1.016E-09	2.216E-12	5.240E-10	1.983E-10	1.284E-09	3.232E-09

TABLE II. β-decay half-lives of the ground-states $T_{1/2}$ g.s. and the first excited 2^+ states $T_{1/2}$ 2^+ in even-even nuclei, together with the excitation energies of the first excited 2^+ states $E(2^+)$. We show experimental data (exp.) as well as the results from shell model calculatons using the code OXBASH (SM), from QRPA calculations, and from the application of the $N_\mathrm{p}N_\mathrm{n}$ correlation scheme. The half-life of ^{84}Mo obtained with OXBASH is highly uncertain due to the large deformation of this nucleus. In the network calculations we used the β-decay half-lives predicted by the shell model when no experimental data were available. Only when shell model calculations were not possible, we used the QRPA results.

Isotope	$T_{1/2}$ g.s. (s)			$T_{1/2}$ 2^+ (s) SM	$E(2^+)$ (keV)		
	exp.	SM	QRPA		exp.	$N_\mathrm{p}N_\mathrm{n}$	SM
^{64}Ge	63.7(25)			12	902 [119,120,67]	938	838
^{68}Se	35.5(7)			9.7	854 [67]	776	727
^{72}Kr	17.2(3)			5.2	709.1 [76]	548	832
^{76}Sr	8.9(3)		9.5		260.9 [67]	254	
^{80}Zr			6.9		290 [70]	254	
^{84}Mo		1.2	1.1	0.71	443.8 [121]	548	1041
^{88}Ru		0.71	0.73	0.45		776	961
^{92}Pd		0.46	0.55	0.30		938	929
^{96}Cd		0.66	0.84	0.40		1036	940

TABLE III. Branching ratios for the even-odd $N = Z + 1$ nuclei for the FRDM1 and FRDM2 reaction rates under XRB-conditions. The branching ratio $F_{\beta+}/F_\mathrm{total}$ is the fraction of the reaction flow in the β-decay branch (the remaining reaction flow is processed via proton capture). A large branching ratio therefore indicates a waiting point, delaying the rp-process with its full β-decay lifetime. Given are also the β-decay half-lives (with error if experimental) and the proton capture Q-values. In addition to the proton capture Q-values calculated with the FRDM (1992) mass model, we also show the experimental proton capture Q-value when available. In the cases, where only the mass of one nucleus is known, we also calculated the proton capture Q-value using the Audi and Wapstra extrapolation for the other nucleus.

Isotope	$T_{1/2}$ (s)	$Q_{(\mathrm{p},\gamma)}$ (MeV)			$F_{\beta+}/F_\mathrm{total}$ [a] (%)
		FRDM (1992) [50]	Exp	Exp+Audi [47]	
^{65}Ge	30.9(5)	1.679	2.949		0.00
^{69}Se	27.4(2)	1.539		2.962	0.00
^{73}Kr	27.0(12)	1.589	2.130		0.00
^{77}Sr	9.0(2)	1.049		1.943	0.03
^{81}Zr	15(5)	0.229		1.407	1.1
^{85}Mo	0.4	0.559			32
^{89}Ru	0.3	0.619			67
^{93}Pd	0.3	0.629			67
^{97}Cd	0.2	0.889			95

[a]With $Q_{(\mathrm{p},\gamma)}$ from FRDM (1992) except for ^{65}Ge, ^{73}Kr and ^{81}Zr

TABLE IV. Branching ratios for the even-even $N = Z$ nuclei for the FRDM1 and HIL reaction rates under XRB-conditions. The branching ratio $F_{\beta^+}/F_{\text{total}}$ is the fraction of the reaction flow in the β-decay branch (the remaining reaction flow is processed via proton capture). A large branching ratio therefore indicates a waiting point, delaying the rp-process with its full β-decay lifetime. Given are also the β-decay half-lives (with error if experimental) and the proton capture Q-values as used in FRDM1 and HIL, respectively.

Isotope	$T_{1/2}$ (s)	FRDM1		HIL	
		$Q_{(\text{p},\gamma)}$ (MeV)	$F_{\beta^+}/F_{\text{total}}$ (%)	$Q_{(\text{p},\gamma)}$ (MeV)	$F_{\beta^+}/F_{\text{total}}$ (%)
^{64}Ge	63.7(25)	0.129	0.15	0.858	0.035
^{68}Se	35.5(7)	-0.450	51	0.821	0.15
^{72}Kr	17.2(3)	-0.311	78	0.748	1.2
^{76}Sr	8.9(3)	-0.261	100	0.645	10
^{80}Zr	6.9	-1.181	100	0.175	66
^{84}Mo	1.1	-0.661	100	0.175	66
^{88}Ru	0.71	-0.511	100	0.044	96
^{92}Pd	0.46	-0.491	100	-0.108	100
^{96}Cd	0.66	-0.341	100	-0.277	100

TABLE V. Branching ratios for the most important waiting points for the FRDM1 reaction rates under ABH-conditions. The branching ratio $F_{\beta^+}/F_{\text{total}}$ is the fraction of the reaction flow in the β-decay branch (the remaining reaction flow is processed via proton capture). A large branching ratio therefore indicates a waiting point, delaying the rp-process with its full β-decay lifetime. Given are also the β-decay half-lives (with error if experimental) and the proton capture Q-values as used in FRDM1.

Isotope	$T_{1/2}$ (s)	$Q_{(\text{p},\gamma)}$ (MeV)	$F_{\beta^+}/F_{\text{total}}$ (%)
^{56}Ni	$5.18(1)\cdot 10^5$	0.767	1.3
^{64}Ge	63.7(25)	0.129	96
^{68}Se	35.5(7)	-0.450	100
^{72}Kr	17.2(3)	-0.311	100
^{76}Sr	8.9(3)	-0.261	100
^{77}Sr	9.0(2)	1.049	74
^{80}Zr	6.9	-1.181	100
^{81}Zr	15(5)	1.407	68
^{83}Nb	4.1(3)	2.241	47
^{85}Mo	0.4	0.559	100
^{86}Mo	19.6(11)	1.101	29
^{90}Ru	13(5)	1.134	30
^{94}Pd	9(5)	0.960	93
^{98}Cd	9.2(3)	1.292	99

TABLE VI. The astrophysical (p,γ) reaction rates $N_A < \sigma v >$ in the reaction rate set FRDM1, calculated with the Hauser-Feshbach code SMOKER. Rates smaller than 10^{-99} have been set to zero. To obtain the number of reactions per second and target nucleus (λ), multiply with ρY_p.

Z	N	A	$N_A < \sigma v >$ (cm^3 s^{-1} mole^{-1})								Q (MeV)
			$T_9 = 0.1$	$T_9 = 0.3$	$T_9 = 0.5$	$T_9 = 1.0$	$T_9 = 1.5$	$T_9 = 2.0$	$T_9 = 2.5$	$T_9 = 3.0$	
32	32	64	1.69E-23	1.09E-11	1.20E-07	1.16E-03 5.40E-02	4.82E-01	2.02E+00	5.64E+00	1.28E-01	
32	33	65	2.64E-23	2.40E-11	3.84E-07	1.21E-02 1.15E+00	1.51E+01	7.78E+01	2.40E+02	2.95E+00	
32	34	66	3.58E-23	2.66E-11	4.06E-07	1.04E-02 8.60E-01	1.05E+01	5.30E+01	1.65E+02	2.31E+00	
32	35	67	7.90E-24	6.82E-12	9.78E-08	3.32E-03 3.96E-01	6.64E+00	4.30E+01	1.62E+02	3.51E+00	
32	36	68	1.43E-23	1.13E-11	1.66E-07	5.38E-03 6.38E-01	1.11E+01	7.62E+01	3.08E+02	3.39E+00	
32	37	69	1.26E-23	1.15E-11	1.69E-07	6.58E-03 9.46E-01	1.90E+01	1.45E+02	6.24E+02	4.54E+00	
32	38	70	1.47E-23	1.24E-11	1.86E-07	6.98E-03 9.94E-01	2.04E+01	1.61E+02	7.30E+02	4.62E+00	
32	39	71	1.49E-23	1.42E-11	2.14E-07	8.90E-03 1.37E+00	2.84E+01	2.18E+02	9.26E+02	5.61E+00	
32	40	72	1.84E-23	1.54E-11	2.34E-07	9.18E-03 1.41E+00	3.14E+01	2.64E+02	1.26E+03	5.66E+00	
32	41	73	2.14E-23	1.71E-11	2.66E-07	6.12E-03 4.34E-01	4.80E+00	2.36E+01	7.62E+01	6.85E+00	
32	42	74	2.52E-22	2.22E-10	3.20E-06	1.09E-01 1.43E+01	2.66E+02	1.87E+03	7.32E+03	6.90E+00	
33	32	65	4.62E-24	6.52E-12	1.22E-07	3.32E-03 2.46E-01	2.60E+00	1.14E+01	3.10E+01	2.59E+00	
33	33	66	5.04E-24	8.62E-12	1.71E-07	6.68E-03 6.48E-01	8.04E+00	3.84E+01	1.09E+02	3.36E+00	
33	34	67	3.30E-24	5.64E-12	1.12E-07	4.76E-03 5.28E-01	7.58E+00	4.14E+01	1.33E+02	4.88E+00	
33	35	68	3.42E-24	5.48E-12	9.36E-08	3.90E-03 5.06E-01	8.64E+00	5.52E+01	2.02E+02	4.71E+00	
33	36	69	2.66E-24	4.26E-12	7.48E-08	3.46E-03 5.36E-01	1.12E+01	8.72E+01	3.82E+02	6.13E+00	
33	37	70	1.36E-24	2.14E-12	3.82E-08	1.86E-03 3.16E-01	7.24E+00	6.08E+01	2.82E+02	6.41E+00	
33	38	71	1.42E-24	2.22E-12	4.02E-08	1.99E-03 3.46E-01	8.24E+00	7.32E+01	3.62E+02	7.29E+00	
33	39	72	1.41E-24	2.44E-12	4.30E-08	1.79E-03 2.44E-01	4.48E+00	3.10E+01	1.23E+02	7.28E+00	
33	40	73	1.69E-24	2.66E-12	4.86E-08	2.50E-03 4.60E-01	1.15E+01	1.04E+02	5.12E+02	8.55E+00	
33	41	74	1.30E-24	1.27E-12	1.89E-08	4.42E-04 4.36E-02	7.08E-01	4.88E+00	2.04E+01	8.60E+00	
33	42	75	5.12E-23	7.70E-11	1.50E-06	5.56E-02 5.08E+00	6.20E+01	3.08E+02	9.52E+02	9.51E+00	
34	35	69	7.32E-25	1.79E-12	3.78E-08	1.50E-03 1.53E-01	2.06E+00	1.09E+01	3.42E+01	1.54E+00	
34	36	70	3.34E-25	7.64E-13	1.66E-08	7.50E-04 9.56E-02	1.65E+00	1.08E+01	4.16E+01	2.21E+00	
34	37	71	4.58E-25	1.41E-12	3.10E-08	1.79E-03 2.90E-01	5.96E+00	4.44E+01	1.86E+02	3.00E+00	
34	38	72	4.72E-25	1.26E-12	2.84E-08	1.58E-03 2.50E-01	5.14E+00	3.94E+01	1.71E+02	2.95E+00	
34	39	73	2.42E-24	7.80E-12	1.82E-07	1.02E-02 1.47E+00	2.68E+01	1.80E+02	6.90E+02	4.38E+00	
34	40	74	2.88E-25	8.62E-13	1.98E-08	1.23E-03 2.22E-01	5.20E+00	4.44E+01	2.12E+02	4.22E+00	
34	41	75	3.22E-25	9.86E-13	2.24E-08	1.49E-03 2.98E-01	7.68E+00	7.08E+01	3.56E+02	5.41E+00	
34	42	76	1.05E-23	2.96E-11	6.66E-07	3.38E-02 4.88E+00	9.54E+01	7.10E+02	3.04E+03	5.27E+00	
34	43	77	1.13E-23	3.00E-11	6.78E-07	3.70E-02 4.90E+00	7.84E+01	4.56E+02	1.50E+03	6.14E+00	
34	44	78	1.20E-23	3.30E-11	6.88E-07	3.30E-02 5.22E+00	1.13E+02	9.18E+02	4.14E+03	6.33E+00	
35	34	69	9.02E-26	4.10E-13	1.01E-08	4.82E-04 5.22E-02	7.10E-01	3.68E+00	1.13E+01	2.43E+00	
35	35	70	1.59E-25	8.14E-13	2.00E-08	1.11E-03 1.43E-01	2.26E+00	1.32E+01	4.40E+01	3.10E+00	
35	36	71	6.74E-26	3.58E-13	9.12E-09	5.70E-04 8.74E-02	1.64E+00	1.12E+01	4.32E+01	4.53E+00	
35	37	72	1.24E-25	6.60E-13	1.64E-08	1.10E-03 1.97E-01	4.40E+00	3.50E+01	1.52E+02	4.99E+00	
35	38	73	7.62E-26	3.74E-13	9.40E-09	6.32E-04 1.21E-01	3.00E+00	2.66E+01	1.31E+02	5.90E+00	
35	39	74	2.20E-24	1.08E-11	2.10E-07	1.01E-02 1.50E+00	2.90E+01	2.08E+02	8.30E+02	6.22E+00	
35	40	75	6.80E-26	3.28E-13	8.86E-09	6.76E-04 1.48E-01	4.12E+00	4.14E+01	2.26E+02	7.13E+00	
35	41	76	4.86E-26	3.00E-13	8.34E-09	6.96E-04 1.48E-01	3.74E+00	3.26E+01	1.53E+02	7.17E+00	
35	42	77	2.08E-24	1.06E-11	2.92E-07	2.10E-02 3.88E+00	9.04E+01	7.60E+02	3.52E+03	8.20E+00	
35	43	78	1.43E-24	2.88E-12	4.42E-08	2.14E-03 3.28E-01	6.40E+00	4.56E+01	1.83E+02	8.28E+00	
35	44	79	2.40E-24	1.13E-11	3.04E-07	2.08E-02 3.40E+00	6.64E+01	4.62E+02	1.79E+03	9.11E+00	
36	37	73	2.92E-26	2.72E-13	1.00E-08	9.00E-04 1.55E-01	3.02E+00	2.06E+01	7.82E+01	2.13E+00	
36	38	74	2.04E-26	1.94E-13	7.04E-09	5.34E-04 8.28E-02	1.58E+00	1.13E+01	4.70E+01	2.34E+00	
36	39	75	6.70E-26	6.16E-13	2.32E-08	2.56E-03 5.32E-01	1.21E+01	9.32E+01	3.88E+02	3.53E+00	
36	40	76	6.22E-26	5.86E-13	2.20E-08	1.96E-03 3.20E-01	6.10E+00	4.24E+01	1.70E+02	3.14E+00	
36	41	77	9.92E-26	8.28E-13	3.04E-08	3.00E-03 5.64E-01	1.19E+01	8.86E+01	3.64E+02	4.05E+00	
36	42	78	1.20E-25	1.03E-12	3.58E-08	3.08E-03 5.32E-01	1.10E+01	8.44E+01	3.72E+02	3.92E+00	
36	43	79	1.33E-25	1.18E-12	3.96E-08	3.38E-03 6.60E-01	1.56E+01	1.31E+02	6.10E+02	5.02E+00	
36	44	80	1.39E-25	1.20E-12	4.02E-08	3.18E-03 5.62E-01	1.26E+01	1.06E+02	5.12E+02	4.85E+00	
36	45	81	1.30E-25	1.16E-12	3.74E-08	3.18E-03 6.66E-01	1.65E+01	1.40E+02	6.26E+02	5.78E+00	
36	46	82	9.56E-26	8.68E-13	3.02E-08	2.70E-03 5.58E-01	1.42E+01	1.30E+02	6.68E+02	5.77E+00	
36	47	83	1.10E-25	9.58E-13	3.36E-08	2.52E-03 3.58E-01	5.98E+00	3.68E+01	1.30E+02	7.06E+00	
36	48	84	1.22E-25	1.08E-12	3.46E-08	2.82E-03 5.80E-01	1.47E+01	1.30E+02	6.16E+02	7.02E+00	

37	37	74	6.86E-27	1.09E-13	4.92E-09	5.60E-04	1.00E-01	1.89E+00	1.21E+01	4.24E+01	3.17E+00
37	38	75	3.70E-27	6.20E-14	2.80E-09	2.82E-04	4.74E-02	8.80E-01	5.76E+00	2.10E+01	4.47E+00
37	39	76	1.13E-26	1.87E-13	8.54E-09	1.17E-03	2.78E-01	6.90E+00	5.62E+01	2.40E+02	4.79E+00
37	40	77	1.21E-26	2.06E-13	9.50E-09	1.23E-03	2.74E-01	6.56E+00	5.24E+01	2.24E+02	5.63E+00
37	41	78	6.78E-26	9.22E-13	3.04E-08	2.62E-03	5.42E-01	1.34E+01	1.15E+02	5.32E+02	5.83E+00
37	42	79	2.70E-26	4.32E-13	1.84E-08	2.06E-03	4.38E-01	1.04E+01	8.42E+01	3.68E+02	6.80E+00
37	43	80	6.16E-26	8.94E-13	3.20E-08	2.88E-03	5.88E-01	1.42E+01	1.21E+02	5.58E+02	6.64E+00
37	44	81	2.74E-26	4.34E-13	1.76E-08	1.88E-03	4.24E-01	1.12E+01	1.03E+02	5.10E+02	7.84E+00
37	45	82	2.38E-26	3.66E-13	1.49E-08	1.46E-03	2.86E-01	6.56E+00	5.36E+01	2.44E+02	7.90E+00
37	46	83	1.94E-26	3.14E-13	1.31E-08	1.58E-03	3.92E-01	1.08E+01	9.98E+01	4.80E+02	8.86E+00
37	47	84	1.38E-26	1.58E-13	3.32E-09	1.89E-04	4.24E-02	1.29E+00	1.36E+01	7.52E+01	8.64E+00
37	48	85	2.46E-26	3.70E-13	1.56E-08	1.61E-03	2.76E-01	5.14E+00	3.34E+01	1.20E+02	9.64E+00
38	39	77	1.49E-27	4.48E-14	2.46E-09	2.78E-04	4.66E-02	8.54E-01	5.52E+00	2.00E+01	1.06E+00
38	40	78	1.52E-27	4.60E-14	2.64E-09	3.44E-04	6.36E-02	1.28E+00	9.14E+00	3.70E+01	2.47E+00
38	41	79	5.34E-27	1.53E-13	8.56E-09	1.36E-03	3.22E-01	7.76E+00	6.16E+01	2.60E+02	5.17E+00
38	42	80	4.16E-27	1.19E-13	6.40E-09	8.82E-04	1.93E-01	4.66E+00	3.96E+01	1.88E+02	3.00E+00
38	43	81	7.16E-27	1.99E-13	9.86E-09	1.30E-03	3.10E-01	8.12E+00	7.14E+01	3.34E+02	3.96E+00
38	44	82	8.76E-27	2.38E-13	1.24E-08	1.60E-03	3.32E-01	7.74E+00	6.44E+01	3.02E+02	3.61E+00
38	45	83	3.24E-27	9.52E-14	5.06E-09	7.92E-04	2.20E-01	6.62E+00	6.66E+01	3.52E+02	4.73E+00
38	46	84	6.90E-27	1.90E-13	1.00E-08	1.48E-03	3.60E-01	9.72E+00	9.12E+01	4.72E+02	4.49E+00
38	47	85	8.34E-27	2.20E-13	1.13E-08	1.71E-03	4.58E-01	1.33E+01	1.29E+02	6.64E+02	5.47E+00
38	48	86	9.12E-27	2.48E-13	1.23E-08	1.61E-03	3.82E-01	1.04E+01	1.01E+02	5.36E+02	5.78E+00
38	49	87	1.07E-26	2.92E-13	1.44E-08	1.69E-03	3.98E-01	1.09E+01	1.02E+02	5.10E+02	6.71E+00
38	50	88	1.09E-26	2.80E-13	1.30E-08	1.51E-03	3.52E-01	9.68E+00	9.40E+01	4.98E+02	7.07E+00
39	39	78	3.64E-28	1.74E-14	1.20E-09	2.10E-04	4.22E-02	8.10E-01	5.10E+00	1.75E+01	2.36E+00
39	40	79	1.99E-28	1.07E-14	7.20E-10	9.80E-05	1.75E-02	3.28E-01	2.14E+00	7.78E+00	3.63E+00
39	41	80	7.32E-28	3.90E-14	2.74E-09	4.30E-04	8.42E-02	1.71E+00	1.22E+01	4.92E+01	2.79E+00
39	42	81	9.38E-28	4.82E-14	3.14E-09	5.72E-04	1.44E-01	3.58E+00	2.86E+01	1.19E+02	5.47E+00
39	43	82	2.10E-27	1.05E-13	6.50E-09	1.19E-03	3.44E-01	1.02E+01	9.76E+01	4.82E+02	5.56E+00
39	44	83	5.90E-28	3.12E-14	2.08E-09	4.16E-04	1.24E-01	3.74E+00	3.64E+01	1.82E+02	6.00E+00
39	45	84	1.43E-27	7.36E-14	4.68E-09	9.22E-04	2.94E-01	9.58E+00	1.01E+02	5.44E+02	6.21E+00
39	46	85	1.43E-27	7.04E-14	4.52E-09	8.60E-04	2.60E-01	8.18E+00	8.44E+01	4.54E+02	7.25E+00
39	47	86	1.66E-27	8.24E-14	5.08E-09	9.32E-04	2.84E-01	9.02E+00	9.30E+01	4.96E+02	7.36E+00
39	48	87	1.91E-27	9.36E-14	5.64E-09	9.44E-04	2.72E-01	8.46E+00	8.70E+01	4.66E+02	7.90E+00
39	49	88	2.16E-27	1.02E-13	5.90E-09	9.32E-04	2.54E-01	7.46E+00	7.24E+01	3.66E+02	7.86E+00
39	50	89	2.36E-27	1.09E-13	6.14E-09	8.84E-04	2.36E-01	7.02E+00	7.04E+01	3.72E+02	8.36E+00
40	41	81	8.42E-28	4.60E-14	3.60E-09	5.00E-04	8.16E-02	1.41E+00	8.74E+00	3.12E+01	1.41E+00
40	42	82	6.22E-29	5.96E-15	4.98E-10	9.04E-05	2.06E-02	4.92E-01	4.00E+00	1.78E+01	2.06E+00
40	43	83	1.36E-28	1.26E-14	1.06E-09	2.80E-04	9.40E-02	2.94E+00	2.84E+01	1.39E+02	1.93E+00
40	44	84	1.02E-28	9.40E-15	8.08E-10	1.71E-04	4.18E-02	1.04E+00	8.76E+00	4.10E+01	2.11E+00
40	45	85	2.36E-28	2.06E-14	1.67E-09	4.40E-04	1.56E-01	5.34E+00	5.72E+01	3.10E+02	3.96E+00
40	46	86	1.65E-28	1.45E-14	1.16E-09	2.46E-04	6.94E-02	2.02E+00	1.99E+01	1.05E+02	3.66E+00
40	47	87	3.52E-28	2.90E-14	2.24E-09	5.30E-04	1.70E-01	5.38E+00	5.50E+01	2.94E+02	4.01E+00
40	48	88	2.30E-28	1.93E-14	1.48E-09	2.90E-04	7.76E-02	2.20E+00	2.12E+01	1.12E+02	4.24E+00
40	49	89	2.62E-28	2.18E-14	1.57E-09	3.12E-04	9.54E-02	3.08E+00	3.28E+01	1.84E+02	5.08E+00
40	50	90	5.04E-28	4.00E-14	2.88E-09	5.12E-04	1.31E-01	3.66E+00	3.58E+01	1.94E+02	5.16E+00
40	51	91	5.42E-28	4.10E-14	2.82E-09	4.84E-04	1.14E-01	2.66E+00	1.99E+01	7.74E+01	5.85E+00
40	52	92	5.56E-28	4.14E-14	2.60E-09	4.46E-04	1.30E-01	3.82E+00	3.46E+01	1.54E+02	6.04E+00
41	41	82	2.06E-29	3.24E-15	3.32E-10	7.06E-05	1.38E-02	2.48E-01	1.47E+00	4.78E+00	1.26E+00
41	42	83	1.20E-29	2.04E-15	2.02E-10	3.66E-05	7.28E-03	1.49E-01	1.05E+00	4.08E+00	2.24E+00
41	43	84	6.84E-29	9.72E-15	8.80E-10	2.32E-04	6.64E-02	1.70E+00	1.35E+01	5.52E+01	3.79E+00
41	44	85	1.99E-29	3.26E-15	3.42E-10	1.10E-04	3.74E-02	1.15E+00	1.08E+01	5.16E+01	4.66E+00
41	45	86	6.36E-29	1.01E-14	9.64E-10	2.98E-04	1.13E-01	4.02E+00	4.36E+01	2.38E+02	5.16E+00
41	46	87	3.36E-29	5.26E-15	5.14E-10	1.51E-04	5.30E-02	1.76E+00	1.83E+01	9.76E+01	5.81E+00
41	47	88	7.50E-29	1.07E-14	9.20E-10	2.36E-04	8.10E-02	2.72E+00	2.88E+01	1.55E+02	6.22E+00
41	48	89	5.78E-29	8.36E-15	7.64E-10	2.02E-04	6.78E-02	2.22E+00	2.30E+01	1.25E+02	6.88E+00
41	49	90	9.66E-29	1.40E-14	1.26E-09	3.28E-04	1.11E-01	3.70E+00	4.00E+01	2.24E+02	6.84E+00
41	50	91	1.08E-28	1.53E-14	1.34E-09	3.18E-04	9.96E-02	3.18E+00	3.38E+01	1.90E+02	7.46E+00
41	51	92	6.04E-29	1.71E-14	1.29E-09	2.82E-04	8.32E-02	2.36E+00	2.10E+01	9.62E+01	7.64E+00
41	52	93	1.29E-28	1.70E-14	1.23E-09	1.79E-04	4.04E-02	9.96E-01	8.26E+00	3.64E+01	8.49E+00

42	43	85	5.20E-30	1.49E-15	1.88E-10	5.70E-05	1.52E-02	3.60E-01	2.68E+00	1.05E+01	5.64E-01
42	44	86	4.06E-30	1.21E-15	1.50E-10	3.74E-05	9.10E-03	2.18E-01	1.78E+00	7.92E+00	1.10E+00
42	45	87	9.22E-30	2.58E-15	3.22E-10	1.18E-04	4.24E-02	1.34E+00	1.29E+01	6.26E+01	1.91E+00
42	46	88	7.08E-30	1.97E-15	2.34E-10	6.56E-05	1.90E-02	5.32E-01	4.92E+00	2.44E+01	2.08E+00
42	47	89	1.88E-29	4.70E-15	5.34E-10	1.88E-04	6.92E-02	2.28E+00	2.32E+01	1.19E+02	3.57E+00
42	48	90	1.05E-29	2.66E-15	3.08E-10	8.88E-05	2.62E-02	7.48E-01	7.08E+00	3.64E+01	3.10E+00
42	49	91	1.20E-29	3.00E-15	3.36E-10	1.05E-04	3.66E-02	1.22E+00	1.31E+01	7.32E+01	4.02E+00
42	50	92	1.35E-29	3.30E-15	3.66E-10	9.88E-05	2.84E-02	8.10E-01	7.84E+00	4.16E+01	4.09E+00
42	51	93	2.68E-29	6.56E-15	6.54E-10	1.73E-04	6.12E-02	2.16E+00	2.40E+01	1.35E+02	4.64E+00
42	52	94	2.88E-29	6.62E-15	6.30E-10	1.53E-04	5.20E-02	1.83E+00	2.10E+01	1.25E+02	4.89E+00
42	53	95	2.84E-29	5.80E-15	5.48E-10	1.47E-04	4.56E-02	1.28E+00	1.09E+01	4.72E+01	5.40E+00
42	54	96	2.74E-29	5.68E-15	4.94E-10	1.20E-04	4.54E-02	1.72E+00	1.99E+01	1.12E+02	5.72E+00
43	43	86	1.04E-30	4.68E-16	7.12E-11	3.08E-05	9.74E-03	2.46E-01	1.84E+00	7.08E+00	2.36E+00
43	44	87	7.80E-31	3.94E-16	6.22E-11	2.34E-05	6.76E-03	1.68E-01	1.30E+00	5.22E+00	3.54E+00
43	45	88	3.24E-30	1.40E-15	2.08E-10	1.01E-04	3.84E-02	1.16E+00	1.02E+01	4.52E+01	3.74E+00
43	46	89	1.42E-30	6.84E-16	1.03E-10	4.46E-05	1.68E-02	5.38E-01	5.24E+00	2.58E+01	4.55E+00
43	47	90	4.02E-30	1.81E-15	2.50E-10	1.10E-04	4.66E-02	1.71E+00	1.89E+01	1.03E+02	4.58E+00
43	48	91	2.26E-30	9.74E-16	1.37E-10	5.60E-05	2.10E-02	6.92E-01	7.06E+00	3.70E+01	5.55E+00
43	49	92	4.54E-30	1.93E-15	2.64E-10	1.09E-04	4.32E-02	1.53E+00	1.68E+01	9.38E+01	5.62E+00
43	50	93	2.90E-30	1.23E-15	1.67E-10	6.28E-05	2.26E-02	7.42E-01	7.76E+00	4.26E+01	6.25E+00
43	51	94	3.34E-30	1.41E-15	1.77E-10	6.10E-05	2.32E-02	8.36E-01	9.42E+00	5.42E+01	6.58E+00
43	52	95	3.70E-30	1.47E-15	1.72E-10	5.58E-05	2.14E-02	7.74E-01	8.74E+00	4.96E+01	7.34E+00
43	53	96	4.76E-30	1.33E-15	1.65E-10	4.12E-05	1.16E-02	3.28E-01	3.04E+00	1.48E+01	7.58E+00
43	54	97	3.64E-30	1.21E-15	1.38E-10	4.10E-05	1.19E-02	3.14E-01	2.62E+00	1.15E+01	8.29E+00
44	45	89	3.70E-31	3.22E-16	6.10E-11	2.66E-05	7.88E-03	1.94E-01	1.48E+00	5.88E+00	6.22E-01
44	46	90	2.96E-31	2.58E-16	4.72E-11	1.73E-05	4.72E-03	1.17E-01	9.62E-01	4.28E+00	1.14E+00
44	47	91	6.86E-31	5.46E-16	9.48E-11	4.62E-05	1.77E-02	5.68E-01	5.50E+00	2.70E+01	1.86E+00
44	48	92	5.02E-31	3.78E-16	6.44E-11	2.50E-05	7.64E-03	2.12E-01	1.92E+00	9.40E+00	2.00E+00
44	49	93	1.01E-30	7.24E-16	1.23E-10	5.90E-05	2.26E-02	7.34E-01	7.38E+00	3.84E+01	2.77E+00
44	50	94	5.06E-31	3.90E-16	6.86E-11	2.60E-05	7.12E-03	1.73E-01	1.40E+00	6.28E+00	3.06E+00
44	51	95	7.28E-31	5.30E-16	8.42E-11	3.60E-05	1.39E-02	4.86E-01	5.36E+00	3.06E+01	3.46E+00
44	52	96	8.26E-31	5.74E-16	8.62E-11	3.18E-05	1.12E-02	3.80E-01	4.26E+00	2.54E+01	3.81E+00
44	53	97	1.59E-30	1.04E-15	1.40E-10	5.10E-05	2.18E-02	8.88E-01	1.13E+01	7.10E+01	4.35E+00
44	54	98	1.57E-30	9.32E-16	1.23E-10	4.36E-05	1.78E-02	7.22E-01	9.42E+00	6.36E+01	4.58E+00
45	45	90	6.46E-32	9.18E-17	2.14E-11	1.37E-05	4.90E-03	1.29E-01	9.92E-01	3.86E+00	2.02E+00
45	46	91	6.00E-32	9.00E-17	2.00E-11	9.68E-06	2.92E-03	7.24E-02	5.64E-01	2.32E+00	3.19E+00
45	47	92	2.14E-31	2.88E-16	5.94E-11	3.68E-05	1.55E-02	5.04E-01	4.74E+00	2.22E+01	3.39E+00
45	48	93	1.12E-31	1.36E-16	2.82E-11	1.69E-05	6.68E-03	2.12E-01	2.02E+00	9.84E+00	4.37E+00
45	49	94	3.06E-31	3.62E-16	7.10E-11	4.36E-05	1.98E-02	7.14E-01	7.58E+00	4.00E+01	4.40E+00
45	50	95	1.42E-31	1.71E-16	3.58E-11	1.99E-05	7.26E-03	2.22E-01	2.14E+00	1.09E+01	5.12E+00
45	51	96	1.60E-31	2.02E-16	3.96E-11	2.26E-05	1.01E-02	3.82E-01	4.32E+00	2.46E+01	5.46E+00
45	52	97	3.22E-31	3.92E-16	7.06E-11	3.42E-05	1.47E-02	5.62E-01	6.68E+00	4.04E+01	6.00E+00
45	53	98	3.66E-31	4.16E-16	6.92E-11	3.22E-05	1.44E-02	5.76E-01	6.94E+00	4.14E+01	6.27E+00
45	54	99	3.82E-31	3.98E-16	6.20E-11	2.70E-05	1.26E-02	5.48E-01	7.24E+00	4.72E+01	7.00E+00
46	47	93	3.66E-32	9.26E-17	2.40E-11	1.26E-05	3.82E-03	9.40E-02	7.20E-01	2.92E+00	6.24E-01
46	48	94	2.38E-32	5.54E-17	1.37E-11	6.56E-06	1.89E-03	4.72E-02	3.90E-01	1.77E+00	9.60E-01
46	49	95	5.14E-32	1.05E-16	2.54E-11	1.64E-05	6.64E-03	2.12E-01	2.06E+00	1.02E+01	1.87E+00
46	50	96	3.30E-32	6.74E-17	1.67E-11	8.94E-06	2.86E-03	7.78E-02	6.90E-01	3.36E+00	1.74E+00
46	51	97	6.34E-32	1.33E-16	3.22E-11	2.00E-05	8.06E-03	2.68E-01	2.74E+00	1.45E+01	2.37E+00
46	52	98	4.10E-32	8.80E-17	2.04E-11	1.08E-05	4.00E-03	1.30E-01	1.36E+00	7.58E+00	2.71E+00
46	53	99	8.30E-32	1.66E-16	3.48E-11	1.96E-05	8.92E-03	3.54E-01	4.34E+00	2.70E+01	3.29E+00
46	54	100	9.04E-32	1.60E-16	3.18E-11	1.62E-05	6.98E-03	2.74E-01	3.46E+00	2.28E+01	3.29E+00
47	47	94	8.86E-33	3.68E-17	1.19E-11	1.01E-05	4.04E-03	1.15E-01	9.56E-01	4.02E+00	1.95E+00
47	48	95	7.56E-32	3.50E-17	1.14E-11	6.32E-06	1.85E-03	4.62E-02	3.70E-01	1.59E+00	3.20E+00
47	49	96	5.92E-33	5.40E-17	1.46E-11	1.23E-05	5.78E-03	1.98E-01	1.93E+00	9.32E+00	3.26E+00
47	50	97	6.44E-32	2.34E-17	7.46E-12	5.98E-06	2.36E-03	7.26E-02	6.78E-01	3.32E+00	4.32E+00
47	51	98	1.51E-32	5.52E-17	1.62E-11	1.43E-05	7.06E-03	2.62E-01	2.84E+00	1.54E+01	4.11E+00
47	52	99	1.75E-32	3.26E-17	9.58E-12	7.14E-06	3.14E-03	1.11E-01	1.19E+00	6.52E+00	4.83E+00
47	53	100	2.54E-32	9.24E-17	2.30E-11	1.52E-05	7.50E-03	3.06E-01	3.74E+00	2.28E+01	4.88E+00
47	54	101	1.25E-32	3.88E-17	9.22E-12	5.98E-06	3.00E-03	1.29E-01	1.67E+00	1.09E+01	5.48E+00

47	55	102	3.40E-32	3.50E-17	8.00E-12	4.94E-06	2.74E-03	1.34E-01	1.94E+00	1.36E+01	5.94E+00
47	56	103	3.06E-32	2.92E-17	6.58E-12	4.14E-06	2.34E-03	1.18E-01	1.81E+00	1.35E+01	6.47E+00
47	57	104	8.62E-33	2.44E-17	5.34E-12	3.52E-06	1.93E-03	9.06E-02	1.26E+00	8.56E+00	6.51E+00
47	58	105	8.20E-33	2.00E-17	4.48E-12	3.20E-06	1.99E-03	1.03E-01	1.50E+00	1.04E+01	7.35E+00
47	59	106	2.60E-33	1.28E-17	2.54E-12	1.59E-06	8.22E-04	3.54E-02	4.48E-01	2.80E+00	7.34E+00
47	60	107	1.85E-33	1.46E-17	3.42E-12	2.58E-06	1.25E-03	4.60E-02	4.94E-01	2.66E+00	8.14E+00
48	49	97	1.74E-32	1.90E-17	6.70E-12	4.42E-06	1.46E-03	3.90E-02	3.28E-01	1.46E+00	8.92E-01
48	50	98	1.00E-32	9.78E-18	3.54E-12	2.60E-06	8.80E-04	2.40E-02	2.08E-01	9.94E-01	1.29E+00
48	51	99	5.78E-33	2.26E-17	7.92E-12	7.34E-06	3.34E-03	1.13E-01	1.13E+00	5.76E+00	2.06E+00
48	52	100	1.57E-32	1.20E-17	4.62E-12	3.64E-06	1.39E-03	4.38E-02	4.44E-01	2.44E+00	1.94E+00
48	53	101	8.12E-33	2.54E-17	8.64E-12	7.10E-06	3.22E-03	1.16E-01	1.27E+00	7.04E+00	2.06E+00
48	54	102	2.44E-32	2.72E-17	8.56E-12	5.60E-06	2.32E-03	8.48E-02	9.92E-01	6.10E+00	2.47E+00
48	55	103	3.44E-33	2.68E-17	7.14E-12	5.42E-06	2.90E-03	1.28E-01	1.68E+00	1.11E+01	2.71E+00
48	56	104	4.02E-33	1.28E-17	3.40E-12	2.48E-06	1.31E-03	5.94E-02	8.34E-01	5.96E+00	2.79E+00
48	57	105	7.20E-33	1.83E-17	4.80E-12	3.78E-06	2.36E-03	1.26E-01	2.00E+00	1.55E+01	3.57E+00
48	58	106	7.30E-33	1.46E-17	3.98E-12	3.26E-06	2.04E-03	1.12E-01	1.86E+00	1.54E+01	3.72E+00
48	59	107	7.02E-33	1.22E-17	3.38E-12	2.94E-06	2.04E-03	1.20E-01	2.04E+00	1.63E+01	4.41E+00
48	60	108	2.96E-32	5.72E-18	1.74E-12	1.49E-06	1.04E-03	6.58E-02	1.23E+00	1.10E+01	4.52E+00
49	49	98	1.51E-33	7.64E-18	3.22E-12	3.08E-06	1.26E-03	3.76E-02	3.32E-01	1.52E+00	1.89E+00
49	50	99	9.98E-33	6.90E-18	2.86E-12	1.95E-06	6.10E-04	1.59E-02	1.34E-01	6.18E-01	2.99E+00
49	51	100	1.63E-33	8.52E-18	3.64E-12	4.58E-06	2.38E-03	8.48E-02	8.62E-01	4.40E+00	3.32E+00
49	52	101	2.42E-33	4.56E-18	2.10E-12	2.42E-06	1.17E-03	4.12E-02	4.32E-01	2.32E+00	4.13E+00
49	53	102	1.43E-33	1.26E-17	5.10E-12	5.62E-06	3.22E-03	1.38E-01	1.72E+00	1.06E+01	4.09E+00
49	54	103	8.70E-34	6.48E-18	2.42E-12	2.26E-06	1.14E-03	4.52E-02	5.36E-01	3.24E+00	4.24E+00
49	55	104	1.07E-33	8.70E-18	2.74E-12	2.34E-06	1.41E-03	7.04E-02	1.04E+00	7.46E+00	4.46E+00
49	56	105	1.36E-33	7.20E-18	2.18E-12	1.84E-06	1.12E-03	5.68E-02	8.54E-01	6.28E+00	5.24E+00
49	57	106	2.08E-33	8.06E-18	2.50E-12	2.32E-06	1.60E-03	9.24E-02	1.56E+00	1.27E+01	5.24E+00
49	58	107	4.56E-33	6.30E-18	2.04E-12	1.94E-06	1.39E-03	8.48E-02	1.53E+00	1.32E+01	5.74E+00
49	59	108	2.86E-33	8.82E-18	2.64E-12	2.32E-06	1.58E-03	9.24E-02	1.59E+00	1.31E+01	5.82E+00
49	60	109	3.20E-33	2.42E-18	8.44E-13	8.84E-07	6.96E-04	4.62E-02	8.84E-01	7.90E+00	6.64E+00

TABLE VII. The astrophysical (p,α) reaction rates $N_A < \sigma v >$ in the reaction rate set FRDM1, calculated with the Hauser-Feshbach code SMOKER. Rates smaller than 10^{-99} have been set to zero. To obtain the number of reactions per second and target nucleus (λ) multiply with ρY_p.

Z	N	A	$N_A < \sigma v >$ (cm^3 s^{-1} mole^{-1})								Q (MeV)
			$T_9 = 0.1$	$T_9 = 0.3$	$T_9 = 0.5$	$T_9 = 1.0$	$T_9 = 1.5$	$T_9 = 2.0$	$T_9 = 2.5$	$T_9 = 3.0$	
32	32	64	0.00E+00	7.46E-71	3.72E-47	9.92E-26	6.08E-17	6.10E-12	1.21E-08	2.84E-06	-2.03E+00
32	33	65	1.87E-53	1.85E-32	1.28E-23	7.04E-14	3.14E-09	2.66E-06	3.22E-04	1.24E-02	4.50E-01
32	34	66	3.62E-76	4.80E-39	6.92E-28	5.76E-16	2.16E-10	5.28E-07	1.14E-04	6.14E-03	-6.80E-02
32	35	67	1.20E-24	8.02E-27	2.96E-20	7.06E-12	8.74E-08	3.34E-05	2.38E-03	6.24E-02	1.05E+00
32	36	68	5.32E-51	9.52E-32	3.54E-23	2.38E-13	1.58E-08	1.73E-05	2.42E-03	9.86E-02	5.40E-01
32	37	69	6.12E-28	3.20E-24	1.36E-18	7.28E-11	6.46E-07	1.83E-04	9.38E-03	1.79E-01	1.49E+00
32	38	70	8.54E-38	1.83E-26	1.86E-19	3.90E-11	5.94E-07	2.94E-04	2.50E-02	7.26E-01	1.18E+00
32	39	71	6.58E-36	6.34E-22	2.30E-16	1.69E-09	4.54E-06	6.00E-04	1.82E-02	2.42E-01	2.04E+00
32	40	72	1.05E-41	2.36E-24	3.62E-18	1.50E-10	1.34E-06	4.30E-04	2.42E-02	4.78E-01	1.60E+00
32	41	73	2.62E-30	2.88E-20	7.82E-16	2.02E-10	2.12E-07	2.82E-05	1.20E-03	2.42E-02	2.47E+00
32	42	74	1.58E-38	1.20E-23	1.92E-17	7.30E-10	3.78E-06	6.74E-04	2.30E-02	3.12E-01	1.58E+00
33	32	65	8.16E-66	8.34E-35	9.66E-26	7.66E-16	5.14E-11	5.56E-08	7.60E-06	3.08E-04	3.49E-01
33	33	66	7.62E-38	9.94E-30	3.36E-22	2.10E-13	3.54E-09	1.69E-06	1.41E-04	4.22E-03	7.49E-01
33	34	67	4.24E-32	9.78E-22	7.86E-16	2.90E-09	3.44E-06	2.92E-04	7.26E-03	8.94E-02	2.65E+00
33	35	68	1.41E-31	5.00E-17	2.60E-13	7.72E-08	1.02E-04	1.00E-02	2.36E-01	2.42E+00	2.40E+00
33	36	69	3.46E-34	5.78E-20	9.82E-15	2.20E-08	3.50E-05	3.48E-03	8.68E-02	9.68E-01	3.40E+00
33	37	70	1.88E-33	6.64E-19	7.64E-14	1.13E-07	1.48E-04	1.27E-02	2.70E-01	2.58E+00	3.18E+00
33	38	71	7.52E-31	2.14E-17	1.33E-12	7.28E-07	5.60E-04	3.54E-02	6.32E-01	5.40E+00	3.95E+00
33	39	72	3.68E-30	2.30E-17	1.63E-12	5.72E-07	2.86E-04	1.40E-02	2.20E-01	1.80E+00	3.73E+00
33	40	73	7.46E-30	2.74E-16	1.13E-11	3.36E-06	1.82E-03	8.88E-02	1.29E+00	9.20E+00	4.47E+00
33	41	74	1.34E-27	6.42E-16	7.64E-12	1.39E-07	2.82E-05	1.33E-03	2.76E-02	3.16E-01	3.91E+00
33	42	75	1.42E-27	7.98E-15	4.22E-10	6.24E-05	1.24E-02	2.72E-01	2.26E+00	1.14E+01	4.41E+00
34	35	69	8.64E-97	1.86E-60	6.66E-43	2.18E-24	2.24E-16	1.01E-11	1.33E-08	2.52E-06	6.38E-01
34	36	70	0.00E+00	7.08E-57	7.56E-39	1.22E-21	2.60E-14	5.26E-10	4.32E-07	5.96E-05	-1.05E+00
34	37	71	4.56E-63	1.14E-38	3.52E-28	9.60E-17	2.08E-11	4.08E-08	8.44E-06	4.70E-04	2.77E-01
34	38	72	4.92E-68	1.98E-41	3.18E-30	1.21E-17	1.07E-11	4.62E-08	1.46E-05	1.02E-03	5.07E-02
34	39	73	3.14E-25	1.29E-29	6.44E-23	5.94E-14	1.64E-09	1.00E-06	9.66E-05	3.18E-03	9.89E-01
34	40	74	3.48E-52	3.86E-36	7.88E-27	9.08E-16	1.93E-10	3.54E-07	6.34E-05	2.96E-03	5.44E-01
34	41	75	1.39E-12	1.97E-31	1.16E-24	8.52E-15	4.76E-10	3.56E-07	3.58E-05	1.18E-03	9.25E-01
34	42	76	1.21E-57	3.94E-34	7.34E-25	3.42E-14	5.04E-09	7.50E-06	1.12E-03	4.38E-02	5.65E-01
34	43	77	6.10E-38	4.58E-23	1.17E-18	1.17E-12	4.92E-09	1.75E-06	1.46E-04	4.68E-03	1.12E+00
34	44	78	5.96E-51	1.98E-31	3.32E-23	2.26E-13	1.39E-08	1.13E-05	1.09E-03	3.10E-02	8.71E-01
35	34	69	3.26E-86	1.32E-46	1.46E-33	5.72E-20	9.66E-14	6.00E-10	2.54E-07	2.30E-05	-3.15E-01
35	35	70	9.06E-79	1.59E-38	1.22E-27	1.12E-16	1.41E-11	2.14E-08	3.72E-06	1.83E-04	3.28E-01
35	36	71	5.98E-34	3.00E-26	1.16E-19	6.26E-12	2.92E-08	5.74E-06	2.52E-04	4.70E-03	1.62E+00
35	37	72	1.81E-34	8.54E-19	1.04E-15	6.02E-11	9.64E-08	1.95E-05	1.03E-03	2.18E-02	1.98E+00
35	38	73	1.12E-29	5.26E-21	3.98E-17	7.64E-11	2.72E-07	4.66E-05	1.62E-03	2.28E-02	2.92E+00
35	39	74	1.27E-29	2.28E-15	4.88E-12	9.84E-08	4.00E-05	2.72E-03	6.22E-02	6.94E-01	3.02E+00
35	40	75	4.82E-29	1.58E-18	2.42E-14	1.39E-08	1.90E-05	1.94E-03	5.00E-02	5.74E-01	3.62E+00
35	41	76	4.62E-40	2.24E-18	1.02E-13	1.60E-08	1.17E-05	1.05E-03	2.92E-02	3.82E-01	2.79E+00
35	42	77	8.76E-31	2.76E-17	2.80E-12	1.86E-06	1.37E-03	8.28E-02	1.43E+00	1.19E+01	3.84E+00
35	43	78	8.44E-32	4.52E-18	4.40E-13	1.83E-07	1.07E-04	6.22E-03	1.16E-01	1.10E+00	3.58E+00
35	44	79	2.08E-30	4.88E-17	5.02E-12	2.84E-06	1.32E-03	4.96E-02	5.86E-01	3.74E+00	4.05E+00
36	36	72	0.00E+00	1.33E-88	7.44E-59	1.55E-32	4.84E-22	4.18E-16	3.50E-12	2.36E-09	-2.81E+00
36	37	73	0.00E+00	3.24E-65	6.10E-45	9.50E-26	1.19E-17	6.46E-13	1.02E-09	2.28E-07	-1.41E+00
36	38	74	0.00E+00	7.06E-62	9.68E-43	1.91E-24	1.21E-16	4.66E-12	6.02E-09	1.17E-06	-1.21E+00
36	39	75	1.57E-87	1.22E-46	2.86E-33	1.19E-19	1.60E-13	8.52E-10	3.30E-07	2.88E-05	-1.96E-01
36	40	76	4.18E-92	1.14E-51	1.04E-36	2.18E-21	1.25E-14	1.36E-10	8.00E-08	9.16E-06	-5.53E-01
36	41	77	2.26E-87	3.30E-44	1.75E-31	1.29E-18	8.54E-13	2.94E-09	8.24E-07	5.52E-05	-1.90E-04
36	42	78	1.25E-96	2.52E-47	2.04E-33	1.83E-19	3.72E-13	2.48E-09	1.08E-06	9.86E-05	-1.55E-01
36	43	79	6.78E+00	6.30E-33	8.72E-27	1.81E-16	2.48E-11	3.32E-08	4.98E-06	2.20E-04	7.11E-01
36	44	80	1.24E-56	2.70E-41	8.08E-31	2.82E-18	2.98E-12	1.43E-08	4.88E-06	3.58E-04	2.05E-01
36	45	81	9.38E-65	1.26E-39	1.07E-28	2.02E-17	2.60E-12	3.98E-09	7.54E-07	4.20E-05	6.23E-01
36	46	82	1.79E-71	2.82E-40	1.14E-29	8.86E-18	3.76E-12	8.90E-09	1.63E-06	6.86E-05	3.44E-01
36	47	83	1.08E-50	1.06E-33	4.26E-27	7.52E-18	7.78E-13	1.11E-09	1.62E-07	6.26E-06	7.72E-01

36	48	84	1.96E-69	7.38E-40	2.10E-29	7.44E-18	1.32E-12	1.53E-09	1.67E-07	5.08E-06	4.08E-01
37	37	74	7.16E-84	1.43E-43	6.84E-32	1.77E-19	1.10E-13	3.68E-10	1.01E-07	6.66E-06	9.60E-02
37	38	75	7.10E-51	1.26E-31	6.22E-24	5.12E-15	1.19E-10	6.38E-08	5.32E-06	1.52E-04	1.10E+00
37	39	76	2.72E-47	6.12E-31	1.42E-23	2.96E-14	1.08E-09	5.96E-07	4.48E-05	1.15E-03	1.27E+00
37	40	77	1.17E-40	3.94E-25	7.74E-19	2.18E-11	9.10E-08	1.51E-05	5.22E-04	7.50E-03	2.21E+00
37	41	78	9.16E-31	1.06E-20	2.92E-17	1.18E-11	4.56E-08	1.41E-05	9.14E-04	2.14E-02	2.17E+00
37	42	79	1.71E-30	4.26E-20	1.51E-15	3.78E-09	9.32E-06	1.16E-03	3.08E-02	3.38E-01	3.05E+00
37	43	80	1.36E-29	7.70E-18	3.60E-14	3.04E-09	3.14E-06	4.26E-04	1.67E-02	2.78E-01	2.86E+00
37	44	81	4.12E-34	8.64E-20	1.28E-14	2.04E-08	3.16E-05	3.38E-03	8.92E-02	1.01E+00	3.57E+00
37	45	82	3.14E-35	3.78E-19	1.46E-14	2.58E-09	2.36E-06	2.86E-04	1.05E-02	1.71E-01	3.12E+00
37	46	83	1.23E-34	1.75E-20	4.14E-15	1.05E-08	1.51E-05	1.31E-03	2.76E-02	2.58E-01	3.69E+00
37	47	84	1.83E-38	1.06E-23	8.34E-18	6.08E-11	1.43E-07	1.90E-05	6.20E-04	9.02E-03	2.81E+00
37	48	85	1.38E-36	3.22E-22	1.21E-16	3.30E-10	3.62E-07	2.54E-05	4.96E-04	4.80E-03	3.29E+00
38	39	77	0.00E+00	3.70E-82	5.48E-55	9.72E-31	8.88E-21	4.28E-15	2.40E-11	1.19E-08	-2.32E+00
38	40	78	0.00E+00	8.26E-62	7.88E-43	1.40E-24	1.40E-16	7.16E-12	1.05E-08	2.14E-06	-1.09E+00
38	41	79	2.06E-97	3.04E-47	6.52E-34	5.32E-20	1.36E-13	9.56E-10	4.02E-07	3.48E-05	-1.34E-01
38	42	80	0.00E+00	1.99E-54	2.86E-38	3.72E-22	6.44E-15	1.28E-10	1.03E-07	1.38E-05	-6.12E-01
38	43	81	4.36E-77	3.10E-41	5.52E-31	9.68E-19	1.15E-12	5.42E-09	1.62E-06	1.02E-04	2.74E-01
38	44	82	0.00E+00	2.24E-50	8.20E-36	1.02E-20	8.56E-14	1.14E-09	6.88E-07	7.38E-05	-3.50E-01
38	45	83	1.76E-78	2.64E-42	7.24E-32	2.22E-19	3.66E-13	2.24E-09	8.32E-07	6.40E-05	2.39E-01
38	46	84	0.00E+00	9.88E-51	2.68E-36	2.96E-21	2.76E-14	4.12E-10	2.68E-07	2.98E-05	-3.23E-01
38	47	85	1.00E-96	4.90E-48	6.56E-35	5.36E-21	1.25E-14	7.34E-11	2.54E-08	1.87E-06	-4.90E-02
38	48	86	0.00E+00	2.46E-55	3.10E-39	8.16E-23	2.08E-15	4.48E-11	3.36E-08	3.90E-06	-5.84E-01
38	49	87	0.00E+00	1.07E-51	1.59E-37	5.56E-23	2.52E-16	2.56E-12	1.44E-09	1.60E-07	-2.64E-01
38	50	88	0.00E+00	2.44E-60	3.08E-42	1.23E-24	4.16E-17	8.62E-13	6.38E-10	8.18E-08	-8.89E-01
39	39	78	0.00E+00	6.22E-49	2.52E-35	3.24E-21	1.29E-14	1.37E-10	8.14E-08	9.30E-06	-2.09E-01
39	40	79	2.00E-49	7.90E-31	1.54E-23	4.80E-15	8.16E-11	3.98E-08	3.28E-06	9.50E-05	1.33E+00
39	41	80	0.00E+00	6.80E-55	2.38E-39	1.00E-23	1.42E-16	2.88E-12	2.50E-09	3.62E-07	-5.18E-01
39	42	81	2.60E-44	8.80E-28	4.92E-21	4.90E-13	4.12E-09	1.11E-06	5.54E-05	1.05E-03	2.02E+00
39	43	82	1.14E-33	2.36E-23	2.02E-19	1.47E-12	2.10E-08	1.04E-05	7.66E-04	1.79E-02	2.15E+00
39	44	83	1.02E-45	7.90E-29	7.02E-22	1.67E-13	2.90E-09	1.38E-06	1.05E-04	2.68E-03	1.89E+00
39	45	84	6.56E-32	1.25E-20	4.72E-17	3.80E-12	7.98E-09	2.74E-06	2.62E-04	9.68E-03	2.15E+00
39	46	85	2.38E-36	2.72E-23	8.78E-18	1.25E-10	6.16E-07	1.22E-04	4.64E-03	6.84E-02	3.02E+00
39	47	86	2.86E-38	4.06E-22	2.74E-17	1.23E-11	2.34E-08	5.34E-06	3.40E-04	9.00E-03	2.38E+00
39	48	87	1.53E-41	9.48E-26	1.85E-19	6.86E-12	4.42E-08	1.08E-05	4.94E-04	8.50E-03	2.49E+00
39	49	88	8.42E-38	2.10E-25	3.36E-21	8.40E-15	9.44E-11	6.68E-08	8.56E-06	3.50E-04	1.67E+00
39	50	89	5.18E-48	1.39E-30	2.60E-23	1.71E-14	3.72E-10	1.67E-07	1.11E-05	2.54E-04	1.68E+00
40	41	81	0.00E+00	6.56E-64	1.06E-44	1.88E-26	1.94E-18	1.08E-13	1.73E-10	3.82E-08	-1.03E+00
40	42	82	0.00E+00	8.94E-63	6.66E-44	9.90E-26	1.16E-17	7.28E-13	1.27E-09	2.94E-07	-9.72E-01
40	43	83	1.09E-39	3.64E-26	1.82E-21	4.34E-14	1.06E-09	7.88E-07	7.96E-05	2.34E-03	1.70E+00
40	44	84	0.00E+00	4.28E-70	5.38E-48	1.51E-27	8.34E-19	1.06E-13	2.76E-10	8.36E-08	-1.46E+00
40	45	85	0.00E+00	3.70E-49	1.19E-35	3.04E-21	1.61E-14	1.69E-10	8.82E-08	8.68E-06	-9.80E-02
40	46	86	0.00E+00	1.30E-57	8.84E-41	4.60E-24	1.68E-16	5.54E-12	6.38E-09	1.12E-06	-6.13E-01
40	47	87	0.00E+00	7.26E-52	2.66E-37	3.16E-22	2.82E-15	4.02E-11	2.58E-08	2.94E-06	-2.52E-01
40	48	88	0.00E+00	8.48E-63	3.76E-44	6.58E-26	1.05E-17	7.96E-13	1.51E-09	3.56E-07	-9.16E-01
40	49	89	0.00E+00	7.38E-60	2.50E-42	4.26E-25	2.12E-17	7.72E-13	9.10E-10	1.58E-07	-7.26E-01
40	50	90	0.00E+00	6.18E-62	1.81E-43	1.80E-25	1.55E-17	6.86E-13	8.60E-10	1.50E-07	-8.91E-01
40	51	91	1.92E-40	2.20E-29	6.56E-25	1.14E-17	3.74E-13	5.26E-10	1.19E-07	8.44E-06	1.27E+00
40	52	92	8.76E-35	4.28E-20	1.42E-14	4.08E-08	4.28E-05	2.46E-03	3.52E-02	2.40E-01	4.11E+00
41	41	82	7.24E-58	2.44E-35	8.98E-27	4.04E-17	2.18E-12	2.26E-09	3.34E-07	1.57E-05	1.02E+00
41	42	83	1.66E-37	3.66E-22	1.79E-16	3.12E-10	3.38E-07	3.00E-05	7.86E-04	1.01E-02	2.81E+00
41	43	84	1.36E-32	2.06E-20	1.25E-15	3.92E-09	6.64E-06	4.98E-04	8.42E-03	6.54E-02	3.62E+00
41	44	85	1.43E-55	2.46E-35	7.26E-27	5.42E-17	3.30E-12	3.00E-09	3.54E-07	1.30E-05	1.26E+00
41	45	86	7.76E-53	2.40E-33	5.08E-26	3.56E-16	3.68E-11	3.90E-08	4.38E-06	1.45E-04	1.50E+00
41	46	87	3.00E-51	4.36E-33	1.80E-25	3.44E-16	1.79E-11	1.72E-08	2.16E-06	8.24E-05	1.71E+00
41	47	88	1.48E-34	3.50E-24	6.64E-21	6.68E-15	8.74E-11	7.36E-08	1.03E-05	4.32E-04	1.94E+00
41	48	89	6.38E-47	8.66E-30	8.52E-23	2.26E-14	3.96E-10	1.88E-07	1.44E-05	3.80E-04	2.09E+00
41	49	90	1.60E-35	8.54E-27	2.42E-23	2.64E-16	1.29E-11	2.04E-08	3.88E-06	1.91E-04	1.55E+00
41	50	91	6.58E-50	4.60E-32	1.41E-24	1.67E-15	5.92E-11	4.18E-08	4.12E-06	1.30E-04	1.85E+00
41	51	92	6.30E-34	1.98E-19	5.96E-15	3.48E-10	1.09E-07	6.26E-06	1.49E-04	2.04E-03	3.28E+00

41	52	93	3.22E-32	1.51E-17	2.04E-12	5.38E-07	1.57E-04	4.96E-03	5.54E-02	3.52E-01	6.42E+00
42	43	85	3.20E-96	4.56E-47	1.29E-34	2.46E-21	6.04E-15	5.40E-11	2.98E-08	3.30E-06	1.39E-01
42	44	86	0.00E+00	3.18E-64	9.70E-45	3.74E-26	6.82E-18	5.90E-13	1.32E-09	3.74E-07	-9.72E-01
42	45	87	0.00E+00	7.40E-80	2.88E-54	2.36E-31	9.02E-22	3.28E-16	1.72E-12	8.70E-10	-1.91E+00
42	46	88	0.00E+00	7.52E-60	4.00E-42	7.30E-25	4.80E-17	2.50E-12	4.18E-09	9.84E-07	-6.85E-01
42	47	89	0.00E+00	1.02E-54	3.96E-39	1.81E-23	2.70E-16	5.66E-12	5.00E-09	7.36E-07	-3.12E-01
42	48	90	0.00E+00	4.38E-68	3.68E-47	2.72E-27	1.47E-18	2.14E-13	6.36E-10	2.12E-07	-1.12E+00
42	49	91	0.00E+00	1.04E-69	3.76E-48	4.38E-28	1.99E-19	2.38E-14	6.12E-11	1.85E-08	-1.27E+00
42	50	92	0.00E+00	5.92E-71	9.14E-49	5.10E-28	4.98E-19	9.32E-14	3.12E-10	1.10E-07	-1.36E+00
42	51	93	7.20E-70	3.04E-38	6.14E-30	1.38E-19	5.88E-14	2.78E-10	1.17E-07	1.12E-05	7.18E-01
42	52	94	8.60E-41	9.42E-25	1.39E-18	3.50E-11	2.30E-07	6.44E-05	3.42E-03	6.56E-02	3.09E+00
42	53	95	2.52E-39	1.04E-23	1.12E-17	1.05E-10	2.06E-07	2.16E-05	5.88E-04	7.88E-03	3.61E+00
42	54	96	2.68E-40	1.36E-24	1.62E-18	3.52E-11	1.45E-07	2.40E-05	8.20E-04	1.17E-02	3.28E+00
43	43	86	3.54E-72	1.19E-40	9.76E-31	1.29E-19	2.84E-14	5.56E-11	1.13E-08	6.44E-07	8.32E-01
43	44	87	0.00E+00	6.42E-49	2.60E-36	5.24E-23	1.25E-16	1.07E-12	5.70E-10	6.14E-08	1.68E-01
43	45	88	0.00E+00	1.63E-51	1.95E-37	5.28E-23	3.06E-16	4.30E-12	3.26E-09	4.60E-07	-3.70E-02
43	46	89	6.14E-84	2.14E-44	2.12E-33	3.66E-21	2.92E-15	1.18E-11	3.56E-09	2.48E-07	5.69E-01
43	47	90	7.50E-59	2.04E-37	1.87E-28	7.88E-18	1.31E-12	2.34E-09	4.56E-07	2.50E-05	1.27E+00
43	48	91	4.18E-73	2.12E-42	2.58E-32	1.86E-20	1.09E-14	3.48E-11	8.80E-09	5.48E-07	8.48E-01
43	49	92	3.42E-63	1.80E-37	4.60E-30	8.78E-20	4.76E-14	2.26E-10	8.18E-08	6.44E-06	9.32E-01
43	50	93	3.18E-57	4.94E-37	1.64E-28	2.98E-18	4.02E-13	6.70E-10	1.24E-07	6.38E-06	1.43E+00
43	51	94	2.80E-35	4.86E-21	5.76E-17	4.72E-12	4.84E-09	8.22E-07	4.54E-05	1.11E-03	2.91E+00
43	52	95	4.56E-35	7.76E-20	2.86E-14	6.04E-08	6.22E-05	4.10E-03	6.98E-02	5.54E-01	5.65E+00
43	53	96	9.34E-35	1.15E-19	3.14E-14	2.18E-08	1.18E-05	5.96E-04	9.92E-03	8.86E-02	5.85E+00
43	54	97	4.64E-35	4.90E-20	1.66E-14	2.24E-08	1.35E-05	6.02E-04	8.32E-03	6.26E-02	6.05E+00
44	45	89	0.00E+00	0.00E+00	7.28E-69	6.14E-39	5.62E-27	3.66E-20	1.19E-15	2.10E-12	-3.21E+00
44	46	90	0.00E+00	3.22E-98	2.82E-65	8.78E-37	3.30E-25	1.44E-18	3.64E-14	5.46E-11	-2.88E+00
44	47	91	0.00E+00	5.24E-92	1.11E-61	2.62E-35	1.63E-24	2.74E-18	3.84E-14	3.84E-11	-2.50E+00
44	48	92	0.00E+00	8.50E-94	1.66E-62	2.20E-35	2.70E-24	6.78E-18	1.26E-13	1.54E-10	-2.60E+00
44	49	93	0.00E+00	6.82E-83	2.48E-56	1.28E-32	1.14E-22	7.18E-17	5.48E-13	3.60E-10	-1.94E+00
44	50	94	0.00E+00	1.51E-79	4.88E-54	5.46E-31	3.30E-21	1.74E-15	1.19E-11	7.12E-09	-1.71E+00
44	51	95	5.34E-95	4.28E-46	9.52E-35	2.84E-22	8.38E-16	1.33E-11	1.27E-08	2.16E-06	3.55E-01
44	52	96	3.28E-47	1.06E-29	9.00E-23	2.48E-14	7.76E-10	7.18E-07	9.90E-05	4.12E-03	2.40E+00
44	53	97	8.82E-43	3.40E-25	1.28E-18	2.48E-11	7.58E-08	1.22E-05	4.88E-04	8.94E-03	2.91E+00
44	54	98	1.97E-45	2.26E-28	1.13E-21	1.83E-13	4.04E-09	2.72E-06	2.80E-04	9.00E-03	2.66E+00
45	45	90	0.00E+00	1.15E-76	7.30E-53	2.30E-31	4.04E-22	1.30E-16	7.32E-13	4.18E-10	-1.47E+00
45	46	91	0.00E+00	2.96E-61	6.92E-44	3.08E-27	1.32E-19	6.58E-15	1.20E-11	3.16E-09	-4.83E-01
45	47	92	0.00E+00	1.15E-62	2.26E-44	5.38E-27	3.76E-19	2.28E-14	4.62E-11	1.31E-08	-6.15E-01
45	48	93	0.00E+00	1.24E-54	6.04E-40	4.48E-25	4.76E-18	1.06E-13	1.10E-10	1.95E-08	-4.70E-02
45	49	94	0.00E+00	2.58E-54	2.14E-39	1.85E-24	1.74E-17	3.34E-13	3.10E-10	5.02E-08	2.00E-03
45	50	95	2.80E-95	3.70E-48	3.40E-36	4.00E-23	8.92E-17	7.44E-13	3.92E-10	4.28E-08	4.88E-01
45	51	96	7.72E-38	1.02E-23	1.24E-19	3.72E-14	1.49E-10	7.34E-08	8.78E-06	3.72E-04	2.51E+00
45	52	97	6.64E-37	4.74E-21	3.04E-15	1.11E-08	1.58E-05	1.39E-03	3.12E-02	3.22E-01	4.84E+00
45	53	98	9.90E-37	4.78E-21	2.46E-15	8.06E-09	1.22E-05	1.16E-03	2.86E-02	3.18E-01	5.14E+00
45	54	99	3.56E-36	1.26E-20	4.50E-15	1.05E-08	1.52E-05	1.47E-03	3.56E-02	3.78E-01	5.42E+00
46	47	93	0.00E+00	0.00E+00	2.62E-72	5.72E-41	1.70E-28	2.22E-21	1.15E-16	2.88E-13	-3.39E+00
46	48	94	0.00E+00	0.00E+00	9.62E-71	7.94E-40	1.99E-27	2.50E-20	1.31E-15	3.28E-12	-3.27E+00
46	49	95	0.00E+00	2.32E-95	9.42E-64	1.18E-36	1.35E-25	3.44E-19	6.72E-15	8.70E-12	-2.54E+00
46	50	96	0.00E+00	1.63E-96	2.48E-64	1.36E-36	2.92E-25	1.05E-18	2.52E-14	3.78E-11	-2.63E+00
46	51	97	0.00E+00	1.01E-62	2.78E-44	1.02E-26	8.18E-19	5.00E-14	9.64E-11	2.56E-08	-5.54E-01
46	52	98	2.26E-53	5.92E-34	4.08E-26	1.25E-16	1.22E-11	2.24E-08	4.96E-06	2.96E-04	1.91E+00
46	53	99	1.08E-43	5.60E-28	3.92E-22	1.31E-14	2.32E-10	1.82E-07	2.48E-05	1.09E-03	2.33E+00
46	54	100	4.62E-50	1.21E-31	3.02E-24	2.86E-15	1.29E-10	1.34E-07	1.99E-05	9.08E-04	2.22E+00
47	47	94	0.00E+00	9.30E-88	1.25E-59	6.64E-35	1.52E-24	1.74E-18	2.04E-14	1.85E-11	-1.99E+00
47	48	95	0.00E+00	2.22E-72	1.68E-50	1.71E-30	1.03E-21	1.90E-16	7.48E-13	3.28E-10	-1.05E+00
47	49	96	0.00E+00	2.54E-71	9.04E-50	6.28E-30	3.36E-21	5.80E-16	2.18E-12	9.22E-10	-1.01E+00
47	50	97	0.00E+00	1.08E-55	1.75E-40	2.50E-25	3.40E-18	8.90E-14	1.05E-10	2.04E-08	-3.89E-03
47	51	98	1.37E-30	1.63E-38	2.32E-30	2.16E-19	5.32E-14	1.25E-10	3.18E-08	2.24E-06	1.47E+00
47	52	99	2.96E-35	3.68E-23	6.16E-17	5.36E-10	8.76E-07	8.10E-05	1.93E-03	2.16E-02	4.32E+00
47	53	100	1.70E-45	2.74E-20	1.01E-15	5.46E-10	1.08E-06	1.63E-04	5.76E-03	8.34E-02	4.50E+00

47	54	101	1.30E-40	2.68E-23	2.44E-17	2.04E-10	5.92E-07	9.12E-05	3.16E-03	4.56E-02	4.93E+00
47	55	102	4.78E-41	4.32E-23	3.92E-17	3.28E-10	1.01E-06	1.63E-04	5.86E-03	8.66E-02	5.01E+00
47	56	103	5.10E-43	4.30E-23	3.18E-17	2.30E-10	7.38E-07	1.21E-04	4.28E-03	6.00E-02	5.30E+00
47	57	104	3.10E-36	2.20E-23	2.66E-17	2.00E-10	3.96E-07	4.68E-05	1.43E-03	1.99E-02	5.18E+00
47	58	105	6.56E-45	2.10E-22	8.52E-17	4.00E-10	9.30E-07	1.06E-04	2.62E-03	2.78E-02	5.72E+00
47	59	106	5.52E-39	3.84E-23	3.16E-17	1.09E-10	1.55E-07	1.59E-05	4.62E-04	6.48E-03	5.41E+00
47	60	107	6.48E-35	1.15E-22	5.80E-17	1.42E-10	1.36E-07	9.56E-06	2.04E-04	2.26E-03	5.85E+00
48	49	97	0.00E+00	0.00E+00	4.80E-76	3.92E-43	6.24E-30	1.69E-22	1.32E-17	4.40E-14	-3.64E+00
48	50	98	0.00E+00	0.00E+00	3.34E-73	1.09E-41	6.96E-29	1.25E-21	7.78E-17	2.26E-13	-3.36E+00
48	51	99	0.00E+00	1.00E-65	1.56E-47	1.12E-28	2.90E-20	3.16E-15	8.98E-12	3.26E-09	-7.69E-01
48	52	100	3.38E-59	5.22E-41	1.38E-31	2.30E-20	1.15E-14	5.78E-11	2.78E-08	3.24E-06	1.11E+00
48	53	101	1.80E-85	9.36E-27	3.02E-22	2.92E-16	6.90E-12	9.78E-09	2.00E-06	1.11E-04	1.99E+00
48	54	102	2.68E-55	2.10E-32	9.14E-25	7.04E-16	2.98E-11	3.32E-08	5.56E-06	2.90E-04	2.17E+00
48	55	103	0.00E+00	2.48E-28	2.38E-21	5.64E-15	4.98E-11	3.56E-08	4.98E-06	2.32E-04	2.37E+00
48	56	104	1.07E-51	1.94E-34	7.44E-27	4.10E-17	6.36E-12	1.47E-08	3.66E-06	2.36E-04	2.11E+00
48	57	105	6.58E-91	4.48E-28	4.76E-22	2.62E-15	4.96E-11	4.54E-08	6.32E-06	2.66E-04	2.54E+00
48	58	106	9.72E-63	6.36E-33	4.10E-25	6.46E-16	5.20E-11	8.44E-08	1.66E-05	8.90E-04	2.52E+00
48	59	107	2.84E-44	5.74E-28	2.84E-22	2.18E-14	4.66E-10	2.80E-07	2.54E-05	7.62E-04	2.99E+00
48	60	108	2.18E-94	2.46E-33	5.62E-25	4.28E-16	2.98E-11	4.06E-08	6.28E-06	2.60E-04	2.68E+00
49	49	98	0.00E+00	1.65E-95	1.09E-66	6.34E-39	2.30E-27	9.80E-21	2.56E-16	4.16E-13	-2.58E+00
49	50	99	0.00E+00	2.42E-77	6.20E-56	8.08E-34	4.18E-24	2.26E-18	1.75E-14	1.27E-11	-1.45E+00
49	51	100	2.26E-87	5.78E-51	4.14E-38	3.36E-24	1.38E-17	1.89E-13	1.58E-10	2.58E-08	4.27E-01
49	52	101	3.46E-34	6.70E-27	4.46E-20	1.44E-12	4.00E-09	5.94E-07	2.28E-05	4.02E-04	3.49E+00
49	53	102	2.48E-40	1.77E-21	5.62E-17	8.98E-12	1.03E-08	1.69E-06	8.76E-05	2.06E-03	3.52E+00
49	54	103	1.98E-40	1.67E-24	2.78E-18	2.86E-11	7.16E-08	9.88E-06	3.32E-04	4.90E-03	4.53E+00
49	55	104	3.54E-47	1.74E-21	1.11E-16	3.92E-11	6.60E-08	1.17E-05	5.48E-04	1.08E-02	4.54E+00
49	56	105	1.68E-36	3.64E-22	1.96E-17	4.86E-11	1.57E-07	2.84E-05	1.11E-03	1.73E-02	4.80E+00
49	57	106	8.94E-40	2.28E-22	1.26E-16	1.34E-10	2.00E-07	3.32E-05	1.62E-03	3.46E-02	4.90E+00
49	58	107	2.38E-39	2.18E-24	2.36E-18	2.12E-11	8.54E-08	2.08E-05	1.14E-03	2.44E-02	5.28E+00
49	59	108	6.80E-39	2.28E-21	6.42E-16	3.52E-10	3.64E-07	4.52E-05	1.76E-03	3.14E-02	5.09E+00
49	60	109	1.46E-43	5.60E-25	5.86E-19	5.08E-12	2.14E-08	5.26E-06	2.82E-04	5.82E-03	5.51E+00

TABLE VIII. The astrophysical (γ,p) photodisintegration rates λ in the reaction rate set FRDM1, calculated from the (p,γ) reaction rates in Tab. VI using detailed balance. Rates smaller than 10^{-99} have been set to zero.

Z	N	A	λ (s^{-1})								Q (MeV)
			$T_9 = 0.1$	$T_9 = 0.3$	$T_9 = 0.5$	$T_9 = 1.0$	$T_9 = 1.5$	$T_9 = 2.0$	$T_9 = 2.5$	$T_9 = 3.0$	
32	32	64	0.00E+00	6.52E-85	1.09E-46	1.63E-16	7.04E-06	2.22E+00	5.24E+03	1.02E+06	-5.03E+00
32	33	65	0.00E+00	1.98E-83	2.54E-46	6.40E-17	1.68E-06	4.30E-01	9.00E+02	1.61E+05	-4.86E+00
32	34	66	0.00E+00	0.00E+00	2.64E-59	7.66E-23	4.98E-10	2.14E-03	2.56E+01	1.51E+04	-6.26E+00
32	35	67	0.00E+00	0.00E+00	8.74E-60	1.78E-23	1.11E-10	4.80E-04	5.84E+00	3.48E+03	-6.22E+00
32	36	68	0.00E+00	0.00E+00	1.37E-70	2.16E-28	1.28E-13	5.64E-06	2.86E-01	4.48E+02	-7.39E+00
32	37	69	0.00E+00	0.00E+00	1.29E-70	4.08E-29	1.34E-14	4.56E-07	2.06E-02	3.10E+01	-7.30E+00
32	38	70	0.00E+00	0.00E+00	6.18E-82	5.68E-34	3.04E-17	1.32E-08	2.62E-03	1.01E+01	-8.53E+00
32	39	71	0.00E+00	0.00E+00	1.90E-81	6.50E-35	1.12E-18	2.86E-10	4.42E-05	1.53E-01	-8.29E+00
32	40	72	0.00E+00	0.00E+00	6.84E-94	2.84E-40	7.12E-22	1.85E-12	1.10E-06	9.42E-03	-9.73E+00
32	41	73	0.00E+00	0.00E+00	0.00E+00	2.90E-44	3.86E-25	2.68E-15	2.94E-09	3.72E-05	-9.99E+00
32	42	74	0.00E+00	0.00E+00	0.00E+00	4.68E-48	3.10E-27	1.53E-16	5.46E-10	1.52E-05	-1.10E+01
33	32	65	9.06E-22	6.10E-05	1.05E+01	1.27E+06	1.78E+08	3.12E+09	2.14E+10	8.62E+10	-1.28E-01
33	33	66	0.00E+00	1.22E-51	2.78E-27	1.82E-07	2.86E+00	1.74E+04	3.84E+06	1.53E+08	-2.95E+00
33	34	67	0.00E+00	3.08E-41	3.46E-21	1.12E-04	1.30E+02	2.14E+05	2.20E+07	5.42E+08	-2.31E+00
33	35	68	0.00E+00	6.14E-62	7.56E-34	3.64E-11	6.36E-03	1.46E+02	7.78E+04	5.86E+06	-3.51E+00
33	36	69	0.00E+00	6.02E-60	1.20E-32	1.38E-10	1.51E-02	2.84E+02	1.40E+05	1.03E+07	-3.39E+00
33	37	70	0.00E+00	1.53E-78	1.47E-43	1.17E-15	1.29E-05	2.58E+00	5.28E+03	1.00E+06	-4.54E+00
33	38	71	0.00E+00	1.54E-80	5.60E-45	1.16E-16	1.75E-06	4.22E-01	9.92E+02	2.10E+05	-4.62E+00
33	39	72	0.00E+00	9.02E-97	1.53E-54	3.52E-21	2.68E-09	4.46E-03	3.22E+01	1.38E+04	-5.61E+00
33	40	73	0.00E+00	1.01E-97	3.56E-55	1.33E-21	1.21E-09	2.34E-03	1.96E+01	9.80E+03	-5.66E+00
33	41	74	0.00E+00	0.00E+00	3.14E-66	6.96E-27	2.94E-13	2.84E-06	5.54E-02	4.72E+01	-6.85E+00
33	42	75	0.00E+00	0.00E+00	1.50E-66	8.70E-27	8.26E-13	1.49E-05	4.38E-01	4.70E+02	-6.90E+00
34	32	66	0.00E+00	2.86E-45	2.78E-23	2.34E-05	7.10E+01	1.72E+05	2.12E+07	5.60E+08	-2.59E+00
34	33	67	0.00E+00	1.58E-58	2.64E-31	2.56E-09	2.04E-01	2.60E+03	8.56E+05	4.32E+07	-3.36E+00
34	34	68	0.00E+00	6.86E-84	1.89E-46	9.14E-17	2.96E-06	8.26E-01	1.82E+03	3.36E+05	-4.88E+00
34	35	69	0.00E+00	2.32E-81	3.78E-45	2.42E-16	4.70E-06	1.12E+00	2.36E+03	4.34E+05	-4.71E+00
34	36	70	0.00E+00	0.00E+00	4.42E-59	4.82E-23	2.78E-10	1.27E-03	1.71E+01	1.14E+04	-6.13E+00
34	37	71	0.00E+00	0.00E+00	8.92E-63	2.58E-25	4.78E-12	4.10E-05	8.22E-01	7.16E+02	-6.41E+00
34	38	72	0.00E+00	0.00E+00	5.32E-71	4.14E-29	2.34E-14	1.15E-06	6.72E-02	1.23E+02	-7.29E+00
34	39	73	0.00E+00	0.00E+00	6.72E-72	3.70E-30	1.54E-15	5.66E-08	2.54E-03	3.70E+00	-7.28E+00
34	40	74	0.00E+00	0.00E+00	9.70E-84	1.65E-35	1.27E-18	7.32E-10	1.88E-04	9.08E-01	-8.55E+00
34	41	75	0.00E+00	0.00E+00	2.30E-85	3.30E-37	1.66E-20	6.94E-12	1.44E-06	6.14E-03	-8.60E+00
34	42	76	0.00E+00	0.00E+00	5.94E-92	5.18E-39	8.14E-21	1.49E-11	6.40E-06	4.08E-02	-9.51E+00
34	43	77	0.00E+00	0.00E+00	2.84E-95	8.00E-42	2.72E-23	9.42E-14	6.72E-08	6.36E-04	-9.60E+00
34	44	78	0.00E+00	0.00E+00	0.00E+00	7.46E-45	4.22E-25	5.98E-15	1.01E-08	1.72E-04	-1.04E+01
35	35	70	4.90E-94	4.04E-29	4.24E-14	2.80E-01	2.04E+04	8.40E+06	3.70E+08	5.04E+09	-1.54E+00
35	36	71	0.00E+00	4.34E-41	1.46E-21	2.60E-05	3.16E+01	6.02E+04	7.22E+06	2.02E+08	-2.21E+00
35	37	72	0.00E+00	1.30E-53	9.44E-29	2.14E-08	7.10E-01	7.50E+03	2.56E+06	1.44E+08	-3.00E+00
35	38	73	0.00E+00	4.66E-53	1.61E-28	1.96E-08	5.24E-01	5.04E+03	1.67E+06	9.32E+07	-2.95E+00
35	39	74	0.00E+00	6.80E-75	9.10E-41	1.69E-13	1.01E-03	1.35E+02	2.04E+05	3.06E+07	-4.38E+00
35	40	75	0.00E+00	1.05E-74	1.10E-41	3.40E-15	1.36E-05	1.70E+00	2.70E+03	4.42E+05	-4.22E+00
35	41	76	0.00E+00	8.50E-94	9.28E-53	3.16E-20	1.42E-08	1.97E-02	1.35E+02	5.86E+04	-5.41E+00
35	42	77	0.00E+00	6.36E-91	8.22E-51	4.38E-19	8.38E-08	6.76E-02	3.20E+02	1.06E+05	-5.27E+00
35	43	78	0.00E+00	0.00E+00	3.70E-59	5.22E-23	2.66E-10	9.48E-04	9.60E+00	4.82E+03	-6.14E+00
35	44	79	0.00E+00	0.00E+00	1.80E-61	1.97E-24	2.48E-11	1.72E-04	3.02E+00	2.40E+03	-6.33E+00
36	35	71	0.00E+00	4.06E-55	1.42E-29	9.12E-09	3.48E-01	3.38E+03	1.00E+06	4.82E+07	-3.10E+00
36	36	72	0.00E+00	4.36E-79	6.10E-44	6.90E-16	7.76E-06	1.42E+00	2.58E+03	4.34E+05	-4.53E+00
36	37	73	0.00E+00	4.58E-87	8.94E-49	2.50E-18	2.02E-07	1.09E-01	3.98E+02	1.09E+05	-4.99E+00
36	38	74	0.00E+00	0.00E+00	4.60E-58	4.64E-23	1.32E-10	4.50E-04	5.26E+00	3.26E+03	-5.90E+00
36	39	75	0.00E+00	0.00E+00	6.68E-61	2.10E-24	1.63E-11	8.26E-05	1.13E+00	7.30E+02	-6.22E+00
36	40	76	0.00E+00	0.00E+00	3.46E-70	6.26E-29	2.36E-14	9.88E-07	5.42E-02	9.66E+01	-7.13E+00
36	41	77	0.00E+00	0.00E+00	1.50E-71	4.88E-30	2.12E-15	8.72E-08	4.38E-03	6.92E+00	-7.17E+00
36	42	78	0.00E+00	0.00E+00	1.65E-79	7.36E-33	1.51E-16	4.22E-08	6.74E-03	2.34E+01	-8.20E+00
36	43	79	0.00E+00	0.00E+00	1.68E-81	1.19E-34	2.70E-18	7.30E-10	1.08E-04	3.42E-01	-8.28E+00
36	44	80	0.00E+00	0.00E+00	1.20E-88	1.94E-37	1.18E-19	1.60E-10	6.08E-05	3.56E-01	-9.11E+00
36	45	81	0.00E+00	0.00E+00	9.12E-91	1.06E-39	5.72E-22	7.64E-13	3.02E-07	1.87E-03	-9.10E+00

36	46	82	0.00E+00	0.00E+00	9.10E-97	7.38E-42	5.88E-23	2.80E-13	2.34E-07	2.46E-03	-9.90E+00
36	47	83	0.00E+00	0.00E+00	1.05E-97	2.34E-43	1.64E-24	8.30E-15	7.58E-09	8.50E-05	-9.77E+00
36	48	84	0.00E+00	0.00E+00	0.00E+00	3.08E-47	8.96E-27	3.00E-16	8.62E-10	2.08E-05	-1.07E+01
37	37	74	0.00E+00	1.32E-39	2.08E-20	2.80E-04	3.34E+02	6.14E+05	6.94E+07	1.79E+09	-2.13E+00
37	38	75	0.00E+00	4.80E-44	2.02E-23	2.74E-06	6.66E+00	1.82E+04	2.76E+06	9.16E+07	-2.34E+00
37	39	76	0.00E+00	3.24E-63	1.31E-34	2.40E-11	7.68E-03	2.46E+02	1.59E+05	1.33E+07	-3.53E+00
37	40	77	0.00E+00	9.28E-57	9.00E-31	1.47E-09	8.22E-02	1.04E+03	3.84E+05	2.30E+07	-3.14E+00
37	41	78	0.00E+00	1.28E-70	1.74E-38	1.31E-12	2.90E-03	2.40E+02	2.74E+05	3.42E+07	-4.05E+00
37	42	79	0.00E+00	6.50E-70	1.14E-38	1.67E-13	2.06E-04	1.30E+01	1.32E+04	1.59E+06	-3.92E+00
37	43	80	0.00E+00	1.30E-87	4.84E-49	2.26E-18	2.18E-07	1.29E-01	5.14E+02	1.53E+05	-5.02E+00
37	44	81	0.00E+00	3.04E-85	8.70E-48	5.46E-18	2.52E-07	1.03E-01	3.38E+02	9.12E+04	-4.85E+00
37	45	82	0.00E+00	0.00E+00	3.48E-56	1.18E-21	2.34E-09	6.44E-03	6.28E+01	3.24E+04	-5.78E+00
37	46	83	0.00E+00	0.00E+00	2.26E-57	7.10E-23	1.33E-10	3.68E-04	3.84E+00	2.26E+03	-5.77E+00
37	47	84	0.00E+00	0.00E+00	3.38E-69	2.64E-28	4.94E-14	1.08E-06	3.34E-02	3.64E+01	-7.06E+00
37	48	85	0.00E+00	0.00E+00	6.22E-70	3.62E-29	8.64E-15	2.68E-07	1.15E-02	1.64E+01	-7.02E+00
38	37	75	0.00E+00	2.16E-57	4.26E-31	1.32E-09	9.22E-02	1.23E+03	4.36E+05	2.34E+07	-3.17E+00
38	38	76	0.00E+00	1.06E-78	1.08E-43	1.00E-15	9.88E-06	1.60E+00	2.60E+03	3.96E+05	-4.47E+00
38	39	77	0.00E+00	8.18E-84	1.17E-46	6.00E-17	2.88E-06	1.15E+00	3.38E+03	7.72E+05	-4.79E+00
38	40	78	0.00E+00	5.64E-98	4.12E-55	3.78E-21	4.54E-09	9.04E-03	7.00E+01	3.08E+04	-5.63E+00
38	41	79	0.00E+00	0.00E+00	8.88E-58	5.24E-23	1.24E-10	3.74E-04	3.90E+00	2.14E+03	-5.83E+00
38	42	80	0.00E+00	0.00E+00	2.26E-66	1.31E-26	1.35E-12	2.52E-05	7.64E-01	8.44E+02	-6.80E+00
38	43	81	0.00E+00	0.00E+00	3.60E-65	2.78E-26	1.50E-12	2.12E-05	5.62E-01	5.82E+02	-6.64E+00
38	44	82	0.00E+00	0.00E+00	4.54E-77	4.46E-32	2.74E-16	4.30E-08	4.94E-03	1.39E+01	-7.84E+00
38	45	83	0.00E+00	0.00E+00	9.36E-79	1.66E-33	1.11E-17	1.68E-09	1.84E-04	4.96E-01	-7.90E+00
38	46	84	0.00E+00	0.00E+00	2.58E-87	4.02E-37	1.41E-19	1.67E-10	6.28E-05	3.78E-01	-8.86E+00
38	47	85	0.00E+00	0.00E+00	9.04E-87	5.16E-38	6.98E-21	5.92E-12	1.98E-06	1.15E-02	-8.64E+00
38	48	86	0.00E+00	0.00E+00	3.94E-95	4.64E-41	2.34E-22	8.44E-13	5.54E-07	4.56E-03	-9.64E+00
38	49	87	0.00E+00	0.00E+00	3.50E-95	2.22E-42	5.54E-24	1.80E-14	1.30E-08	1.25E-04	-9.42E+00
38	50	88	0.00E+00	0.00E+00	0.00E+00	9.42E-47	1.12E-26	2.02E-16	3.72E-10	6.60E-06	-1.06E+01
39	39	78	3.26E-72	1.66E-22	2.48E-10	1.68E+01	3.10E+05	6.76E+07	2.08E+09	2.24E+10	-1.06E+00
39	40	79	0.00E+00	2.16E-46	1.09E-24	1.15E-06	5.60E+00	2.06E+04	3.64E+06	1.31E+08	-2.47E+00
39	41	80	0.00E+00	3.02E-91	2.00E-51	1.03E-19	2.18E-08	1.77E-02	7.90E+01	2.40E+04	-5.17E+00
39	42	81	0.00E+00	7.52E-55	1.26E-29	6.48E-09	2.86E-01	3.52E+03	1.36E+06	8.64E+07	-3.00E+00
39	43	82	0.00E+00	4.88E-71	2.00E-39	6.46E-14	1.25E-04	1.06E+01	1.29E+04	1.68E+06	-3.96E+00
39	44	83	0.00E+00	7.90E-65	1.66E-35	9.70E-12	4.32E-03	1.67E+02	1.29E+05	1.30E+07	-3.61E+00
39	45	84	0.00E+00	8.84E-84	5.82E-47	1.69E-17	7.52E-07	3.26E-01	1.10E+03	2.96E+05	-4.73E+00
39	46	85	0.00E+00	1.03E-79	1.81E-44	3.30E-16	5.18E-06	1.28E+00	3.08E+03	6.76E+05	-4.49E+00
39	47	86	0.00E+00	1.07E-95	6.58E-54	1.02E-20	7.70E-09	1.35E-02	1.04E+02	4.84E+04	-5.47E+00
39	48	87	0.00E+00	0.00E+00	2.10E-57	1.10E-22	2.50E-10	7.60E-04	8.44E+00	5.18E+03	-5.78E+00
39	49	88	0.00E+00	0.00E+00	2.66E-66	5.66E-27	4.54E-13	8.28E-06	2.62E-01	3.08E+02	-6.71E+00
39	50	89	0.00E+00	0.00E+00	2.26E-70	3.30E-29	1.08E-14	3.98E-07	1.98E-02	3.28E+01	-7.07E+00
40	39	79	0.00E+00	2.02E-44	2.10E-23	7.96E-06	2.70E+01	7.60E+04	1.03E+07	2.88E+08	-2.36E+00
40	40	80	0.00E+00	8.22E-66	2.78E-36	2.02E-12	8.16E-04	2.62E+01	1.61E+04	1.27E+06	-3.63E+00
40	41	81	0.00E+00	1.68E-51	1.36E-27	6.76E-08	1.17E+00	8.00E+03	2.02E+06	9.26E+07	-2.79E+00
40	42	82	0.00E+00	4.54E-96	3.46E-54	6.26E-21	4.42E-09	6.60E-03	4.18E+01	1.57E+04	-5.47E+00
40	43	83	0.00E+00	1.18E-97	3.62E-55	1.97E-21	2.30E-09	4.90E-03	4.16E+01	1.99E+04	-5.56E+00
40	44	84	0.00E+00	0.00E+00	9.06E-60	9.04E-24	6.02E-11	3.08E-04	4.42E+00	3.02E+03	-6.00E+00
40	45	85	0.00E+00	0.00E+00	1.14E-61	1.25E-24	1.97E-11	1.63E-04	3.22E+00	2.78E+03	-6.21E+00
40	46	86	0.00E+00	0.00E+00	5.14E-72	9.62E-30	8.08E-15	4.82E-07	3.12E-02	6.04E+01	-7.25E+00
40	47	87	0.00E+00	0.00E+00	2.08E-73	1.33E-30	1.71E-15	1.27E-07	9.38E-03	1.95E+01	-7.36E+00
40	48	88	0.00E+00	0.00E+00	1.81E-78	5.60E-33	5.54E-17	1.15E-08	1.58E-03	5.02E+00	-7.90E+00
40	49	89	0.00E+00	0.00E+00	2.00E-78	3.82E-33	3.10E-17	5.66E-09	7.02E-04	2.04E+00	-7.86E+00
40	50	90	0.00E+00	0.00E+00	4.86E-83	2.60E-35	1.40E-18	6.72E-10	1.53E-04	6.84E-01	-8.36E+00
41	41	82	3.60E-90	1.92E-28	9.26E-14	4.50E-01	3.10E+04	1.26E+07	5.58E+08	7.76E+09	-1.41E+00
41	42	83	0.00E+00	4.68E-41	5.86E-22	7.26E-06	8.82E+00	1.74E+04	2.16E+06	6.22E+07	-2.06E+00
41	43	84	0.00E+00	2.08E-37	3.56E-19	1.44E-03	1.55E+03	3.14E+06	3.98E+08	1.14E+10	-1.93E+00
41	44	85	0.00E+00	1.11E-41	3.06E-22	7.78E-06	1.22E+01	2.76E+04	3.76E+06	1.18E+08	-2.11E+00
41	45	86	0.00E+00	1.17E-71	8.96E-40	6.44E-14	1.93E-04	2.18E+01	3.24E+04	4.96E+06	-3.96E+00
41	46	87	0.00E+00	1.34E-67	9.48E-38	1.65E-13	1.22E-04	6.52E+00	6.28E+03	7.46E+05	-3.66E+00
41	47	88	0.00E+00	2.96E-72	4.72E-40	5.44E-14	1.78E-04	2.04E+01	3.08E+04	4.84E+06	-4.01E+00

41	48	89	0.00E+00	3.22E-77	1.72E-43	2.32E-16	1.53E-06	2.44E-01	4.56E+02	8.40E+04	-4.24E+00
41	49	90	0.00E+00	2.00E-90	4.14E-51	9.10E-20	1.73E-08	1.59E-02	8.58E+01	3.22E+04	-5.08E+00
41	50	91	0.00E+00	2.72E-92	1.97E-52	9.92E-21	2.16E-09	2.00E-03	1.09E+01	4.20E+03	-5.16E+00
41	51	92	0.00E+00	0.00E+00	8.94E-59	1.28E-23	3.68E-11	1.08E-04	9.96E-01	4.70E+02	-5.85E+00
41	52	93	0.00E+00	0.00E+00	2.18E-61	3.02E-25	2.30E-12	1.24E-05	1.74E-01	1.09E+02	-6.04E+00
42	41	83	6.16E-84	8.68E-27	5.34E-13	6.96E-01	3.22E+04	1.02E+07	3.62E+08	4.10E+09	-1.26E+00
42	42	84	0.00E+00	1.46E-42	3.56E-22	3.60E-05	7.64E+01	1.84E+05	2.44E+07	7.08E+08	-2.24E+00
42	43	85	0.00E+00	5.60E-69	2.98E-38	2.64E-13	3.16E-04	1.89E+01	1.70E+04	1.70E+06	-3.79E+00
42	44	86	0.00E+00	4.76E-83	2.34E-46	6.70E-17	2.88E-06	1.12E+00	3.30E+03	7.64E+05	-4.66E+00
42	45	87	0.00E+00	8.78E-92	7.94E-52	6.64E-20	2.12E-08	2.48E-02	1.49E+02	5.74E+04	-5.16E+00
42	46	88	0.00E+00	0.00E+00	9.58E-58	1.53E-22	5.68E-10	2.20E-03	2.72E+01	1.71E+04	-5.81E+00
42	47	89	0.00E+00	0.00E+00	1.46E-62	2.30E-25	4.08E-12	3.52E-05	7.06E-01	6.14E+02	-6.22E+00
42	48	90	0.00E+00	0.00E+00	2.40E-68	8.42E-28	1.86E-13	5.62E-06	2.40E-01	3.50E+02	-6.88E+00
42	49	91	0.00E+00	0.00E+00	1.83E-68	3.82E-28	7.20E-14	2.04E-06	8.62E-02	1.26E+02	-6.84E+00
42	50	92	0.00E+00	0.00E+00	6.42E-74	1.63E-30	3.16E-15	2.84E-07	2.42E-02	5.70E+01	-7.46E+00
42	51	93	0.00E+00	0.00E+00	2.02E-76	4.12E-32	1.55E-16	1.78E-08	1.58E-03	3.50E+00	-7.64E+00
42	52	94	0.00E+00	0.00E+00	2.22E-84	5.64E-36	4.30E-19	2.20E-10	4.86E-05	2.00E-01	-8.49E+00
42	53	95	0.00E+00	0.00E+00	7.70E-87	5.24E-38	4.40E-21	2.26E-12	5.18E-07	2.34E-03	-8.63E+00
42	54	96	0.00E+00	0.00E+00	2.48E-93	4.38E-41	5.78E-23	1.20E-13	6.22E-08	4.66E-04	-9.30E+00
43	43	86	1.09E-49	1.44E-15	2.42E-06	1.44E+03	6.28E+06	6.80E+08	1.37E+10	1.08E+11	-5.64E-01
43	44	87	8.06E-78	1.22E-25	8.22E-13	2.04E-01	6.48E+03	2.02E+06	8.20E+07	1.13E+09	-1.10E+00
43	45	88	0.00E+00	5.90E-38	1.09E-19	4.86E-04	5.20E+02	1.02E+06	1.26E+08	3.52E+09	-1.91E+00
43	46	89	0.00E+00	6.62E-42	1.65E-22	4.08E-06	6.82E+00	1.65E+04	2.40E+06	7.84E+07	-2.08E+00
43	47	90	0.00E+00	1.28E-65	3.22E-36	3.26E-12	2.22E-03	1.13E+02	1.01E+05	1.09E+07	-3.57E+00
43	48	91	0.00E+00	5.80E-59	1.06E-32	3.86E-11	3.46E-03	6.14E+01	3.00E+04	2.24E+06	-3.10E+00
43	49	92	0.00E+00	1.64E-73	4.08E-41	6.56E-15	2.38E-05	2.92E+00	4.66E+03	7.66E+05	-4.02E+00
43	50	93	0.00E+00	2.30E-75	1.60E-42	4.84E-16	1.87E-06	2.24E-01	3.46E+02	5.72E+04	-4.09E+00
43	51	94	0.00E+00	9.36E-84	3.04E-47	5.52E-18	2.24E-07	9.62E-02	3.26E+02	8.76E+04	-4.64E+00
43	52	95	0.00E+00	1.23E-88	1.97E-50	6.34E-20	6.62E-09	4.66E-03	2.20E+01	7.52E+03	-4.89E+00
43	53	96	0.00E+00	1.47E-96	5.70E-55	7.04E-22	4.70E-10	6.96E-04	4.38E+00	1.62E+03	-5.40E+00
43	54	97	0.00E+00	0.00E+00	7.64E-59	3.50E-24	9.86E-12	3.66E-05	4.52E-01	2.78E+02	-5.72E+00
44	43	87	0.00E+00	3.28E-46	8.14E-25	8.10E-07	4.38E+00	1.64E+04	2.68E+06	8.42E+07	-2.36E+00
44	44	88	0.00E+00	4.58E-65	9.28E-36	6.68E-12	3.12E-03	1.12E+02	7.32E+04	6.00E+06	-3.54E+00
44	45	89	0.00E+00	8.12E-69	3.38E-38	3.16E-13	4.18E-04	2.68E+01	2.52E+04	2.64E+06	-3.74E+00
44	46	90	0.00E+00	6.80E-82	8.78E-46	9.68E-17	3.00E-06	9.88E-01	2.64E+03	5.80E+05	-4.55E+00
44	47	91	0.00E+00	6.94E-83	1.26E-46	1.92E-17	7.42E-07	2.96E-01	9.26E+02	2.30E+05	-4.58E+00
44	48	92	0.00E+00	1.53E-98	9.74E-56	1.11E-21	1.64E-09	3.84E-03	3.44E+01	1.75E+04	-5.55E+00
44	49	93	0.00E+00	0.00E+00	6.94E-57	1.70E-22	3.44E-10	9.86E-04	1.03E+01	5.82E+03	-5.62E+00
44	50	94	0.00E+00	0.00E+00	1.20E-62	3.94E-25	8.18E-12	7.34E-05	1.51E+00	1.37E+03	-6.25E+00
44	51	95	0.00E+00	0.00E+00	1.46E-66	2.06E-27	1.63E-13	3.02E-06	9.88E-02	1.21E+02	-6.58E+00
44	52	96	0.00E+00	0.00E+00	1.28E-73	1.13E-30	1.68E-15	1.37E-07	1.08E-02	2.36E+01	-7.34E+00
44	53	97	0.00E+00	0.00E+00	1.20E-76	1.30E-32	3.60E-17	3.62E-09	3.10E-04	6.96E-01	-7.58E+00
44	54	98	0.00E+00	0.00E+00	2.48E-83	1.29E-35	5.86E-19	2.20E-10	3.86E-05	1.36E-01	-8.29E+00
45	45	90	9.24E-54	3.32E-17	2.04E-07	3.42E+02	2.08E+06	2.62E+08	5.74E+09	4.86E+10	-6.22E-01
45	46	91	8.98E-81	6.50E-27	1.12E-13	6.24E-02	2.54E+03	8.78E+05	3.76E+07	5.32E+08	-1.14E+00
45	47	92	0.00E+00	1.00E-37	1.13E-19	3.56E-04	3.30E+02	5.92E+05	6.92E+07	1.88E+09	-1.86E+00
45	48	93	0.00E+00	2.80E-41	2.92E-22	3.94E-06	5.10E+00	1.04E+04	1.35E+06	4.10E+07	-2.00E+00
45	49	94	0.00E+00	7.42E-53	1.03E-28	1.21E-08	3.74E-01	3.94E+03	1.37E+06	7.98E+07	-2.77E+00
45	50	95	0.00E+00	4.50E-59	6.44E-33	1.87E-11	1.31E-03	1.83E+01	7.24E+03	4.56E+05	-3.06E+00
45	51	96	0.00E+00	5.88E-65	3.82E-36	1.32E-12	6.20E-04	2.72E+01	2.32E+04	2.56E+06	-3.46E+00
45	52	97	0.00E+00	2.10E-71	2.56E-40	4.06E-15	6.52E-06	5.36E-01	6.94E+02	1.03E+05	-3.81E+00
45	53	98	0.00E+00	3.46E-79	1.68E-44	1.43E-16	2.28E-06	6.42E-01	1.77E+03	4.26E+05	-4.35E+00
45	54	99	0.00E+00	1.76E-83	2.94E-47	3.54E-18	1.31E-07	5.74E-02	2.12E+02	6.52E+04	-4.58E+00
46	45	91	0.00E+00	5.50E-41	1.07E-21	3.00E-05	4.90E+01	9.94E+04	1.12E+07	2.74E+08	-2.02E+00
46	46	92	0.00E+00	3.30E-60	4.14E-33	6.50E-11	8.14E-03	1.48E+02	6.50E+04	4.14E+06	-3.19E+00
46	47	93	0.00E+00	1.54E-63	4.38E-35	9.80E-12	3.82E-03	1.36E+02	9.16E+04	7.78E+06	-3.39E+00
46	48	94	0.00E+00	8.06E-80	7.78E-45	1.31E-16	2.06E-06	4.64E-01	9.80E+02	1.84E+05	-4.37E+00
46	49	95	0.00E+00	2.72E-80	4.06E-45	1.02E-16	2.08E-06	5.70E-01	1.39E+03	2.90E+05	-4.40E+00
46	50	96	0.00E+00	4.50E-92	5.50E-52	5.78E-20	1.57E-08	1.49E-02	7.66E+01	2.72E+04	-5.12E+00
46	51	97	0.00E+00	2.96E-98	6.40E-56	3.52E-22	4.38E-10	9.88E-04	8.86E+00	4.54E+03	-5.46E+00

46	52	98	0.00E+00	0.00E+00	6.44E-61	1.53E-24	1.45E-11	9.38E-05	1.65E+00	1.36E+03	-6.00E+00
46	53	99	0.00E+00	0.00E+00	2.44E-64	1.29E-26	3.66E-13	4.16E-06	1.01E-01	1.02E+02	-6.27E+00
46	54	100	0.00E+00	0.00E+00	2.30E-71	5.42E-30	2.70E-15	1.37E-07	8.56E-03	1.65E+01	-7.00E+00
47	47	94	4.56E-55	5.56E-18	4.82E-08	1.00E+02	6.24E+05	7.88E+07	1.74E+09	1.50E+10	-6.24E-01
47	48	95	1.51E-72	3.30E-24	4.96E-12	4.62E-01	1.00E+04	2.46E+06	8.66E+07	1.09E+09	-9.60E-01
47	49	96	0.00E+00	8.60E-39	1.55E-20	7.20E-05	7.26E+01	1.33E+05	1.56E+07	4.32E+08	-1.87E+00
47	50	97	0.00E+00	3.84E-37	9.30E-20	7.82E-05	3.78E+01	4.54E+04	4.22E+06	1.03E+08	-1.74E+00
47	51	98	0.00E+00	1.74E-47	7.60E-26	1.18E-07	8.36E-01	4.18E+03	9.36E+05	4.08E+07	-2.37E+00
47	52	99	0.00E+00	1.78E-53	1.54E-29	1.10E-09	2.72E-02	2.60E+02	8.82E+04	5.28E+06	-2.71E+00
47	53	100	0.00E+00	2.16E-62	1.31E-34	8.26E-12	2.36E-03	8.40E+01	6.56E+04	6.86E+06	-3.29E+00
47	54	101	0.00E+00	3.08E-63	1.61E-35	8.58E-13	2.26E-04	7.86E+00	6.26E+03	6.90E+05	-3.29E+00
48	47	95	0.00E+00	3.42E-40	3.10E-21	5.06E-05	7.02E+01	1.35E+05	1.50E+07	3.74E+08	-1.95E+00
48	48	96	0.00E+00	7.48E-61	1.69E-33	3.62E-11	4.66E-03	8.70E+01	4.00E+04	2.68E+06	-3.20E+00
48	49	97	0.00E+00	5.56E-62	2.52E-34	1.58E-11	4.08E-03	1.17E+02	7.00E+04	5.54E+06	-3.26E+00
48	50	98	0.00E+00	8.28E-80	5.98E-45	7.94E-17	1.04E-06	2.08E-01	4.08E+02	7.42E+04	-4.32E+00
48	51	99	0.00E+00	6.40E-76	1.63E-42	2.06E-15	1.49E-05	2.40E+00	4.28E+03	7.32E+05	-4.11E+00
48	52	100	0.00E+00	2.66E-88	5.06E-50	2.44E-19	2.60E-08	1.62E-02	6.62E+01	2.00E+04	-4.83E+00
48	53	101	0.00E+00	3.06E-89	1.08E-50	8.38E-20	1.21E-08	9.70E-03	4.78E+01	1.67E+04	-4.88E+00
48	54	102	0.00E+00	1.07E-98	3.66E-56	2.80E-22	4.14E-10	1.10E-03	1.15E+01	6.86E+03	-5.48E+00
48	55	103	0.00E+00	0.00E+00	1.47E-61	2.12E-25	2.04E-12	1.48E-05	2.94E-01	2.68E+02	-5.94E+00
48	56	104	0.00E+00	0.00E+00	2.04E-66	1.53E-27	1.19E-13	2.54E-06	9.92E-02	1.46E+02	-6.47E+00
49	49	98	6.76E-69	3.58E-23	2.68E-11	1.57E+00	3.00E+04	6.92E+06	2.28E+08	2.68E+09	-8.92E-01
49	50	99	1.18E-89	1.54E-30	5.78E-16	3.88E-03	3.58E+02	1.83E+05	9.98E+06	1.70E+08	-1.29E+00
49	51	100	0.00E+00	9.40E-43	5.10E-23	3.30E-06	8.02E+00	2.26E+04	3.46E+06	1.14E+08	-2.06E+00
49	52	101	0.00E+00	2.98E-41	2.48E-22	3.12E-06	3.90E+00	8.02E+03	1.07E+06	3.46E+07	-1.94E+00
49	53	102	0.00E+00	1.93E-42	9.40E-23	5.08E-06	1.21E+01	3.60E+04	5.94E+06	2.14E+08	-2.06E+00
49	54	103	0.00E+00	2.48E-50	6.94E-28	3.74E-09	4.06E-02	2.74E+02	7.88E+04	4.32E+06	-2.47E+00
49	55	104	0.00E+00	1.62E-53	1.42E-29	1.32E-09	4.56E-02	5.80E+02	2.48E+05	1.75E+07	-2.71E+00
49	56	105	0.00E+00	1.18E-55	4.00E-31	9.96E-11	4.78E-03	7.44E+01	3.74E+04	3.04E+06	-2.79E+00

TABLE IX. The astrophysical (γ,α) photodisintegration rates λ in the reaction rate set FRDM1, calculated from (α,γ) reaction rates using detailed balance. Not included are γ-induced α-decay rates of α-unbound nuclei. Rates smaller than 10^{-99} have been set to zero.

Z	N	A	λ (s^{-1})								Q (MeV)
			$T_9 = 0.1$	$T_9 = 0.3$	$T_9 = 0.5$	$T_9 = 1.0$	$T_9 = 1.5$	$T_9 = 2.0$	$T_9 = 2.5$	$T_9 = 3.0$	
32	32	64	0.00E+00	1.72E−71	2.67E−43	2.25E−18	1.33E−08	2.60E−03	5.39E+00	1.01E+03	−2.67E+00
32	33	65	0.00E+00	5.36E−69	5.24E−42	6.09E−18	2.23E−08	3.91E−03	8.14E+00	1.59E+03	−2.49E+00
32	34	66	0.00E+00	1.47E−75	4.95E−46	4.85E−20	1.12E−09	6.46E−04	3.22E+00	1.23E+03	−2.88E+00
32	35	67	0.00E+00	3.78E−75	1.20E−45	1.16E−19	2.55E−09	1.39E−03	6.65E+00	2.47E+03	−2.87E+00
32	36	68	0.00E+00	2.47E−84	2.87E−51	1.19E−22	2.24E−11	4.08E−05	4.45E−01	3.05E+02	−3.40E+00
32	37	69	0.00E+00	7.79E−88	2.36E−53	1.16E−23	4.42E−12	1.03E−05	1.22E−01	8.47E+01	−3.61E+00
32	38	70	0.00E+00	8.37E−96	3.99E−58	4.70E−26	1.28E−13	9.23E−07	2.36E−02	2.92E+01	−4.09E+00
32	39	71	0.00E+00	0.00E+00	3.06E−61	2.07E−27	1.29E−14	9.62E−08	2.03E−03	1.97E+00	−4.45E+00
32	40	72	0.00E+00	0.00E+00	2.84E−67	1.40E−30	1.29E−16	4.93E−09	3.15E−04	6.62E−01	−5.00E+00
32	41	73	0.00E+00	0.00E+00	2.83E−71	7.68E−34	6.22E−20	1.91E−12	1.10E−07	2.35E−04	−5.30E+00
32	42	74	0.00E+00	0.00E+00	3.43E−80	5.42E−37	3.60E−21	7.41E−13	1.06E−07	3.68E−04	−6.29E+00
33	32	65	0.00E+00	1.85E−66	6.66E−42	3.98E−20	2.51E−11	2.41E−06	4.60E−03	1.05E+00	−2.19E+00
33	33	66	0.00E+00	1.41E−69	1.36E−42	3.44E−19	3.79E−10	3.19E−05	4.55E−02	7.78E+00	−2.50E+00
33	34	67	0.00E+00	2.96E−68	6.28E−42	1.48E−18	2.94E−09	3.94E−04	7.46E−01	1.45E+02	−2.38E+00
33	35	68	0.00E+00	1.44E−70	1.47E−43	1.34E−19	3.86E−10	5.83E−05	1.17E−01	2.40E+01	−2.47E+00
33	36	69	0.00E+00	2.59E−76	8.62E−47	7.43E−21	1.32E−10	5.94E−05	2.47E−01	8.40E+01	−2.85E+00
33	37	70	0.00E+00	2.70E−80	2.25E−49	2.34E−22	8.38E−12	4.83E−06	2.24E−02	8.18E+00	−3.04E+00
33	38	71	0.00E+00	3.38E−86	1.18E−52	1.32E−23	3.05E−12	5.38E−06	5.24E−02	3.18E+01	−3.44E+00
33	39	72	0.00E+00	2.20E−88	5.89E−54	3.10E−24	1.09E−12	2.13E−06	2.08E−02	1.23E+01	−3.57E+00
33	40	73	0.00E+00	2.43E−96	1.19E−58	1.81E−26	5.20E−14	3.44E−07	7.64E−03	8.09E+00	−4.06E+00
33	41	74	0.00E+00	0.00E+00	4.66E−62	2.25E−28	9.41E−16	4.92E−09	8.07E−05	6.86E−02	−4.38E+00
33	42	75	0.00E+00	0.00E+00	2.55E−71	1.01E−32	3.26E−18	1.91E−10	1.40E−05	3.05E−02	−5.32E+00
34	29	63	0.00E+00	7.48E−80	8.22E−51	2.16E−25	3.21E−15	1.89E−09	1.23E−05	6.79E−03	−2.77E+00
34	30	64	0.00E+00	3.27E−58	3.30E−38	6.64E−19	6.61E−11	2.70E−06	3.42E−03	6.41E−01	−1.70E+00
34	31	65	0.00E+00	3.35E−59	8.17E−38	2.24E−18	2.00E−10	7.92E−06	9.99E−03	1.89E+00	−1.77E+00
34	32	66	0.00E+00	6.26E−67	1.98E−41	1.80E−19	6.82E−11	4.91E−06	8.48E−03	1.93E+00	−2.24E+00
34	33	67	0.00E+00	1.96E−72	1.09E−44	2.76E−20	3.90E−11	3.87E−06	7.21E−03	1.70E+00	−2.61E+00
34	34	68	0.00E+00	1.60E−76	4.83E−48	1.96E−21	3.99E−11	1.29E−05	3.71E−02	9.21E+00	−2.94E+00
34	35	69	0.00E+00	2.13E−67	1.81E−41	5.33E−18	1.75E−08	3.38E−03	8.09E+00	1.81E+03	−2.31E+00
34	36	70	0.00E+00	1.43E−88	7.70E−54	3.22E−24	1.37E−12	3.88E−06	5.40E−02	4.23E+01	−3.51E+00
34	37	71	0.00E+00	2.53E−84	2.84E−52	4.51E−24	7.23E−13	1.32E−06	1.51E−02	1.10E+01	−3.24E+00
34	38	72	0.00E+00	1.46E−86	6.99E−53	4.68E−24	1.10E−12	2.81E−06	4.38E−02	4.16E+01	−3.34E+00
34	39	73	0.00E+00	8.82E−92	3.57E−55	3.14E−25	1.57E−13	5.18E−07	8.13E−03	6.89E+00	−3.55E+00
34	40	74	0.00E+00	6.47E−99	3.53E−60	1.25E−27	5.01E−15	5.35E−08	2.00E−03	3.40E+00	−4.08E+00
34	41	75	0.00E+00	0.00E+00	1.05E−66	4.10E−31	1.00E−17	1.60E−10	5.78E−06	8.59E−03	−4.69E+00
34	42	76	0.00E+00	0.00E+00	2.05E−70	1.18E−32	2.46E−18	1.73E−10	1.71E−05	4.84E−02	−5.09E+00
34	43	77	0.00E+00	0.00E+00	2.37E−76	1.51E−35	8.03E−21	7.87E−13	1.05E−07	4.24E−04	−5.73E+00
34	44	78	0.00E+00	0.00E+00	4.87E−79	1.24E−36	6.81E−21	1.50E−12	2.35E−07	8.64E−04	−6.03E+00
35	35	70	0.00E+00	2.48E−82	1.59E−51	7.66E−25	2.26E−14	1.64E−08	1.12E−04	6.19E−02	−3.03E+00
35	36	71	0.00E+00	4.66E−85	3.75E−52	3.45E−24	1.77E−13	1.49E−07	1.07E−03	6.09E−01	−3.26E+00
35	37	72	0.00E+00	5.67E−76	3.50E−47	1.24E−21	9.98E−12	2.93E−06	1.06E−02	3.68E+00	−2.73E+00
35	38	73	0.00E+00	6.14E−78	1.56E−48	6.21E−22	1.92E−11	1.15E−05	6.04E−02	2.54E+01	−2.90E+00
35	39	74	0.00E+00	5.92E−85	3.68E−53	4.62E−24	7.21E−13	9.41E−07	8.58E−03	5.75E+00	−3.39E+00
35	40	75	0.00E+00	1.15E−91	1.00E−56	3.96E−26	3.15E−14	1.10E−07	1.77E−03	1.61E+00	−3.67E+00
35	41	76	0.00E+00	0.00E+00	1.16E−64	4.16E−30	6.50E−17	8.31E−10	2.68E−05	3.79E−02	−4.48E+00
35	42	77	0.00E+00	0.00E+00	2.52E−67	2.11E−31	1.00E−17	2.98E−10	1.76E−05	3.73E−02	−4.70E+00
35	43	78	0.00E+00	0.00E+00	8.04E−70	1.90E−32	1.06E−18	1.99E−11	9.02E−07	1.92E−03	−5.02E+00
35	44	79	0.00E+00	0.00E+00	4.35E−74	2.14E−34	2.14E−19	2.46E−11	2.83E−06	8.61E−03	−5.46E+00
36	31	67	0.00E+00	1.33E−66	4.53E−43	1.82E−21	1.54E−12	2.13E−07	5.74E−04	1.77E−01	−1.93E+00
36	32	68	0.00E+00	9.50E−70	1.91E−44	7.73E−22	1.24E−12	2.33E−07	7.44E−04	2.51E−01	−2.17E+00
36	33	69	0.00E+00	8.86E−73	3.01E−46	1.41E−22	4.76E−13	1.22E−07	4.60E−04	1.74E−01	−2.41E+00
36	34	70	0.00E+00	3.14E−78	1.14E−48	1.85E−23	1.77E−13	7.86E−08	4.12E−04	1.90E−01	−2.75E+00
36	35	71	0.00E+00	5.74E−79	5.73E−48	2.78E−22	1.78E−12	5.18E−07	2.01E−03	7.66E−01	−2.77E+00
36	36	72	0.00E+00	9.62E−78	2.92E−49	2.05E−22	5.99E−12	2.74E−06	1.10E−02	3.69E+00	−2.90E+00
36	37	73	0.00E+00	1.04E−80	7.81E−51	1.69E−23	1.39E−12	1.38E−06	9.21E−03	4.31E+00	−3.01E+00

36	38	74	0.00E+00	8.81E−92	1.35E−55	8.70E−26	4.71E−14	1.66E−07	2.69E−03	2.33E+00	−3.42E+00
36	39	75	0.00E+00	4.63E−85	1.33E−52	3.22E−24	6.88E−13	1.37E−06	1.54E−02	1.04E+01	−3.20E+00
36	40	76	0.00E+00	3.16E−93	2.13E−56	4.08E−26	3.16E−14	1.58E−07	3.70E−03	4.53E+00	−3.51E+00
36	41	77	0.00E+00	0.00E+00	3.89E−64	3.51E−29	5.75E−16	1.22E−08	5.94E−04	1.06E+00	−4.38E+00
36	42	78	0.00E+00	0.00E+00	7.40E−65	2.48E−30	5.51E−17	1.42E−09	8.77E−05	2.05E−01	−4.37E+00
36	43	79	0.00E+00	0.00E+00	2.73E−67	1.85E−31	1.02E−17	2.58E−10	1.30E−05	2.57E−02	−4.70E+00
36	44	80	0.00E+00	0.00E+00	5.84E−71	5.43E−33	1.93E−18	1.63E−10	1.74E−05	5.34E−02	−5.06E+00
36	45	81	0.00E+00	0.00E+00	3.81E−76	3.66E−36	3.68E−21	4.15E−13	4.87E−08	1.56E−04	−5.52E+00
36	46	82	0.00E+00	0.00E+00	2.86E−80	1.33E−37	1.34E−21	4.89E−13	1.12E−07	5.43E−04	−5.99E+00
36	47	83	0.00E+00	0.00E+00	1.30E−85	2.45E−41	5.63E−25	3.40E−16	1.25E−10	9.58E−07	−6.49E+00
36	48	84	0.00E+00	0.00E+00	1.47E−91	1.27E−43	4.20E−26	7.41E−17	4.43E−11	4.38E−07	−7.09E+00
37	37	74	0.00E+00	1.31E−84	2.09E−53	1.14E−26	4.31E−16	4.05E−10	3.46E−06	2.29E−03	−3.00E+00
37	38	75	0.00E+00	3.61E−91	8.70E−57	8.46E−28	1.58E−16	3.18E−10	4.25E−06	3.78E−03	−3.42E+00
37	39	76	0.00E+00	1.51E−96	4.52E−59	3.43E−28	1.92E−16	6.26E−10	1.11E−05	1.21E−02	−3.72E+00
37	40	77	0.00E+00	7.79E−95	6.22E−59	3.96E−28	1.96E−16	5.11E−10	7.46E−06	7.00E−03	−3.69E+00
37	41	78	0.00E+00	1.65E−98	1.76E−60	4.70E−28	1.03E−15	5.60E−09	1.19E−04	1.34E−01	−4.05E+00
37	42	79	0.00E+00	0.00E+00	1.59E−62	1.39E−29	6.80E−17	6.00E−10	1.70E−05	2.29E−02	−4.08E+00
37	43	80	0.00E+00	0.00E+00	7.82E−65	1.44E−30	1.79E−17	2.38E−10	8.61E−06	1.38E−02	−4.31E+00
38	33	71	0.00E+00	2.02E−68	7.66E−45	2.73E−23	3.00E−14	4.93E−09	1.43E−05	4.52E−03	−1.83E+00
38	34	72	0.00E+00	2.21E−71	3.28E−46	1.05E−23	1.92E−14	3.84E−09	1.24E−05	4.21E−03	−2.10E+00
38	35	73	0.00E+00	4.06E−73	4.75E−47	5.86E−24	1.69E−14	4.39E−09	1.69E−05	6.50E−03	−2.25E+00
38	36	74	0.00E+00	7.07E−77	1.38E−48	5.09E−24	2.90E−14	8.53E−09	3.24E−05	1.20E−02	−2.55E+00
38	37	75	0.00E+00	2.12E−85	2.00E−53	3.85E−26	1.64E−15	1.29E−09	8.78E−06	4.82E−03	−3.08E+00
38	38	76	0.00E+00	2.02E−90	5.33E−56	8.45E−27	1.37E−15	1.70E−09	1.30E−05	7.01E−03	−3.36E+00
38	39	77	0.00E+00	2.45E−92	4.61E−57	1.19E−26	6.74E−15	1.53E−08	1.58E−04	1.00E−01	−3.51E+00
38	40	78	0.00E+00	2.06E−91	1.12E−56	1.64E−26	1.27E−14	4.34E−08	6.23E−04	4.83E−01	−3.43E+00
38	41	79	0.00E+00	2.25E−95	3.66E−59	7.63E−28	1.55E−15	9.27E−09	2.02E−04	2.21E−01	−3.66E+00
38	42	80	0.00E+00	1.62E−96	1.16E−59	7.99E−28	2.73E−15	2.23E−08	5.85E−04	7.06E−01	−3.75E+00
38	43	81	0.00E+00	1.93E−96	2.20E−59	1.89E−27	7.48E−15	7.28E−08	2.31E−03	3.38E+00	−3.78E+00
38	44	82	0.00E+00	0.00E+00	8.05E−65	2.32E−30	6.83E−17	1.89E−09	1.15E−04	2.62E−01	−4.27E+00
38	45	83	0.00E+00	0.00E+00	1.79E−70	1.64E−33	2.68E−19	1.54E−11	1.37E−06	3.95E−03	−4.78E+00
38	46	84	0.00E+00	0.00E+00	7.04E−74	6.90E−35	6.66E−20	1.04E−11	1.73E−06	7.36E−03	−5.17E+00
38	47	85	0.00E+00	0.00E+00	1.20E−80	1.91E−38	1.56E−22	5.95E−14	1.57E−08	9.03E−05	−5.83E+00
38	48	86	0.00E+00	0.00E+00	5.32E−86	5.58E−41	4.14E−24	4.40E−15	1.98E−09	1.52E−05	−6.36E+00
38	49	87	0.00E+00	0.00E+00	3.59E−96	4.10E−46	3.05E−28	1.31E−18	1.99E−12	4.48E−08	−7.32E+00
38	50	88	0.00E+00	0.00E+00	0.00E+00	4.27E−49	5.29E−30	6.24E−20	1.22E−13	2.74E−09	−7.91E+00
39	39	78	0.00E+00	4.99E−93	8.41E−59	2.05E−29	8.55E−18	2.68E−11	4.50E−07	4.55E−04	−3.38E+00
39	40	79	0.00E+00	1.79E−94	5.96E−59	1.60E−28	1.36E−16	5.56E−10	1.04E−05	1.08E−02	−3.56E+00
39	41	80	0.00E+00	0.00E+00	3.02E−76	4.70E−37	2.77E−22	2.76E−14	3.40E−09	1.23E−05	−5.30E+00
39	42	81	0.00E+00	9.05E−95	1.16E−58	1.24E−27	2.10E−15	1.32E−08	3.23E−04	4.01E−01	−3.61E+00
39	43	82	0.00E+00	3.91E−96	1.04E−59	1.44E−28	1.74E−16	7.05E−10	1.15E−05	1.03E−02	−3.68E+00
39	44	83	0.00E+00	0.00E+00	9.41E−62	6.72E−29	4.74E−16	5.99E−09	2.19E−04	3.47E−01	−3.96E+00
39	45	84	0.00E+00	0.00E+00	1.59E−67	3.10E−32	1.10E−18	2.77E−11	1.45E−06	2.85E−03	−4.49E+00
39	46	85	0.00E+00	0.00E+00	9.14E−71	1.95E−33	5.16E−19	4.30E−11	4.97E−06	1.69E−02	−4.82E+00
39	47	86	0.00E+00	0.00E+00	1.40E−78	8.83E−38	2.05E−22	3.80E−14	6.33E−09	2.69E−05	−5.52E+00
39	48	87	0.00E+00	0.00E+00	1.77E−86	2.97E−41	3.12E−24	4.64E−15	2.82E−09	2.78E−05	−6.37E+00
39	49	88	0.00E+00	0.00E+00	3.02E−93	1.50E−45	6.38E−28	1.78E−18	1.67E−12	2.44E−08	−6.97E+00
39	50	89	0.00E+00	0.00E+00	0.00E+00	2.13E−49	4.82E−30	7.98E−20	1.92E−13	4.91E−09	−7.96E+00
40	36	76	0.00E+00	1.31E−79	1.40E−51	6.79E−27	9.05E−17	5.89E−11	4.22E−07	2.51E−04	−2.42E+00
40	37	77	0.00E+00	7.39E−85	1.02E−54	2.84E−28	1.59E−17	2.10E−11	2.27E−07	1.78E−04	−2.74E+00
40	38	78	0.00E+00	2.10E−81	4.53E−52	1.23E−26	2.36E−16	1.76E−10	1.35E−06	8.44E−04	−2.63E+00
40	39	79	0.00E+00	1.06E−78	7.52E−50	6.27E−25	6.50E−15	3.31E−09	2.00E−05	1.07E−02	−2.57E+00
40	40	80	0.00E+00	2.07E−74	6.15E−47	1.57E−22	7.50E−13	1.50E−07	3.63E−04	8.73E−02	−2.30E+00
40	41	81	0.00E+00	1.99E−92	3.12E−58	7.36E−29	2.84E−17	7.96E−11	1.19E−06	1.06E−03	−3.31E+00
40	42	82	0.00E+00	1.04E−93	2.55E−58	1.25E−27	1.31E−15	4.80E−09	6.98E−05	5.50E−02	−3.44E+00
40	43	83	0.00E+00	2.43E−93	2.76E−58	1.01E−27	1.33E−15	6.64E−09	1.29E−04	1.30E−01	−3.41E+00
40	44	84	0.00E+00	0.00E+00	6.90E−65	7.71E−31	1.36E−17	2.51E−10	1.09E−05	1.89E−02	−4.12E+00
40	45	85	0.00E+00	0.00E+00	2.23E−64	1.01E−30	1.42E−17	2.38E−10	9.75E−06	1.60E−02	−4.05E+00
40	46	86	0.00E+00	0.00E+00	4.85E−66	2.29E−31	8.74E−18	2.79E−10	1.82E−05	4.28E−02	−4.22E+00
40	47	87	0.00E+00	0.00E+00	7.16E−74	2.05E−35	1.45E−20	2.02E−12	3.21E−07	1.36E−03	−4.98E+00

40	48	88	0.00E+00	0.00E+00	4.40E−78	1.88E−37	7.22E−22	2.33E−13	6.10E−08	3.58E−04	−5.41E+00
40	49	89	0.00E+00	0.00E+00	5.43E−86	1.86E−41	1.17E−24	1.35E−15	7.11E−10	6.48E−06	−6.19E+00
40	50	90	0.00E+00	0.00E+00	8.96E−91	8.50E−44	3.76E−26	1.20E−16	1.15E−10	1.53E−06	−6.67E+00
40	51	91	0.00E+00	0.00E+00	4.01E−78	1.23E−37	2.39E−22	4.23E−14	6.88E−09	2.85E−05	−5.44E+00
40	52	92	0.00E+00	1.72E−85	2.44E−53	5.33E−25	1.52E−13	3.89E−07	5.08E−03	3.70E+00	−2.96E+00
41	41	82	0.00E+00	3.51E−60	4.04E−39	3.82E−20	3.52E−12	1.53E−07	1.97E−04	3.64E−02	−1.33E+00
41	42	83	0.00E+00	6.34E−89	4.96E−56	1.31E−27	2.08E−16	3.30E−10	3.24E−06	2.12E−03	−3.03E+00
41	43	84	7.31E−99	1.13E−40	1.31E−26	6.63E−12	4.17E−05	4.34E−01	2.02E+02	1.64E+04	−2.34E−01
41	44	85	0.00E+00	4.44E−97	9.79E−61	9.26E−30	1.24E−17	6.39E−11	1.35E−06	1.49E−03	−3.56E+00
41	45	86	0.00E+00	0.00E+00	8.34E−65	3.99E−31	4.52E−18	6.65E−11	2.59E−06	4.29E−03	−4.06E+00
41	46	87	0.00E+00	0.00E+00	3.10E−68	2.38E−33	9.36E−20	3.12E−12	2.17E−07	5.45E−04	−4.28E+00
41	47	88	0.00E+00	0.00E+00	4.52E−67	2.78E−32	8.46E−19	2.22E−11	1.27E−06	2.75E−03	−4.27E+00
41	48	89	0.00E+00	0.00E+00	4.34E−77	1.30E−37	1.86E−22	3.72E−14	7.42E−09	3.69E−05	−5.16E+00
41	49	90	0.00E+00	0.00E+00	3.65E−83	1.85E−40	2.62E−24	1.46E−15	5.06E−10	3.57E−06	−5.80E+00
41	50	91	0.00E+00	0.00E+00	4.41E−86	4.81E−42	2.29E−25	2.41E−16	1.21E−10	1.05E−06	−6.05E+00
41	51	92	0.00E+00	0.00E+00	1.41E−70	1.60E−34	8.53E−21	2.77E−13	1.86E−08	4.80E−05	−4.58E+00
41	52	93	0.00E+00	6.66E−70	1.86E−44	3.75E−21	2.99E−11	1.56E−05	8.71E−02	3.76E+01	−1.93E+00
42	38	80	0.00E+00	9.72E−70	6.88E−46	1.69E−24	1.81E−15	3.58E−10	1.31E−06	5.05E−04	−1.73E+00
42	39	81	0.00E+00	2.87E−65	5.51E−43	1.02E−22	5.02E−14	6.90E−09	2.03E−05	6.83E−03	−1.49E+00
42	40	82	0.00E+00	5.45E−42	2.64E−28	7.54E−15	1.04E−08	5.58E−05	2.04E−02	1.60E+00	−1.93E−01
42	41	83	0.00E+00	8.85E−43	2.14E−28	2.25E−14	3.04E−08	1.17E−04	3.06E−02	1.85E+00	−2.34E−01
42	43	85	0.00E+00	5.35E−41	5.55E−27	2.36E−12	2.07E−05	3.16E−01	1.92E+02	1.75E+04	−1.64E−01
42	44	86	0.00E+00	2.50E−95	1.52E−59	7.49E−29	8.90E−17	4.24E−10	8.21E−06	8.32E−03	−3.40E+00
42	45	87	0.00E+00	1.88E−99	3.81E−62	4.83E−30	2.03E−17	1.89E−10	5.56E−06	7.62E−03	−3.66E+00
42	46	88	0.00E+00	0.00E+00	1.79E−66	4.04E−32	9.56E−19	2.39E−11	1.37E−06	2.98E−03	−4.10E+00
42	47	89	0.00E+00	0.00E+00	2.04E−68	4.34E−33	2.66E−19	1.03E−11	7.28E−07	1.77E−03	−4.28E+00
42	48	90	0.00E+00	0.00E+00	1.99E−73	1.49E−35	6.77E−21	8.03E−13	1.21E−07	5.01E−04	−4.79E+00
42	49	91	0.00E+00	0.00E+00	2.34E−78	4.98E−38	1.49E−22	4.48E−14	1.17E−08	7.18E−05	−5.28E+00
42	50	92	0.00E+00	0.00E+00	1.30E−81	1.26E−39	1.43E−23	8.86E−15	3.60E−09	2.88E−05	−5.61E+00
42	51	93	0.00E+00	0.00E+00	7.71E−69	4.67E−33	4.02E−19	1.88E−11	1.48E−06	3.77E−03	−4.36E+00
42	52	94	0.00E+00	1.38E−72	6.54E−46	9.92E−22	1.65E−11	1.40E−05	1.17E−01	7.21E+01	−2.07E+00
42	53	95	0.00E+00	1.26E−75	1.23E−47	1.52E−22	3.88E−12	3.27E−06	2.33E−02	1.17E+01	−2.24E+00
42	54	96	0.00E+00	2.77E−84	7.32E−53	3.96E−25	8.82E−14	2.29E−07	3.18E−03	2.42E+00	−2.76E+00
42	55	97	0.00E+00	7.33E−86	1.11E−53	1.40E−25	2.24E−14	3.59E−08	3.34E−04	1.97E−01	−2.85E+00
43	44	87	0.00E+00	2.59E−74	3.10E−47	3.22E−23	2.01E−13	6.99E−08	2.90E−04	1.08E−01	−2.07E+00
43	45	88	0.00E+00	0.00E+00	8.97E−65	1.70E−32	6.60E−20	6.98E−13	2.62E−08	4.71E−05	−3.82E+00
43	46	89	0.00E+00	7.71E−86	4.97E−54	2.20E−26	2.73E−15	4.73E−09	5.44E−05	4.13E−02	−2.77E+00
43	47	90	0.00E+00	0.00E+00	9.32E−65	1.10E−31	7.56E−19	9.68E−12	3.78E−07	6.64E−04	−3.88E+00
43	48	91	0.00E+00	0.00E+00	1.23E−68	1.66E−33	7.71E−20	2.86E−12	2.18E−07	5.97E−04	−4.23E+00
43	49	92	0.00E+00	0.00E+00	6.51E−79	9.10E−39	1.56E−23	3.21E−15	6.61E−10	3.50E−06	−5.29E+00
43	50	93	0.00E+00	0.00E+00	1.20E−80	2.12E−39	1.09E−23	4.11E−15	1.22E−09	8.10E−06	−5.45E+00
43	51	94	0.00E+00	0.00E+00	2.25E−65	1.18E−31	1.85E−18	3.65E−11	1.77E−06	3.42E−03	−3.92E+00
43	52	95	0.00E+00	2.94E−69	4.72E−44	5.45E−21	4.07E−11	2.40E−05	1.66E−01	9.28E+01	−1.80E+00
43	53	96	0.00E+00	2.88E−69	6.18E−44	7.02E−21	4.22E−11	1.89E−05	9.76E−02	4.09E+01	−1.79E+00
43	54	97	0.00E+00	7.08E−80	2.21E−50	4.21E−24	3.55E−13	6.40E−07	7.72E−03	5.64E+00	−2.44E+00
43	55	98	0.00E+00	5.71E−81	7.22E−51	2.49E−24	1.40E−13	1.48E−07	1.14E−03	6.10E−01	−2.49E+00
44	42	86	0.00E+00	5.54E−62	4.52E−40	1.21E−20	1.39E−12	6.54E−08	8.63E−05	1.57E−02	−1.28E+00
44	43	87	0.00E+00	4.61E−66	2.63E−42	2.90E−21	1.12E−12	8.91E−08	1.60E−04	3.73E−02	−1.53E+00
44	44	88	0.00E+00	3.88E−97	3.35E−61	3.19E−31	1.61E−19	6.04E−13	1.21E−08	1.41E−05	−3.37E+00
44	45	89	0.00E+00	0.00E+00	4.35E−65	7.85E−33	2.76E−20	2.88E−13	1.13E−08	2.19E−05	−3.77E+00
44	46	90	0.00E+00	0.00E+00	5.67E−67	3.75E−33	4.65E−20	8.37E−13	4.12E−08	8.53E−05	−3.98E+00
44	47	91	0.00E+00	4.42E−96	3.44E−60	1.83E−29	2.41E−17	1.35E−10	3.16E−06	3.90E−03	−3.31E+00
44	48	92	0.00E+00	0.00E+00	3.40E−74	1.77E−36	5.66E−22	5.19E−14	6.56E−09	2.47E−05	−4.71E+00
44	49	93	0.00E+00	0.00E+00	1.07E−73	5.11E−36	1.56E−21	1.32E−13	1.54E−08	5.38E−05	−4.69E+00
44	50	94	0.00E+00	0.00E+00	3.00E−75	7.54E−37	4.87E−22	7.10E−14	1.25E−08	5.92E−05	−4.82E+00
44	51	95	0.00E+00	0.00E+00	2.09E−63	1.00E−30	1.06E−17	2.26E−10	1.29E−05	2.86E−02	−3.67E+00
44	52	96	0.00E+00	1.52E−68	1.19E−43	5.51E−21	3.18E−11	1.86E−05	1.38E−01	8.57E+01	−1.69E+00
44	53	97	0.00E+00	3.37E−69	4.66E−44	3.83E−21	2.48E−11	1.36E−05	8.77E−02	4.56E+01	−1.73E+00
44	54	98	0.00E+00	1.06E−77	3.55E−49	1.04E−23	5.25E−13	8.53E−07	1.09E−02	9.07E+00	−2.24E+00
44	55	99	0.00E+00	1.63E−79	3.99E−50	4.04E−24	1.97E−13	2.34E−07	2.14E−03	1.33E+00	−2.33E+00

44	56	100	0.00E+00	4.07E−88	2.47E−55	1.06E−26	5.00E−15	2.10E−08	4.21E−04	4.32E−01	−2.85E+00
45	45	90	0.00E+00	0.00E+00	1.61E−68	2.11E−35	2.03E−22	4.05E−15	2.46E−10	6.38E−07	−3.83E+00
45	46	91	0.00E+00	0.00E+00	2.67E−69	2.50E−35	5.02E−22	1.49E−14	1.14E−09	3.39E−06	−4.02E+00
45	47	92	0.00E+00	0.00E+00	4.63E−72	1.93E−36	1.22E−22	6.13E−15	6.35E−10	2.33E−06	−4.35E+00
45	48	93	0.00E+00	0.00E+00	1.99E−74	2.14E−37	3.72E−23	3.04E−15	4.23E−10	1.89E−06	−4.60E+00
45	49	94	0.00E+00	0.00E+00	2.87E−75	9.11E−38	2.65E−23	2.81E−15	4.48E−10	2.16E−06	−4.70E+00
45	50	95	0.00E+00	0.00E+00	9.81E−76	1.31E−37	5.55E−23	6.76E−15	1.15E−09	5.71E−06	−4.78E+00
45	51	96	0.00E+00	1.12E−93	1.23E−58	1.28E−28	1.43E−16	9.51E−10	2.74E−05	3.96E−02	−3.11E+00
45	52	97	0.00E+00	3.16E−65	1.30E−41	3.87E−20	8.35E−11	3.06E−05	1.77E−01	9.34E+01	−1.41E+00
45	53	98	0.00E+00	2.63E−65	1.92E−41	8.57E−20	2.05E−10	7.57E−05	4.18E−01	2.05E+02	−1.44E+00
45	54	99	0.00E+00	2.80E−73	4.53E−46	5.34E−22	8.51E−12	8.85E−06	9.75E−02	7.93E+01	−1.92E+00
45	55	100	0.00E+00	1.43E−77	1.08E−48	3.46E−23	1.29E−12	1.51E−06	1.46E−02	9.70E+00	−2.20E+00
45	56	101	0.00E+00	8.47E−85	5.51E−53	2.09E−25	4.82E−14	1.75E−07	3.71E−03	4.29E+00	−2.61E+00
46	44	90	0.00E+00	9.42E−99	4.54E−63	4.49E−33	4.64E−21	3.75E−14	1.48E−09	2.98E−06	−3.21E+00
46	45	91	0.00E+00	0.00E+00	1.03E−65	2.62E−34	8.06E−22	1.06E−14	5.52E−10	1.35E−06	−3.49E+00
46	46	92	0.00E+00	0.00E+00	1.87E−66	1.75E−34	5.81E−22	7.48E−15	3.81E−10	9.22E−07	−3.67E+00
46	47	93	0.00E+00	0.00E+00	2.15E−69	1.36E−35	2.04E−22	5.64E−15	4.53E−10	1.49E−06	−4.01E+00
46	48	94	0.00E+00	0.00E+00	5.57E−73	9.02E−37	8.18E−23	4.86E−15	5.64E−10	2.24E−06	−4.41E+00
46	49	95	0.00E+00	0.00E+00	1.12E−72	1.51E−36	1.15E−22	6.03E−15	6.45E−10	2.48E−06	−4.39E+00
46	50	96	0.00E+00	0.00E+00	4.92E−75	2.45E−37	8.10E−23	1.01E−14	1.89E−09	1.03E−05	−4.64E+00
46	51	97	0.00E+00	1.74E−92	6.04E−58	2.04E−28	1.50E−16	8.12E−10	2.07E−05	2.74E−02	−2.95E+00
46	52	98	0.00E+00	2.77E−62	6.28E−40	1.82E−19	1.93E−10	5.39E−05	2.77E−01	1.37E+02	−1.16E+00
46	53	99	0.00E+00	1.65E−61	1.32E−39	2.66E−19	2.65E−10	7.00E−05	3.41E−01	1.64E+02	−1.13E+00
46	54	100	0.00E+00	2.86E−69	3.52E−44	1.42E−21	8.13E−12	5.28E−06	4.54E−02	3.23E+01	−1.58E+00
46	55	101	0.00E+00	5.81E−72	8.49E−46	2.42E−22	2.36E−12	1.76E−06	1.50E−02	9.96E+00	−1.75E+00
46	56	102	0.00E+00	1.87E−78	1.33E−49	3.04E−24	1.38E−13	2.46E−07	3.74E−03	3.72E+00	−2.13E+00
47	47	94	0.00E+00	0.00E+00	3.66E−72	1.58E−37	5.56E−24	2.41E−16	2.50E−11	9.62E−08	−4.01E+00
47	48	95	0.00E+00	0.00E+00	2.89E−73	1.20E−37	9.43E−24	6.43E−16	8.99E−11	4.25E−07	−4.24E+00
47	49	96	0.00E+00	0.00E+00	2.22E−74	5.96E−38	7.51E−24	6.11E−16	9.34E−11	4.68E−07	−4.40E+00
47	50	97	0.00E+00	0.00E+00	9.63E−74	2.06E−37	2.34E−23	1.72E−15	2.41E−10	1.13E−06	−4.37E+00
47	51	98	0.00E+00	6.58E−93	1.42E−58	1.69E−29	5.88E−18	2.01E−11	4.02E−07	4.85E−04	−2.92E+00
47	52	99	0.00E+00	9.39E−57	8.68E−37	7.28E−18	2.29E−09	2.99E−04	8.72E−01	2.82E+02	−8.09E−01
47	53	100	0.00E+00	7.05E−60	2.22E−38	1.14E−18	5.32E−10	8.30E−05	2.69E−01	9.25E+01	−9.58E−01
47	54	101	0.00E+00	6.22E−62	3.97E−40	5.29E−20	4.07E−11	9.07E−06	4.00E−02	1.82E+01	−1.07E+00
47	55	102	0.00E+00	1.92E−65	5.15E−42	8.51E−21	1.48E−11	4.79E−06	2.47E−02	1.15E+01	−1.26E+00
47	56	103	0.00E+00	1.01E−72	1.05E−46	2.77E−23	3.20E−13	2.95E−07	3.11E−03	2.53E+00	−1.70E+00
48	44	92	0.00E+00	0.00E+00	3.61E−66	3.32E−35	8.69E−23	1.24E−15	7.28E−11	1.98E−07	−3.25E+00
48	45	93	0.00E+00	0.00E+00	2.06E−68	5.91E−36	4.30E−23	9.23E−16	6.62E−11	2.04E−07	−3.52E+00
48	46	94	0.00E+00	0.00E+00	4.05E−69	2.59E−36	2.51E−23	6.89E−16	6.08E−11	2.20E−07	−3.63E+00
48	47	95	0.00E+00	0.00E+00	1.28E−71	2.77E−37	7.85E−24	3.34E−16	3.65E−11	1.50E−07	−3.94E+00
48	48	96	0.00E+00	0.00E+00	2.00E−74	7.58E−39	4.86E−25	3.28E−17	5.01E−12	2.68E−08	−4.25E+00
48	49	97	0.00E+00	0.00E+00	1.89E−73	9.51E−38	6.34E−24	3.88E−16	5.17E−11	2.43E−07	−4.27E+00
48	50	98	0.00E+00	0.00E+00	1.09E−73	1.82E−37	1.60E−23	1.03E−15	1.38E−10	6.49E−07	−4.32E+00
48	51	99	0.00E+00	2.56E−89	1.91E−56	3.61E−28	7.69E−17	1.94E−10	3.02E−06	2.97E−03	−2.63E+00
48	52	100	0.00E+00	6.76E−54	4.37E−35	2.36E−17	2.97E−09	2.79E−04	7.23E−01	2.22E+02	−5.18E−01
48	53	101	0.00E+00	1.53E−51	8.02E−34	8.36E−17	6.63E−09	5.06E−04	1.17E+00	3.36E+02	−3.79E−01
48	54	102	0.00E+00	1.82E−54	2.54E−35	1.86E−17	2.74E−09	3.11E−04	9.88E−01	3.66E+02	−5.46E−01
48	55	103	0.00E+00	1.66E−60	4.35E−39	2.36E−19	1.57E−10	3.82E−05	1.94E−01	1.01E+02	−9.25E−01
48	56	104	0.00E+00	7.94E−65	1.38E−41	1.42E−20	2.42E−11	9.53E−06	6.58E−02	4.24E+01	−1.17E+00
49	49	98	0.00E+00	0.00E+00	2.02E−78	2.60E−41	1.53E−26	2.56E−18	5.79E−13	3.83E−09	−4.54E+00
49	50	99	0.00E+00	0.00E+00	7.98E−79	6.84E−41	5.31E−26	1.08E−17	2.87E−12	2.16E−08	−4.66E+00
49	51	100	0.00E+00	1.07E−91	3.73E−59	2.16E−30	7.52E−19	2.74E−12	6.17E−08	8.72E−05	−2.83E+00
49	52	101	5.76E−70	4.84E−53	5.28E−38	1.49E−19	6.61E−11	8.16E−06	2.01E−02	5.57E+00	−8.25E−01
49	53	102	1.06E+39	1.80E−34	4.57E−29	3.76E−15	1.04E−07	2.73E−03	2.34E+00	2.97E+02	−6.60E−02
49	54	103	5.20E−74	3.10E−48	6.92E−34	3.26E−17	2.15E−09	1.06E−04	1.63E−01	3.47E+01	−3.05E−01
49	55	104	1.78E+15	3.27E−40	3.61E−32	8.74E−17	6.73E−09	3.37E−04	4.68E−01	8.71E+01	−3.39E−01
49	56	105	0.00E+00	2.48E−55	5.24E−37	3.34E−18	1.21E−09	1.77E−04	5.90E−01	2.23E+02	−6.81E−01

TABLE X. The β^+-decay half-lives $T_{1/2}$ used in the reaction rate set FRDM1. Shown are experimental data and theoretical values calculated with the QRPA or the shell model code OXBASH. When both, QRPA and shell model calculations are available, we used the shell model calculations.

Z	N	A	$T_{1/2}$ (s) exp.	QRPA	SM	Q_{EC} (MeV)
32	28	60		8.20E−02		1.17E+01
32	29	61	4.00E−02			1.30E+01
32	30	62	1.10E−01			8.47E+00
32	31	63	9.50E−02			9.56E+00
32	32	64	6.37E+01			4.41E+00
32	33	65	3.09E+01			6.24E+00
32	34	66	8.14E+03			2.10E+00
32	35	67	1.13E+03			4.22E+00
32	36	68	2.34E+07			1.06E−01
32	37	69	1.41E+05			2.23E+00
32	39	71	9.88E+05			2.32E−01
33	32	65	1.90E−01			9.70E+00
33	33	66	9.58E−02			9.55E+00
33	34	67	4.25E+01			6.01E+00
33	35	68	1.52E+02			8.10E+00
33	36	69	9.12E+02			4.01E+00
33	37	70	3.16E+03			6.22E+00
33	38	71	2.35E+05			2.01E+00
33	39	72	9.36E+04			4.36E+00
33	40	73	6.94E+06			3.41E−01
33	41	74	1.54E+06			2.56E+00
34	30	64		9.17E−02		1.17E+01
34	31	65		4.27E−02		1.33E+01
34	32	66		6.49E−01		8.80E+00
34	33	67	1.07E−01			9.91E+00
34	34	68	3.55E+01			4.91E+00
34	35	69	2.74E+01			6.78E+00
34	36	70	2.47E+03			2.46E+00
34	37	71	2.84E+02			4.43E+00
34	38	72	7.26E+05			3.35E−01
34	39	73	2.57E+04			2.74E+00
34	41	75	1.03E+07			8.63E−01
35	34	69		2.02E−01		9.98E+00
35	35	70	7.91E−02			1.13E+01
35	36	71	2.14E+01			6.60E+00
35	37	72	7.86E+01			8.74E+00
35	38	73	2.04E+02			4.68E+00
35	39	74	1.52E+03			6.91E+00
35	40	75	5.80E+03			3.03E+00
35	41	76	5.83E+04			4.96E+00
35	42	77	2.05E+05			1.36E+00
35	43	78	3.88E+02			3.57E+00
36	31	67		2.15E−02		1.60E+01
36	32	68		7.26E−02		1.20E+01
36	33	69		3.32E−02		1.38E+01
36	34	70		3.90E−01		9.09E+00
36	35	71	9.70E−02			1.04E+01
36	36	72	1.72E+01			5.04E+00
36	37	73	2.70E+01			6.65E+00
36	38	74	6.90E+02			3.14E+00
36	39	75	2.58E+02			4.90E+00
36	40	76	5.33E+04			1.31E+00
36	41	77	2.68E+05			3.06E+00
36	43	79	1.26E+05			1.63E+00
36	45	81	6.62E+12			2.80E−01

37	37	74	6.50E−02		1.04E+01
37	38	75	1.90E+01		7.02E+00
37	39	76	3.65E+01		8.50E+00
37	40	77	2.27E+02		5.34E+00
37	41	78	1.06E+03		7.22E+00
37	42	79	1.37E+03		3.65E+00
37	43	80	3.34E+01		5.72E+00
37	44	81	1.65E+04		2.24E+00
37	45	82	7.64E+01		4.40E+00
37	46	83	7.45E+06		9.09E−01
37	47	84	2.83E+06		2.68E+00
38	32	70		1.59E−02	1.47E+01
38	33	71		1.40E−02	1.67E+01
38	34	72		5.64E−02	1.22E+01
38	35	73		4.40E−02	1.41E+01
38	36	74		3.86E−01	9.54E+00
38	37	75		1.97E−01	1.07E+01
38	38	76	8.90E+00		6.04E+00
38	39	77	9.00E+00		6.85E+00
38	40	78	1.50E+02		3.76E+00
38	41	79	1.32E+02		5.32E+00
38	42	80	6.38E+03		1.87E+00
38	43	81	1.34E+03		3.93E+00
38	44	82	2.21E+06		1.80E−01
38	45	83	1.17E+05		2.28E+00
38	47	85	5.60E+06		1.06E+00
39	39	78		2.15E−01	1.19E+01
39	40	79	1.48E+01		7.12E+00
39	41	80	3.50E+01		6.95E+00
39	42	81	7.24E+01		5.51E+00
39	43	82	9.50E+00		7.82E+00
39	44	83	4.25E+02		4.47E+00
39	45	84	4.60E+00		6.49E+00
39	46	85	9.65E+03		3.25E+00
39	47	86	5.31E+04		5.24E+00
39	48	87	2.87E+05		1.86E+00
39	49	88	9.21E+06		3.62E+00
40	35	75		1.68E−02	1.69E+01
40	36	76		5.25E−02	1.31E+01
40	37	77		2.94E−02	1.47E+01
40	38	78		2.61E−01	1.03E+01
40	39	79		1.65E−01	1.13E+01
40	40	80		6.85E+00	6.60E+00
40	41	81	1.50E+01		7.16E+00
40	42	82	3.20E+01		4.00E+00
40	43	83	4.40E+01		5.87E+00
40	44	84	1.55E+03		3.40E+00
40	45	85	4.72E+02		4.69E+00
40	46	86	5.94E+04		1.48E+00
40	47	87	6.05E+03		3.67E+00
40	48	88	7.21E+06		6.73E−01
40	49	89	2.82E+05		2.83E+00
41	41	82		6.63E−01	1.07E+01
41	42	83	4.10E+00		7.50E+00
41	43	84	1.20E+01		1.13E+01
41	44	85	2.09E+01		6.00E+00
41	45	86	8.80E+01		7.98E+00
41	46	87	2.22E+02		5.17E+00
41	47	88	8.70E+02		7.55E+00
41	48	89	4.25E+03		4.29E+00

41	49	90	5.26E+04			6.11E+00
41	50	91	2.20E+10			1.25E+00
41	51	92	1.10E+15			2.01E+00
42	37	79		5.02E−03		1.83E+01
42	38	80		4.79E−02		1.33E+01
42	39	81		3.17E−02		1.48E+01
42	40	82		6.65E−02		1.14E+01
42	41	83		8.82E−02		1.19E+01
42	42	84		1.07E+00		4.84E+00
42	43	85		3.68E−01		9.60E+00
42	44	86	1.96E+01			5.27E+00
42	45	87	1.45E+01			6.49E+00
42	46	88	4.80E+02			3.37E+00
42	47	89	1.22E+02			5.58E+00
42	48	90	2.04E+04			2.49E+00
42	49	91	9.29E+02			4.43E+00
42	51	93	1.26E+11			4.05E−01
43	43	86		8.83E−02	2.40E−02	1.37E+01
43	44	87		3.38E−01	5.82E−01	9.49E+00
43	45	88		2.73E−01	1.70E+00	1.16E+01
43	46	89	1.28E+01			7.51E+00
43	47	90	8.70E+00			8.96E+00
43	48	91	1.88E+02			6.22E+00
43	49	92	7.38E+01			7.87E+00
43	50	93	9.90E+03			3.20E+00
43	51	94	1.76E+04			4.25E+00
43	52	95	7.20E+04			1.69E+00
43	53	96	5.27E+06			2.97E+00
44	39	83		3.74E−03		1.84E+01
44	40	84		1.09E−02		1.45E+01
44	41	85		2.01E−02		1.50E+01
44	42	86		8.63E−02	4.39E−02	1.06E+01
44	43	87		6.51E−02	2.36E−02	1.24E+01
44	44	88		7.27E−01	7.07E−01	7.86E+00
44	45	89		2.94E−01	4.25E−01	9.92E+00
44	46	90	1.30E+01			6.31E+00
44	47	91	9.00E+00			7.41E+00
44	48	92	2.19E+02			4.51E+00
44	49	93	5.97E+01			6.34E+00
44	50	94	3.11E+03			1.59E+00
44	51	95	5.92E+03			2.57E+00
44	53	97	2.51E+05			1.11E+00
45	45	90		8.31E−02	2.19E−02	1.39E+01
45	46	91		4.10E−01	4.93E−01	9.88E+00
45	47	92		3.47E−01	1.00E+00	1.19E+01
45	48	93		2.65E+00	5.09E+00	8.20E+00
45	49	94		4.73E+00		1.02E+01
45	50	95	3.01E+02			5.11E+00
45	51	96	5.94E+02			6.45E+00
45	52	97	1.84E+03			3.52E+00
45	53	98	5.22E+02			5.06E+00
45	54	99	1.39E+06			2.04E+00
45	55	100	7.49E+04			3.63E+00
45	56	101	1.04E+08			5.42E−01
46	42	88		1.96E−02		1.32E+01
46	43	89		1.59E−02		1.55E+01
46	44	90		6.83E−02	3.13E−02	1.12E+01
46	45	91		5.25E−02		1.30E+01
46	46	92		5.46E−01	4.61E−01	8.55E+00
46	47	93		2.24E−01	2.83E−01	1.05E+01

46	48	94	9.00E+00			6.60E+00
46	49	95	9.21E−01			8.70E+00
46	50	96	1.22E+02			3.45E+00
46	51	97	1.86E+02			4.79E+00
46	52	98	1.06E+03			1.87E+00
46	53	99	1.28E+03			3.39E+00
46	54	100	3.14E+05			3.62E−01
46	55	101	3.05E+04			1.98E+00
47	47	94		7.71E−02	6.00E−02	1.43E+01
47	48	95	2.00E+00			1.00E+01
47	49	96	5.10E+00			1.21E+01
47	50	97	1.90E+01			6.70E+00
47	51	98	4.67E+01			8.42E+00
47	52	99	1.24E+02			5.43E+00
47	53	100	1.20E+02			7.05E+00
47	54	101	6.66E+02			4.20E+00
47	55	102	7.74E+02			5.96E+00
47	56	103	3.94E+03			2.69E+00
47	57	104	4.15E+03			4.28E+00
47	58	105	3.57E+06			1.35E+00
47	59	106	1.44E+03			2.96E+00
48	43	91		7.24E−03		1.75E+01
48	44	92		2.02E−02		1.35E+01
48	45	93		1.61E−02		1.56E+01
48	46	94		7.31E−02		1.16E+01
48	47	95		5.12E−02		1.33E+01
48	48	96		8.38E−01	6.57E−01	8.70E+00
48	49	97		2.03E−01	3.44E−01	1.06E+01
48	50	98	9.20E+00			5.20E+00
48	51	99	1.60E+01			6.89E+00
48	52	100	4.91E+01			3.88E+00
48	53	101	7.20E+01			5.48E+00
48	54	102	3.30E+02			2.59E+00
48	55	103	4.38E+02			4.14E+00
48	56	104	3.46E+03			1.14E+00
48	57	105	3.33E+03			2.74E+00
48	59	107	2.34E+04			1.42E+00
49	49	98		9.30E−02		1.40E+01
49	50	99		3.94E+00		8.02E+00
49	51	100	6.10E+00			1.02E+01
49	52	101	1.60E+01			6.41E+00
49	53	102	2.40E+01			9.25E+00
49	54	103	6.50E+01			6.05E+00
49	55	104	1.08E+02			7.91E+00
49	56	105	3.04E+02			4.85E+00
49	57	106	3.72E+02			6.52E+00
49	58	107	1.94E+03			3.43E+00
49	59	108	2.38E+03			5.16E+00
49	60	109	1.51E+04			2.02E+00
49	61	110	6.91E+01			3.88E+00
49	62	111	2.45E+05			8.65E−01

FIG. 1. The proton bound nuclei on the neutron-deficient side of stability for the FRDM (1992) and Hilf et al. mass models. An open square indicates that not all proton bound nuclei are displayed. Indicated is also the level density of the compound nucleus in the Gamow window for proton capture at $T = 1$ GK. Also marked are nuclei, where no calculation of the level density has been performed since they are not important for the rp-process. Generally, the number of levels will be larger for higher temperatures. About 10 levels are considered to be sufficient for the application of the statistical model (Hauser-Feshbach approach). Also shown are the nuclei for which experimental mass and β-decay half-life data exist (enclosed by solid line) as well as the nuclei for which only β-decay half-life measurements exist (enclosed by dashed line).

FIG. 2. Theoretical error (see Eq. 1) (in MeV) for the predictions of experimentally known masses in the isotopic chains $Z = 32$ to $Z = 49$ based on various global mass models. Only nuclei on the neutron-deficient side of the valley of stability are considered.

FIG. 3. The microscopic mass correction $\mathrm{M}_{\mathrm{mic}} = M_{\mathrm{exp}} - M_{\mathrm{drop}}$ plotted for experimentally known isotopes in the range $Z = 32$ to $Z = 45$ as a function of the P-factor. The solid line indicates the best fit through the data which was used for interpolating the microscopic mass corrections of experimentally unknown isotopes.

FIG. 4. Mass difference between experimental masses and mass model predictions as a function of the mass number A for the isotopic chains with $Z = 34$, 36, 38, 39, 40, 42, 44, 46, and 48. The solid horizontal line indicates zero mass difference.

FIG. 5. The theoretical errors (see Eq. 1)(in MeV) for the various mass model predictions for nuclei with $32 \leq Z \leq 48$ on the neutron deficient side of β-stability (excluding stable nuclei). Shown are the errors for the predictions of masses (148 nuclei), of masses measured since 1988 (23 nuclei), of proton separation energies S_{p}, of α-separation energies S_{α} and of electron capture Q-values Q_{EC}.

FIG. 6. The quadrupole deformation parameter ϵ_2 predicted by theoretical models and derived from experimental data, plotted as function of mass number A for the isotopic chains with $Z = 34$, 36, 38, 40, 42, and 44. Note that only the data derived from measurements of the quadrupole moment and the FRDM (1992) mass model predictions yield the sign of ϵ_2. For all other values the positive sign was chosen arbitrarily.

FIG. 7. $\epsilon_2/\epsilon_{2(sp)}$ values derived from experimental data as function of $N_{\mathrm{p}}N_{\mathrm{n}}$ in the range $28 \leq Z, N \leq 50$. The data yield from measurements of the reduced transition strength $B(E2: 2^+ \rightarrow 0^+)$ and of ground-state electric quadrupole moments. The solid line indicates the best fit to the data which is used for interpolating the deformation of unknown isotopes in this mass range. Since some of the data obtained from spectroscopic ground-state quadrupole moments are not correlated at all, only the $B(E2: 2^+ \rightarrow 0^+)$ data are used for the fit (see text).

FIG. 8. The predicted proton separation energies $S_{\rm p}$ for the four discussed mass models as a function of mass number A for the odd Z isotopic chains from $Z = 33$ to $Z = 49$. Shown are also the experimental data available. The solid horizontal line indicates $S_{\rm p} = 0$ (proton drip line).

FIG. 9. The proton drip line in the element range $36 \leq Z \leq 48$ as predicted by various mass models. Shown is the mass number of the first proton unbound nucleus along each isotopic chain as a function of charge number Z.

FIG. 10. The predicted α-separation energies S_α for the four discussed mass models as a function of mass number A for the even Z isotopic chains from $Z = 34$ to $Z = 48$. Shown are also the experimental data available.

FIG. 11. The predicted electron capture Q-values $Q_{\rm EC}$ for the four discussed mass models as a function of mass number A for the even Z isotopic chains from $Z = 34$ to $Z = 48$. Shown are also the experimental data available.

FIG. 12. The deviations between experimental masses and the masses predicted by the various mass models as a function of the z-component of the isospin T_z. The comparison is limited to nuclei on the proton rich side of stability with $32 \leq Z \leq 50$. The errorbars reflect the uncertainty in the mass measurements. To improve the clarity of the plot we omitted errorbars of less than 100 keV and added a small offset (< 0.5) to some of the T_z values to avoid overlapping error bars.

FIG. 13. The reaction rates calculated with SMOKER using the Hauser-Feshbach method for the reactions (a) ^{64}Ge(p,γ)^{65}As, (b) ^{81}Zr(p,γ)^{82}Nb, (c) ^{89}Ru(p,γ)^{90}Rh, and (d) ^{83}Nb(p,α)^{80}Zr as a function of temperature. The ^{64}Ge(p,γ)^{65}As reaction rate is calculated using the FRDM (1992) mass model with a level density of only 7 states in the Gamow window. For comparison the Van Wormer et al. reaction rate based on 3 experimentally known levels is shown (recalculated for the FRDM (1992) Q-value). For the other reaction rates (a), (b) and (c) we show the results as predicted with the mass and deformation parameters of the Hilf et al. and the FRDM (1992) models respectively. For the FRDM (1992) mass model we show the rate obtained with a simple level density treatment (old $\rho(E, J, \pi)$), which is identical to the one used with the Hilf et al. mass formula. In addition we show the rection rates resulting from the application of the Ignatyuk level density treatment (new $\rho(E, J, \pi)$).

FIG. 14. The Hauser-Feshbach rate for the ^{81}Zr(p,γ)^{82}Nb-reaction as a function of the proton separation energy. Indicated are also the Q-values predicted by several mass models. The rate is calculated with the ground-state deformations as predicted by the FRDM (1992) mass model for a temperature of 1.5 GK.

FIG. 15. The Hauser-Feshbach rate for the ^{81}Zr(p,γ)^{82}Nb-reaction as a function of the quadrupole deformation parameter ϵ_2 of the initial and the final nucleus respectively. The rates are calculated for a temperature of 1.5 GK with Q-values as predicted by the FRDM (1992) mass model.

FIG. 16. The ratio of the $i(\gamma,\mathrm{p})j$ photodisintegration rate to the proton capture rate $j(\mathrm{p},\gamma)i$ for various proton capture Q-values. A density of 10^6 g/cm^3 and a proton abundance of $\mathrm{Y_p}$=0.77 was assumed.

FIG. 17. The ratio of the (γ,α) photodisintegration rate to the proton capture rate as a function of temperature for the nuclei that are predicted to have low α-separation energies. Shown are the results for the mass and deformation parameters predicted by the (a) FRDM (1992) mass model, and (b) by the Hilf et al. mass formula. A density of 10^6 g/cm^3 and a proton abundance of $\mathrm{Y_p}$=0.77 was assumed.

FIG. 18. The Gamow-Teller strength function for the decay of ^{76}Sr, ^{80}Zr, ^{84}Mo, ^{88}Ru, ^{92}Pd, and ^{96}Cd from QRPA calculations using a Folded-Yukawa potential and Lipkin-Nogami pairing. For details see [46]. The strength (as defined in Eqs. 39 and 40) is summed in 0.25 MeV energy bins and then given per MeV. Q-values and deformations were taken from the FRDM (1992) mass model.

FIG. 19. Gamow-Teller strength function for the decay of ^{88}Ru, ^{92}Pd, and ^{96}Cd from shell model calculations using the code OXBASH. The strength (as defined in Eqs. 39 and 40) is summed in 0.25 MeV energy bins and then given per MeV. Q-values were obtained from the FRDM (1992) mass model.

FIG. 20. Gamow-Teller strength function from QRPA calculations for the β^+-decay of ^{80}Zr for different deformation parameters. Shown are the results for a spherical nucleus ($\epsilon_2 = \epsilon_4 = \epsilon_6 = 0$) and for a negative quadrupole deformation ϵ_2 instead of the positive value predicted by the FRDM (1992) mass model. The strength function for the deformation predicted by the FRDM (1992) mass model can be found in Fig. 18.

FIG. 21. The β-decay half-lives of ^{80}Zr and ^{84}Mo calculated using the QRPA as a function of quadrupole ground-state deformation ϵ_2. These calculations were preformed using a version of the QRPA-code that is based on Nielsson potentials, which gives slightly different half-life predictions. Indicated are also the values obtained from the deformation ϵ_2 predicted by the FRDM (1992) mass model. The β-decay Q-value was calculated with the FRDM (1992) mass model.

FIG. 22. The β-decay half-lives of (a) ^{80}Zr and (b) ^{92}Pd calculated using the QRPA as a function of the (experimentally unknown) electron capture Q-value Q_{EC}. Indicated are also the results for the Q-values obtained from various mass models.

FIG. 23. The β-decay half-lives for the isotopic chains $Z = 38$, 40, 42, 44, 46, and 48 as functions of mass number. Shown are experimental values and the data obtained from various theoretical approaches.

FIG. 24. The experimentally determined energies of the first excited 2^+ state in even-even nuclei with $28 < Z < 50$ and $28 < N < 50$ as a function of $N_{\rm p}N_{\rm n}$. The solid line shows the linear fit through the data that is used to predict unknown $E(2^+)$ values. The Zn, Ge and Se isotopes are not used in the fit and shown only for illustration.

FIG. 25. The calculated β-decay half-lives as functions of temperature for the even-even waiting point isotopes. The temperature dependence is the result of the thermal population of the first excited 2^+ state.

FIG. 26. The rp-proces reaction flow integrated over 1000 s. The solid line indicates a flow of more than 10%, the dashed line of 1–10% and the thinner dotted line for 0.1–1% of the total reaction flow in that region, which was $6.7{\cdot}10^{-3}$ mole/g in all cases. Each square resembles an isotope predicted to be proton bound by the respective mass model. Shaded isotopes are important waiting points at the shown conditions while filled squares indicate stable isotopes. In addition, p-only nuclei are marked with a "p". Shown is the flow for XRB-conditions (10^6 g/cm^3, 1.5 GK) with reaction rates FRDM1, for XRB-conditions with reaction rates FRDM2, for XRB-conditions with reaction rates HIL, and for ABH-conditions (10^4 g/cm^3, 1.5 GK) with reaction rates FRDM1.

FIG. 27. The stellar half-life of ^{68}Se including β-decay and proton capture processes as a function of the proton capture Q-value under XRB-conditions (10^6 g/cm^3, 1.5 GK). Indicated are also the various proton capture Q-value predictions by mass models as well as the estimated upper limit from the Blank et al. experiment.

FIG. 28. The density required for equal net proton capture and β-decay rate as a function of temperature for the important even-even $N = Z$ nuclei and the even-odd $N = Z + 1$ nuclei in the FRDM1 caluclations as well as for the even-even $N = Z$ nuclei in the HIL calculations. For densities above the line net proton capture is the dominant destruction mechanism, either by sequential (p,γ) reactions or via 2p-capture. The total stellar lifetime will then be greatly reduced and the waiting point will be bridged. Below the curve the nucleus will be a waiting point with a lifetime given by the β-decay rate.

FIG. 29. The abundance versus processing time for H, He and some important waiting point isotopes, assuming a solar system distribution of seed nuclei. Shown are the abundances for XRB-conditions (10^6 g/cm^3, 1.5 GK) from the calculations with the reaction rate sets FRDM1, FRDM2, and HIL.

FIG. 30. The mass number of the isotope that has built up its peak abundance as a function of processing time. The slope of the curve therefore indicates the processing speed in a certain mass region. All the calculations were done at a temperature of 1.5 GK. Shown are the results for the FRDM1, FRDM2 and HIL calculations at a density of 10^6 g/cm^3 (XRB-conditions) as well as for FRDM1 at a density of $5{\cdot}10^6$ g/cm^3 (these conditions can also occur in X-ray bursts) and 10^4 g/cm^3 (ABH-conditions).

FIG. 31. The abundance versus processing time for H, He and some important waiting point isotopes, assuming a solar system distribution of seed nuclei. Shown are the abundances for ABH-conditions (10^4 g/cm^3, 1.5 GK) from the calculations with the reaction rate set FRDM1.

FIG. 32. The contours in the temperature density plane for the processing time required to reach the maximum $A = 92$ abundance and for the overproduction factor obtained at that maximum. Plots are shown for the three sets of reaction rates FRDM1, FRDM2, and HIL. Also indicated are the temperatures and densities for XRB- and ABH-conditions that were used in this paper to simulate nucleosynthesis in X-ray bursts and accretion disks around black holes, respectively.

FIG. 33. The total abundance in each mass chain (neutron-deficient isotopes only) at the time of maximum $A = 92$ abundance. Shown are the results for XRB-conditions (10^6 g/cm^3, 1.5 GK) with reaction rates FRDM1, for XRB-conditions with reaction rates FRDM2, for XRB-conditions with reaction rates HIL, and for ABH-conditions (10^4 g/cm^3, 1.5 GK) with reaction rates FRDM1.

FIG. 34. The overproduction factor for the most neutron-deficient stable isotope in each isobaric chain in respect to the solar abundance. The overproduction factors are calulated assuming the complete β-decay of all nuclei produced at the time the $A = 92$ abundance has reached its maximum (see the corresponding abundance distribution in Fig. 33). β-delayed proton decay was neglected (see text). Shown are the overproduction factors for XRB-conditions (10^6 g/cm^3, 1.5 GK) with reaction rates FRDM1, for XRB-conditions with reaction rates FRDM2, for XRB-conditions with reaction rates HIL, and for ABH-conditions (10^4 g/cm^3, 1.5 GK) with reaction rates FRDM1.

FIG. 35. The overproduction factor for ^{92}Mo as a function of the β-decay half-life of the corresponding waiting point nucleus ^{92}Pd. The calculations were done for XRB-conditions (10^6 g/cm^3, 1.5 GK) with reaction rates FRDM1 under the assumptions for Fig. 34.

FIG. 36. The neutrino loss corrected nuclear energy produced at a density of 10^6 g/cm^3 and a temperature of 1.5 GK (XRB-conditions) as a function of processing time. Shown are the data for the 3 different reaction libraries FRDM1, FRDM2 and HIL as well as for a small network including nuclei between H and Kr and for the 3α-reaction only. These results do not reflect the energy produced in a specific astrophysical site but are rather illustrations of how changes in the reaction rate database and in the network size affect energy production.

FIG. 37. The time integrated reaction flow for the Zr-Nb cycle at a density of 10^6 g/cm^3 and a temperature of 2 GK.

FIG. 38. The branching ratio for the (p,α)-branch at ^{83}Nb as a function of temperature. A timescale faster than 0.1 s was assumed as requirement to establish equilibrium between ^{83}Nb and ^{84}Mo. This condition is fulfilled above about 1.5 GK, which is the reason for the discontinuity at that temperature.

FIG. 39. The abundance as a function of time for a density of 10^6 g/cm^3 and a temperature of 2 GK. Under these conditions the rp-process ends in the Nb-Zr cycle.