NUCLEOSYNTHESIS IN FAST EXPANSIONS OF HIGH-ENTROPY, PROTON-RICH MATTER

G. C. JORDAN IV AND B. S. MEYER

Department of Physics and Astronomy, Clemson University, 118 Kinard Laboratory, Clemson, SC 29634-0978; gjordan@ces.clemson.edu

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ABSTRACT

We demonstrate that nucleosynthesis in rapid, high-entropy expansions of proton-rich matter from high temperature and density can result in a wider variety of abundance patterns than heretofore appreciated. In particular, such expansions can produce iron-group nuclides, $p$-process nuclei, or even heavy, neutron-rich isotopes. Such diversity arises because the nucleosynthesis enters a little explored regime in which the free nucleons are not in equilibrium with the abundant $^4$He. This allows nuclei significantly heavier than iron to form in the presence of abundant free nucleons early in the expansion. As the temperature drops, nucleons increasingly assemble into $^4$He and heavier nuclei. If the assembly is efficient, the resulting depletion of free neutrons allows disintegration flows to drive nuclei back down to iron and nickel. If this assembly is inefficient, then the large abundance of free nucleons prevents the disintegration flows and leaves a distribution of heavy nuclei after reaction freezeout. For cases in between, an intermediate abundance distribution, enriched in $p$-process isotopes, is frozen out. These last expansions may contribute to the solar system’s supply of the $p$-process nuclides if mildly proton-rich, high-entropy matter is ejected from proto-neutron stars winds or other astrophysical sites. Also significant is the fact that, because the nucleosynthesis is primary, the signature of this nucleosynthesis may be evident in metal-poor stars.

Subject headings: nuclear reactions, nucleosynthesis, abundances — supernovae: general

1. INTRODUCTION

Since the early attempts to understand the $p$-process nuclei in terms of hydrogen burning in supernovae (Burbidge et al. 1957; Audouze & Truran 1975), astrophysicists have focused relatively little attention on explosive nucleosynthesis in neutron-rich stellar environments. One reason is that after hydrogen burning, mainline stellar evolution proceeds at conditions of equal numbers of neutrons and protons or at a slight degree of neutron richness. Accordingly, if the star subsequently explodes, the resulting nucleosynthesis typically occurs in a neutron-rich environment. Proton-rich, explosive nucleosynthesis in stars thus tends to occur only in rather rare settings such as when hydrogen-rich material accretes onto a white dwarf or neutron star. In this case, the nucleosynthesis can proceed in a series of proton-capture and $\beta^-$-decay reactions known as the rp-process (Wallace & Woosley 1981). Alternatively, neutrino interactions in material ejected from some astrophysical setting, such as winds from a proto-neutron star or a collapsar disk, can make the matter proton-rich prior to or during the nucleosynthesis.

A second reason little attention has focused on proton-rich, explosive nucleosynthesis is that it is thought to be already relatively well understood. The details of the rp-process have been well-studied (e.g., Schatz et al. 2001) since it was first delineated. It is generally imagined that proton-rich matter that achieves higher temperatures would freeze out from an equilibrium in which iron-group isotopes would dominate the abundances.

The purpose of this Letter is to show that explosive nucleosynthesis in proton-rich environments can be much more complex than previously thought. While it is indeed the case that the final abundances for many conditions are dominated by iron-group isotopes, for sufficiently fast, high-entropy expansions of proton-rich matter, isotopes considerably heavier than iron can form. These other distributions of nuclei can include interesting quantities of light and heavy $p$-process isotopes or even neutron-rich species usually ascribed to the $\beta^+$-process.

We have not identified an astrophysical site in which such fast, high-entropy, proton-rich expansions may occur. A possible setting, however, is a neutrino-driven wind from a high-mass proto-neutron star early in its epoch of Kelvin-Helmholtz cooling by neutrino emission. Some general relativistic calculations find wind entropies, $s/k_B$, of order 100–200 and expansion timescales (which we here define as the radial expansion timescale, $\tau = r/v_r$, where $r$ is the radial coordinate and $v_r$ is the radial velocity of a parcel of wind matter) as short as a few milliseconds (e.g., Otsuki et al. 2000; Thompson et al. 2001). The $Y_e$, that is, the electron-to-baryon ratio, in these winds is set by the interaction of neutrinos and antineutrinos with free nucleons. Most wind calculations studied to date have shown ejected matter to be neutron-rich, but recent one-dimensional supernova models show development of $Y_e > 0.5$ neutrino-heated matter when the treatment of neutrino transport is made increasingly realistic (e.g., Ramp & Janka 2000; Liebendörfer et al. 2001). Other sites may also be possible. For example, neutrino interactions also give fast, proton-rich expansions in winds from the inner regions of collapsar models (e.g., Pruet et al. 2004a, 2004b; Surman & McLaughlin 2004). So far those models give only modest entropies, and the resulting nucleosynthesis yields do not show the variety we find. However, as future models become more realistic (such as the inclusion of strong magnetic fields by Kohri et al. 2004, for example), they may yield conditions more similar to those we study here.

2. CALCULATIONS

The model of the expansion of matter occurring in the fast winds was adopted from the fast expansion calculations of Meyer (2002) and similarly uses the Clemson Nucleosynthesis Code (Meyer 1995) with NACRE (Angulo et al. 1999) and NON-SMOKE (Rauscher & Thielemann 2000) rate compilations. The nuclear reaction network contains 3942 isotopes of all elements up to polonium ($Z = 84$) and follows 30,318 strong, electromagnetic, and weak interactions among these species. Neutrino reactions on nucleons and heavier nuclei are not included in the calculations. The abundance of heavy nuclei is quite small throughout the calculation, so neutrino-nucleus interactions will have little effect on the electron-to-baryon ratio.
Moreover, because of the very rapid expansions we consider, we expect neutrino-nucleon interactions to change \( Y_e \) very little once the material is actually ejected, although we plan to consider such reactions in future work. In our calculations, material expands at constant entropy from high temperature and density. Since the calculations model the acceleration phase of the winds, \( r \) grows exponentially on a timescale \( \tau \), so that the density drops exponentially on the timescale \( \tau/5 \).

We computed several fast expansions. Table 1 presents the final production factors, defined as \( Q = X/X_G \), where \( X \) and \( X_G \) are the final mass fractions of species \( i \) in the expansion and in the solar system, respectively. For all calculations, the initial composition of the starting material was a mixture of protons and neutrons that gave the desired initial value of \( Y_e \). This material was given an initial temperature of \( T_0 = 7\times10^9 \) K and an initial density consistent with the chosen value of entropy. The calculations were halted after the temperature had dropped below \( T_0 = 0.01 \), by which time all capture reactions had long since ceased and nuclei were only beta-decaying back to stability.

We focus on a particularly interesting expansion with \( s/k_G = 145 \), \( \tau = 0.003 \) s, and \( Y_e = 0.510 \), which produces a significant quantity of the light \( p \)-process nuclei \(^{56}\)Ni and \(^{60}\)Ru. Figure 1 shows the nuclear abundances as a function of atomic number for 4 times in the expansion.

In the initial stages of the expansion, nuclear reactions are sufficiently fast for the system to maintain nuclear statistical equilibrium (NSE). Table 2 shows that, at \( T_0 = 9.5 \), the network and NSE mass fractions of neutrons, protons, and \(^4\)He are identical. We note that other species contribute much smaller mass fractions at this time (and throughout the expansion). As the temperature falls, the neutrons and protons begin to undergo \(^4\)He. By \( T_0 = 7.5 \), the network mass fractions of neutrons and protons are greater than their NSE values, while the network mass fraction of \(^4\)He is less than its NSE value. The network abundances are clearly no longer in NSE at this stage in the expansion.

This remarkable result is due to the fact that the flows from neutrons and protons into \(^4\)He are slower than NSE demands. At high temperature, \(^2\)H, \(^3\)H, and \(^4\)He are in equilibrium with the free nucleons, but their abundances are low because of the high entropy. Since these nuclei carry the flows that assemble \(^4\)He from the free nucleons, their low abundances render them unable to assemble \(^4\)He as rapidly as NSE demands. Quantitatively, at \( T_0 = 7.5 \) neutrons are disappearing (because of assembly through \(^2\)H and \(^4\)He into \(^4\)He) on a timescale of 0.43 ms; however, the NSE neutron abundance is changing on a faster timescale of 0.33 ms.

This disequilibrium between the nucleons and the \(^4\)He grows as the temperature and density fall further. By \( T_0 = 6 \), for example, the neutrons, protons, and \(^4\)He are even farther from their NSE values. At this point in the expansion, the network is consuming neutrons on a 0.77 ms timescale, but the NSE neutron abundance is falling on a more rapid 0.15 ms timescale. The network is not keeping pace with the NSE.

While the free nucleons and \(^4\)He dominate the mass fractions in this calculation, some heavy nuclei do form. As evident from Table 2, the number of heavy nuclei at \( T_0 = 6 \) is only \(~10^{-8} \) nucleon\(^{-1}\). That number increases to \(~10^{-6} \) nucleon\(^{-1}\) by \( T_0 = 5 \), to \(~10^{-5} \) nucleon\(^{-1}\) by \( T_0 = 4.39 \), and freezes out at \( 2.90 \times 10^{-7} \) nucleon\(^{-1}\). The critical point is that the disequilibrium between free nucleons and \(^4\)He allows a large abundance of free nucleons to exist at temperatures that would normally see almost all the free nucleons locked into \(^4\)He. The few heavy nuclei that have been forming at these times thus coexist with a large overabundance of free nucleons (Meyer 2002). Through a rapid sequence of neutron and proton captures, these few heavy nuclei are driven to quite high nuclear mass, as shown in the \( T_0 = 4.39 \) panel in Figure 1.

As the temperature continues to fall, new heavy nuclei assemble from the abundant \(^4\)He. The growing abundance of heavy nuclei is then increasingly able to catalyze the synthesis of \(^4\)He from free neutrons and protons through reaction cycles such as \(^{20}\)Ne(\(n, \gamma\))\(^{21}\)Ne(\(n, \gamma\))\(^{22}\)Ne(\(p, \gamma\))\(^{23}\)Ne(\(p, \alpha\))\(^{20}\)Ne. Owing to the proton excess, these reaction cycles lead to a rapid disappearance of neutrons but a residual abundance of free protons (compare the proton and neutron mass fractions at \( T_0 = 4.39 \) and 3.8 in Table 2). The neutrons helped hold the nuclear abundances in a high-mass, nonequilibrium distribution. When they begin to disappear, the heavy nuclei begin to disintegrate toward the favored iron-group nuclei. The \( T_0 = 3.93 \) and 3.80 panels of Figure 1 show the abundance distribution during disintegration. Notice the abundance peaks corresponding to closed nuclear shells at \( Z = 50 \) and \( N = 50 \). The disintegration ends on the \( N = 50 \) closed shell, predominantly at \(^{92}\)Mo.
the temperature continues to decline, proton captures deplete $^{92}$Mo and create a large abundance of $^{96}$Ru and $^{96}$Pd. These isotopes subsequently decay to $^{94}$Mo and $^{94}$Ru after reaction freezeout.

It is essential to note the large final abundances of protons and $^4$He. These species, and even a residual quantity of neutrons, are abundant throughout the disintegration epoch. This high abundance of neutrons and $^4$He nuclei hinders the disintegration flow, despite the high temperatures at the time the disintegrations are occurring ($T_\nu \approx 4$). As mentioned, the high abundance of protons also leads to proton-capture reactions along the $N = 50$ closed shell that modify the abundances at late times.

Table 1 presents the top 10 production factors from six calculations of nucleosynthesis in fast, high-entropy, proton-rich expansions. Calculation B is the model discussed above ($s = 145k_B$, $Y_e = 0.510$, and $\tau = 0.003$ s). The most produced isotopes are light $p$-nuclei, which are made as $N = 50$ progenitors. At the top of the list are $^{96}$Ru (made as $^{96}$Pd) and $^{94}$Mo (made as $^{94}$Ru) at production factors of $\sim 10^6$. Also significantly produced are $^{92}$Mo (made as $^{92}$Rh) and $^{94}$Mo, which, although depleted late in the expansion by proton captures, is ultimately produced off of the $N = 50$ shell as $^{94}$Ru.

Calculations A and F show results for expansions with lower entropy ($s = 140k_B$) or lower $Y_e$ ($Y_e = 0.505$) than in our reference calculation, B. In both cases, the disintegration flow began earlier in the expansion than in calculation B and allowed the nuclei to return to the iron-group isotopes before freezeout. On the other hand, the larger entropy in calculation C delayed the onset of disintegration and thereby strongly produced higher mass $p$-process isotopes. In calculation D, still heavier proton-rich isotopes formed. Rare isotopes such as $^{160}$Ta are particularly strongly produced in this expansion. Finally, in calculation E, the entropy was sufficiently high that the neutrons disappeared only at low temperature, so the disintegration flow never occurred. This proton-rich expansion was able to produce heavy, neutron-rich isotopes whose production is normally associated with rapid neutron-capture nucleosynthesis (see Meyer 2002). These results show the remarkable sensitivity of the nucleosynthesis to small changes in the expansion parameters.

Because of this sensitivity and the nonequilibrium aspect of the nuclear flows, we expect that the nucleosynthetic yields of such proton-rich expansions will be sensitive to nuclear reaction rates on a large suite of isotopes. Table 3 shows results for calculations identical to calculation C, but with all charged-particle and electromagnetic (EM) rates on nuclei with $Z \geq 27$ increased or decreased by a factor of 2. In the former case, the larger reaction cross sections enhance the assembly of nucleons into $^4$He and cause earlier and more efficient disintegration of the heavy isotopes into iron-group nuclei. In the latter case, the assembly of nucleons and hence the disintegration flows are less efficient, and a heavier distribution of nuclei results than in calculation C. Reaction surveys will be needed to further clarify this issue.

3. IMPLICATIONS

The precise mechanism for the production of the light $p$-process nuclei, $^{92,94}$Mo and $^{96,98}$Ru, remains mysterious. The "$\gamma$-process" in core-collapse supernovae (Woosley & Howard 1978) accounts for the heavy $p$-process nuclei but underproduces the light $p$-isotopes (for a review, see Arnould & Goriely 2003). This suggests some other process than the $\gamma$-process may be responsible for their origin (although questions remain about the need for an exotic site; Costa et al. 2000). Other suggested sites or processes are thermonuclear supernovae (Howard et al. 1991), $\alpha$-rich freezeouts in core-collapse supernovae (Fuller & Meyer 1995; Hoffman et al. 1996), or the rp-process (e.g., Schatz et al. 2001). Each of these processes or sites has its own difficulties in either producing the right isotopes or actually ejecting them into interstellar space, and the question of the site of origin of the light $p$-process nuclei remains. The fairly large production factors we find for certain expansions suggest that neutrino-heated ejecta from core-collapse supernovae may account for a significant fraction of the solar system’s supply of these isotopes, if the winds achieve high enough entropies or fast enough ex-
pansions. Alternatively, it may be that some other as yet not fully characterized astrophysical setting, such as a gamma-ray burst fireball or wind from a collapsar (MacFadyen & Woosley 1999) or neutron star merger (e.g., Eichler et al. 1989), may achieve the necessary conditions.

Finally, we note that comparison of observations of abundances of Sr, Y, and Zr in low-metallicity stars and models of the chemical evolution of the Galaxy hint at a "lighter element primary process" (LEPP) that may contribute to the synthesis of $A \approx 90$ nuclei (Travaglio et al. 2004). The process we describe is primary and could therefore be an interesting component of this inferred LEPP.

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### TABLE 2

| t (ms) | $T_1$ | $X_1$ | $X_2$ | $X_3$ | $Y_1^*$ |
|-------|-------|-------|-------|-------|--------|
| 0.16  | ...   | 9.50  | 4.69  | $10^{-3}$ | $4.89 \times 10^{-1}$ | $4.24 \times 10^{-2}$ | $3.60 \times 10^{-1}$ |
| 0.94  | ...   | 7.48  | 1.36  | $10^{-2}$ | $1.56 \times 10^{-1}$ | $7.08 \times 10^{-1}$ | $8.38 \times 10^{-1}$ |
| 1.64  | ...   | 6.00  | 4.18  | $10^{-2}$ | $6.18 \times 10^{-2}$ | $8.96 \times 10^{-1}$ | $1.13 \times 10^{-1}$ |
| 2.22  | ...   | 5.00  | 1.44  | $10^{-2}$ | $3.45 \times 10^{-2}$ | $9.51 \times 10^{-1}$ | $1.55 \times 10^{-1}$ |
| 2.64  | ...   | 4.39  | 3.45  | $10^{-2}$ | $2.37 \times 10^{-2}$ | $9.71 \times 10^{-1}$ | $1.04 \times 10^{-1}$ |
| 3.00  | ...   | 3.93  | 5.83  | $10^{-2}$ | $2.03 \times 10^{-2}$ | $9.77 \times 10^{-1}$ | $2.01 \times 10^{-1}$ |
| 3.11  | ...   | 3.80  | 3.68  | $10^{-2}$ | $2.02 \times 10^{-2}$ | $9.78 \times 10^{-1}$ | $2.20 \times 10^{-1}$ |
| 3.90  | ...   | 3.00  | 3.21  | $10^{-3}$ | $2.01 \times 10^{-2}$ | $9.78 \times 10^{-1}$ | $2.77 \times 10^{-5}$ |
| 6.56  | ...   | 1.43  | 2.42  | $10^{-2}$ | $2.01 \times 10^{-2}$ | $9.78 \times 10^{-1}$ | $2.90 \times 10^{-5}$ |
| 14.11 | ...   | 0.12  | ...   | ...   | $2.01 \times 10^{-2}$ | $9.78 \times 10^{-1}$ | $2.90 \times 10^{-5}$ |

$a$ $Y_1$ is the number of heavy nuclei (i.e., nuclei with charge greater than or equal to 6) per nucleon.

### TABLE 3

| Rank | $A^\prime$ | $O$ | $A^\prime$ | $O$ |
|------|------------|-----|------------|-----|
| 1    | $^{45}$Sc  | $4.19 \times 10^3$ | $^{133}$Xe | $1.66 \times 10^7$ |
| 2    | $^{60}$Ni  | $2.96 \times 10^3$ | $^{133}$Xe | $7.62 \times 10^7$ |
| 3    | $^{44}$Ti  | $9.22 \times 10^3$ | $^{133}$Xe | $4.29 \times 10^7$ |
| 4    | $^{46}$Ti  | $1.38 \times 10^3$ | $^{133}$Xe | $3.53 \times 10^10$ |
| 5    | $^{46}$Co  | $1.27 \times 10^1$ | $^{133}$Xe | $3.53 \times 10^10$ |
| 6    | $^{46}$Ti  | $1.20 \times 10^1$ | $^{133}$Xe | $2.44 \times 10^7$ |
| 7    | $^{46}$Ca  | $9.89 \times 10^6$ | $^{133}$Xe | $1.96 \times 10^7$ |
| 8    | $^{47}$Ca  | $7.42 \times 10^6$ | $^{133}$Xe | $1.15 \times 10^7$ |
| 9    | $^{46}$Ti  | $5.41 \times 10^6$ | $^{133}$Xe | $9.30 \times 10^7$ |
| 10   | $^{50}$Cr  | $4.79 \times 10^6$ | $^{133}$Xe | $9.07 \times 10^7$ |

$a$ Calculation A: $\tau = 0.003$, $s = 150$, initial $Y_1 = 0.510$; (charged-particle and EM rates for $Z \geq 27$) $\times 2$.

$b$ Calculation B: $\tau = 0.003$, $s = 150$, initial $Y_1 = 0.510$; (charged-particle and EM rates for $Z \geq 27$) $\times 0.5$. 