Enormous explosion energy of Type IIP SN 2017gmr with bipolar $^{56}\text{Ni}$ ejecta

Victor P. Utrobin$^{1,2,3}$, * Nikolai N. Chugai$^2$, Jennifer E. Andrews$^4$, Nathan Smith$^4$, Jacob Jencson$^4$, D. Andrew Howell$^{5,6}$, Jamison Burke$^{5,6}$, Daichi Hiramatsu$^{5,6}$, Curtis McCully$^{5,6}$, and K. Azalee Bostroem$^7$

$^1$NRC ‘Kurchatov Institute’ – Institute for Theoretical and Experimental Physics, B. Cheremushkinskaya St. 25, 117218 Moscow, Russia
$^2$Institute of Astronomy, Russian Academy of Sciences, Pyatnitskaya St. 48, 119017 Moscow, Russia
$^3$Max-Planck-Institut für Astrophysik, Karl-Schwarzschild-Str. 1, 85748 Garching, Germany
$^4$Steward Observatory, University of Arizona, 933 North Cherry Avenue, Tucson, AZ 85721-0065, USA
$^5$Department of Physics, University of California, Santa Barbara, CA 93106-9530, USA
$^6$Las Cumbres Observatory, 6740 Cortona Dr, Suite 102, Goleta, CA 93117-5575, USA
$^7$Department of Physics and Astronomy, University of California, 1 Shields Avenue, Davis, CA 95616-5270, USA

Accepted 2021 May 9. Received 2021 April 19; in original form 2021 March 1

ABSTRACT

The unusual Type IIP SN 2017gmr is revisited in order to pinpoint the origin of its anomalous features, including the peculiar light curve after about 100 days. The hydrodynamic modelling suggests the enormous explosion energy of $\approx 10^{52}$ erg. We find that the light curve with the prolonged plateau/tail transition can be reproduced either in the model with a high hydrogen abundance in the inner ejecta and a large amount of radioactive $^{56}\text{Ni}$, or in the model with an additional central energy source associated with the fallback/magnetar interaction in the propeller regime. The asymmetry of the late Hz emission and the reported linear polarization are reproduced by the model of the bipolar $^{56}\text{Ni}$ ejecta. The similar bipolar structure of the oxygen distribution is responsible for the two-horn structure of the [O i] $\lambda\lambda 6360, 6364$ Å emission. The bipolar $^{56}\text{Ni}$ structure along with the high explosion energy are indicative of the magneto-rotational explosion. We identify narrow high-velocity absorption features in Hz and He i $10830$ Å lines with their origin in the fragmented cold dense shell formed due to the outer ejecta deceleration in a confined circumstellar shell.

Key words: hydrodynamics – methods: numerical – supernovae: general – supernovae: individual: SN 2017gmr

1 INTRODUCTION

Type IIP supernovae (SNe IIP) originate from a core collapse of massive stars ($>9 M_\odot$) that retain a significant fraction of the hydrogen envelope until the explosion. The general paradigm is that the SN IIP light curve at the plateau stage is maintained by the release of the internal energy deposited during the propagation of the shock wave through the presupernova (pre-SN) envelope (Grassberg et al. 1971), whereas the luminosity tail is powered by the radioactive decay $^{56}\text{Co} \rightarrow ^{56}\text{Fe}$ (Weaver & Woosley 1980). The debatable explosion mechanisms of SNe IIP or, in a broad sense, of core-collapse supernovae include the neutrino-driven explosion (Colgate & White 1966; Janka 2017; Burrows & Vartanyan 2021), the magneto-rotational explosion related to the magnetar formation (LeBlanc & Wilson 1970; Bisnovatyi-Kogan 1971; Khokhlov et al. 1999), and the jet-powered supernovae related to the collapsar (rotating black hole plus disk) formation (MacFadyen et al. 2001).

In most observed cases one cannot distinguish between different options, because the outcome of the SN explosion with the typical energy of $\sim 10^{51}$ erg is not sensitive to the explosion mechanism. An exception is the case when the explosion energy inferred from the hydrodynamic modelling significantly exceeds the upper limit for the neutrino-driven mechanism $E_{\nu\nu} \approx 2 \times 10^{51}$ erg (Janka 2017). Among eleven SNe IIP explored so far via the uniform hydrodynamic modelling (cf. Utrobin & Chugai 2019) only two events can be attributed to the category of high-energy SNe IIP with $E > E_{\nu\nu}$: SN 2009kf with $E = 2.15 \times 10^{52}$ erg (Utrobin et al. 2010) and SN 2000cb with $E = 4.4 \times 10^{51}$ erg (Utrobin & Chugai 2011).

The recent Type IIP SN 2017gmr is another candidate for high-energy SNe IIP, because it exhibits both high luminosity and high expansion velocities (Andrews et al. 2019). The hydrodynamic modelling of SN 2017gmr (Goldberg & Bildsten 2020) prefers a high explosion energy of $\approx 5 \times 10^{51}$ erg despite of the claimed model degeneracy with respect to SN parameters. It is noteworthy that SN 2017gmr shows signatures of asymmetry in the Hz and [O i] $\lambda\lambda 6360, 6364$ Å emission lines on day 312 (Andrews et al. 2019) and in the polarization (Nagao et al. 2019), which indicate a non-spherical explosion.

Even more remarkable feature of SN 2017gmr is a steeper luminosity decline of the post-plateau tail compared to the $^{56}\text{Co}$ decay luminosity (Andrews et al. 2019). The hydrodynamic modelling of
Goldberg & Bildsten (2020) is unable to account for this behavior that is interpreted as an excess in the observed luminosity over the computed luminosity at the radioactive tail. Possible explanations for the fast tail decline include the early escape of gamma-quanta due to the decay of the radioactive $^{56}\text{Co}$, the circumstellar (CS) interaction, or the early dust formation (Andrews et al. 2019), although the latter seems unlikely because of the high gas temperature at this stage. The CS interaction might be relevant, since early spectra show the narrow Hz emission with broad wings indicating the dense CS shell. Another possibility might be the high velocity of radioactive $^{56}\text{Ni}$, favoring the early escape of the gamma rays similar to the case of SN 2013ej (Utrobin & Chugai 2017). The applicability of this conjecture for SN 2017gmr at the moment is unclear and must be examined.

Here we address unusual features of SN 2017gmr with an emphasis on the fast decline of the luminosity tail. To this end we revisit the construction of the bolometric light curve (Section 2) and then perform the hydrodynamic modelling and examine a number of possibilities for the origin of the fast tail decline (Section 3). The asymmetry effects in emission lines at the nebular stage and the polarization are explored in order to constrain the extent of mixing of radioactive $^{56}\text{Ni}$ (Section 4). We then address the effects of the CS interaction at the photospheric stage (Section 5). Finally, results are summarized and discussed in Section 6.

Below we adopt the distance of 19.6 Mpc and the reddening of $E(B-V) = 0.3$ mag (Andrews et al. 2019). The explosion date is set to be 2017 September 1.88, MJD 57997.89, which is recovered from the fit of the earliest $r$ magnitudes by the hydrodynamic modelling. This moment is 1.20 day earlier compared to that adopted by Andrews et al. (2019).

## 2 BOLOMETRIC LIGHT CURVE AND PHOTOSPHERIC VELOCITIES

Additional late-time photometry presented here and not published in Andrews et al. (2019) was obtained via the Las Cumbres Observatory 1-m telescope network (Brown et al. 2013) with the Sinistro cameras in the framework of the Global Supernova Project. The data were reduced using the Beautiful Algorithms to Normalize Zillions of Astronomical Images (BANZAI) pipeline (McCully et al. 2018). PSF-fitting photometry was then performed using lcofgsnpipe (Valenti et al. 2016), a PyRAF-based photometric reduction pipeline. $UBV$-band data were calibrated to Vega magnitudes (Stetson 2000) using standard fields observed on the same night with the same telescope, while $gri$-band data were calibrated to AB magnitudes using the Sloan Digital Sky Survey (SDSS Collaboration; Albareti et al. 2017).

As described in Andrews et al. (2019), a quasi-bolometric light curve was created using the routine superbol (Nicholl 2018), where the reddening and redshift corrected photometry in each band (from UV to IR) was interpolated with the $g$-band as reference, then converted to a spectral luminosity ($L_{\lambda}$). The bolometric luminosity was then computed from the integration of the SED for each epoch. While uncertainties in the reddening ($E(B-V) = 0.30 \pm 0.06$ mag) and distance ($19.6 \pm 1.4$ Mpc) introduce uncertainties in the total bolometric luminosity, the bolometric light curve falls much more rapidly than expected for a fully-trapped $^{56}\text{Co}$ decay.

To get an idea of the late-time ($t > 170 $days) bolometric luminosity, we convert $V$ magnitudes into bolometric fluxes using a relation between $V$ and bolometric magnitudes for SN 1987A (Catchpole et al. 1988; Whitelock et al. 1988). This procedure suggests that spectral energy distributions of SN 2017gmr and SN 1987A are similar and differ only by a constant factor that can be fixed by matching bolometric luminosities of SN 2017gmr around day 170 (Fig. 1).

The velocities at the photosphere (Fig. 2) are estimated from absorption minima of H$\beta$, Fe $\beta$ 5169 Å, He $\alpha$ 5876 Å, and Na $\alpha$ 5892 Å doublet in spectra of SN 2017gmr (Andrews et al. 2019). In the case of Na $\alpha$ we treat photon scatterings in the doublet by the Monte Carlo technique to recover the photospheric velocity from the absorption minimum. According to the Fe $\beta$ 5169 Å absorption the photospheric velocity is lower by $\sim 500 \text{ km s}^{-1}$ compared to the velocities from H$\beta$ and Na $\alpha$ 5892 Å doublet (Fig. 2). The possible reason for that might be the presence of Mg $\delta$ 5167, 5173, 5184 Å triplet. The optical depth of the strongest triplet line Mg $\delta$ 5184 Å is comparable to that of Fe $\beta$ 5169 Å for the ionization fraction Mg $\delta$/Mg $\beta \sim 0.025$. We performed Monte Carlo simulations of the radiation transfer for the blend of Fe $\beta$ and Mg $\delta$ triplet assuming that the Sobolev optical depth of the strongest Mg $\delta$ line 5184 Å line is equal to that of the Fe $\beta$ 5169 Å line. The absorption minimum in this case turns out shifted towards red by about $500 \text{ km s}^{-1}$, which could explain the lower velocity according to the Fe $\beta$ 5169 Å line.

Noteworthy, the radial velocity of an absorption minimum ($|v_{\text{min}}|$)
of the P Cygni profile may differ from the photospheric velocity \(v_{\text{ph}}\). The equality \(|v_{\text{min}}| = v_{\text{ph}}\) takes place only for the scattering lines of a moderate strength. For a strong line even with the conservative scattering the minimum is displaced towards blue due to the scattered emission, so in this case \(|v_{\text{min}}| > v_{\text{ph}}\). This effect is strengthened by a net emission likewise in the case of H\(_2\). For a weak absorption, on the contrary, \(|v_{\text{min}}| < v_{\text{ph}}\), because in this case the absorption is formed by a narrow layer close to the photosphere and projection effects shift the absorption minimum towards zero radial velocity. We therefore admit that the velocity recovered from the shallow H\(\beta\) and He\(\iota\) absorptions at the early stage \((t < 20 \text{ days})\) likely underestimate the photospheric velocity by about 20\%. At the later stage \((t > 20 \text{ days})\) the H\(\beta\) absorption is moderately strong and the net and scattered emissions are suppressed because of the H\(\beta\) quanta conversion into P\(\tau\) and H\(\alpha\). We thus believe that the H\(\beta\) absorption provides a reliable measure of the photospheric velocity at this stage.

3 HYDRODYNAMIC MODELLING

3.1 Model overview

To hydrodynamically model the Type IIP SN 2017gmr, we use the time-implicit Lagrangian radiation hydrodynamics code \(\text{Cradle}\) which integrates the set of spherically symmetric hydrodynamic equations including self-gravity, and a radiation transfer equation in the gray approximation (Utrobin 2004, 2007). The explosion of the pre-SN model is initiated by a supersonic piston applied to the stellar envelope at the boundary with the collapsing 1.6\(M_\odot\) core. The pre-SN model is the hydrostatic non-evolutionary model of a red super-giant (RSG) star. It should be emphasized that the choice of the non-evolutionary model is motivated by the following arguments. First, the description of the light curve and photospheric velocities of SNe IIP by the spherically symmetric hydrodynamics cannot be attained based on the current evolutionary pre-SN models (Utrobin & Chugai 2008): the modification of the density distribution of the hydrogen envelope and the smoothing density and composition gradients at the metals/He and He/H composition interfaces is needed. Second, the RSG explosion in three-dimensional (3D) hydrodynamic simulations indeed results in smoothing the density and composition gradients at the metals/He and He/H composition interfaces (Utrobin et al. 2017). Third, at the final stage of the pre-SN evolution the density distribution in the hydrogen envelope can be modified by the acoustic waves excited by the vigorous convection at the Ne-burning stage (Shiode & Quataert 2014). These arguments unavoidably compel us to choose a non-evolutionary RSG model in which the structure modification can be implemented by hand to make the spherically symmetric hydrodynamic model appropriate for the description of the observational data (Figs. 3 and 4). It is noteworthy that the hydrogen abundance in the inner layers of the mixed ejecta \(X_\text{e} \approx 0.05\) is comparable to the value of about 0.03 produced by the 3D hydrodynamic model of SN 1999em (Utrobin et al. 2017). However, mixing induced by the SN explosion can be affected by particular features of the explosion mechanism, e.g., explosion asymmetry, so the higher hydrogen abundance in the central zone of the expanding ejecta cannot be ruled out.

We do not solve the complicated optimization problem to construct the optimal pre-SN model, instead we rely on the well-studied effects of model parameters on the light curve and the photospheric velocity (e.g., Grassberg et al. 1971; Woosley 1988; Utrobin 2007). The optimal pre-SN model that reproduces the major observational data of SN 2017gmr (i.e., the light curve and the evolution of the photospheric velocity) is found by means of hydrodynamic simulations for an extended set of SN parameters and a several options for the energy source at the plateau/tail transition, the distribution of radioactive \(^{56}\text{Ni}\) being fixed by the nebular spectra.

Along with the pre-SN radius, the ejecta mass, the explosion energy, and the total amount of radioactive \(^{56}\text{Ni}\), the extent of its mixing in velocity space affects the light curve at the plateau as well. For SN 2017gmr the extent of \(^{56}\text{Ni}\) mixing is constrained by the H\(\alpha\) emission on day 312 (Section 4). We find that the outer boundary of \(^{56}\text{Ni}\) ejecta should lie at about 3300 km\(\text{s}^{-1}\) in the freely ex-
Table 1. Basic properties of hydrodynamic models$^a$.

| Model         | $M_{ej}$ | $E$  | $M_{Ni}$ | $V_{Ni}^{max}$ | $X_c$ | Note          |
|---------------|----------|------|----------|----------------|-------|---------------|
|               | $(M_\odot)$ | $(10^{51}$ erg) | $(M_\odot)$ | $(km/s)$      |       |               |
| HM-refm       | 22.0     | 10.2 | 0.16     | 3300          | 0.05  | reference     |
| HM-hmix       | 22.0     | 10.2 | 0.18     | 3300          | 0.20  | H mixing      |
| HM-Imas       | 14.0     | 10.0 | 0.20     | 3400          | 0.04  | low $M_{ej}$  |
| HM-enni       | 22.0     | 10.2 | 0.23     | 7900          | 0.05  | outer $^{56}$Ni |
| HM-magn       | 22.0     | 10.2 | 0.01     | 3300          | 0.05  | magnetar      |
| HM-optm       | 22.0     | 10.2 | 0.11     | 3300          | 0.05  | FM-mechanism  |

$^a$ In all models the pre-SN radius $R_0$ is 525 $R_\odot$.

3.2 Plateau/tail transition and luminosity tail

The hydrodynamic modelling leads us to the reference model HM-refm (Table 1) that fits in general the initial luminosity peak, the plateau, the late-time radioactive tail, and the photospheric velocities of SN 2017gmr (Fig. 6). However, this model demonstrates the essential deficit in the luminosity at the plateau/tail transition in the range of 90 – 140 days. The impression is that the plateau/tail transition of SN 2017gmr is overlong compared to ordinary SNe IIP. Note that the recent hydrodynamic modelling of SN 2017gmr by Goldberg & Bildsten (2020) faces the same problem of the model drawback at the plateau/tail transition.

Increasing both the total amount of radioactive $^{56}$Ni and the opacity of the inner layers of the ejecta might compensate for the luminosity deficit. The higher opacity could be related to a more intense inward mixing of hydrogen-rich matter as a result of a strong explosion asymmetry compared to ordinary SNe IIP. The study of this possibility results in model HM-hmix (Table 1) whose $^{56}$Ni mass of 0.182 $M_\odot$ and hydrogen abundance in the central zone $X_c = 0.2$ are larger than 0.158 $M_\odot$ and $X_c = 0.05$ of the reference model, respectively. The excellent fit to the observations in the range $t \leq 130$ days lends credibility to this model, although the later steep tail decline is not fully reproduced (Fig. 7).

Among other possibilities to resolve the issue of the steep luminosity tail one can admit the low ejecta mass that favors a more efficient escape of gamma rays from the $^{56}$Co decay. The appropriate low-mass model HM-Imas (Table 1) that fits the tail requires the ejecta mass of 14 $M_\odot$ with the $^{56}$Ni mass of 0.2 $M_\odot$ (Fig. 8, blue line). This model expectedly produces too short plateau and should be rejected.

Another possibility is prompted by the model with the external $^{56}$Ni ejecta that was successfully applied to SN 2013ej (Utrobin & Chugai 2017). The similar hydrodynamic model HM-exni (Table 1) for SN 2017gmr with the external $^{56}$Ni fits the radioactive tail (Fig. 8, red line), if all the 0.23 $M_\odot$ of $^{56}$Ni resides in the velocity range of $6000 - 8000$ km s$^{-1}$. The drawback of this model is a pronounce luminosity excess at the plateau in the interval of 30–80 days and the luminosity deficit at the plateau/tail transition. The external $^{56}$Ni thus cannot resolve the problems of the anomalous light curve of SN 2017gmr.

One can consider somewhat exotic model HM-magn in which the magnetar, not radiative $^{56}$Ni, determines the luminosity tail. With a standard magnetar luminosity evolution $L = L_0/(1 + t/t_0)^2$ the tail can be described with the following parameter values: $L_0 = 4.47 \times 10^{42}$ erg s$^{-1}$ and $t_0 = 19.07$ days. The rest of model parameters are in Table 1. The apparent drawback of the magnetar model is a large luminosity excess at the end of the plateau (Fig. 8, green line). The magnetar model in a simple version thus cannot resolve simultaneously the issue of the plateau/tail transition and of the steep luminosity tail.

We briefly address a possible role of the CS interaction. Although this mechanism potentially is able to provide the steep luminosity tail; however, in this case it is highly unlikely. The point is that the interaction power of $\sim 10^{42}$ erg s$^{-1}$ required at about 100 days is released in the outer layers of the ejecta that would be inevitably accompanied by a strong broad H$\alpha$ emission with the luminosity of $\sim 10^{44}$ erg s$^{-1}$ and...
a specific line profile lacking the absorption component. This is not the case for SN 2017gmr (cf. Andrews et al. 2019), so the significant contribution of the spherical CS interaction is ruled out. To overcome this problem, one might admit that the CS matter is arranged in the form of a dense equatorial ring that is overtaken by the SN envelope (e.g., Chugai & Danziger 1994; Smith et al. 2015). In that case one expects that the forward and reverse shocks are submerged inside the envelope, so the released radiation creates a quasi-spherical photosphere. This scenario, however, should reveal strong effects of the non-spherical interaction in the early Hz profile. Instead, the spectra show the usual steady evolution — a characteristic of spherical SNe IIP. The CS interaction as a major energy source at the early luminosity tail thus should be rejected.

3.3 Central energy source?

An interesting possibility to resolve the issue of the post-plateau light curve involves an additional energy source related to the fallback on the magnetar. The conjecture looks quite sensible by two reasons: first, the fallback accretion flow always accompanies the core-collapse supernovae (Colgate 1971; Chevalier 1989), and second, the magnetar formation could be a natural outcome of the magneto-rotational explosion mechanism that is required to provide the enormous explosion energy of \( \sim 10^{52} \) erg. The extra energy source involving both the fallback and the magnetar we dub “FM-mechanism”, for short. The dense environment associated with the fallback inhibits the magnetic dipole radiation of the magnetar, whereas the FM-mechanism could release rotational energy in the propeller regime that operates when the Alfvén radius \( R_{\text{m}} \) exceeds the corotation radius \( R_c \) and both of these radii are less than the radius of the light cylinder (Illarionov & Sunyaev 1975; Shakura 1975). Conditions required for the FM-mechanism to operate successfully can be illustrated by the case of the unusual supernova ASASSN-15nx with the luminosity decreasing in the range of \( 5 \times 10^{42} \) erg s\(^{-1} \) to \( 5 \times 10^{41} \) erg s\(^{-1} \) between days 100 and 200 (Chugai 2019). In that case the light curve has been modelled in terms of the FM-mechanism assuming a spherical fallback with the accretion rate \( M \sim 10^{-3} M_\odot \) yr\(^{-1} \) onto the magnetosphere of the neutron star with the magnetic moment \( \mu \sim 5 \times 10^{31} \) G cm\(^3\), and the rotation period \( p \sim 10^{-2} \) s. These values could be applicable for the FM-mechanism in the case of SN 2017gmr, although some complications could arise due to the unknown specific angular momentum of the fallback material.

We will constrain the rate of the power release by the FM-mechanism, \( L_C(t) \), using the hydrodynamic modelling. To this end we impose the power \( L_C(t) \) at the inner boundary of the computational domain immediately after the explosion. The model HM-optm (Table 1) with the \( ^{56}\text{Ni} \) mass of 0.11 \( M_\odot \) reproduces both the observed plateau/tail transition and the luminosity tail decline (Fig. 9) for the adopted luminosity \( L_C(t) \) shown in the same plot. The required evolution of the additional energy source is characterized by a slow rise toward the maximum at about 80 days and an exponential decline later on. In fact, the exponential behavior is constrained only to the stage of \( t < 180 \) days; we retain the later exponential behavior simply to minimize a number of parameters, although the preferred option is zero contribution of the extra energy source to the luminosity on day 312.

At first glance the adopted luminosity evolution of the FM-mechanism looks highly artificial. In fact, it is not, because the fallback accretion rate at the early stage is high enough to shrink the magnetosphere so strongly that the inequality \( R_{\text{m}} < R_c \) is fulfilled thus turning off the propeller regime. In this case the gravitational energy of the fallback accretion flow is released in the close vicinity of neutron star producing the neutrino luminosity (Chevalier 1989).

It should be emphasized that the conjecture about the extra energy source is viable so long as the conclusion on the fast decline of the observed luminosity tail remains valid.

3.4 Constraining the ejecta mass

The adequate hydrodynamic model is constructed by fitting the photometric and spectroscopic observations of the object under study. The important physical parameters — the initial radius \( R_0 \), the ejecta mass \( M_{ej} \), and the explosion energy \( E \) — can be reliably estimated from the detailed observations of the whole outburst, particularly, at both the shock breakout and the plateau/tail transition. Photometric data of SN 2017gmr are well defined at both epochs and are sufficiently comprehensive to construct the bolometric light curve (Andrews et al. 2019). The rising part of the \( r \)-band light curve after the shock breakout is strengthened by the upper limit in \( r \)-band.

Figure 8. The bolometric light curves of three models (Table 1) that are able to reproduce the observed luminosity tail: model HM-Imas with the low ejecta mass (blue line), model HM-exni with the external \( ^{56}\text{Ni} \) (red line), and model HM-magn with the standard magnetar luminosity (green line).

Figure 9. The bolometric light curve of model HM-optm (blue line, Table 1) overlaid on the observational data (black and magenta circles). A good fit to the observations is obtained with the additional energy source applied to the internal boundary of the ejecta (green line) and the energy deposition of gamma rays by radioactive decay of \( ^{56}\text{Ni} \) (solid red line). Dotted red line is the total power of the radioactive decay. Inset shows the evolution of the photospheric velocity in the same way as in Fig. 6.
obtained two days before the SN discovery (Andrews et al. 2019) in constraining the basic parameters of hydrodynamic model.

Model HM-optm excellently fits the early $M_r$ magnitudes and satisfies the upper limit provided the explosion occurs at MJD 57997.89 (Fig. 10). It is noteworthy that during the first several days the SN spectrum is essentially the Planck function that permits us to adequately calculate the r-band magnitude with the radiation hydrodynamics code CRAB. Interestingly, the flux in r-band shows a double peak structure that is to our knowledge has never been noticed either observationally or numerically. In this particular model the double peak phenomenon is related to the formation of the opaque thin dense shell at about 0.38 days after the collapse and ~1 hour after the shock breakout due to the sweeping of the external layers into the thin shell by the pressure of the escaping radiation flux. To illustrate the point we show the evolution of related values, viz., the absolute bolometric magnitude $M_{bol}$, the absolute r-band magnitude $M_r$, the effective temperature $T_{eff}$, the photospheric velocity $v_{ph}$, and the photospheric radius $R_{ph}$ (Fig. 11). It is noteworthy that the r-band flux shows a jump at ~1 hour after the shock breakout, whereas the bolometric flux does not. This suggests that the variation of the r-band flux is essentially related to that in the color temperature (Fig. 11a). The formation of the thin opaque shell is accompanied by its radiative acceleration clearly seen in the behavior of the velocity at the photosphere (Fig. 11b). The double peak structure in the r (in other bands as well) light curve is an interesting phenomenon that could be used in future as an additional tool in order to constrain SN IIP parameters.

The plateau length $t_p$ of hydrodynamic models of SNe IIP is known to be almost independent of the ejecta mass given the constant ratio $E/M_{ej}$. For example, in the case of the normal IIP SN 1999em it shows only weak dependence on the explosion energy $t_p \propto E^{-0.18}$ and the pre-SN radius $t_p \propto R_{0.10}$ (Urobin 2007). Such a behavior of hydrodynamic models produces the impression of a degeneracy with respect to the model parameters. However, in the rigorous sense the degeneracy on the ejecta mass is absent. We demonstrate this fact for SN 1999em with auxiliary hydrodynamic models in which $E/M_{ej} = \text{const}$ and the calculated bolometric luminosity is secured at the observed plateau luminosity by the appropriate choice of the pre-SN radius. These models with the different ejecta masses deviated from the value of 19 $M_\odot$ by $\pm 2 M_\odot$ show significant difference at the plateau/tail transition (Fig. 12). This numerical experiment suggests that the distributions of the major elements and $^{56}$Ni are fixed in the ejecta (Figs. 4 and 5). Any variation of these distributions results in the deformation of the model light curve at the end of the plateau and the plateau/tail transition, which is inconsistent with observations. It is noteworthy that the uncertainty of the ejecta mass becomes significant (>10%), if the date of the explosion is fixed with an error worse than ~2 days.

In the special case of SN 2017gmr with the additional central energy source the dependence on the ejecta mass is weaker, yet even in this case we are able to distinguish between the models with the ejecta masses of 21 $M_\odot$, 22 $M_\odot$, and 24 $M_\odot$ (Fig. 12, inset). It is remarkable that the upper limit in r-band taken two days before the SN discovery imposes serious constraint on the ejecta mass. Its physical meaning can be explained by the following reasonings. Using the approximate formulae relating the physical parameters to the observable properties of hydrodynamic models (Urobin 2007) with $E/M_{ej} = \text{const}$ and the fixed bolometric luminosity at the plateau, we find that $R_0 \propto M_{ej}^{-1.51}$. In other words, the lower the ejecta mass is, the larger the pre-SN radius is, and, consequently, the larger the characteristic expansion time is. The latter results in the longer time interval between the shock breakout and the epoch of the SN discovery at $M_r = -17.18$ mag. This implies that there is a lower limit of the ejecta mass determined by the interval between the time of the upper limit in r-band and the discovery epoch. This interval is exactly realized in the hydrodynamic model with the ejecta mass of 21 $M_\odot$ (Fig. 10, inset). Model HM-optm with the ejecta mass of
22 $M_\odot$ excellently fits both the upper limit in $r$-band and the early $M_\star$ magnitudes (Fig. 10), and the observed bolometric light curve as a whole (Fig. 9).

In addition to the lower limit of the ejecta mass of 21 $M_\odot$, we can estimate the uncertainty in the derived SN parameters by varying the model parameters around the optimal model. The uncertainty of 1.4 Mpc in the SN distance and the uncertainty of about 0.06 mag in the reddening $E(B-V)$ (Andrews et al. 2019) imply nearly 20 per cent uncertainty in the bolometric luminosity. The scatter in the plot of the photospheric velocity versus time (Fig. 9, inset) suggests the uncertainty of about 7 per cent in the photospheric velocity. We estimate the maximal uncertainty of the plateau duration as 4 days, i.e. 4 per cent of the plateau duration. With these uncertainties of observables, we have the initial radius of $525 \pm 176 R_\odot$, the ejecta mass of $22^{+5}_{-3} M_\odot$, the explosion energy of $(10.2 \pm 0.83) \times 10^{51}$ erg, and the total $^{56}$Ni mass of $0.110 \pm 0.023 M_\odot$.

Table 2. Parameters of $^{56}$Ni components.

| Component | $v_\infty$ | $v_r$ | $M_{Ni}$ |
|-----------|------------|-------|----------|
| front     | 2400       | 900   | 0.059    |
| rear      | 2100       | 900   | 0.025    |
| central   | 0          | 700   | 0.026    |


4 ASYMMETRY OF $^{56}$NI EJECTA

The triple-peaked Hα profile in the spectrum on day 312 (Andrews et al. 2019) suggests a non-spherical line-emitting region. Following a concept employed for the Type IIP SN 2004dj (Chugai et al. 2005), the Type IIP SN 2016X (Utrobin & Chugai 2019), and the Type II SN 2010ip (Smith et al. 2012), we attribute the Hα asymmetry to the bipolar $^{56}$Ni ejecta embedded into a spherical envelope. For SN 2017gmr the $^{56}$Ni distribution is represented by three collinear homogeneous spheres: central, front, and rear. The radii, the shifts, and the masses of components are found via the fit of the Hα line profile for the optimal inclination angle. The latter is constrained by relying on the linear polarization $p = 0.37 \pm 0.04\%$ on day 136 (Nagao et al. 2019). The density distribution in the envelope is set analytically as $p(\nu) = p_0(\nu_0/\nu)^{0.5}/[1+(\nu/\nu_0)^{7.5}]$ with $p_0$ and $\nu_0$ specified by the ejecta mass of 22 $M_\odot$ and the explosion energy of $10^{51}$ erg. The adopted density distribution in the outer layers $\rho \propto \nu^{-8}$ is well consistent with the hydrodynamic model HM-optm (Fig. 5).

The energy deposition by gamma-rays is treated in a single flight approximation with the absorption coefficient $k_\gamma = 0.03(1 + X) \text{ cm}^2 \text{ g}^{-1}$ (Kozma & Fransson 1992), where $X$ is the hydrogen mass fraction assumed to be uniform in the ejecta. We adopt $X = 0.5$ to allow for the synthesized helium. Positrons from $e^- \rightarrow \text{capture}$ channel deposit their kinetic energy on-the-spot. The ionization rate by the Compton and secondary electrons is calculated with the energy fractions spent on ionization, excitation, and heating according to Xu et al. (1992). The recombination rate corresponds to the total recombinations on levels $n > 2$ assuming the electron temperature of 5000 K. The adopted recombination regime approximately allows for the ionization from the second level by the recombinasion Balmer continuum and by the two-photon hydrogen continuum. The recombination Hα emissivity corresponds to case C (opaque Balmer lines) which implies that each recombination onto levels $n > 2$ ends up with the Hα quanta emission.

The inclination angle is found using the iterative procedure starting with a certain $^{56}$Ni configuration: $^{56}$Ni $\rightarrow$ Hα $\rightarrow$ polarization $\rightarrow$ $^{56}$Ni etc. The polarization computation (Chugai 2006) is based on the Monte Carlo technique that follows the history of photons created by a central spherical source with a subsequent Thomson multiple scattering in the non-spherical distribution of electrons produced by the bipolar $^{56}$Ni ejecta. The found inclination angle is $\theta = 40^\circ$ (Fig. 13a) with the radii $v_r$, the shifts $v_\infty$, and the masses of $^{56}$Ni components given in Table 2. The bipolar components are not completely symmetrical: rear component has the lower bulk velocity and the lower mass compared to the front component. For the $^{56}$Ni mass of 0.11 $M_\odot$ the model Hα luminosity of $3 \times 10^{50}$ erg s$^{-1}$ on day 312 coincides with the observed value of $3.2 \times 10^{50}$ erg s$^{-1}$ estimated from the flux calibrated spectrum. This choice of the $^{56}$Ni mass is consistent also with the light curve of model HM-optm (Fig. 9).

A reliable modelling of the thermal state of the [O i] doublet-emitting region is beyond reach given the significant role of the cooling by CO and SiO in the oxygen-rich matter (Liu & Dalgarno 1995). To get idea of the asymmetry of the oxygen line-emitting region, we decompose [O i] 6300, 6364 Å doublet using spherical Gaussian components: central, front, and rear assuming the inclination angle of 40°. The components are specified by the normalized emissivity $j = A \exp[-(\nu/\nu_0)^2]$, where $\nu$ is a velocity distance from the center of the component with a certain velocity shift $v_\infty$ (Table 3, Fig. 13b). The bipolar components of the oxygen emissivity in [O i] 6300, 6364 Å doublet are rather similar to those of $^{56}$Ni (Fig. 13b, inset). This could be interpreted in two ways: either oxygen components reflect the distribution of the electron temperature due to the bipolar $^{56}$Ni in a spherical oxygen-rich gas, or the oxygen ejecta have essentially bipolar structure. The absence in the oxygen 6300 Å line...
narrow counterpart related to the central $^{56}$Ni component favors the bipolar oxygen distribution.

A conjecture that the $[\text{O} \, \text{i}]$ doublet originates from the toroidal oxygen distribution might be conceivable. However, the modelling shows that for any parameters the toroidal structure fails to reproduce the observed $[\text{O} \, \text{i}]$ doublet profile. The toroidal oxygen distribution is thus ruled out.

## 5 CIRCUMSTELLAR INTERACTION

The narrow Hα emission revealed by the early spectra of SN 2017gmr (Andrews et al. 2019) on days 1.5 and 2.3 and disappeared 3 days later suggests the presence of a dense confined CS shell similar to that of SN 2013fs (Yaron et al. 2017; Bullivant et al. 2018). The hydrodynamic interaction of the ejecta with the CS shell should result in the additional optical luminosity powered by the forward and reverse shocks. Another consequence of the CS interaction is the formation of a thin cold dense shell (CDS) between two shocks. The latter is observed in some SNe IIP as a narrow high-velocity ($\sim 10^5 \text{km s}^{-1}$) absorption (NHVAs) in the Hα, He i 10830 Å, and possibly Hβ lines at about 50 – 100 days (Chugai et al. 2007). The optical and near infrared spectra of SN 2017gmr indeed reveal distinctive NHVAs in the Hα and He i 10830 Å lines with velocities decreasing in the range $14000 – 12000 \text{km s}^{-1}$ between days 50 and 100. We identify this feature with the CDS partially fragmented due to Rayleigh-Taylor instability (Chugai et al. 2007). The manifestation of the CS interaction as the luminosity excess and the NHVAs can be used to constrain the parameters of the SN ejecta and CS shell.

We treat the CS interaction based on a thin shell approximation (Chevalier 1982). The interaction model was described earlier (Chugai 2001) and here we recap only the essential points. The CDS dynamics is computed using Runge-Kutta 4-th order solver (Press et al. 2007). The shock radiative cooling time $\tau_c$ at a certain moment is determined assuming the electron-ion equilibrium with the postshock density four times greater than the preshock density. This description is not valid at the very early stage, $t < 10$ hours after the shock breakout, when the radiative precursor strongly accelerates the preshock gas thus diminishing the viscous shock. The forward and reverse shock luminosity is approximated as $L_k/(1 + 2/v)$, where $L_k$ is the shock kinetic luminosity. The interaction optical luminosity is equal to the X-ray luminosity absorbed by the unshocked ejecta and the CDS. The density of the homologously expanding ejecta is set as $\rho = \rho_0 (v_0/v)^{0.5} [1 + v/v_0]^{-5.5}$. The density distribution of the confined CS shell is adopted to be uniform in the range of $r < 5 \times 10^{14} \text{cm}$ with a drop $\rho \propto 1/r^2$ in the range of $(1-2) \times 10^{15} \text{cm}$ and the steady wind $\rho \propto 1/r^2$ in the outer zone $r > 2 \times 10^{15} \text{cm}$ (Fig. 14b, inset). The adopted CS shell extent is in line with that for SN 2013fs, $(0.4 \pm 1) \times 10^{15} \text{cm}$ (Yaron et al. 2017).

Our strategy is to find a model that minimizes the CS luminosity and meets the kinematic requirements imposed by the NHVAs. The luminosity and the kinematic properties of the CDS depend on the CS density and the density $\rho (v)$ of the external layers of the SN ejecta. For a given ejecta mass of $22 M_\odot$ the latter is determined by the explosion energy that can be found from the luminosity and the kinematic constraints. We find that the preferred explosion energy is of $6 \times 10^{51} \text{erg}$, whereas the mass of the confined CS shell is of $4 \times 10^{-3} M_\odot$. Remarkably the latter value is comparable to the mass estimate for the confined dense CS shell in SN 2013fs (Yaron et al. 2017). This case produces a moderate interaction luminosity with the maximal contribution of about 18% in the range of 2 – 7 days (Fig. 14a) that compensates a small deficit in the bolometric luminosity of the hydrodynamic model HM-optm at this stage (Fig. 9).

Simultaneously, the CS interaction model meets the kinematic constraints imposed by the early He i 10830 Å absorption and the NHVAs of Hα and He i 10830 Å (Fig. 14b).

Table 3. Parameters of $[\text{O} \, \text{i}]$ doublet components.

| Component | $v_0$ ($\text{km s}^{-1}$) | $b$ | $A$ |
|-----------|-------------------------|----|----|
| front     | 2000                    | 500| 1.0|
| rear      | 2000                    | 600| 0.3|
| central   | 600                     | 600| 0.047|

Figure 13. Asymmetry in the Hα and $[\text{O} \, \text{i}]$ 6300, 6364 Å line profiles taken on day 312 (Andrews et al. 2019, gray lines). Panel (a): the model Hα line is given by red line. Inset shows the computed polarization as a function of inclination angle (blue line) and the observed polarization degree (horizontal magenta line) along with the $\sigma$ lines (horizontal gray lines) (Nagao et al. 2019). Panel (b): thick blue line represents the resultant model $[\text{O} \, \text{i}]$ 6300, 6364 Å line and thin blue line corresponds to the extended spherical emissivity component. Cartoon shows the spherical and bipolar components of the radioactive $^{56}$Ni distribution (red circles) and the bipolar components of the $[\text{O} \, \text{i}]$ emissivity (blue circles). Arrow indicates the direction towards observer.

Figure 14. Observational effects of the CS interaction for the ejecta mass of $22 M_\odot$ and the explosion energy of $6 \times 10^{51} \text{erg}$. Left panel shows the computed bolometric luminosity due to only the CS interaction (red line) compared to the observed bolometric luminosity (gray crosses). Inset shows a zoom in of the first ten days that demonstrates a small contribution of the CS interaction to the SN luminosity at the stage $t < 2$ days. Right panel shows the evolution of the CDS velocity (blue line) and the maximal velocity of the unshocked ejecta (red line). Filled circle corresponds to the lower limit of the maximal ejecta velocity inferred from the He i 10830 Å absorption on day 14.48. Crosses correspond to the NHVA inferred from the Hα profile, while open circles are the NHVA inferred from the He i 10830 Å. Inset shows the corresponding CS density distribution.
The tension between the explosion energy of $6 \times 10^{51}$ erg suggested by the CS interaction model and that of $10^{52}$ erg implied by the hydrodynamic model should not be considered as an irreducible one. Given the simplicity of the CS interaction model, this disparity indicates that the realistic hydrodynamic model should include the shock wave propagation in the CS shell, which however would require a more complex hydrodynamic approach in order to treat an essentially 3D physics related to the CDS formation and the Rayleigh-Taylor instability.

6 DISCUSSION AND CONCLUSIONS

The paper has been aimed at the study of the non-standard Type IIP SN 2017gmr with the emphasis on the unusual light curve at the plateau/tail transition and the steep tail decline. We find that a standard hydrodynamic model of the RSG star exploding with the $^{56}$Ni ejecta is not able to describe the plateau/tail transition and the very tail for any model parameter set. In that sense we confirm the conclusion of Goldberg & Bildsten (2020) who demonstrate that their hydrodynamical model of SNe IIP fails to account for the plateau/tail transition and the luminosity tail of the SN 2017gmr light curve. The model with the high hydrogen abundance in the central zone of the ejecta and the high amount of $^{56}$Ni is able to fit the light curve in the range of $t < 130$ days, although later on the model tail is somewhat less steep compared to the observed one. Yet this kind of hydrodynamic model should be considered as a viable contender for the hydrodynamic model with the additional central energy source.

We find also that the light curve of SN 2017gmr can be reproduced in the framework of the FM-mechanism which implies a central energy source with the specific temporal behavior of the power release. The extra energy source is attributed in this case to the fallback interaction with the magnetar in the propeller regime. This ad hoc scenario is in line with the extremely high explosion energy of $\approx 10^{52}$ erg and the bipolar $^{56}$Ni asymmetry, both indicative of the magneto-rotational explosion. At the moment we are not aware of other SNe IIP that would show similar behavior of the bolometric light curve at the plateau/tail transition and the early stage of the luminosity tail. The another SN IIP that demonstrates a steep decline of the early tail is SN 2013ej (Dhungana et al. 2016). Although this behavior has been explained in the model with the external $^{56}$Ni (Utrobin & Chugai 2017), we do not rule out the FM-mechanism as a viable alternative. One should keep eye open on the possibility that the FM-mechanism could sometimes manifest itself at the radioactive tail of core-collapse SNe including SNe II varieties and SNe Ib/c.

The unusual bolometric light curve of SN 2017gmr raises a serious caveat. The point is that at the plateau/tail transition of SNe IIP the radiative cooling regime changes from the photospheric to the nebular one, which is accompanied by the corresponding spectrum transformation. This poses a question whether the technique for the bolometric flux reconstruction that is appropriate for the photospheric stage preserves the same accuracy at the nebular stage. The only case when we have no doubts in this regard is SN 1987A. In other cases of SNe IIP some degree of doubt remains.

The hydrodynamic modelling confirms the earlier suggestion that the high luminosity and the fast expansion of the ejecta indicate the high explosion energy (Andrews et al. 2019). The explosion energy of $10^{52}$ erg inferred in our model is twice as large compared to $4.6 \times 10^{51}$ erg, the value preferred by Goldberg & Bildsten (2020). Note, however, that the degeneracy of the light curves for SN 2017gmr admits also a model with the larger explosion energy of $\approx 10^{52}$ erg and, consequently, the ejecta mass larger than $22 M_\odot$ (Goldberg & Bildsten 2020). Despite both hydrodynamic codes STELLA (Blinnikov et al. 1998; Blinnikov & Sorokina 2004; Baklanov et al. 2005; Blinnikov et al. 2006) used by Goldberg & Bildsten (2020) and our code CRAB (Utrobin 2004, 2007) produce the similar results for the normal Type IIP SN 1999em (Baklanov et al. 2005; Utrobin 2007), in the case of SN 2017gmr some inconsistency is unavoidable because of the different adopted $^{56}$Ni distribution and the different pre-SN configuration, namely evolutionary and non-evolutionary pre-SN models, respectively.

We have identified the NHVA in the H$\alpha$ and He$\alpha$ 10830 Å lines.

---

Table 4. Hydrodynamic models of Type IIP supernovae.

| SN    | $R_{97}$ | $M_{ej}$ | $E$  | $M_{Ni}$ | $v_{Ni}^{max}$ | $v_{Ni}^{min}$ |
|-------|----------|----------|------|----------|----------------|----------------|
| 1987A | 35       | 18       | 1.5  | 7.65     | 3000           | 600            |
| 1999em| 500      | 19       | 1.3  | 3.6      | 660            | 700            |
| 2000cb| 35       | 22.3     | 4.4  | 8.3      | 8400           | 440            |
| 2003Z | 230      | 14       | 0.245| 0.63     | 535            | 360            |
| 2004et| 1500     | 22.9     | 2.3  | 6.8      | 1000           | 300            |
| 2005cs| 600      | 15.9     | 0.41 | 0.82     | 610            | 300            |
| 2008in| 570      | 13.6     | 0.505| 1.5      | 770            | 490            |
| 2009kF| 2000     | 28.1     | 2.15 | 40.0     | 7700           | 410            |
| 2012A | 715      | 13.1     | 0.525| 1.16     | 710            | 400            |
| 2013ej| 1500     | 26.1     | 1.4  | 3.9      | 6500           | 800            |
| 2016X | 436      | 28.0     | 1.73 | 2.95     | 4000           | 760            |
| 2017gmr| 525     | 22.0     | 10.2 | 11.0     | 3300           | 640            |

---

Figure 15. Explosion energy (Panel (a)) and $^{56}$Ni mass (Panel (b)) versus ejecta mass for SN 2017gmr and eleven other core-collapse SNe (Utrobin & Chugai 2019). Dotted line in Panel (a) is the upper limit of the explosion energy of $2 \times 10^{51}$ erg for the neutrino-driven mechanism (Janka 2017) with the uncertainty of about $\pm 10^{51}$ erg$^2$ shown by the shaded green band.

---

$^2$ H.-Th. Janka, private communication.
in the spectra between 50 days and 100 days and use the velocities of these features to constrain the mass of the confined CS shell of $\sim 4 \times 10^{-3} M_\odot$ which turns out to be comparable to that of SN 2013fs (Yaron et al. 2017). The requirement of the low contribution of the CS interaction luminosity combined with the kinematic constraints from the NHV A implies the preferred explosion energy of $6 \times 10^{51}$ erg that is lower than the energy implied by the hydrodynamic model. The contradiction casts a shadow on the thin shell model and suggests the need for the full radiation hydrodynamics treatment of the shock wave propagation in the dense CS shell, which however cannot be done by the available hydrodynamic code.

The explosion energy of SN 2017gmr is indeed enormous for SNe IIP and places this event to the category of high-energy SNe IIP with other two cases of SN 2000cb and SN 2009kf (Table 4, Fig. 15). The explosion energies of these three supernovae exceed the upper limit for the neutrino-driven explosion mechanism which implies that their explosions could be related to the rotational energy of the collapsing core. Unfortunately, spectra of SN 2009kf and SN 2000cb at the late nebular stage are lacking, so one cannot say anything about possible asymmetry of the $^{56}$Ni ejecta in these supernovae. It is noteworthy that in both preferred models for SN 2017gmr, HMXB and HP-optm (Table 1), the $^{56}$Ni mass exceeds the amount of $^{56}$Ni typical for SNe IIP (Table 4, Fig. 15), which is in line with the unusually high explosion energy of SN 2017gmr.

ACKNOWLEDGEMENTS

VPU is partially supported by Russian Scientific Foundation grant 19-12-00229. This work makes use of observations from the Las Cumbres Observatory global telescope network. The LCO team is supported by NSF grants AST-1911225 and AST-1911151.

DATA AVAILABILITY

The data underlying this article will be shared on reasonable request to the corresponding author.

REFERENCES

Albareti F. D., et al., 2017, ApJS, 233, 25
Andrews J. E., et al., 2019, ApJ, 885, 43
Baklanov P. V., Blinnikov S. I., Pavlyuk N. N., 2005, Astronomy Letters, 31, 429
Bisnovatyi-Kogan G. S., 1971, Soviet Astr., 14, 652
Blinnikov S., Sorokina E., 2004, Ap&SS, 290, 13
Blinnikov S. I., Eastman R., Bartunov O. S., Popolitov V. A., Woosley S. E., 1998, ApJ, 496, 454
Blinnikov S. I., Röpke F. K., Sorokina E. I., Gieseler M., Reinecke M., Travaglio C., Hillebrandt W., Stritzinger M., 2006, A&A, 453, 229
Brown T. M., et al., 2013, PASP, 125, 1031
Bullivant C., et al., 2018, MNRAS, 476, 1497
Burrows A., Vartanyan D., 2021, Nature, 589, 29
Catchpole R. M., et al., 1988, MNRAS, 231, 75P
Chevalier R. A., 1982, ApJ, 259, 302
Chevalier R. A., 1989, ApJ, 346, 847
Chugai N. N., 2001, MNRAS, 326, 1448
Chugai N. N., 2006, Astronomy Letters, 32, 739
Chugai N. N., 2019, Astronomy Letters, 45, 427
Chugai N. N., Danziger I. J., 1994, MNRAS, 268, 173
Chugai N. N., Fabrika S. N., Sholukhova O. N., Goranskij V. P., Abolmasov P. K., Vlasyuk V. V., 2005, Astronomy Letters, 31, 792
Chugai N. N., Chevalier R. A., Utrobin V. P., 2007, ApJ, 662, 1136
Colgate S. A., 1971, ApJ, 163, 221
Colgate S. A., White R. H., 1966, ApJ, 143, 626
Dhungana G., et al., 2016, ApJ, 822, 6
Elmhamdi A., et al., 2003, MNRAS, 338, 939
Goldberg J. A., Bildsten L., 2020, ApJ, 895, L45
Grassberg E. K., Imshennik V. S., Nadyozhin D. K., 1971, Ap&SS, 10, 28
Hillionon A. F., Sunyaev R. A., 1975, Soviet Astronomy Letters, 1, 73
Janka H.-T., 2017, Neutrino-Driven Explosions. p. 1095, doi:10.1007/978-3-319-21846-5_109
Khokhlov A. M., Höffich P. A., Oran E. S., Wheeler J. C., Wang L., Chchelkchening A. Y., 1999, ApJS, 524, L107
Kozma C., Fransson C., 1992, ApJ, 390, 602
LeBlanc J. M., Wilson J. R., 1970, ApJ, 161, 541
Liu W., Dalgarano A., 1995, ApJ, 454, 472
MacFadyen A. I., Woosley S. E., Heger A., 2001, ApJ, 550, 410
McCully C., Volgenau N. H., Harbeck D.-R., Lister T. A., Saunders E. S., Turner M. L., Siuverd R. J., Bowman M., 2018, in Guzman J. C., Ibsen J., eds, Society of Photo-Optical Instrumentation Engineers (SPIE) Conference Series Vol. 10707, Software and Cyberinfrastructure for Astronomy V. p. 107070K (arXiv: 1811.04163), doi:10.1117/12.2314340
Nagao T., et al., 2019, MNRAS, 489, L69
Nicholl M., 2018, Research Notes of the American Astronomical Society, 2, 230
Press W. H., Teukolsky S. A., Vetterling W. T., Flannery B. P., 2007, Numerical Recipes 3rd Edition: The Art of Scientific Computing, 3 edn. Cambridge University Press
Shakura N. I., 1975, Soviet Astronomy Letters, 1, 223
Shiode J. H., Quataert E., 2014, ApJ, 780, 96
Smith N., et al., 2012, MNRAS, 420, 1135
Smith N., et al., 2015, MNRAS, 449, 1876
Stetson P. B., 2000, PASP, 112, 925
Utrobin V. P., 2004, Astronomy Letters, 30, 293
Utrobin V. P., 2007, A&A, 461, 233
Utrobin V. P., Chugai N. N., 2008, A&A, 491, 507
Utrobin V. P., Chugai N. N., 2011, A&A, 532, 1100
Utrobin V. P., Chugai N. N., 2017, MNRAS, 472, 5004
Utrobin V. P., Chugai N. N., 2019, MNRAS, 490, 2042
Utrobin V. P., Chugai N. N., Botticella M. T., 2010, ApJ, 723, L89
Utrobin V. P., Wongwathanarat A., Janka H.-T., Müller E., 2017, ApJ, 846, 37
Valenti S., et al., 2016, MNRAS, 459, 3939
Weaver T. A., Woosley S. E., 1980, in Ninth Texas Symposium on Relativistic Astrophysics. pp 335–357, doi:10.1111/j.1749-6632.1980.tb15942.x
Whitelock P. A., et al., 1988, MNRAS, 234, 5P
Woosley S. E., 1988, ApJ, 330, 218
Xu Y., McCray R., Oliiva E., Rand Icch S., 1992, ApJ, 386, 181
Yaron O., et al., 2017, Nature, Physics, 13, 510

This paper has been typeset from a TeX/\LaTeX file prepared by the author.