Constraining Magnetic Field Amplification in SN Shocks Using Radio Observations of SNe 2011fe and 2014J

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Abstract

We modeled the radio non-detection of two Type Ia supernovae (SNe), SN 2011fe and SN 2014J, considering synchrotron emission from the interaction between SN ejecta and the circumstellar medium. For ejecta whose outer parts have a power-law density structure, we compare synchrotron emission with radio observations. Assuming that 20% of the bulk shock energy is being shared equally between electrons and magnetic fields, we found a very low-density medium around both the SNe. A less tenuous medium with particle density \(\sim 1 \text{ cm}^{-3}\), which could be expected around both SNe, can be estimated when the magnetic field amplification is less than that presumed for energy equipartition. This conclusion also holds if the progenitor of SN 2014J was a rigidly rotating white dwarf (WD) with a main-sequence (MS) or red giant companion. For a He star companion, or a MS for SN 2014J, with 10% and 1% of bulk kinetic energy in magnetic fields, we obtain mass-loss rates of \(<10^{-9}\) and \(<4 \times 10^{-9}\) \(M_\odot\) yr\(^{-1}\) for a wind velocity of 100 km s\(^{-1}\). The former requires a mass accretion efficiency of >99% onto the WD, but is less restricted for the latter case. However, if the tenuous medium is due to a recurrent nova, it is difficult from our model to predict synchrotron luminosities. Although the formation channels of SNe 2011fe and 2014J are not clear, the null detection in radio wavelengths could point toward a low amplification efficiency for magnetic fields in SN shocks.

Key words: circumstellar matter – ISM: magnetic fields – supernovae: general – supernovae: individual (SN 2011fe, SN 2014J)

1. Introduction

Astrophysical shocks are suitable places for high-energy particle acceleration and magnetic field amplification. The accelerated ions, via plasma instabilities, can amplify magnetic fields (Bykov et al. 2013). As particle acceleration is not very clear, how much magnetic field is generated in these shocks still remains an open question. Recently, Caprioli & Spitkovsky (2014a, 2014b) tried to investigate these mechanisms through hybrid simulations. However, their effort was restricted by computational limitations, which do not allow simulations in which a shock has a very high Mach number like supernova (SN) shocks. Another way to understand these complex phenomena is through observing radiation emitted by the shock-accelerated particles. A highly relativistic particle loses a fraction of its energy in the form of synchrotron radiation. The power of this radiation is proportional to the energy density of the ambient magnetic field. Therefore, a proper modeling of synchrotron emission from these shocks, along with observations, could give an estimate of the ambient magnetic field, and this can then shed light on the complex mechanisms of particle acceleration.

The interaction of high-velocity SN ejecta with the ambient medium launches a strong shock into the circumstellar medium (CSM). Particles accelerated in these shocks can emit synchrotron radiation in radio wavelengths. When these SNe are nearby the probability of detection increases. In this regard, the two nearest young SNe Ia, SN 2011fe, in M101, and SN 2014J, in M82, at distances of 6.4 Mpc (Shappee & Stanek 2011) and 3.4 Mpc (Marion et al. 2015), respectively, provide a unique opportunity to detect the radiation and understand the circumstances which led to that emission. Several observational attempts were made to detect radio emission from these two SNe, but nothing was detected within the first 150 days of explosion (Horesh et al. 2012; Chomiuk et al. 2012; Pérez-Torres et al. 2014; Chomiuk et al. 2016). As reported here, observations show null detection even around 1.5 and 4 years after the explosions of SNe 2014J and 2011fe, respectively.

The CSM plays an important role in shock dynamics. This medium is shaped by the pre-SN mass lost history of the progenitor system. Both the aforementioned SNe are of Type Ia, which are widely accepted as the thermonuclear explosion of a CO white dwarf (WD) (Hoyle & Fowler 1960). Before the explosion, this WD accretes matter from a companion star, which could be a main-sequence (MS) or an asymptotic giant branch star, and explodes when it reaches the Chandrasekhar (CH) mass limit (Whelan & Iben 1973). This is called the single degenerate (SD) channel. A second possibility is to have another WD as a companion, with the total mass of the system being more than the CH limit. In this case, known as the double degenerate (DD) scenario, the two spiraling WDs merge due to the emission of gravitational waves, which reduces the orbital separation, and, under the right physical conditions, gives rise to a successful thermonuclear runaway (Iben & Tutukov 1984; Webbink 1984). The companion star, in the SD scenario, loses a fraction of its initial mass either due to a wind or Roche Lobe overflow, which results in a high-density circumbinary medium. On the contrary, the merger model predicts a very clean environment around the progenitor system.

Previous modelings of null detection of radio emission from SNe 2011fe and 2014J (Chomiuk et al. 2012; Pérez-Torres...
et al. 2014; Chomiuk et al. 2016) prefer to put stringent upper limits on the ambient medium density, $n_{\text{ISM}}$, and the mass-loss rate, $\dot{M}$, of the progenitor system, assuming equipartition of energy, in the post-shock region, between electric and magnetic fields. This assumption suggests a very tenuous medium around both the SNe. The radio-silent feature of these SNe could also be due to a lower strength of magnetic fields in these shocks. In the present work we examine both equipartition and a lower-energy density in magnetic fields compared to that in electric fields, and try to constrain the magnetic field amplification in SN shocks.

The paper is organized as follows. In the next section we discuss the explosion models of the WDs used and SN ejecta profile. Along with this we discuss the synchrotron radio emission model. Radio observations are given in Section 3. We present our results in Section 4 and discuss them in Section 5. Conclusions are drawn in Section 6.

2. Explosion Models and Radio Emission

Both SNe 2011fe and 2014J are normal SNe Ia with $^{56}$Ni masses of $\sim 0.6 M_\odot$. (Röpke et al. 2012; Churazov et al. 2014).

Within the framework of the SD scenario, the N100 model (Röpke et al. 2012; Seitenzahl et al. 2013) is capable of reproducing the observables of normal SNe Ia with reasonable accuracy. This is a delayed detonation model where the central region is ignited by 100 sparks. Another channel that could account for normal SNe Ia is the merger of two sub-CH WDs (Pakmor et al. 2012). Here, two C/O degenerate stars with masses of 1.1 and 0.9 $M_\odot$ merge and lead to a successful SN explosion. This model probes the DD formation channel. The total masses and asymptotic kinetic energies of ejecta for N100 and the violent merger model are 1.4 $M_\odot$, 1.95 $M_\odot$, and 1.45 $\times 10^{51}$ erg, 1.7 $\times 10^{51}$ erg, respectively.

For both models, numerical simulations provide the density of the ejecta as a function of velocity, in the homologous phase, up to $\sim 2.5 \times 10^4$ km s$^{-1}$. The ejecta structure of the extreme outer part of the expanding WD is not available from these simulations. This part of the ejecta interacts with the CSM, and hence plays an important role in radio emission. Therefore, beyond the above velocity range we have extrapolated the density ($\rho$) profile considering a power law $\rho \propto r^{-n}$, where $r$ represents the radius of the SN and $n$ is the power-law index. The initial structure of the outer part of the WD is uncertain.

The accretion rate and subsequent burning phases play a vital role in determining the properties of the outer edge. We therefore allow $n$ to vary in the range 12–14. For convenience we call the N100 + power-law structure the N100 model and similarly refer to the violent merger + power-law profile as the merger model.

In the SD scenario the SN ejecta most likely interact with a wind medium characterized by a wind velocity ($v_w$) and mass-loss rate $M$ of the progenitor system. When $M$ and $v_w$ both are constants, the density of this surrounding medium is expressed as $\rho_{\text{CSM}} = M / (4 \pi r^2 v_w)$. On the other hand, in the DD formation channel, the ejecta launching a strong shock into a presumably constant density medium with $\rho_{\text{CSM}} = n_{\text{ISM}} \mu$, where $\mu$ represents the mean atomic weight of the surrounding medium. Therefore, the density of the CSM can be written as $\rho_{\text{CSM}} = A r^{-s}$, with $s = 2$ or 0 for a wind or a constant density medium, respectively. The interaction of the ejecta, with power-law structure, and this surrounding medium, gives rise to a self-similar structure (Chevalier 1982). In this case, the forward-shock radius, $r_\text{f}$, varies with $A$ and time, $t$, as $r_\text{f} \propto A^{(s-n)/(n-3)} t^{(3-n)/(n-3)}$, which implies a shock speed of $v_\text{f} = \frac{d r}{d t} = \frac{(n-3)}{(n-3) r_\text{f}}$. The synchrotron power is inversely proportional to the square of the mass of the emitting particle. Therefore, in a medium with low magnetic field strength, electrons rather than ions are more likely to be responsible for this radiation. It is expected that along with ions electrons are also accelerated in shocks. For a test particle, in a non-relativistic shock with a compression factor of 4, diffusive shock acceleration predicts a power-law energy spectrum with an index $p = 2$. When synchrotron or inverse Compton losses are important, the index of the integrated particle spectrum becomes $(p + 1)$. The presence of a sub-shock along with the main shock predicts a particle spectrum with index $> 2$ for low-energy particles (Ellison & Reynolds 1991). For the case of SNe Type Ib/c, which originate from compact progenitors like SNe Type Ia, it is observed that the spectrum of relativistic particles is steeper, with an index $p = 3$, with the loss mechanisms having no role to play (Chevalier & Fransson 2006). These loss mechanisms are found to be irrelevant for both the SNe we are concerned with here. The insufficient knowledge of particle acceleration, at shocks, does not allow a unique choice of $p$. We therefore assume that the energy spectrum of the electrons obeys $dN/dE = N_0 E^{-p}$ with $p = 3$, considering the above similarities with SNe Type Ib/c. Here $N_0$ and $E$ represent the normalization constant of distribution and electron energy, respectively. We assume that all electrons in the post-shock region are accelerated. For a power-law energy distribution the minimum energy of this population varies as $E_{\text{min}} \propto r_\text{f}^2 T^{-2} \epsilon_{\text{r}}$, for a constant value of $\epsilon_{\text{r}}$, where $\epsilon_{\text{r}} = \epsilon_{\text{rel}} + \epsilon_{\text{rad}} = u_e / u_\text{th}$ represents the fraction of post-shock energy density, $u_\text{th} = \frac{9}{8} \rho (f) v_\text{r}^2$, that goes into electrons. $u_e$ is the energy density in electric fields. This energy is shared by relativistic electrons with $E > m_e c^2 (\equiv E_{\text{rel}})$, where $m_e$ and $c$ are the electron mass and velocity of light in a vacuum, and non-relativistic electrons with a kinetic energy of the electron $E_{\text{r}} < E_{\text{rel}}$. The two quantities $\epsilon_{\text{rel}}$ and $\epsilon_{\text{rad}}$ give an account of the shared energy between the two components. Initially, when the shock velocity is very high, the entire electron population is relativistic. This implies that $\epsilon_{\text{r}} = \epsilon_{\text{rel}}$ and $\epsilon_{\text{rel}} > 0.16 \left(\frac{v_r}{10^4 \text{ km s}^{-1}}\right)^2$ for $p = 3$ (Chevalier & Fransson 2006). When the above criterion is not fulfilled, there exists a non-relativistic electron population. The presence of this component implies that $\epsilon_{\text{rad}} = \epsilon_{\text{r}} \left(\frac{E_{\text{rad}}}{E_{\text{rel}}}\right)^{(p-2)}$. We consider electrons with $E > E_{\text{est}}$ to contribute to the synchrotron emission, as only relativistic particles radiate part of their energy via synchrotron. We note that the assumption of a power-law spectrum for non-relativistic electrons may not be realistic; however, this allows us to estimate the energy density in relativistic electrons as a function of time. In previous studies (Pérez-Torres et al. 2014; Chomiuk et al. 2016) it was assumed that all the electrons, in the post-shock region, will be relativistic at any point in time. This introduces an artificial kink in their synchrotron radio light curve, and beyond this kink the light curve rises rather steeply. Therefore, when the radio emissions predicted from such models are compared with observations, they suggest a rather low-density medium around 2.
the progenitor system. As we consider the effect of shock deceleration in estimating $\xi_{b}$, and hence in the synchrotron light curve, our model predicts a more feasible density of the circumbinary medium where the shock was launched. Caprioli & Spitkovsky (2014a) show that at non-relativistic shocks, energetic particles could acquire 10%–20% of the post-shock energy. As particle acceleration is a complex mechanism and depends on the orientation of the background magnetic fields (see Caprioli & Spitkovsky 2014a for details) it is reasonable to consider the lower limit of acceleration efficiency. Therefore, for electrons we assumed that 10% of the bulk kinetic energy is transferred to them, i.e., $\epsilon_{e} = 0.1$, and kept it fixed at this value for our entire study.

An SN will be in the free-expansion phase until the reverse-shock remains in the outer power-law part of the ejecta. The unshocked ejecta profiles, as a function of radius at 20 days after the explosion, for both the N100 and merger models, are shown in Figure 1. The left panel depicts the ejecta structure when the SN plows through a constant density medium with $n_{ISM} = 0.2 \text{ cm}^{-3}$, and the right panel shows the ejecta profile (for N100) in a wind medium, characterized by $\dot{M} = 8 \times 10^{-10} \left( \frac{v_{s}}{100 \text{ km s}^{-1}} \right)^{-1} \text{ M}_{\odot} \text{ yr}^{-1}$. In contrast to Pérez-Torres et al. (2014), where a thin-shell approximation was used for the shock structure, here we take into account the actual structure of the shocked gas from the similarity solutions (Chevalier 1982) for the various values of $s$ and $n$ investigated.

The intensity of the synchrotron radiation, with synchrotron self-absorption (SSA) as the main absorption mechanism, from a shell of thickness $\Delta r$ can be written as

$$I_{s}(h) = \frac{2kT_{\text{bright}}}{c^2 f\left(\frac{\nu_{\text{peak}}}{\nu_{\text{abs}}},\frac{\nu_{s}}{\nu_{\text{abs}}}\right)} \nu_{s}^{5/2} \left[ 1 - \exp\left( -\xi_{b} \tau_{e} \right) \right]$$

(1)

with

$$f(x) = x^{1/2} \left[ 1 - \exp\left( -x^{-(p+4)/2} \right) \right]$$

(2)

(Pérez-Torres et al. 2014; Björnsson & Lundqvist 2014). Here $\xi_{b} = \Delta s(h)/(2\Delta r)$ gives the normalized path length traversed by the radiation along the line of sight and $0 < h \leq 1$ takes into account emission from different parts of the shell. $\nu_{\text{abs}}$ is the absorption frequency for which optical depth $\tau_{\text{abs}} = 1$ and for $h = 0$ $\nu_{\text{abs}} = \nu_{\text{abs,0}}$ and $\tau_{e} = \tau_{e,0}$, respectively. $T_{\text{bright}}$ and $\nu_{\text{peak}}$ represent the brightness temperature and peak frequency of the radiation. $k$ is the Boltzmann constant.

According to the models, SNe 2011fe and 2014J are in the optically thin regime when the radio observations are done (see Section 3 for radio observations). Therefore, the luminosity for $\tau_{e} < 1$ from a shell with an outer radius $r_s$ can be written as:

$$L_{s,\text{thin}} = \frac{8\pi^2 k T_{\text{bright}}}{c^2 f(\nu_{\text{peak}}/\nu_{\text{abs}}, \nu_{s}/\nu_{\text{abs}})} \nu_{s}^{(p+3)/2} \nu_{e}^{-(p-1)/2},$$

(3)

with

$$\nu_{\text{abs,0}} = (2\Delta r \tau(p) N_0 B^{(p+2)/2})^{2/(p+4)}$$

(4)

where $\nu_{e} = \frac{L_s}{L_{e,0}}$ with $L_{e,0} = 4\pi^2 \nu_{s}^2 \epsilon_{e}$. $B$ represents the magnetic field strength and $\tau(p)$ is the SSA coefficient.

Observational evidence suggests that the brightness temperature, $T_{\text{bright}}$, is $\sim 10^{11}$ K (Readhead 1994). In our modeling we consider a value of $5 \times 10^{10}$ K and keep it constant throughout the evolution, which is also what was assumed in Pérez-Torres et al. (2014).

3. Radio Data

In this paper, we modeled the publicly available radio data for both SN 2011fe and SN 2014J, as well as radio observations obtained by us. Namely, for the modeling of SN 2014J, we used data published in Pérez-Torres et al. (2014), complemented by recently obtained data with the Jansky Very Large Array (JVLA) at 3.0 GHz, and by the European VLBI Network (EVN) at 1.66 GHz. For the modeling of the radio emission from SN 2011fe, we used data published by Chomiuk et al. (2012), also complemented by recently obtained 3.0 GHz JVLA observations, and by 1.66 GHz EVN observations.

3.1. New EVN Observations

We observed SN 2014J on 2015 February 28–March 1 (project EP092A; PI: Pérez-Torres), and SN 2011fe on 2015 June 14 (project EP092B; PI: Pérez-Torres), using the EVN at 1.66 GHz, for 12 hr in each case (with an integrated on-source time of about 8.0 hr). Each time, we used a sustained data
recording rate of 1024 Mbit s⁻¹, in dual-polarization mode and with 2 bit sampling. Each frequency band was split into 8 intermediate subbands of 16 MHz bandwidth each, for a total synthesized bandwidth of 128 MHz. Each subband was in turn split into 32 spectral channels of 500 kHz bandwidth each.

The EVN observations of SN 2014J (EP092A) included the following 13-antenna array: Effelsberg, Westerbork, Jodrell Bank, Onsala, Medicina, Torun, Urumqi, Sheshan, Svetloe, Zelenchuk, Badary, Sardinia, and the DSN antenna of Robledo (Madrid), which joined the observations for about half of the total time, although it was only recorded in LCP. We note that Onsala could not take data due to heavy winds, and, most importantly, Effelsberg, our most sensitive antenna, could not observe for 5 hr. Fortunately, we had additional large dishes for this experiment, which allowed us to obtain a final rms value close to the expected one (see below). We observed our target source, SN 2014J, phase-referenced to the core of the nearby galaxy M81, known to be very compact at VLBI scales, with a typical duty cycle of 5 minutes. We used the strong QSO 1150 +812 as a fringe finder and bandpass calibrator. For our observations of SN 2011fe (EP092B), the array included 10 antennas: Effelsberg, Westerbork, Jodrell Bank, Onsala, Medicina, Torun, Svetloe, Zelenchuk, Badary, and Tianma. Unfortunately, the data for Jodrell Bank and Onsala were of very bad quality, so we had to remove them completely, thus resulting in an eight-antenna array. Furthermore, the performances of Badary and especially Tianma (the second most sensitive antenna) were rather poor. As a result, the attained rms was far from the initially expected one, and significantly worse than that attained for EP092A (see below). We observed our target source, SN 2011fe, phase-referenced to the strong, nearby source J1359+5544, known to be very compact at VLBI scales, with a typical duty cycle of 5 minutes. We used the strong source 3C309.1 as a fringe finder and bandpass calibrator.

All the data were correlated at the EVN MkIV data processor of the Joint Institute for VLBI in Europe (JIVE, the Netherlands), using an averaging time of 4 s for SN 2014J, and 2 s for SN 2011fe.

We used AIPS for calibration, data inspection, and flagging of our eEVN data, using standard procedures. Those steps included a priori gain calibration (using the measured gains and system temperatures of each antenna), parallactic angle correction, and correction for ionosphere effects. We then aligned the visibility phases in the different subbands, i.e., "fringe-fitted" the data, solved for the residual delays and delay rates, and interpolated the resulting gains into the scans of SN 2014J. We then imaged a field of view of 3′ × 3′, centered at the position given by Smith et al. (2014), and applied standard imaging procedures using AIPS, without averaging the data either in time, or frequency, to prevent time and bandwidth smearing of the images. We used natural UV-weighting to maximize the signal-to-noise ratios in our final images, which were 9.8 μJy/beam and 34.3 μJy/beam for SN 2014J and SN 2011fe, respectively.

3.2. Jansky VLA Observations

We also observed SN 2014J and SN 2011fe with the Jansky VLA on 2015 August 15–16 and 2015 August 31 (project 15A-076; PI: Pérez-Torres), respectively, while the VLA was in its most extended configuration. Each target was observed for 2 hr, yielding an effective on-target time of ~87 minutes. 3C286 was used as a flux and bandpass calibrator. We centered our observations at the frequency of 3.0 GHz, covering the whole S-band (from 2.0 to 4.0 GHz), and recorded in full polarization, using a dump time of 2 s. The nearby, bright, point-like sources J0954+7435 (for SN 2014J) and J1419+5423 (for SN 2011fe) had been considered for phase calibration. We used the Common Astronomy Software Applications (McMullin et al. 2007) package for data reduction purposes, which consisted of standard amplitude and phase calibration. Since our target sources were far from any bright point-like source, a natural weighting scheme was used to attain the best possible sensitivity in our images. In the case of SN 2011fe, we essentially reached the expected off-source rms value (4.0 μJy/beam). This was possible because the VLA observations were carried out with the array in its most extended configuration, i.e., with the smallest possible synthesized beam at a given frequency. In this way, most diffuse emission was expected to be filtered out for M101, as was indeed the case. However, for SN 2014J, the host galaxy, M82, a starburst galaxy, is extremely bright at low frequencies in the few central kiloparsecs (see Figure 2). In fact, the 3.0 GHz VLA emission is not limited by the thermal noise, since very far from the galaxy we actually reach values of about 7μJy/beam, very close to the nominal values. However, we are limited by the dynamic range and the intrinsic background radio emission in the surroundings of SN 2014J, which are significantly higher, ~32μJy/beam.

Tables 1 and 2 show the summary of our new radio observations of SN 2011fe and SN 2014J, respectively, along with other existing radio data.

4. Results

When the SN plows through a constant density medium the radio luminosity evolution, for an optically thin shell, can be
Table 1

| Starting UTC | Epoch days | $t_{int}$ hours | Array | $\nu$ GHz | $S_\nu$ $\mu$Jy | $L_{\nu,23}$ |
|--------------|------------|----------------|-------|------------|----------------|-------------|
| 2011 Aug 25.77 | 2.1 ... | JVLA$^a$ | 5.90 | 17.4 | 8.53 |
| 2015 Jan 14 | 1390 | 8.0 | EVN$^b$ | 1.66 | 105 | 50.5 |
| 2015 Aug 31 | 1468 | 1.45 | JVLA$^b$ | 3.00 | 12.0 | 5.88 |

Notes. The columns, from left to right, are as follows: starting observing time, UTC; mean observing epoch (in days since explosion, assumed to be on 2011 August 23.7); integration time, in hr; facility; central frequency in GHz; $3\sigma$ flux density upper limits, in $\mu$Jy; the corresponding $3\sigma$ spectral luminosity, assuming a distance of 6.4 Mpc, in units of $10^{23}$ erg s$^{-1}$ Hz$^{-1}$.  

$^a$ Data from Chomiuk et al. (2012).  

$^b$ This work.

Table 2

| Starting UTC | Epoch days | $t_{int}$ hours | Array | $\nu$ GHz | $S_\nu$ $\mu$Jy | $L_{\nu,23}$ |
|--------------|------------|----------------|-------|------------|----------------|-------------|
| 2014 Jan 23.2 | 8.2 ... | JVLA$^a$ | 5.50 | 12.0 | 1.67 |
| 2014 Jan 24.4 | 9.4 ... | JVLA$^a$ | 22.0 | 24.0 | 3.31 |
| 2014 Jan 28.8 | 13.8 13.6 | JVLA$^a$ | 1.55 | 37.2 | 5.15 |
| 2014 Jan 29.5 | 14.5 14.0 | eMERLIN$^b$ | 6.17 | 40.8 | 5.63 |
| 2014 Feb 4.0 | 20.0 11.0 | eEVN$^b$ | 1.66 | 32.4 | 4.47 |
| 2014 Feb 19.1 | 35.1 10.0 | eEVN$^b$ | 1.66 | 28.5 | 3.94 |
| 2014 Apr 11.2 | 86.2 ... | JVLA$^c$ | 5.90 | 21.0 | 2.92 |
| 2014 Jun 12.0 | 148 ... | JVLA$^c$ | 7.40 | 41.0 | 5.71 |
| 2015 Mar 1 | 410 8.0 | EVN$^d$ | 1.66 | 29.4 | 4.10 |
| 2015 Aug 16 | 578 1.45 | JVLA$^d$ | 3.00 | 96.0 | 13.4 |

Notes. The columns, from left to right, are as follows: starting observing time, UTC; mean observing epoch (in days since explosion, assumed to be on 2014 January 15.0); integration time, in hr; facility; central frequency in GHz; $3\sigma$ flux density upper limits, in $\mu$Jy; the corresponding $3\sigma$ spectral luminosity, assuming a distance of 3.4 Mpc to M82, in units of $10^{23}$ erg s$^{-1}$ Hz$^{-1}$.  

$^a$ Data from Chandler & Marvil (2014).  

$^b$ Pérez-Torres et al. (2014).  

$^c$ Chomiuk et al. (2012).  

$^d$ This work.

written as follows for $p = 3$:  

$$L_{\nu,\text{thin}} \propto \left( \frac{n-3}{n} \right)^{3.86} \frac{1}{T_{\text{bright}}} \frac{1}{\epsilon_{\text{rel}}} \left( \frac{1}{\eta_{\text{ISM}}} \right) \frac{T_{\nu}}{T_{\nu}}$$  

and in the case of a wind stratified medium, $s = 2$,  

$$L_{\nu,\text{thin}} \propto \left( \frac{n-3}{n-2} \right)^{3.86} \frac{1}{T_{\text{bright}}} \frac{1}{\epsilon_{\text{rel}}} \left( \frac{1}{\eta_{\text{ISM}}} \right) \frac{T_{\nu}}{T_{\nu}}$$  

Here $\epsilon_B = \frac{n_B}{\eta_{\text{ISM}} B}$, where $\eta_B = B^2/8\pi$ is the amount of post-shock energy density in magnetic fields.  

We consider two cases: (1) equipartition of energy between electric and magnetic fields, i.e., when $\epsilon_B = \epsilon_B = 0.1$. This implies that each field carries 10% of the bulk shock energy; and (2) $\epsilon_B = 0.1$ and $\epsilon_B = 0.01$, i.e., when 10% and 1% of post-shock energy are in electric and magnetic fields, respectively. Therefore, in the second situation the amplification of the magnetic field, in the post-shock region, is less than that in the former one. In the paper when we mention a low amplification efficiency of magnetic fields, or the magnetic field amplification is less than that presumed for energy equipartition, or $\epsilon_B = 0.01$, we refer to the second case. The first case where energy equipartition is assumed is cited as $\epsilon_B = 0.1$ throughout the paper.

For $\epsilon_B = 0.1$ and $n = 13$ the light curves predicted by the merger model are shown in the left panel of Figure 3 for SN 2014J. As can be seen, the density of the CSM is constrained by the upper limit at 1.66 GHz measured around 400 days after the explosion. This enforces the density of the surroundings to be $<0.32$ cm$^{-3}$ for a constant density medium. For the same parameters the radio light curves calculated using the N100 model in a wind medium are depicted in the right panel of Figure 3. In this case, the mass-loss rate of the progenitor system is restricted by the observational upper limit obtained at 5.0 GHz, around 8 days after the explosion of SN 2014J (Pérez-Torres et al. 2014). According to our model this implies $M < 6.9 \times 10^{-10} (v_{\nu}/100$ km s$^{-1})^{-1} M_\odot$ yr$^{-1}$ (right panel of Figure 3).

As for SN 2014J, SN 2011fe was also thoroughly searched for in the radio just after the explosion (Horesh et al. 2012; Chomiuk et al. 2012; Chomiuk et al. 2016). Among the early observations, we found that the upper limit on the luminosity at 5.9 GHz, measured around 2 days after the explosion, (Chomiuk et al. 2012; tabulated in Table 1) constrains the maximum possible mass-loss rate from the system. This was also noted by Chomiuk et al. (2012) and Pérez-Torres et al. (2014). For a case with a constant density medium, the radio luminosity increases with time (Equation (5)) when $12 \leq n \leq 14$, in contrast to the wind medium in which the luminosity decreases with time (Equation (6)) (see the trends of the light curves of SN 2014J displayed in Figure 3 for wind and constant density medium $n = 13$). It is therefore found that for a constant density circumbinary medium, the radio upper limit measured around 4 years after the explosion, at 3 GHz, limits $\eta_{\text{ISM}}$. Table 1 contains the details of these observations, which are relevant to derive the upper limits on the CSM density, together with the hitherto unpublished limits measured at 1390 days at 1.66 GHz.

In Figure 4, we show the dependence of the CSM density on $n$ using N100 (blue) and the merger model (red). The solid lines with crosses (filled circles) represent the upper limits on $\eta_{\text{ISM}}$ or $M$ for $v_{\nu} = 100$ km s$^{-1}$ around SN 2011fe (SN 2014J), when $\epsilon_B = 0.1$. The dashed lines show the effect of a low magnetic field energy density, $\epsilon_B = 0.01$. As is clear from the figure, there is little variation on the upper limits of $\eta_{\text{ISM}}$ when $n$ varies in range 12–14. In case of a wind medium, $M/v_{\nu}$, calculated from N100, has a rather strong dependence on $n$. The upper limits on $\eta_{\text{ISM}}$, from both models and for both SNe, are $\sim 0.35$ cm$^{-3}$ when equipartition of energy between electric and magnetic fields is considered, i.e., $\epsilon_B = \epsilon_B = 0.1$. For a low magnetic field energy density ($\epsilon_B = 0.01$), which implies a low amplification efficiency for magnetic fields in the shocked region, the upper limits are $\sim 2$ cm$^{-3}$. The upper limits on the mass-loss rate predicted by the N100 model, for SNe 2011fe and 2014J, are $<10^{-9} M_\odot$ yr$^{-1}$ and $<4 \times 10^{-9} M_\odot$ yr$^{-1}$ for $\epsilon_B = 0.1$ and 0.01, respectively, assuming $v_{\nu} = 100$ km s$^{-1}$.

Several authors have calculated the upper limits on $M$ and $\eta_{\text{ISM}}$, for both SNe 2011fe and SN 2014J, considering...
\( \epsilon_e = \epsilon_B = 0.1 \) and \( v_e = 100 \text{ km s}^{-1} \). From a very early non-detection in radio, Horesh et al. (2012) constrained \( \dot{M} \) of SN 2011fe to be \(< 10^{-8} M_\odot \text{ yr}^{-1} \) assuming \( v_e = 4 \times 10^4 \text{ km s}^{-1} \). However, just after the explosion, the shock velocity could be much higher than \( 4 \times 10^4 \text{ km s}^{-1} \) (Berger et al. 2002). Chomiuk et al. (2012) estimated \( \dot{M} < 6 \times 10^{-10} M_\odot \text{ yr}^{-1} \) (\( \eta_{\text{ISM}} < 6 \text{ cm}^{-3} \)) using their upper limits on radio luminosity, which were deeper than those reported in Horesh et al. (2012).

For SN 2014J, Pérez-Torres et al. (2014) found \( \dot{M} < 7 \times 10^{-10} M_\odot \text{ yr}^{-1} \) (\( \eta_{\text{ISM}} < 1.3 \text{ cm}^{-3} \)) from their radio analysis. From X-ray observations, where \( \epsilon_B \) plays no role, Margutti et al. (2012, 2014) obtained \( \dot{M} < 2 \times 10^{-9} M_\odot \text{ yr}^{-1} \) and \(< 1.2 \times 10^{-9} M_\odot \text{ yr}^{-1} \) (\( \eta_{\text{ISM}} < 166 \text{ cm}^{-3} \)) and \(< 3.5 \text{ cm}^{-3} \)) for SNe 2011fe and 2014J, respectively. In our present analysis we used an ejecta structure and a radio emission model similar to those presented in Chomiuk et al. (2012) and Pérez-Torres et al. (2014). We estimate \( \dot{M} \sim 7 \times 10^{-10} M_\odot \text{ yr}^{-1} \) (\( \eta_{\text{ISM}} < 0.35 \text{ cm}^{-3} \)) (see Figure 4) for both SNe, when \( \epsilon_e = \epsilon_B = 0.1 \) and \( v_e = 100 \text{ km s}^{-1} \). This shows that our upper limits on \( \dot{M} \) are consistent with the previous findings by Chomiuk et al. (2012) and Pérez-Torres et al. (2014). However, our upper limit on \( \eta_{\text{ISM}} \) is almost an order of magnitude smaller than those reported from previous radio analyses. The reason is that, to derive the limits, Chomiuk et al. (2012) and Pérez-Torres et al. (2014) used the upper limits on radio luminosity measured around 20 (for SN 2011fe) and 35 (for SN 2014J) days after the explosion, which were the best possible options at that time. In our present study we obtained the upper limits from radio observations carried out around 4 and 1.5 years after the explosions of SNe 2011fe and 2014J, respectively. Therefore, this study provides a deeper and more confident limit compared to the previous ones. The limits on \( \dot{M} \) reported from X-ray analyses are midway between our predictions from \( \epsilon_B = 0.1 \) and 0.01.

5. Discussion

SN 2006X, a normal SN Ia, showed a notable presence of circumstellar material in the form of time evolution of a blueshifted absorption feature in the Na I D line (Patat et al. 2007). This SN was radio-silent despite prompt attempts to observe the SN in radio wavelengths (Stockdale et al. 2006).

We note that the optical and radio observations were not conducted simultaneously. However, given the fairly large distance to SN 2006X (\(~16 \text{ Mpc})\), it could be difficult to detect radio emission from it if magnetic field amplification was low in SN shocks.

No variations in Na I D were reported for SN 2011fe (Patat et al. 2013) or for SN 2014J (Goobar et al. 2014). The lack of narrow-line absorption could be due to a viewing angle effect, allowing for incomplete shells to possibly exist around those SNe. How frequently one could expect to detect such shells cannot be predicted in advance, unless proper hydrodynamical modeling of the progenitor evolution and shock CSM-shell interaction are done. With our observations we are probing a region with a radius of a few \( 10^{17} \text{ cm} \) around both the SNe. The radio silence feature of these SNe could be due to no shells existing around them within the above-mentioned radius. Or it may be due to low amplification efficiency of magnetic fields in the shocks. We try to explore the second possibility with our modeling of the radio observations, tabulated in Tables 1 and 2 for SNe 2011fe and 2014J, respectively.

As discussed in Section 4 the evolutions of radio luminosities (Equations (5) and (6), and Figure 4) depend on the nature of the circumbinary medium that carries the signature of the progenitor system. The progenitors of both SNe 2011fe and 2014J are not clear. However, observational evidence has narrowed down the possible formation channels for both of them. From pre-explosion HST images Kelly et al. (2014) have excluded a bright red giant (RG) companion and a symbiotic recurrent nova system, possessing similar luminosity of RS Oph, as a progenitor system for SN 2014J. For SN 2011fe, Li et al. (2011) excluded a RG/helium (He) star/main-sequence (MS) with a mass of more than 3.5 \( M_\odot \) as a companion. From UV observations Brown et al. (2012) precluded a solar mass MS companion with a radius of 0.1 \( R_\odot \) (Bloom et al. 2012), although their model suffers from a viewing angle effect. Combining these findings with non-detections of binary companion materials in optical nebular spectra of the two SNe, Lundqvist et al. (2015) concluded that hydrogen-rich main-sequence donor systems, with separations

\[ \text{From the NASA/IPAC Extragalactic Database (NED)}. \]
less than $\sim 7 R_\odot$ between the binary systems, could be ruled out for SN 2011fe, while there is still some room for this for SN 2014J. In addition, He star donors are constrained to those with large separations (more than $\sim 1 R_\odot$) between themselves and the WD.

These studies favor a DD formation channel for both the SNe. However, a few SD channels are still allowed. For these formation channels we will now examine what upper limits on $n_{\text{ISM}}$ and $\dot{M}$ our models, with $\epsilon_B = 0.1$ and 0.01, predict, and what is expected for that kind of progenitor.

5.1. DD Scenario

The merger of two WDs predicts a low gas density, which is characteristic of the interstellar medium, around the SN site. The H I column density, $N_{\text{H I}}$, around SN 2011fe, is $3 \times 10^{20}$ cm$^{-2}$, as estimated by Chomiuk et al. (2012). With a path length ($l$) of $\sim 100$ pc and mean atomic weight $\mu \approx 1.24$ in M101 (Stoll et al. 2011; Mazzali et al. 2014), the particle density around the location of SN 2011fe is $\mu N_{\text{H I}}/l \sim 1$ cm$^{-3}$. From our modeling with the merger model, the null detection gives an upper limit on the particle density $n_{\text{ISM}} \sim 1$ cm$^{-3}$ (left panel of Figure 4) when $\eta_B < u_e$ with $\epsilon_B = 0.01$. For $\epsilon_B = 0.1$ the upper limit on the particle density is $\sim 0.25$ cm$^{-3}$.

SN 2014J occurred in M82. Ritchey et al. (2015) estimated $N_{\text{HI}}$, in the direction of SN 2014J from 21 cm line observations. They obtained $N_{\text{H I}} = 2.9 \times 10^{21}$ cm$^{-2}$ for the gas around the SN site and compared it with the $N_{\text{HI}}$ computed from the equivalent width, contributed by M82 along the SN line of sight, of the blended $\lambda 5780.5$ diffuse interstellar band (DIB). From the DIB they derived $N_{\text{H I}} = 2.24 \times 10^{21}$ cm$^{-2}$, which is close to that derived from the 21 cm emission. The projected distance of SN 2014J from the center of M82 is $\approx 913$ pc, for a distance to M82 of 3.4 Mpc (see Figure 2). Combining this with the H I column density map shown in the upper left panel of Figure 7 of Leroy et al. (2015), we estimate $l \sim 1$ kpc for the atomic hydrogen gas around SN 2014J in M82. Here we assume that the SN is located at the middle of the disk. With $N_{\text{H I}} \approx 2.9 \times 10^{21}$ cm$^{-2}$, along with the slightly sub-solar metallicity of M82 (Origlia et al. 2004), this implies a particle density of $\sim 1$ cm$^{-3}$ around the SN location. The left panel of Figure 4 shows that the radio non-detection can be well accounted for with the above expected ISM density by assuming $\epsilon_B = 0.01$. The effect of a higher value of $\epsilon_B (=0.1)$ on $n_{\text{ISM}}$ is similar to that for SN 2011fe.

5.2. SD Channels

Considering the observational constraints on the progenitor systems of both the SNe, a He star with large separation ($\gtrsim 1 R_\odot$) could be suggested as a binary companion to the WD. For SN 2014J a hydrogen-rich MS may also be possible, although the separation to the WD is likely to have been $\gtrsim 5 R_\odot$ (Lundqvist et al. 2015). In these cases it is expected that before explosion mass has been transferred from the secondary to the WD via Roche Lobe overflow. As mentioned in Section 4 and shown in the right panel of Figure 4, for SNe 2011fe and 2014J, the upper limit on $M$ is $< 10^{-9}$ ($v_w/100$ km s$^{-1}$)$^{-1}$ $M_\odot$ yr$^{-1}$ for $\epsilon_B = 0.1$. This implies that before the explosion the WD accreted matter with an efficiency of more than 99%. However, after studying radio observations of 27 SNe Ia, Panagia et al. (2006) suggested that in cases with Roche Lobe overflow, the accretion efficiency could be 80% at most. Therefore, for this kind of companion, a rather dense medium is expected with a higher value of $M/v_w$. From our model with $\epsilon_B = 0.01$ we estimate $M < 4 \times 10^{-9}$ ($v_w/100$ km s$^{-1}$)$^{-1}$ $M_\odot$ yr$^{-1}$ for both the SNe. This implies a less restricted mass accretion efficiency of the progenitor WD in comparison to the case with $\epsilon_B = 0.1$.

Within the SD scenario two other possibilities exist, namely the spin-up/down model and recurrent novae. We discuss them next.

5.2.1. Spin-up/down Model

As discussed in Hachisu et al. (2012a, 2012b), rigidly rotating WDs within the mass range $1.38 M_\odot < M_{\text{WD}} \lesssim 1.5 M_\odot$, and with a MS or RG companion, can result in normal SNe Ia. For such WDs the spin-down time, determined by the magnetodipole radiation, can be $> 10^9$ years when they have low magnetic fields and explode with low angular momentum (Ilkov & Soker 2011). Therefore, the MS or the RG companion may evolve into a helium WD (in cases with RG, a C/O WD could also possibly form, as the evolution timescale for RG to form a C/O degenerate core is $< 10^8$ years). This spin-down time is enough to allow the CSM, created by the mass-loss from the progenitor...
system, to diffuse into the ISM. Therefore, the primary WD explodes in an ISM-like environment. In such a system, the companion could also have a negligible imprint on optical nebular spectra (Jaschek 2011), even if it did not yet evolve into a WD (Lundqvist et al. 2015).

We examine the N100 model for SN 2014J, in a constant density medium, considering $\epsilon_B = 0.1$ and 0.01. N100 is an explosion model for a non-rotating WD, in contrast to the spin-up/down model that takes into account WD rotation. However, we expect rotating WDs in the above-mentioned mass range to have an overall similar ejecta structure to that of N100, as both of them could account for normal Type Ia. Therefore, it can be expected that the prediction from N100 could be close to that from a spin-up/down model. From Figure 4 (left panel) the upper limit on the surrounding density is $\sim 3 \text{ cm}^{-3}$ for $\epsilon_B = 0.01$, which is similar to what could be expected around SN 2014J (see Section 5.1). For $\epsilon_B = 0.1$, the limit is $\sim 0.4 \text{ cm}^{-3}$.

A similar spin-up/down model is applicable to SN 2011fe (Hachisu et al. 2012b). For $\epsilon_B = 0.1$, and using our radio observations from 1468 days, the upper limit on $n_{\text{ISM}}$ is $\sim 0.3 \text{ cm}^{-3}$. However, from SN evolution with $\epsilon_B = 0.01$, we note that, in our models, around 3 years after the explosion the reverse-shock starts to move inside the steep outer density profile of the ejecta. This implies that the self-similar solution, which is applicable as long as the SN is in the free-expansion phase, is no longer valid. The SN then evolves into the Sedov–Taylor (ST) phase. The duration of free expansion can be written as $T_{\text{FE}} = \tau_{\text{ST}}^2 \frac{\rho_{\text{ej}}}{\rho_{\text{IS}}} \left( \frac{\zeta}{\zeta_{\text{ec}}} \right)^{\frac{3}{2}}$, where $\beta = \frac{4\pi}{\theta} \left( \frac{3}{5} - \frac{2}{7} \right)$. Here, $\beta = \dot{M}/\nu_w (n_{\text{ISM}})$ for a wind (constant density) medium, $\theta$ is the ratio between swept-up ejecta and circumstellar mass, $\rho_{\text{ej}}$ and $\rho_{\text{ej}}$ represent the density and velocity of the extreme outer part of the inner ejecta. $\zeta = r_e/r_c$ and $\zeta = r_{\text{IS}}/r_{\text{IS}}$, with $r_e$ and $r_c$ being the contact discontinuity and reverse-shock radii. $r_{\text{IS}}$ represents a reference radius and this is $10^{15} \text{ cm}$ in our case. It is found that for a wind medium $T_{\text{FE}} > 100 \text{ years}$ for the range of mass-loss rates we consider here. However, in case of a constant density medium, the ST phase starts much earlier. Using the merger model we estimate that $T_{\text{FE}}$ is around 8 and 5 years for SN 2014J (SN 2011fe) when $\epsilon_B = 0.1$ and 0.01, respectively. For N100, $T_{\text{FE}}$ is comparably shorter for SN 2014J (SN 2011fe), around 5 (6) years for $\epsilon_B = 0.1$. When $\epsilon_B = 0.01$, the ST phase starts $\sim 3$ years after the explosion of both the SNe with $n_{\text{ISM}} \sim 1 \text{ cm}^{-3}$.

For SN 2011fe, $n_{\text{ISM}}$ is constrained by the radio luminosity measured at 1468 days, during which the SN is in the ST phase for $\epsilon_B = 0.01$. Modeling synchrotron emission from this phase is beyond the scope of this paper. This therefore limits our conclusions for the spin-up/down scenario for SN 2011fe when $\epsilon_B = 0.01$.

### 5.2.2. Recurrent Novae

A near-Chandrasekhar mass WD, accreting at a rate between $10^{-7}$–$10^{-12} \text{ M}_\odot \text{ yr}^{-1}$, experiences nova eruptions with a recurrent time, $t_{\text{rec}}$, of a few to a hundred-thousand years (Yaron et al. 2005). The maximum expansion velocity of the ejected shells is in the range $1000$–$5000 \text{ km s}^{-1}$. This nova shell sweeps up the CSM material and a cavity is formed around the progenitor system. If $t_{\text{rec}}$ is long, the cavity is refilled by the new stellar wind. In this case one would expect Equation (6) to govern the radio luminosity evolution. For short

$r_{\text{rec}}$, the cavity will be partially refilled. Therefore, the SN ejecta first encounter a wind medium, then move into a cavity, and finally interact with a nova shell. There could be multiple nova shells present in the circumbinary medium. The evolution of radio synchrotron spectra in this scenario can only be predicted by hydrodynamical modeling, which takes into account the interaction of SN ejecta with different kinds of ambient media, formed at different times of the progenitor evolution.

Recently, Harris et al. (2016) have performed hydrodynamical modeling of an SN shock interacting with CSM shells, to predict the synchrotron emission from such a situation. The authors consider a vacuum between the exploding WD and the interacting CSM shells. In their model the radio light curve has a plateau when the SN ejecta with a velocity $2 \times 10^3 \text{ km s}^{-1}$ interact with two nova shells. This feature is distinctly different from that of the SN interacting with a single nova shell, for which a peak in the light curve is predicted. Harris et al. (2016) show that when the ejecta velocity is increased to $3 \times 10^4 \text{ km s}^{-1}$, the peak radio luminosity decreases. Along with shock velocity, the density and the thickness of the shell, which are both poorly constrained, play important roles in predicting the synchrotron emission from the SN. As several less constrained physical parameters are involved here, from our models it is difficult to predict the expected radio luminosity for this kind of system. However, from very early X-ray non-detection Margutti et al. (2014) have precluded a recurrent nova system, with a recurrence time of 100 yr (Dimitriadis et al. 2014), for SN 2014J. This is also true for SN 2011fe, as this SN was not seen in early X-ray observations as well (Margutti et al. 2012). Moreover, given the constraints on possible SD binary companions (cf. Lundqvist et al. 2015), a nova-shell-like scenario may not be applicable for these SNe, as is also indicated by the absence of NaI D variations (Putat et al. 2013; Goobar et al. 2014). The only absorption-line feature detected for these SNe is in K I λ7665 (Graham et al. 2015; Maeda et al. 2016), but the absorbing gas is at distances of at least $10^{19} \text{ cm}$ from the SN, and could be of interstellar origin (Maeda et al. 2016).

### 6. Conclusions

We have observed the SNe Ia 2011fe and 2014J in the radio at late epochs, up to 1468 and 578 days post-explosion, respectively. We do not detect the SNe, and model the radio non-detections assuming synchrotron emission from shock-accelerated relativistic electrons to be the origin of the radiation. We consider that all the electrons in the post-shock region are accelerated and their energy distributions follow a power-law profile with $p = 3$. It is assumed that 10% of bulk shock energy is converted to electrons, i.e., $\epsilon_e = 0.1$. After a certain time, as the shock slows down, a part of the electron population will have kinetic energy less than the rest-mass energy of electrons. We calculate the fraction of energy that goes into these non-relativistic electrons and estimate the synchrotron flux from the energy available to relativistic electrons. This keeps us from underestimating the density of the ambient medium.

Both SNe 2011fe and 2014J are normal Type Ia SNe. We have therefore used the inner density structure of the N100 model (for the SD channel) and a violent merger model (for the DD channel), which could both account for normal SNe Ia, and added a power-law density profile, with an index in the range 12–14, for the outer part of the ejecta. Considering that the
interaction of SN ejecta with the CSM follows a self-similar structure, we compare the radio light curves predicted by our models with observational upper limits. This gives us upper limits on the ambient medium density, i.e., $\eta_{\text{ISM}}$ for a constant density medium and $M/\nu_w$ for a wind medium. Since synchrotron power is proportional to the magnetic field energy density, the amplification of this field, in the shock, plays an important role in determining the synchrotron luminosity, and hence on the prediction of the ambient medium density.

The circumbinary medium is shaped by the mass-loss history of the progenitor system. Although the kind of binary systems that result in these two SNe are not clear, studies based on observational evidence (Li et al. 2011; Bloom et al. 2012; Brown et al. 2012; Kelly et al. 2014; Lundqvist et al. 2015) have given indications about the formation channels of these two SNe. Among these, the DD scenario is one of the possibilities. From the H I column density observed around both the SNe, $\eta_{\text{ISM}}$ could be expected to be $\sim1\,\text{cm}^{-3}$. To estimate $\eta_{\text{ISM}}$ from the H I surface density, we assume path lengths of 100 pc and 1 kpc for the atomic hydrogen gas in M101 and M82 around SNe 2011fe and 2014J, respectively. From our modeling we found that (left panel of Figure 4) the radio non-detections post one year can be explained with $\eta_{\text{ISM}} \sim 1\,\text{cm}^{-3}$ when the magnetic field amplification is low ($\epsilon_B = 0.01$), compared to what is expected for equipartition of energy (i.e., $\epsilon_B = \epsilon_e = 0.1$).

As discussed in Section 5.2, an MS companion may be possible for SN 2014J in the SD scenario, whereas this is less likely for SN 2011fe. A He star companion is easier to accommodate for both SNe. In this channel, we found that the upper limits on $M/\nu_w$ predicted by our model, considering a low value of $\epsilon_B$, could allow for significant mass-loss (see right panel of Figure 4). For $\epsilon_B = 0.1$ and 0.01, the upper limits on $M$ for both the SNe are $\sim 7 \times 10^{-10}$ and $\sim 4 \times 10^{-9} M_\odot\,\text{yr}^{-1}$ when $\nu_w = 100\, \text{km}\,\text{s}^{-1}$. This suggests that in the former case the mass accretion efficiency of the WD needs to be more than 99%, whereas the latter puts a less constraining limit on accretion efficiency. From a radio study of 27 Type Ia SNe, Panagia et al. (2006) found a maximum accretion efficiency of 80% for this kind of progenitor system.

Another possibility within the SD channel is recurrent novae. In this case, the wind may be swept up in shell-like structures prior to the SN explosion. Our models cannot predict the radio emission for this scenario, as the SN could interact with different kinds of ambient media depending on its evolution. To predict the synchrotron luminosity from this kind of system, proper hydrodynamical modeling is required. However, considering the absence of Na I D variation, early non-detections in X-rays, as well as no indications of hydrogen and metals from an SD companion in optical nebular spectra of the two SNe, recurrent novae may be less likely progenitor systems for these SNe.

A particularly interesting case for the SD scenario is instead the spin-up/down model, which suggests a density around the progenitor system similar to the ISM density. We have tried to explore this channel, for SN 2014J, by considering an SN with a density profile of the N100 model for the inner part of the ejecta that expands into a constant density medium. The upper limit on $\eta_{\text{ISM}}$ predicted by this model (left panel of Figure 4) is $\sim 3\,\text{cm}^{-3}$ when $\eta_B < \eta_e$ with $\epsilon_B = 0.01$.

Modeling of radio and X-ray emissions from core-collapse SNe exhibit a broad range for both $\epsilon_B$ and $\epsilon_e$. From detailed radio and X-ray modeling of SN 2002ap, Björnsson & Fransson (2004) found that the energy density in the electric field is comparable to that of magnetic fields. This was not true for SN 1993J, for which Fransson & Björnsson (1998) deduced $\epsilon_B/\epsilon_e \sim 280$, with $\epsilon_B \sim 0.14$ and $\epsilon_e \sim 5 \times 10^{-3}$. In the case of SN 2011dh it is found that the radio emission can be interpreted by SSA, assuming $\epsilon_B = \epsilon_e = 0.1$ (Soderberg et al. 2012; Krauss et al. 2012). However, to explain the X-ray emission from the same radio-emitting electron population one requires $\epsilon_B = 0.01$ and $\epsilon_e = 0.3$ (Soderberg et al. 2012). Similarly, for SN 2013df, Kamble et al. (2016) estimated a much smaller energy density in magnetic fields compared to that in electric fields ($\epsilon_B/\epsilon_e \sim 5 \times 10^{-3}$).

We found that both SNe 2011fe and 2014J the upper limit on $M/\nu_w$ has a rather strong dependence on $n$, which varies in the range 12–14, compared to that of $\eta_{\text{ISM}}$ (Figure 4). However, it is clear from Figure 4 that $\epsilon_B$ plays the most important role in determining the upper limits on the CSM density ($\eta_{\text{ISM}}$ or $M/\nu_w$) around both the SNe. From particle in cell (PIC) simulations, Caprioli & Spitkovsky (2014b) showed that for shocks with a Mach number ($M_A$) up to 100, the magnetic field amplification is proportional to the square root of $M_A$ (see their Equation (2) and the bottom panel of Figure 5). As $u_\text{e} \propto M_A^2$, for shocks with $M_A \sim 100$, $u_\text{e} \propto M_A$. If this is true for shocks with $M_A \sim$ few $\times 1000$, which is applicable to SN shocks, then this implies that although the magnetic fields could be amplified very efficiently in the post-shock region, $u_B$ will be much smaller than $u_\text{e}$.

In summary, our study shows that the null detection of radio emission from SNe 2011fe and 2014J could be due to the fact that $u_B < u_\text{e}$ with $\epsilon_B < 0.1$ in the post-shock region. There is evidence of circumbinary material around a few Type Ia SNe (e.g., SNe 2002ic, 2005gj, 2006X; Hamuy et al. 2003; Aldering et al. 2006; Patat et al. 2007), however, no radio emission has been detected from them (Berger et al. 2003; Stockdale et al. 2003; Soderberg & Frail 2005; Stockdale et al. 2006). This may indicate that $u_B < u_\text{e}$ with $\epsilon_B < 0.1$ could be a general feature of SN shocks. Future detailed periodic radio observations from nearby SNe, together with results from PIC simulations for $M_A \sim 1000$, could shed light on this.

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