Novae as Tevatrons: Prospects for CTA and IceCube

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ABSTRACT
The discovery of novae as sources of ~ 0.1 – 1 GeV gamma-rays highlights the key role of shocks and relativistic particle acceleration in these transient systems. Although there is evidence for a spectral cut-off above energies ~ 1 – 100 GeV at particular epochs in some novae, the maximum particle energy achieved in these accelerators has remained an open question. The high densities of the nova ejecta (~ 10 orders of magnitude larger than in supernova remnants) render the gas far upstream of the shock neutral and shielded from ionizing radiation. The amplification of the magnetic field needed for diffusive shock acceleration requires ionized gas, thus confining the acceleration process to a narrow photo-ionized layer immediately ahead of the shock. Based on the growth rate in this layer of the hybrid non-resonant cosmic ray current-driven instability (considering also ion-neutral damping), we quantify the maximum particle energy, \( E_{\text{max}} \), across the range of shock velocities and upstream densities of interest. We find values of \( E_{\text{max}} \sim 10 \) GeV - 10 TeV, which are broadly consistent with the inferred spectral cut-offs, but which could also in principle lead to emission extending to higher energies \( \sim 100 \) GeV accessible to atmosphere Cherenkov telescopes, such as the planned Cherenkov Telescope Array (CTA). Detecting TeV neutrinos with IceCube in hadronic scenarios appears to be more challenging, although the prospects are improved for a particularly nearby event (distance \( \lesssim \) kpc) or if the shock power during the earliest, densest phases of the nova outburst is higher than implied by the observed GeV light curves, due to downscattering of the gamma-rays by electrons within the ejecta. Novae provide ideal nearby laboratories to study magnetic field amplification and the onset of cosmic ray acceleration, because other time-dependent sources (e.g. radio supernovae) typically occur too distant to detect as gamma-ray sources.

Key words: keywords

1 INTRODUCTION

Novae are sudden visual outbursts powered by runaway nuclear burning on the surface of a white dwarf accreting from a stellar companion (e.g. Starrfield et al. 1998). They reach peak luminosities of \( \sim 10^5 L_\odot \) and eject moderate quantities of mass \( \sim 10^{-5} - 10^{-4} M_\odot \) at velocities of hundreds to thousands of km s\(^{-1}\) (e.g., Shore 2002). Although in standard models novae are thought to be powered directly by the energy released from nuclear burning (e.g. Hillman et al. 2014), growing evidence suggests that shock interaction plays an important role in powering nova emission across the electromagnetic spectrum. Evidence for shocks includes multiple velocity components in the optical spectra (e.g. Williams & Mason 2010), hard X-ray emission starting weeks to years after the outburst (e.g. Mukai et al. 2008; Osborne 2015), and an early sharp maximum in the radio light curve on timescales of months, in excess of that expected from freely-expanding photo-ionized ejecta (e.g., Chomiuk et al. 2014; Weston et al. 2015).

The most striking indicator of shocks in novae is the recent discovery by Fermi LAT of \( \gtrsim 100 \) MeV gamma-rays, observed at times coincident within a few days of the optical peak and lasting a few weeks (Ackermann et al. 2014). The first nova with detected gamma-rays occurred in the symbiotic binary V407 Cyg (Abdo et al. 2010), suggesting that shocks were produced by the interaction between the nova outflow and the dense wind of the companion red giant. However, gamma-rays have now been detected from at least five ordinary classical novae with main sequence companions (Ackermann et al. 2014). Remarkably, this demonstrates that the nova outflow runs into dense gas even in systems not embedded in the wind of an M giant or associated with recurrent novae. This dense gas instead likely represents lower velocity mass ejected earlier in the outburst (‘internal shocks’; Mukai & Ishida 2001; Metzger et al. 2014). Current observations are consistent with many, and possibly all, novae producing shocks and \( \gtrsim 100 \) MeV gamma-ray emission. Indeed, the LAT-detected novae appear to be distin-
Figure 1. Regimes of shock conditions in different astrophysical environments, in the space of shock velocity $v_{sh}$ and density $n$. Shocks in novae occupy a distinct portion of parameter space of relatively low shock velocities and high densities. Shown for comparison are supernova shock break-out from red supergiant (RSG) and blue supergiant (BSG) progenitor stars, gamma-ray burst (GRB) afterglows, radio supernovae, supernova remnants, the intergalactic medium (IGM), interplanetary medium (IPM) on scales of 1 AU, and shocks in the outflows from binary white dwarf mergers (Beloborodov 2014).

Nova gamma-rays are produced by the decay of neutral pions created by proton-proton collisions, or via Inverse Compton scattering or bremsstrahlung emission from energetic electrons. As both of these scenarios require a population of relativistic particles, these events provide a real-time probe of relativistic particle acceleration in non-relativistic shocks, complementary to those found in other astrophysical environments such as supernova remnants (e.g. Martin & Dubus 2013; Metzger et al. 2015). Constraints on the efficiency and energy spectra of particle acceleration as probed by novae provides a new and promising avenue to address old questions, such as the sources of Galactic cosmic rays and magnetic field amplification at shocks.

Although the velocities of shocks in novae are similar to those in older supernova remnants, the density of the shocked gas in novae is typically 10 orders of magnitude higher (Fig. 1). Such high densities result in several novel physical effects, such as the importance of radiative cooling on the post-shock dynamics (‘radiative’ shocks). Furthermore, the gas both upstream well ahead of the shock and in the downstream cooling layer is largely neutral and hence will absorb thermal soft X-ray/UV radiation from the shock, reprocessing the shock power to optical frequencies (Metzger et al. 2014). Indeed, a comparison between the minimum shock power needed to explain the observed gamma-ray luminosity and the optical light curve shows that a significant fraction of the latter is shocked-powered (Metzger et al. 2015), analogous to interacting supernovae (e.g. Chevalier & Irwin 2011).

The high densities of nova shocks also result in several physical simplifications. The timescale for energy exchange between electrons and protons behind the shock via Coulomb scattering is short compared to the dynamical or thermal cooling rate, justifying a single temperature plasma (Metzger et al. 2015). The high matter and photon densities near the shock furthermore ensure that relativistic particles radiate their energy rapidly compared to the outflow expansion time (fast cooling regime). Gamma-ray producing novae therefore act as “calorimeters” for probing the efficiency and spectrum of relativistic particle acceleration, in ways comple-
mentary to studies of supernova remnants (e.g. Morlino & Caprioli 2012). Exploiting the radiative nature of nova shocks, Metzger et al. (2015) place a minimum of \( f_{\text{in}} \geq 10^{-3} - 10^{-2} \) on the fraction of the kinetic power of the shock placed into relativistic particles. Based on observations of supernova remnant shocks (e.g. Volk et al. 2005), and the results of MHD kinetic hybrid (Caprioli & Spitkovsky 2014a) and particle-in-cell (Kato 2014; Park et al. 2014) simulations, these high efficiencies appear to favor a hadronic origin for the GeV gamma-ray emission.

This paper addresses a key question: what is the highest energy particle accelerated in nova shocks? This issue is crucial to assessing novae as potential targets for present and future TeV telescopes, such as the Cherenkov Telescope Array (CTA; Actis et al. 2011), and current neutrino observatories, such as IceCube (Karle et al. 2003). Although the statistics are poor, the time-integrated spectral energy distribution measured by Fermi LAT shows hints for a spectral cut-off or break in the photon energy range \( E_{\gamma} \sim 1-10 \) GeV (Ackermann et al. 2014). MAGIC Collaboration et al. (2015) presented upper limits on \( \Delta f_{\text{in}} = \text{eqs. [12]} \) and temperature \( T_{\text{sh}} \sim 1 - 2 \times 10^4 \) K in which \( f_{\text{in}} \approx 10^{-4} - 10^{-3} \) (eq. [14]), the latter set by the balance between ionization and recombination. Diffusive particle acceleration is limited to a fraction of the ionized layer \( \Delta \Omega_{\text{ion}} \), which is itself generally much narrower than the ejecta thickness \( R_{\text{ej}} \) or post-shock cooling length \( L_{\text{cool}} \).

A naive application of the Hillas (1984) confinement criterion, using the entire radial extent of the nova ejecta as the extent of the accelerator, results in a maximum particle energy of \( \sim 1 \) PeV (Metzger et al. 2015). However, this estimate is unrealistic because the magnetic field required for particle acceleration via diffusive shock acceleration cannot be supported in neutral regions (comprising the bulk of the ejecta) due to ion-neutral damping. This limits the effective size of the acceleration zone to at most a narrow photo-ionized layer ahead of the shock, substantially reducing the maximum particle energy to a lower value (consistent with current observations). Nevertheless, here we will show that particle energies exceeding \( \sim 1 \) TeV can readily be achieved under physically reasonable conditions during the nova outburst.

This paper is organized as follows. In §2 we briefly review the properties of nova shocks to motivate the relevant physical conditions, including the radial extent of the photo-ionized layer. In §3 we estimate the maximum particle energy in nova shocks by means of a simple stability analysis, based on the growth of the non-resonant (Bell 2004) modes, considering also the role of ion-neutral damping. In §4 we discuss the implications of our results for novae as TeV photon (§4.1) and neutrino (§4.2) sources and for magnetic field amplification in non-relativistic shocks (§4.3). We briefly summarize our results and conclude in §5.

2 Nova Shocks

2.1 Shock Dynamics

Shocks were unexpected in classical novae because the pre-explosion environment surrounding the white dwarf is occupied only by the low density wind of the main sequence companion star, requiring a different source of matter into which the nova outflow collides. One physical picture, consistent with both optical, Schaefer et al. (2014) and radio imaging (e.g. Chomiuk et al. 2014), and the evolution of optical spectral lines (e.g. Ribeiro et al. 2013, Shore et al. 2013), is that the thermonuclear runaway is first accompanied by a slow ejection of mass with a toroidal geometry, the shape of which may be influenced by the binary companion (e.g., Livio et al. 1990; Lloyd et al. 1997). This slow outflow is then followed by a second discrete ejection or continuous wind (e.g., Bath & Shaviv 1976) with a higher velocity and more spherical geometry. The subsequent collision between the fast and slow components produces strong “internal” shocks within the ejecta which are concentrated in the equatorial plane. The fast component continues to expand freely along the polar direction, creating a bipolar morphology (see Fig. 1 of Metzger et al. 2015 for a schematic diagram).

The slow outflow of velocity \( v_{\text{ej}} \) \( \approx 10^2 v_{\text{ej}} \) km s\(^{-1}\) expands to a radius \( R_{\text{ej}} = v_{\text{ej}} t \approx 6 \times 10^3 f_{\text{sh}} v_{\text{ej}} \) cm by a time \( t = f_{\text{sh}} \) week. The characteristic density of the ejecta of assumed thickness \( R_{\text{ej}} \) can be estimated as

\[
n_{\text{ej}} \approx \frac{M_{\text{ej}}}{4 \pi R_{\text{ej}}^3 f_{\text{sh}} m_p} \approx 9 \times 10^{10} M_{\odot} f_{\text{sh}}^{-1} v_{\text{ej}}^{-3} \text{ cm}^{-3},
\]

where \( f_{\text{sh}} \approx 0.5 \) is the fraction of the total solid-angle subtended by the outflow and \( M_{\text{ej}} = 10^{-4} M_{\odot} M_{\odot} \) is the ejecta mass, normalized to a characteristic value (e.g. Seagrass et al. 1980).

A faster outflow (“wind”) of mass loss rate \( M \) and velocity \( v_1 \approx 2 v_{\text{ej}} \) collides with the ejecta from behind. The density of the
wind at the collision radius (\(r \approx r_{\text{sh}}\)) is given by
\[
n_{\text{sh}} = \frac{\dot{M}}{4\pi R_{\text{sh}}^2 \rho v_{\text{sh}}} \sim 2 \times 10^8 M_{-3} \nu_{10}^2 r_{\text{sh}}^{-2} \text{cm}^{-3},
\]
where \(\dot{M} = 10^{-5} M_{-5} M_{\odot} \text{yr}^{-1}\) is normalized to a value resulting in the ejection of \(\sim 10^{-3} M_{\odot}\) over a week. For instance, a total mass \(4 \times 10^{-3} M_{\odot}\) was ejected in the "fast" component of V959 Mon (Chomiuk et al. 2014).

The shock properties are constrained to produce the range observationally unresolved at the time of the gamma-ray emission. The discussion above provides estimates for the characteristic value of \(r_{\text{sh}} \sim v_{\text{sh}} t_{\text{rec}}\).

### 2.2 Ionized Layer

Absent external sources of photo-ionization, the unshocked upstream gas is neutral. The recombination timescale \(t_{\text{rec}} \sim 1/\sigma_{\text{rec}}^2\),
\[
t_{\text{rec}} \sim 4 \times 10^{-10} \nu_{10}^2 \nu_{\text{us}}^{-2} t_{\text{sh}}^{-1}
\]
is short compared to the expansion time, where \(\sigma_{\text{rec}}^2 \approx 4 \times 10^{-13} Z^2 (T/10^7 \text{K})^{-0.8} \text{cm}^{-1} \text{s}^{-1}\) is the approximate radiative recombination rate for hydrogen-like atomic species of charge \(Z\) and \(t_{\text{sh}} = T_{\text{sh}} / n_{\text{sh}}\) is the temperature of the unshocked gas. Absent photo-ionization heating, the shielded neutral gas will cool to a temperature \(T_{\text{sh}} < 10^3 \text{K}\), comparable or less than the effective temperature of the nova emission, for which \(f_{\nu} \sim \Omega(1)\) in thermal ionization balance (Saha equilibrium).

The upstream is subject to ionizing UV and X-ray radiation from the shocks, which penetrates the upstream medium to a depth \(\Delta^2\) (see eq. 12 below) that depends on the element \(Z\) dominating the bound-free opacity at the frequency of relevance. The neutral fraction of this exposed layer, \(f_{\text{us}} = (1 + \lambda_{\text{us}})^{-1}\), is set by the balance between ionization and recombination, where
\[
\lambda_{\text{Z}} = \frac{4\pi}{v_{\text{sh}}^2} \int \frac{\sigma_{\text{rec}}^2}{\nu^2} d\nu \approx \frac{L_{\text{ion}}}{2 \pi \nu^2 n_{\text{sh}}^2 r_{\text{sh}}^2} \frac{t_{\text{sh}}}{J_{\nu}}
\]
\[
\approx \frac{\sigma_{\text{rec}}^2 v_{\text{sh}}}{2 \pi \nu^2 (1 + 5/2)} \left( \frac{J_{\nu}}{J_{\nu}^*} \right) \cdot \frac{8 \times 10^{-18} Z^2}{1 + 5/2} \text{cm}^{-2} Z^2
\]
\[
J_{\nu} = 13.6 Z^2 \text{eV}^{-1}
\]

Here \(J_{\nu}^*\) is the average intensity of ionizing radiation, \(n_{\text{sh}} \sim n\) is the free electron density in the ionized layer (which is nearly fully ionized; see eq. 14 below), \(\sigma_{\text{rec}}^2 \approx \sigma_{\text{ REC}}^2 (\nu/\nu_{\text{th}})^{-3}\) is the approximate bound-free cross section, and \(\sigma_{\text{rec}}^2 \approx 6 \times 10^{-18} Z^2 \text{cm}^{-2}\) is the approximate cross section at the ionization threshold frequency \(h\nu_{\text{th}} = 13.6 eV\).

Equation 14 shows that \(J_{\nu} \gg 1\) (\(J_{\nu} \ll 1\)) for light elements with \(Z \lesssim 5\). Radiation of frequency \(\nu \sim v_{\text{sh}}\) penetrates the neutral gas to the depth \(\Delta^2\) where the absorbptive optical depth equals unity,
\[
(n A) \sim 2 \times 10^{18} n_{\text{sh}} \nu_{10}^2 \nu_{\text{us}}^{-2} \text{cm}^{-2}
\]
\[
\Delta_{\text{abs}} = \Delta_{\text{abs}} \sim 4 \times 10^5 \text{cm}^{-2} \left( \frac{v_{\text{sh}}}{1 + 5/2} \right) \frac{1}{n_{10}^2}
\]
\[
A_{\text{Z}} = A_{\text{Z}} \sim 10 \text{erg s}^{-1}
\]
\[
A_{\text{Z}} = \frac{A_{\text{Z}}}{\text{erg s}^{-1}} \text{erg cm}^{-2}
\]
where we have taken $X_{\text{ion}} = 0.5$. This is typically 2–3 orders of magnitude smaller than the characteristic radius of the ejecta at the time of the shocks, $R_0 \sim 10^{13}$ cm. The residual neutral fraction in the ionizing layer is approximately

$$f_n \approx \frac{X_{\text{ion}}}{X_{\text{ion}} + 1} \approx 1.3 \times 10^{-3} \left(1 + 5\zeta/2\right) \approx 8 \frac{L_{\text{UV}}}{0.1} r_{\text{sh}}^2.$$  

(14)

Direct free-free shock luminosity sets the minimum ionizing intensity, but $J_\nu$ can exceed $J_{\text{eff}}$ if line cooling becomes more important than free-free emission or due to ionization by secondary photons produced by the absorption and reprocessing of higher frequency radiation by the neutral gas. Such reprocessing is challenging to determine accurately without a detailed photo-ionization calculation, but a simple estimate suggests that the extent of the ionizing layer could be expanded by an order of magnitude or more. If a fraction $f_{\text{eff}}$ of the total shock luminosity is placed into H-ionizing photons of energy $\sim 2h\nu_{\text{He}}^{\text{z+1}}$, then the maximum thickness of the ionized layer is larger than the minimum value of $\Delta_{\text{ion}}$ (equation [12]) by a factor of

$$J_\nu / J_{\text{eff}} \sim 1 + f_{\text{eff}} (kT_{\text{sh}} / 2h\nu_{\text{He}}^{\text{z+1}}) \approx 8 (f_{\text{eff}} / 0.1) r_{\text{sh}}^2.$$  

(15)

### 3 PARTICLE ACCELERATION

In the standard theory of Diffusive Shock Acceleration (DSA; e.g., Blandford & Ostriker 1978), particles are accelerated to energies exceeding that of the thermal plasma by diffusing back and forth across the shock front via interaction with magnetic turbulence upstream and downstream. The momentum distribution, $f(p) \propto p^{-4}$, predicted by the simplest DSA models corresponds to an energy distribution

$$\frac{dN}{dE} \propto \left\{ \begin{array}{ll} E^{1/2}, & kT_{\text{sh}} \leq E \ll m_{\text{p}}c^2 \\
\text{constant}, & m_{\text{p}}c^2 \ll E \ll E_{\text{max}}. \end{array} \right.$$  

(16)

that concentrates the non-thermal energy in relativistic particles. Equation [16] is valid for strong shocks, which are justified because the post-shock temperature $T_{\text{sh}}$ greatly exceeds the temperature of the upstream ionized gas, $T_{\text{u}} \sim 10^4$ K. Because of the high degree of ionization of the upstream gas (equation [14]), the effects of a neutral return flux (e.g., Blasi et al. 2012) do not appear to be relevant to nova shocks.

The magnetic field of the unshocked nova ejecta will likely be weak due to flux freezing dilution from the white dwarf surface (Metzger et al. 2013). The confinement needed for relativistic particle acceleration in nova shocks thus requires significant magnetic field amplification. A promising candidate for such amplification on small scales in non-relativistic shocks is the hybrid non-resonant (NRH) cosmic ray current-driven instability identified by Bell (2004) (see also Reville et al. 2006, Blasi & Anton 2008, Zirakashvili et al. 2008, Reuel & Spitkovsky 2009, Caprioli & Spitkovsky 2014b). However, the high densities in novae imply that ion-neutral collisions can potentially damp the NRH (Reville et al. 2007), which could suppress particle acceleration in the neutral or quasi-neutral regions ahead of the shock (Fig. 2).

The growth rate of the NRH modes, accounting for ion-neutral damping, is given by (Reville et al. 2007)

$$\gamma \approx \delta \left( \frac{\sqrt{2} \rho_{\text{He}}^{1/4} + \sqrt{3} \gamma_0 - \nu_{\text{He}} / 2}{\gamma_0 \nu_{\text{He}}} \right) \left( \frac{\gamma_0}{\nu_{\text{He}}} \right)^{1/2} \approx 1,$$

(17)

where $\gamma_0 \equiv (kT_{\text{He}} / e)_{\text{He,min}}^{1/2}$ is the growth rate for fully ionized gas and $\nu_{\text{He}} \approx 90 f_\nu \Pi_0 r_{\text{sh}}^{-2} \approx 0.1$ is the ion-neutral momentum exchange frequency (Kulsrud & Cesarsky 1971). Here $\zeta \approx \epsilon_\nu \nu_{\text{He}} / c$ is a dimensionless parameter characterizing the strength of the ion current driving term, which is proportional to the fraction $\epsilon_\nu$ of the shock power used to accelerate relativistic protons, and $r_{\text{He,min}} = \Pi_0 r_{\text{sh}}$ is the gyroradius of the minimum energy cosmic rays, where $r_{\text{He,min}}$ is the minimum momentum of the current-carrying ions. The simplification made in the top line of equation (17) that $\gamma_0 \gg \gamma$ has been checked after the fact.

The minimum cosmic ray momentum $p_{\text{min}} \approx E_{\text{min}} / c$ increases with distance ahead of the shock because particles with larger energy (larger gyroradii) can diffuse further in the face of downstream advection. The minimum ion energy $E_{\text{min}}(z)$ at distance $z$ ahead of the shock is estimated by equating the upstream diffusion timescale $t_{\text{diff}} \sim D / v_{\text{sh}}^2$ to the advection time $t_{\text{adv}} \approx z / v_{\text{sh}}$, from $z$ to the shock, where $D$ is the diffusion coefficient, taking $D = r_{\text{He}}^2 / 3$, corresponding to the limit of Bohm diffusion (Caprioli & Spitkovsky 2014b), where $r_{\text{He}} = E / eB$ is the gyroradius and $B$ is the magnetic field strength, gives

$$E_{\text{min}}(z) \approx 3eBv_{\text{sh}} z / c \approx 170 \text{ GeV}$$

\begin{equation}
\frac{1}{1 + 8\zeta/2} \left( \frac{\epsilon_\nu}{\epsilon_{\text{He}}} \right)^{1/2} (T_{\text{He}} / T_{\text{sh}})^{1/2} \left( \frac{L_{\text{UV}} / 0.1 \text{ cm}}{v_{\text{sh}}^2} \right),
\end{equation}

(18)

where in the final line of equation (18), $z$ is normalized to the thickness of the H-ionized layer (eq. [15]). We have normalized the magnetic field strength, $B$, to its post-shock value

$$B_{\text{sh}} = \sqrt{8\epsilon_\nu m_{\text{He}}^2 / \rho_{\text{He}}} = 5.6v_{\text{He}} T_{\text{He}}^{1/2} / B_{\text{sh}},$$

(19)

which is parameterized as a fraction $\epsilon_{\text{He}} / \epsilon_\nu = \epsilon_{\text{He}} / 0.01$ of the post-shock pressure $P = (3/4)\mu m_{\text{He}}^2$. Strictly speaking, the growth rate should be calculated using the initial, unamplified magnetic field strength $B_0$, which corresponds to $\epsilon_{\text{He}} \approx 0.01$. However, a larger value of $\epsilon_{\text{He}} \sim 10^{-3}$ is motivated by simulations of the NRH, which show that in the non-linear stage the growth of the instability is reduced by a compensating factor of $(B / B_0)^2$ (Caprioli & Spitkovsky 2014b). Thus, although our estimate may be wrong in the first fold of the growth, it becomes better and better and converges to the growth rate in the amplified field once saturation is achieved.

Using equation (18), the growth rate (equation [17]) is now written

$$\gamma = \left\{ \begin{array}{ll} k_{\text{He}} \epsilon_{\nu \text{He}} / 3eB (3\epsilon_\nu / E_{\text{He}}) \approx 1, & f_\nu \approx 1, \\
k_{\text{He}} \epsilon_{\nu \text{He}} / 3eB (3\epsilon_\nu / E_{\text{He}})^{1/2} \approx 1, & f_\nu \approx 1. \end{array} \right.$$  

(20)

Because the growth rate increases with $k$ and particles of energy $E$ scatter most effectively off modes of wavenumber $k \leq 1 / r_{\text{He}}$, the modes with $k_{\text{He}} \sim 1$ are the most relevant ones. Such modes only have time to act if their growth time, $t_{\text{He}} = 1 / \gamma$, is less than the time it takes the disturbance to be advected downstream, $t_{\text{adv}} = z / v_{\text{sh}}$.

This ratio is given by

$$t_{\text{He}} \left|_{f_\nu \approx 1} \right. \approx \left\{ \begin{array}{ll} \frac{3\epsilon_\nu E / (\epsilon_{\nu \text{He}} eB)}{3} \approx 16 f_\nu \epsilon_{\nu \text{He}} c^2 / (8^2), & T_{\text{He}} \approx 1, \\
3E / (3eB) \approx 1, & f_\nu \approx 1. \end{array} \right.$$  

(21)

Equation (21) shows that in the far upstream neutral layer ($f_\nu = 1$), we have $t_{\text{He}} \gg t_{\text{adv}}$, making the acceleration of particles with energies $E \gg \text{GeV}$ unlikely in this region.

However, the lower value of $f_\nu \sim 10^{-3} \ll 1$ in the H-ionized layer (equ. [14]) allows for more efficient particle acceleration. Because $t_{\text{He}} / t_{\text{adv}}$ decreases with distance from the shock, the maximum
energy particle is limited by the growth rate at $z \sim \Delta_{\text{ion}}$. At this location, we have $t_{\gamma\gamma} \lesssim t_{\text{adv}}$ for particles up to an energy

$$E_{\max,1} \approx \frac{e_\gamma e_B}{3} \left(\frac{v_{\text{sh}}}{c}\right) \approx 2.2 \text{ TeV} \frac{e_\gamma}{e_B} \frac{3^{1/2}}{n_{10}^{1/2}(1 + 5x/2)} \frac{J_f}{J_{\text{f}}}.$$  \hfill (22)

Equation (22) sets the maximum particle energy only when $E_{\max,1}$ is less than the value limited by the size of the ionized layer (diffusive confinement criterion), which from equation (18) is given by

$$E_{\max,2} = E_{\text{min}}(z = \Delta_{\text{ion}}) \approx \frac{700 \text{ GeV}}{(1 + 5x/2)} 3^{1/2} \frac{v_{\text{sh}}}{n_{10}^{1/2}} \frac{J_f}{J_{\text{f}}}.$$  \hfill (23)

The fact that $E_{\max,2} = [9v_{\text{sh}}/(e_Bc)]E_{\max,1} \approx 0.3v_{\text{sh}}e_\gamma J_{\text{max}}/J_{\text{f}}$, shows that the thickness of the ionized region is generally more important than the finite NRH growth rate for fiducial parameters. $E_{\max,1}$ and $E_{\max,2}$ are increasing functions of both the shock velocity $v_{\text{sh}}$ and the amount of ionizing radiation $J_f/J_{\text{f}}$, the latter of which is expected to depend inversely on $v_{\text{sh}}$ due to the greater importance of line emission and radiation trapping/downscattering for lower shock velocities.

Finally, we check that the cosmic ray acceleration time $t_{\text{acc}} \approx 8D/c_{\text{sh}}$ (e.g., Caprioli & Spitkovsky 2014),

$$t_{\text{acc}} \approx \frac{E_c}{3eBc_{\text{sh}}^{3/2} v_{\text{sh}}} \approx 10^{-3} \text{ week} \frac{v_{\text{sh}}^{3/2}}{n_{10}^{1/2}} \frac{J_f}{E_{\text{TeV}}},$$  \hfill (24)

is much less than expansion time $t$. The acceleration time is also shorter than the timescale for pionic losses, $t_{\pi} \sim 1/(4\sigma_{\pi}\gamma_c)$,

$$t_{\text{acc}} \approx 0.014v_{\text{sh}}^{3/2} \frac{n_{10}^{1/2}}{E_{\text{TeV}}} \frac{J_f}{J_{\text{f}}}.$$  \hfill (25)

where we have taken $\sigma_{\pi} \approx 20 \text{ mbar}$ as an estimate of the p-p inelastic cross section (Kamae et al. 2006). A similar calculation shows that Coulomb losses are unimportant on the acceleration timescale. The maximum particle energy is thus set by the confinement condition, rather than the finite acceleration time.

4 IMPLICATIONS

4.1 Maximum Particle Energy

Figure 3 shows contours of the maximum particle energy $E_{\max} = \min\{E_{\max,1}, E_{\max,2}\}$ in the space of the shock velocity $v_{\text{sh}}$ and upstream density $n$. The top panel is calculated assuming photoionization ahead of the shock exclusively from the shock’s free-free emission $J_f = J_{\text{f}}$, while the bottom panel is calculated assuming that an additional fraction $f_{\text{ion}} = 0.1$ of the total shock power is available to ionize hydrogen, either via direct line emission or reprocessing of hard radiation absorbed by neutral gas (eq. (15)). A dashed black line separates the regime of adiabatic versus radiative shocks ($t_{\text{cool}} = t_{\text{exp}}$), where $t_{\text{exp}} \approx 2$ weeks is the characteristic timescale of the LAT emission and $t_{\text{cool}}$ is calculated using a standard cooling function for solar metallicity gas (Schure et al. 2009). Dashed blue lines show the gamma-ray luminosity $L_{\gamma}$ (eq. (1)), calculated assuming a shock radius $R_{\text{sh}} = v_{\text{sh}}t$, $\epsilon_{\text{sh}} = 0.1$, and $\epsilon_{\gamma} = 0.1$ (Metzger et al. 2015), bracketing the range $L_{\gamma} \sim 10^{35} - 10^{36} \text{ erg s}^{-1}$ measured by Fermi/LAT. The largest values of $E_{\max}$ at high $v_{\text{sh}}$ in the bottom panel may be unphysical because the $\gamma$ keV thermal X-rays from such high velocity shocks will directly escape instead of ionizing the upstream gas.

Figure 3. Maximum particle energy (solid red lines) accelerated at the shock as a function of the shock velocity $v_{\text{sh}}$ and upstream density $n$, calculated as the minimum of $E_{\max,1}$ (eq. 22) and $E_{\max,2}$ (eq. 23) for $e_B = 10^{-2}$ and shown for two cases: assuming (1) Top Panel: photoionization exclusively from the shock’s free-free emission $J_f = J_{\text{f}},$ (2) Bottom Panel: assuming a fraction $f_{\text{ion}} = 0.1$ of the total shock power is available to ionize hydrogen via direct line emission or reprocessing of hard radiation absorbed by neutral gas (eq. 15). A dashed black line separates the regime of adiabatic versus radiative shocks ($t_{\text{cool}} = t_{\text{exp}}$), where $t_{\text{exp}} \approx 2$ weeks is the characteristic timescale of the LAT emission and $t_{\text{cool}}$ is calculated using a standard cooling function for solar metallicity gas (Schure et al. 2009). Dashed blue lines show the gamma-ray luminosity $L_{\gamma}$ (eq. 1), calculated assuming a shock radius $R_{\text{sh}} = v_{\text{sh}}t$, $\epsilon_{\text{sh}} = 0.1$, and $\epsilon_{\gamma} = 0.1$ (Metzger et al. 2015), bracketing the range $L_{\gamma} \sim 10^{35} - 10^{36} \text{ erg s}^{-1}$ measured by Fermi/LAT. The largest values of $E_{\max}$ at high $v_{\text{sh}}$ in the bottom panel may be unphysical because the $\gamma$ keV thermal X-rays from such high velocity shocks will directly escape instead of ionizing the upstream gas.

Photons of energy 100 GeV can pair create off target photons of energy $E_{\text{opt}} \sim 1$ eV, resulting in an optical depth

$$\tau_{\gamma\gamma} \approx n_{\gamma} \sigma_{\gamma\gamma} R_{\text{sh}} \approx 3(\tau_{\text{opt}} + 1) \frac{E_{\text{opt}}}{eV} \frac{L_{\text{opt}}}{10^{39} \text{ erg s}^{-1}} v_{\text{sh}}^{-1} t_{\text{cool}}^{-1}$$  \hfill (26)

near unity, where $\sigma_{\gamma\gamma} = 10^{-35} \text{ cm}^2$ is the photon-photon absorption cross section near threshold and $n_{\gamma} = L_{\text{opt}}(\tau_{\text{opt}} + 1)/(4\pi c R_{\text{sh}}^2 E_{\text{opt}})$.
is the energy density of the target optical photons near the shock and \( \tau_{\text{opt}} \approx m_\gamma n R_{\text{sh}} k_{\text{opt}} \approx 0.1 v_\gamma \text{g/cm}^2 \) is the optical depth of the shock at optical frequencies, where \( k_{\text{opt}} \approx 0.1 \text{ cm}^2 \text{ g}^{-1} \) is the optical opacity. We find that \( \tau_{\gamma,\gamma} \ll \tau_{\gamma,\gamma} \) few on a timescale of weeks for the relevant range of \( n \sim v_\gamma \) parameter space where \( E_{\max} \) is highest, suggesting that \( \gamma \sim \gamma \) may be relevant. However, due to the decreasing target photon number density from nova in the infrared (the Rayleigh Jeans tail), neglecting dust formation, such suppression may not extend to \( > 1 \text{ TeV} \) energies.

Although the Fermi LAT spectra hint at a \( \sim 10 \text{ GeV} \) spectral cut-off (Ackermann et al. 2014), which was confirmed with high significance at a single observing epoch by MAGIC (MAGIC Collaboration et al. 2015), our results in Figure 3 suggest that this does not exclude novae as TeV sources. The shock properties (\( n, v_\gamma, \tau_{\gamma,\gamma} \)) are likely to vary in time during the nova outburst, so a spectral cut-off measured at one epoch does not exclude a higher value of \( E_{\max} \) at other times.

CTA will reach a 5\( \sigma \) sensitivity of \( \sim 3 \times 10^{-11} \sim 10^{-10} \text{ erg cm}^{-2} \text{ s}^{-1} \) in the energy range of \( \sim 0.1 \sim 1 \text{ TeV} \) for a half-hour integration, and a factor 3 times deeper for a 5 hour integration (Actis et al. 2011). Figure 5 shows the spectral energy distribution from pion decay, assuming the DSA particle spectrum (eq. 16) for different values of the proton spectral cut-off and normalized to the characteristic value of the 0.1 \( \sim 1 \text{ GeV} \) nova LAT flux. We find that a LAT-detected nova with \( E_{\max} \geq 1 \text{ TeV} \) could easily be detected by CTA in a half-hour integration, while a nova with the same \( E_{\max} \) even an order of magnitude dimmer could be detected in 5 hours. We strongly encourage future additional TeV follow-up of novae near and after optical peak, even if not first detected by Fermi. Detectors or upper limits can be used to place meaningful constraints on the maximum accelerated particle energy and the location of the gamma-ray emission.

Although roughly 10 Galactic novae are discovered by optical surveys per year, only \( \sim 1 \) per year has been detected by Fermi-LAT. The main distinguishing feature of LAT-detected nova appears to be their close distances, although selection effects related to which nova receive pointed observations cannot be excluded. The deeper sensitivity of CTA suggests that, if novae spectra indeed routinely extend to sufficiently high energies, a much larger fraction of the optically-discovered events are potentially detectable as TeV sources. Indeed, the total Galactic nova rate is estimated to be \( \sim 40 \) per year, but the majority go undiscovered due to extinction in the Galactic plane. Since novae remains gamma-ray luminous for a few weeks, on average at least one nova is gamma-ray active in the Galaxy at any given time. Nova thus represent a potentially promising transient source for a future CTA survey of the Galactic plane, even in the absence of an optical trigger.

### 4.2 Implications for Neutrino Detection

An additional prediction of the favored hadronic model is a GeV neutrino flux comparable to the gamma-ray flux that may be detected by future experiments (Razzaque et al. 2010). This neutrino energy range is especially interesting with the recent and near-future upgrades of large scale neutrino facilities. Here we consider IceCube with its DeepCore extension (Abbasi et al. 2012), which was designed for detection in the \( \sim 10 \sim 100 \text{ GeV} \) energy range. Future upgrades, such as the Precision IceCube Next Generation Upgrade (PINGU; The IceCube-PINGU Collaboration 2014) or KM3NeT-ORCA (Katz 2014), can further enhance sensitivity in the relevant energy range.

We begin by considering the detection prospects of novae with
Given the significant background rate and the very high cutoff energy required for $N_{\nu} \gg 1$, detecting a neutrino signal from the current sample of LAT-detected novae seems unlikely. Detection prospects may, however, be improved by two factors. First, note that the LAT emission is typically delayed with respect to the peak of the optical emission by a few days (Ackermann et al. 2014), a fact which is puzzling if a significant fraction of the nova optical emission is indeed shock-powered during the LAT phase (Metzger et al. 2015). Metzger et al. (2015) hypothesize that this delay is due to the attenuation of the LAT emission by (inelastic) electron scattering, at early times when the shocks occur at the smallest radii and are passing through the densest material. If this is true, then the true peak of the shock power could be significantly larger, by up to an order of magnitude, than that inferred from the LAT light curve. To suppress 100 MeV gamma-rays by inelastic scattering requires a Thomson optical $\tau_B \sim n\sigma_T R_\odot \gtrsim 100$, i.e. $n \gtrsim 10^{11}$ cm$^{-3}$ for $R_\odot \sim v_{ej} \times 1$ day and $v_{ej} \sim 1000$ km s$^{-1}$. From equation (23), the maximum particle energy is $E_{\max} \sim 400$ GeV for $n = 10^{11}$ cm$^{-3}$, $v_8 \sim 1$, $T_a = 2 \times 10^4$ K and $J_r \sim 10 J_0$. So, in principle $\sim 1$ TeV energy particles can be accelerated at times when $\gtrsim 100$ MeV gamma-rays are blocked.

Considering a scenario in which the neutrino luminosity is $\sim 10$ times higher than in our fiducial estimate, but lasting only the first $\sim 1$ day of the nova outburst, then the total number of detected neutrinos would be similar to our estimate in equation (27) (for the same $E_{\max}$), but the neutrino background would be reduced to $N_{bg} \sim 0.1$ over this shorter time interval. The prospects for detecting a nova during the earliest phases of the nova outburst would therefore be considerably greater.

The probability of detecting individual novae would also be higher for a particularly nearby event, e.g. at a distance of 1 kpc instead of the $\sim 4-8$ kpc typical of the current LAT sample. Finally, given the relatively high rate of Galactic novae ($\sim 10$ yr$^{-1}$), it may be possible to stack multiple novae in one observation, improving the probability of a statistically significant detection.

### 4.3 Magnetic Field Amplification at Non-Relativistic Shocks

The minimum gamma cut-off of $\sim 3$ GeV allowed by LAT measurements requires particles to be accelerated to $E_{\max} \sim 30$ GeV in hadronic scenarios. From equation (23) this requires a minimum magnetic energy fraction of

$$\epsilon_B \gtrsim 2 \times 10^{-3} \epsilon_{p-1} \epsilon_{v-8}^{-2} (J_r/J_0)^{-2} \approx 10^{-3} \left( \frac{L_\nu}{10^{35} \text{erg s}^{-1}} \right) \epsilon_{p-1} \epsilon_{v-8}^{-1} v_8^2 J_r (J_r/J_0)^{-2},$$

where in the final line we have normalized the density using the observed gamma-ray luminosity (eq. 7) assuming $R_\odot \sim v_{ej} a$ and we have assumed that the gamma-ray emission is hadronic.

For $v_8 \lesssim 1$ this minimum value of $\epsilon_B$ is generally larger than the value that would be achieved from flux freezing alone at the shock given the estimated value for the pre-shock field (Metzger et al. 2015). This provides evidence for magnetic field amplification in nova shocks, likely driven by cosmic rays (Bell 2004). There is other evidence for magnetic field amplification in supernova remnant shocks (Molina et al. 2010; Ressler et al. 2014), for instance from fitting the combined radio and X-ray data (e.g., Molina & Caprioli 2012 for Tycho). However, the evidence for similar amplification is less clear in younger radio supernovae (Thompson et al. 2009).

Novae provide ideal nearby laboratories to study the onset of CR acceleration, because other time-dependent sources, such as radio supernovae, typically occur too distant to detect as gamma-ray sources.

### 5 CONCLUSIONS

Shocks in novae occupy a relatively unique regime of low velocities and high gas densities (Fig. 1). Such high densities imply that gas well upstream of the shock is neutral and opaque to ionizing radiation. Because the magnetic field amplification needed for diffusive shock acceleration requires ionized gas, the acceleration process is confined to a narrow photo-ionized layer just ahead of the shock. This limits particle energies to orders of magnitude below the naive Hillas criterion estimate. By comparing the downstream advection rate to the growth rate of the NRH Cosmic Ray Current-Driven Instability in the ionized layer (accounting for ion-neutral damping), we have quantified the maximum particle energy, $E_{\max}$, across the range of shock velocities and upstream densities constrained independently by observations of nova spectra, characteristic ejecta mass, and LAT gamma-ray emission (Fig. 5).

We find values of $E_{\max} \sim 10$ GeV–10 TeV, which are broadly consistent with those needed to produce the spectral cut-offs constrained by LAT observations and observed conclusively at a single epoch by the MAGIC telescope. We also find, however, that in principle acceleration up to $E_{\max} \gg 1$ TeV is achievable for low densities and high shock velocities. Furthermore, the observed spectral cut-offs might not be intrinsic to the accelerator, but instead result from $\gamma-\gamma$ absorption. The latter is likely to become less important at the highest photon energies (at least until dust formation occurs, producing infrared target photons).

Novae gamma-ray emission could therefore extend up to the energy range $\gtrsim 100$ GeV accessible to atmosphere Cherenkov telescopes, such as the planned CTA. Detecting high energy neutrinos by IceCube appears less promising, unless (1) $E_{\max} \gtrsim 10$ TeV and (2) the shock power is higher during the earliest phases of the nova outburst than implied by the LAT emission, or (3) in the case of a particularly nearby nova of distance $\lesssim 1$ kpc, which, however, occur perhaps only once per decade.

Novae are unlikely to be an important source of galactic cosmic rays (CRs) compared to supernova remnants (SNRs). For canonical Galactic rates of supernovae, $\mathcal{R}_{SN} \approx 0.03$ yr$^{-1}$, and novae, $\mathcal{R}_N \approx 40$ yr$^{-1}$, and blast wave energies of $E_{SN} \approx 10^{51}$ erg and $E_N \approx M_4 v_8^2/2 \approx 10^{49} M_4 v_8^2$ erg, respectively, the relative CR flux ratio is approximately $(E_N/R_\odot)/(E_{SN}/R_{SN}) \approx 2 \times 10^{-3}$, i.e., the contribution of novae to the diffuse CR spectrum is negligible. This estimate also assumes equal escape probability, while in fact the nova ejecta is probably more effective than SNRs at trapping cosmic rays due to the higher gas densities and short evolution times.

Interestingly, the CR spectrum measured by PAMELA and AMS–02 shows a hardening around 200 GeV/nucleon, where the proton spectral slope changes from $\sim 2.85$ to 2.7 (Adriani et al. 2011; Aguilar et al. 2015). Such a feature could also be explained by an excess of low-energy CRs, whose spectrum is softer than that produced by SNRs (which dominates at larger energies) and cuts off above $\sim 100$ GeV/nucleon. At first glance, novae represent an intriguing source for this low energy CR component, since their maximum CR energies that fall naturally in the $\sim 100$ GeV range. Producing such a large distortion of the CR spectrum would, however, require a comparable nova and SNR contribution to the CR energy flux in this energy range, which as discussed above is only possible in the chance coincidence of a recent very nearby nova.

The physics of relativistic particle acceleration and magnetic
field amplification as probed by nova shocks represent important open problems across many areas of high energy astrophysics. Our results show that novae provide robustness for the robustness of both processes in non-relativistic shocks, even under novel conditions of high densities and in the presence of upstream neutral gas. Given the sensitive dependence of $E_{\text{max}}$ on the extent of the ionized layer, future work is needed to better constrain the structure of the upstream gas with a more accurate photo-ionization model, accounting for line emission from behind the radiative shock and also for downscattering and reprocessing of high frequency radiation trapped by the neutral gas.

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