Gamma-ray novae as probes of relativistic particle acceleration at non-relativistic shocks

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1 INTRODUCTION

As “guest stars” to our ancestors, novae have been observed since antiquity (Zwicky 1936). Powered by runaway nuclear burning on the surface of white dwarfs (e.g., Starrfield et al. 1972), nova eruptions are usually accompanied by the expulsion of matter outwards at high velocities \( \gtrsim 10^3 \) km s\(^{-1}\). (Shore 2013 and references therein). Most of the radiated energy in a nova outburst occurs at optical and ultraviolet wavelengths and is generally attributed to the outwards transport of thermal energy produced near the white dwarf surface (Prialnik 1986; Prialnik & Kovetz 1995; Starrfield et al. 2000). Dynamical stellar evolution and hydrodynamic calculations are used to interpret basic features of nova light curves, such as how their rate of evolution depends on the white dwarf mass, central temperature, and accretion rate (e.g., Yaron et al. 2005; Hillman et al. 2014). Despite notable successes, models of nova outbursts remain plagued by uncertainties such as the efficiency of convective mixing and the assumed prescription for mass loss. This makes it challenging to accurately predict the quantity, velocity, and time history of matter ejected from the white dwarf surface.

Evidence for shocks in novae has existed well before Fermi in the form of hard X-ray emission (e.g., Mukai & Ishida 2001; Mukai et al. 2008), early peaks in the radio emission which are
inconsistent with thermal emission from freely expanding photoionized gas (e.g., Taylor et al. 1987, Krauss et al. 2011, Weston et al. 2013), and coronal line emission during the nebular phase requiring a harder source of ionizing radiation than expected from the cooling white dwarf (e.g., Shields & Ferland 1978). However, since most of these signatures are observed on timescales of months or later after the eruption, shocks appear to have been relegated to a mere side feature of the main thermonuclear event. The Fermi discovery of luminous shocks in coincidence with the optical peak unambiguously establishes their importance to the qualitative picture of nova eruptions. The dearth of X-ray and radio shock signatures at times coincident with the gamma-rays is unsurprising in retrospect due to the high bound-free and free-free optical depths created, respectively, by large columns of neutral and ionized gas at early times (Hillman et al. 2014, Metzger et al. 2014).

Shocks were unexpected in classical novae because the pre-eruption 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 (e.g., 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 ejection or a continuous wind (e.g., Bath & Shaviv 1976) with a higher velocity and more spherical geometry.

Assuming the expanding ejecta cools sufficiently, 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 (Fig. 1). Interestingly, modeling of the symbiotic gamma-ray nova V407 Cyg (Abdo et al. 2010) also appeared to require the presence of a dense equatorial torus (Martin & Dubus 2013), similar to that inferred in classical novae (see also Sokoloski et al. 2008, Orlando et al. 2009, Drake et al. 2009, Orlando & Drake 2012 for other evidence for bipolar ejecta in symbiotic novae).

Metzger et al. (2014) present a semi-analytic model for nova shocks and their thermal emission, which they fit to the radio and X-ray data of the gamma-ray nova V1324 Sco. A key finding was that the shocks responsible for the radio maximum seen months after the eruption could, at smaller radii and earlier times, also power the nova optical emission. One can generalize the basic argument as follows. Gas ahead of a shock with power sufficient to create the observed gamma-ray emission is necessarily dense. Most of the kinetic energy dissipated by shocks moving through a dense medium is radiated as thermal X-rays (a so-called “radiative” shock). These X-rays are absorbed by neutral gas ahead or behind the shock, reprocessing their energies to lower, optical frequencies, where the lower opacity readily allows their energy to escape. The observed gamma-ray luminosity is typically only a fraction $\sim 10^{-4} - 10^{-2}$ of that emitted as optical/UV radiation (§1). However, since only a fraction of the shock power is used to accelerate relativistic particles and only a fraction of that is radiated in the LAT bandpass, one concludes that a significant fraction of the nova optical emission is shock-powered.

The above result also implies that combined optical and gamma-ray data from novice can be used to probe particle acceleration at non-relativistic shocks. The basic concepts of diffusive shock acceleration (DSA) were developed almost forty years ago (e.g., Axford et al. 1977, Bell 1978, Blandford & Ostriker 1978). However, only recently have plasma kinetic simulations seen the self-consistent development of the DSA cycle in collision-less non-relativistic shocks (e.g., Caprioli & Spitkovsky 2014, Kato 2014, Park et al. 2014). As we shall discuss, gamma-ray novae provide a real-time laboratory for testing the predictions of DSA theory in ways complementary to traditional methods, such as the modeling of supernova remnant emission.

This paper solidifies the above theoretical arguments and puts them into practice using data from the classical novae V1324 Sco and V339 Del (hereafter, ‘nova Sco’ and ‘nova Del’, respectively). Though well-studied at many wavelengths, we do not consider the nova V959 Mon because of the lack of early optical coverage concurrent with the gamma-ray detections. The salient properties of nova shocks are reviewed in §2, including their radiative nature (§2.1), the high efficiency with which shock power is radiated at optical/UV frequencies (§2.2), and the efficiency of particle acceleration and gamma-ray production (§2.3). A combined analysis of the optical and gamma-ray data of V1324 Sco and V339 Del is presented in §3 using our theoretical framework to constrain the minimum fraction of the shock power placed into relativistic particles and the minimum fraction of nova optical light curves that are shock powered. In §4 we discuss our results and their implications for particle acceleration at non-relativistic shocks and for the power source behind nova optical light curves.

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Figure 1. Schematic diagram of the proposed geometry (side view) of shock interaction in classical nova outflows (see text). A slow outflow with velocity $v_{\text{ej}} < 10^3$ km s$^{-1}$ is ejected first, its geometry shaped into an equatorially-concentrated torus, possibly due to interaction with the binary companion of the white dwarf (e.g., Livio et al. 1990). This outflow is followed within a few days by a faster outflow or continuous wind with a higher velocity $v_f \sim 2v_{\text{ej}}$. The fast and slow components collide, produce a forward-reverse shock structure. Kinetic energy dissipated by the shocks is radiated as thermal X-rays, which are absorbed by neutral gas ahead or behind the shock and re-radiated as thermal optical/UV emission (Metzger et al. 2014). A small fraction, $\epsilon_{\text{sh}} \ll 1$, of the shock power is used to accelerate non-thermal ions or electrons, which radiate gamma-rays by interacting with ambient gas or radiation, respectively.
2 SHOCKS IN GAMMA-RAY NOVAE

2.1 The shocks are probably radiative

The slow outflow from a nova with velocity \( v_{ej} = 10^3 v_h \text{ km s}^{-1} \) expands to a radius

\[
R_g = v_{ej} t \approx 6 \times 10^{13} t_{ab} v_h \text{ cm}
\]  

(1)

by a time \( t = t_{ab} \) week. The density in the slow ejecta of assumed thickness \( \sim R_g \) and hydrogen-dominated composition is given by

\[
n_{ej} \approx \frac{M_{ej}}{4\pi R_g^2 m_p \rho_{ej}} \sim 9 \times 10^9 M_{-5,5} t_{ab}^{-1} v_h^{3/2} \text{ cm}^{-3},
\]  

(2)

where \( f_{\text{solid}} \approx 0.5 \) is the fraction of the total solid-angle subtended by the outflow (Fig. 1) and \( M_e = 10^4 M_{-5} M_{ej} \) is the ejecta mass, normalized to a value characteristic of those measured by late thermal radio emission (Hjellming et al. 1979; Sequist et al. 1980). We assume a radially-thick outflow because the resulting low density makes this the most conservative assumption when arguing for radiative shocks.

A faster outflow (hereafter “wind”) of mass loss rate \( \dot{M} \) and velocity \( v_w \sim 2 v_{ej} \) collides with the ejecta from behind. The density of the wind at the collision radius (\( \sim \) radius of the slow ejecta) is given by

\[
n_w \approx \frac{\dot{M}}{4\pi R_{w}^2 m_p v_w} \sim 2 \times 10^3 M_{-5} v_{8}^2 t_{ab}^{-1} \text{ cm}^{-3},
\]  

(3)

where \( \dot{M} = 10^3 M_{-5} M_{ej} \text{ wk}^{-1} \) is normalized to a value resulting in the ejection of \( \sim 10^5 M_{ej} \) over a week. For instance, a total mass \( 4 \times 10^5 M_{ej} \) was ejected in the “fast” component of V959 Mon (Chomiuk et al. 2014; L. Chomiuk, private communication). Our assumption that the fast outflow is a steady wind is also the most conservative one when arguing for radiative shocks; for a fixed fast ejecta mass, the wind density at the reverse shock is smaller than if the ejection had occurred over a shorter duration.

The interaction drives a forward shock (FS) through the slow shell and a reverse shock (RS) back through the wind (Fig. 1). Assuming the shocks are radiative, the post shock material is compressed and piles up in a central cold shell sandwiched by the ram pressure of the two shocks. The FS propagates at a velocity \( v_s \approx v_w - v_{ej} \ll v_{ej} \), while the velocity of the reverse shock is \( v_{rs} = v_w - v_c \), where \( v_c \) is the velocity of the cold central shell, which only moderately exceeds the velocity of the slow ejecta due to the large inertia of the latter (see Metzger et al. 2014), although be wary of notational differences). Hereafter, the velocities of both shocks are parametrized as \( v_{sh} = \eta v_{ej} \), where \( \eta < 1 \) for the FS and \( \eta \sim 1 \) for RS. Also note that for parameters of interest, the rate of energy dissipation by the RS is greater than that by the FS, potentially favoring the former as the site of particle acceleration.

The shocks heat the gas to a temperature

\[
T_{sh} \approx \frac{3}{16k} m_p v_{sh}^2 \sim 2.3 \times 10^7 \eta^2 \text{K}
\]  

(4)

and compresses it to a density \( n_{sh} = 4 n_{ej} \) (FS) or \( n_{in} \) (RS).

Gas cools behind the shock on a characteristic timescale

\[
\tau_{cool} = \frac{3k T_{sh} / 2}{\Lambda (n_{sh})} \approx \left\{ \begin{array}{ll}
3.1 \times 10^5 \eta^2 M_{-5,5} t_{ab}^{-1} \text{ s, } & \text{FS,} \\
6.0 \times 10^5 \eta^2 M_{-5,5} t_{ab}^{-1} \text{ s, } & \text{RS}
\end{array} \right.
\]  

(5)

where \( \Lambda = 2 \times 10^{-7} T^{1/2} \text{ erg cm}^{-3} \text{ s}^{-1} \) is the cooling function and we have assumed \( \eta = 1 \) in the RS case. Our cooling function includes just free-free emission, which is conservative from the standpoint of radiative shocks because line cooling contributes a comparable or greater cooling rate at temperatures \( \leq 3 \times 10^7 \text{ K} \) (e.g. Schure et al. 2009). Here we have assumed that electrons and protons are efficiently coupled behind the shock, as is justified because the timescale for Coulomb energy exchange \( \tau_{ep} = 147 \times 10^{3} / n_{sh} \) s (NRL Plasma Formulary; Huba 2007) is short compared to the expansion timescale \( t_{exp} \sim R_g / v_{ej} \sim t \).

\[
\frac{t_{ep}}{\tau_{cool}} \approx \left\{ \begin{array}{ll}
3 \times 10^3 M_{-5,5}^{1/4} \eta^2 T_{sh}^{1/2} \text{ FS, } & \\
1.2 \times 10^3 M_{-5,5}^{1/4} \text{ RS, } &
\end{array} \right.
\]  

(6)

where we have (conservatively) assumed a pure hydrogen composition.

Whether the shocks are radiative depend on the ratio of the cooling timescale to the expansion timescale,

\[
\chi \equiv \frac{\tau_{cool}}{t} \approx \left\{ \begin{array}{ll}
2.1 \times 10^{-3} \eta^2 M_{-5,5}^{1/4} \text{ FS, } & \\
0.10 \eta^2 M_{-5,5}^{1/4} \text{ RS, } &
\end{array} \right.
\]  

(7)

For the characteristic range of parameters \( M_{-5} \sim 0.3 \sim \chi \) (Sequist & Bode 2008), \( v_h \lesssim 3 \), \( \eta < 1 \) and timescales corresponding to the gamma-ray emission \( t \lesssim \text{few weeks} \), one concludes that the FS is likely to be radiative (\( \chi > 1 \)). It is less clear that the RS is radiative, since in principle for \( v_h \gtrsim 2 \) and a very low wind mass loss rate \( M_{-5} \lesssim 1 \) one can have \( \chi > 1 \).

The argument for radiative shocks can be strengthened by considering additional constraints. First, the optical depth of the outer unshocked ejecta at optical frequencies,

\[
\tau_{opt} \approx \frac{m_p n_{ej} v_{sh} \sqrt{\Lambda (n_{sh})}}{8 \sqrt{3} A_0 (n_{sh} m_p \rho_{sh} \tau_{cool})} \approx \left\{ \begin{array}{ll}
2 \times 10^{-3} \eta^2 \rho_{sh}^{1/4} \text{ FS, } & \\
0.08 \eta^2 M_{sh}^{1/4} t_{opt}^{1/2} \text{ RS, } &
\end{array} \right.
\]  

(8)

must exceed unity, where \( \rho_{sh} \sim 0.1 \text{ cm}^{-3} \) is the optical opacity set by Doppler-broadened iron lines (Pinto & Eastman 2000). This constraint is motivated by the lack of clear evidence (e.g. emission lines) for hot, energetic shocks in optically thin regions during the optical peaks of nova and during most epochs of gamma-ray detection. Over the first few weeks or longer following outburst (i.e., prior to the drop of the "iron curtain"), nova spectra are characterized by broad P-Cygni line profiles, indicating the presence of a pseudo-photosphere.

Equation (7) can be rewritten in terms of \( \tau_{opt} \) as

\[
\chi = \frac{9 \eta^2 k^{3/2} v_{sh} \sqrt{\Lambda (n_{sh})}}{8 \sqrt{3} A_0 (n_{sh} m_p \rho_{sh} \tau_{cool})} \approx \left\{ \begin{array}{ll}
2 \times 10^{-3} \eta^2 \rho_{sh}^{1/4} \text{ FS, } & \\
0.08 \eta^2 M_{sh}^{1/4} t_{opt}^{1/2} \text{ RS, } &
\end{array} \right.
\]  

(9)

For the forward shock to occur below the photosphere (\( \tau_{opt} > 1 \)) yet not be radiative (\( \chi > 1 \)) would thus require unphysically high ejecta velocities, \( v_{ej} \gtrsim 10,000 \text{ km s}^{-1} \). The RS is also radiative if \( \tau_{opt} > 1 \) on timescales of a couple weeks for \( M_{-5} \lesssim 3 \) as long as \( M_{-5} > 0.2 \).

A second requirement is that power dissipated by the shocks

\[
L_{sh} = \frac{9 \eta^2 k^{3/2} n_{sh} \rho_{sh} v_{sh}^3}{32} \text{eV/cm} \text{s} \left| \text{per unit } \frac{1}{4} \text{ } \text{erg/cm} \text{s} \right|
\]  

(9)

is bounded by a factor \( 1/\epsilon_{sh} \leq 1 \) accounting for the fraction of the shock power radiated as gamma-rays, where \( \epsilon_{sh} \) accounts for the efficiency of the shock in accelerating non-thermal particles and \( \epsilon_r \) accounts for the energy of the latter radiated in the LAT bandpass (see Fig. 2.3). Equation (7) can be written in terms of \( L_{sh} = L_{\gamma} / \epsilon_{sh} \) as

\[
\frac{1}{\epsilon_{sh}} \approx \frac{0.02 \eta^2 v_{sh}^2 \left( \frac{L_{\gamma}}{10^{46} \text{ erg s}^{-1}} \right)}{0.01}.
\]  

(10)
For measured values $L_{\gamma} \sim 3 \times 10^{49}$ erg s$^{-1}$ \citep{ackermann2014}, the shocks are thus radiative ($\chi < 1$) on a timescale of weeks for $\epsilon_{\text{th}} \epsilon_\gamma < 0.01$ if the velocity of the shocks is $\lesssim 2,000$ km s$^{-1}$.

In conclusion, both forward and reverse shocks are likely to be radiative at times corresponding to the observed gamma-ray emission, although this statement is the most secure at the earliest times. One can nevertheless keep the following arguments fully general by introducing the shock radiative efficiency $f_{\text{rad}} = (1 + 5\chi_\gamma \chi)^{-1}$, which equals unity for $\chi \ll 1$ but scales $\propto 1/\chi$ for $\chi \gg 1$ \citep{metzger2014}.

\subsection*{2.2 The shock power will mostly emerge as optical radiation}

Absent the immediate presence of a shock, the bulk of the nova ejecta is neutral at early times because the timescale for radiative recombination, $t_{\text{rec}} \sim 1/n_p m_\text{p} c^{-1} \sim 11 \nu_{\text{opt}} M_{\odot}^{-1} Z^{-2}$ s, is extremely short compared to the evolution timescale $\sim$ weeks, where $\nu_{\text{opt}} \sim 10^{-12} Z^2$ cm s$^{-1}$ is the approximate radiative recombination rate for hydrogen-like species of charge $Z$ \citep{osterbrock2006}. X-rays of temperature $T_{\text{sh}} \lesssim 10^4$ K (eq. [11]) are thus absorbed by high columns of neutral gas ahead or behind the shock, before being re-radiated as line emission at lower frequencies. If the shock is radiative then the timescale for reprocessing to lower temperatures is also necessarily short because the line cooling function $\chi(T)$ increases rapidly with decreasing temperature down to $T \sim 10^4$ K (e.g. \citealt{schure2009}).

Radiation escapes efficiently once it reaches the optical/near-UV band due to the much lower opacity at these frequencies compared to the UV/X-ray opacity of neutral gas. Reprocessed optical radiation will furthermore not be degraded by PdV losses provided that the photon diffusion timescale $t_d \sim c / r_{\text{diff}}$ is shorter than the expansion timescale $t_{\text{exp}} \sim R_0 / v_\gamma \approx t_s$, where $t_{\text{exp}}$ is the optical depth \citep{arnett1982}. Equating these two, optical radiation escapes without adiabatic losses after a time \citep{arnett1982}

$$t_{\text{opt}} \approx 0.4 M_{\odot}^{1/2} Z^{-1/2} d, \quad \text{(11)}$$

If the energy released by shocks or diffusion from the central source peaks at times $\lesssim t_{\text{opt}}$, then $t_{\text{opt}}$ also sets the timescale for the optical lightcurve to peak (rise time).

The onset of gamma-ray emission from V1324 Sco and V339 Del was delayed with respect to the peak of the optical radiation by a few days \citep{ackermann2014}. The dominant process by which gamma-rays of energy $\epsilon_\gamma \sim 0.1$–3 GeV are attenuated is inelastic electron scattering, for which the Klein-Nishina opacity is $\kappa_{\text{en}} \sim 10^{-3} (\epsilon_\gamma / \text{GeV})^{-1} \kappa_{\text{opt}} \sim 4 \times 10^{-3} (\epsilon_\gamma / \text{GeV})^{-1} \kappa_{\text{opt}}$. The ejecta will thus remain opaque to gamma-rays of energy $\epsilon_\gamma$ until $t_{\text{opt}} \lesssim 250 (\epsilon_\gamma / \text{GeV})$, as occurs after a time $t_s$ which exceeds that of the nominal optical peak (eq. [11]) by the ratio

$$t_s / t_{\text{opt}} \approx \left( \frac{c \kappa_{\text{en}}}{v_\gamma \kappa_{\text{opt}}} \right)^{1/2} \approx 1.1 \left( \frac{\epsilon_\gamma}{\text{GeV}} \right)^{-1/2} v_8^{-1/2}, \quad \text{(12)}$$

i.e. $t_s \sim$ couple days for $\epsilon_\gamma \sim 0.1$ GeV.

\footnote{Although the bulk of the ejecta is neutral at early times, X-rays from the shocks fully ionize a thin layer of gas just ahead of the shocks \citep{metzger2014}. Thus, we do not expect the shock to be modified by the presence of a neutral upstream medium, as may occur in some supernova remnants \citep[e.g.,][]{blasi2012}.}

\footnote{Photo-pion production is the dominant source of opacity for higher energy gamma-rays $\epsilon_\gamma \gtrsim 3$ GeV, for which the opacity is $\kappa_{\pi} \sim 3 \times 10^{-4} \kappa_{\text{en}}$ \citep[e.g.][]{anchordoqui2002, montanet1994}.}

From the above we can draw two key conclusions: (i) a shock that produces gamma-ray emission which is not absorbed and hence observable (gamma-ray optical depth $\tau_{\gamma} < 1$) necessarily radiates the bulk of its dissipated thermal energy without adiabatic losses at optical frequencies ($\tau_{\text{opt}} \ll c / v_\gamma$); (ii) gamma-ray absorption ($\tau_{\gamma} > 1$) at early times may explain the delayed onset of the gamma-ray emission. If true, the latter implies that the optical emission near peak can be shock-powered, even if gamma-ray emission is suppressed at this time.

\subsection*{2.3 Partitioning the shock energy}

A large fraction $f_{\text{rad}} \sim 1$ of the total power $L_{\gamma}$ dissipated by shocks goes into thermal X-rays, which are absorbed and re-radiated as optical emission \citep{osterbrock2006}. A much smaller fraction $\epsilon_{\text{th}} \ll 1$ goes into accelerating non-thermal ions or electrons. The fraction of non-thermal power $\epsilon_{\text{th}}$ radiated as gamma-rays, $L_{\gamma} = \epsilon_{\text{th}} E_{\gamma} L_{\gamma}$, depends also on a factor $\epsilon_{\gamma} < 1$ accounting for the radiative efficiency of the accelerated particles and the fraction of the total gamma-ray emission emitted in the LAT bandpass. Combining these expressions, the fraction of the total nova optical luminosity powered by shocks can be written

$$f_{\text{sh}} \equiv \frac{L_{\text{opt,sh}}}{L_{\text{opt}}} = \frac{f_{\text{rad}}}{\epsilon_{\text{th}}} \frac{E_{\gamma}}{L_{\gamma}} \quad \text{(13)}$$

Once $\epsilon_{\gamma}$ is specified based on the assumed emission process (hadronic or leptonic), the observed ratio $L_{\gamma} / L_{\text{opt}}$ sets a lower limit on the value of $\epsilon_{\text{th}},$ i.e.

$$f_{\text{sh}} < 1 \Rightarrow \epsilon_{\text{th}} > \epsilon_{\text{th, min}} = \frac{1}{L_{\gamma}} \frac{L_{\gamma}}{L_{\gamma}} \frac{L_{\gamma}}{L_{\gamma}} \quad \text{(14)}$$

assuming a radiative shock ($f_{\text{rad}} = 1$). Alternatively, if the value of $\epsilon_{\text{th}}$ is assumed, then the measured value of $L_{\gamma} / L_{\text{opt}}$ determines $f_{\text{sh}}$. We now consider what values of $\epsilon_{\text{sh}}$ and $\epsilon_{\gamma}$ are expected in hadronic and leptonic scenarios, respectively.

\subsubsection*{2.3.1 Hadronic scenario}

The momentum distribution $dN_p / dE_p \propto E_p^{-4}$ of non-thermal protons accelerated via Diffusive Shock Acceleration \citep[e.g.,][]{blanford1978} corresponds to an energy distribution

$$dN_p / dE_p \propto \begin{ cases} E_p^{1/2}, & kT_{\text{sh}} \ll m_p c^2 \text{ constant, } m_p c^2 \ll E_p < E_{\max}, \end{ cases} \quad \text{(15)}$$

that concentrates most of the non-thermal energy in relativistic particles. The \citep{hillas1984} criterion sets an upper limit on the maximum proton energy

$$E_{\max} = \frac{B_0 v_{\text{sh}} R_0}{c} \lesssim 7 \times 10^{33} \text{ eV} \left( \frac{\epsilon_{\gamma}}{10^{-5}} \right)^{1/2} M_{\odot}^{-1/2} v_8^{-1/2}, \quad \text{(16)}$$

where $B_0$ is the magnetic field strength near the shock, which is estimated by assuming that magnetic pressure $B_0^2 / 8\pi$ is a fraction $\epsilon_B$ of the thermal pressure of the post-shock gas. A high value of $\epsilon_B \gtrsim 10^{-2}$ is expected if ion acceleration is efficient, because cosmic-ray induced instabilities amplify the magnetic field for some distance behind the shock to a significant fraction of its equipartition value \citep[e.g.,][]{caprioli2014}.

Neglecting magnetic field amplification, the magnetic field strength and resulting equipartition factor $\epsilon_B$ within the nova ejecta can be estimated as follows. Energy released by nuclear burning is initially carried outwards from the white dwarf surface by convection. In the convective region, magnetic pressure $B_0^2 / 8\pi$ \citep[© RAS, MNRAS 000:1-??]{2014}.
will be in approximate equipartition with the convective turbulent kinetic energy, which is itself presumably comparable to the kinetic energy density of the outflowing matter \( U_e = \rho v^2/2 \) of velocity \( v \), where \( \rho \) is the outflow density. If the outflow occurs in the form of a steady wind \( (\rho \propto r^{-2}; \text{eq. } [3]) \), then the frozen-in toroidal component of the magnetic field will decrease \( \propto 1/r \) from sideways expansion in two dimensions, resulting in the ratio \( U_B/U_e \) remaining fixed at close to its value at the outflow surface, i.e. \( \epsilon_B \sim O(1) \). However, if (as seems likely) the fast ejecta possesses a gradient in its radial velocity, then radial expansion becomes important at some radius \( R_{opt} \). Flux-freezing in the 3D expanding flow results in \( B \propto 1/r^2 \) at \( r > R_{opt} \). The ratio of \( U_B/U_e \propto \epsilon_B \) will in this case be smaller by a factor \( (R_{opt}/R_{sh})^2 \) than its value at the white dwarf surface, resulting in a value \( \epsilon_B \gtrsim 10^{-6} \) for characteristic values \( R_{sh} \sim 10^{14} \text{ cm} \) and \( R_{opt} > 10^{11} \text{ cm} \), the latter representing the characteristic maximum size of inflated white dwarf envelope (e.g. Hillman et al. 2014 their Fig. 2). Even \( \epsilon_B \sim 10^{-8} \) corresponds to a maximum proton energy \( E_{\text{max}} \) \( \gtrsim 10^{12} \text{ eV} \) eq. [16] which is sufficiently large to explain the highest LAT-detected photon energies via \( \gamma \)-ray creation (Ackermann et al. 2014). We caution, however, that because the bulk of the ejecta is neutral and hence may not be able to support the strong turbulent magnetic field needed to accelerate particles, an additional constraint on \( E_{\text{max}} \) is that the particle Larmor radius \( \gamma_1 \sim E_p/eB_0 \) must be smaller than the width of the X-ray ionized layer ahead of the shock, which Metzger et al. (2014) estimate to be \( \Delta_{\text{kin}} \sim 3 \times 10^{13} (n_{sh}/10^3 \text{ cm}^{-3})^{-1} \text{ cm} \) (their eq. [46]); for typical values of the unshocked gas density \( n_{sh} \sim 10^{-6} \text{ cm}^{-3} \) (eq. [2], [3]), one finds that \( \epsilon_B \gtrsim 0.1 \) is required to have \( \gamma_1 \sim \Delta_{\text{kin}} \) for protons of energy \( E_p \gtrsim 10^{18} \text{ eV} \).

The fraction of the shock power placed into non-thermal ions \( \epsilon_{\text{th}} \) can be estimated from observations of other non-relativistic shocks, such as gamma-ray emission in supernova remnants (e.g. Ackermann et al. 2014). In Tycho, for instance, Morlino & Caprioli (2012) infer \( \epsilon_{\text{th}} \approx 0.1 \) if the gamma-rays are hadronic in origin. Hybrid (kinetic ions - fluid electron) simulations of non-relativistic shocks also find \( \epsilon_{\text{th}} \approx 0.1 \) – 0.2 for cases in which the upstream magnetic field is quasi-parallel to the shock normal (Caprioli & Spitkovsky 2014). However, if the nova ejecta is characterized by a phase of approximately steady-state outflow as described above, then radial expansion will produce a magnetic field dominated by its toroidal component, i.e. perpendicular to the outflow and hence shock direction (Fig. 1). For such quasi-perpendicular magnetic field geometries, Caprioli & Spitkovsky (2014) infer much lower proton acceleration efficiencies (consistent with zero).

Relativistic protons accelerated at the shocks produce pions by colliding with effectively stationary protons in the ejecta. The mean time between interactions is given by \( t_{\text{coll}} \approx (k_{\pi-p} m_{p} m_{\pi})^{-1} \), where \( k_{\pi-p} \approx 0.025 \text{ cm}^2 \text{ g}^{-1} \approx k_{\gamma-p}/4 \) Kamae et al. 2009 is the opacity for inelastic proton collisions of energy \( E_p \geq \text{GeV} \). The number of collisions a proton experiences over an expansion timescale \( t/t_{\text{coll}} \) thus exceeds unity until after a time \( t_{\text{coll}} \approx 6.2 M_{14}^{1/2} \nu_{3/2}^{1/2} \) weeks. (17)

At times \( t \ll t_{\text{coll}} \) the ejecta thus acts as an efficient “calorimeter” for converting relativistic protons into gamma-rays. The fraction of \( p - p \) collisions producing gamma-rays \( \epsilon_{\gamma} \) is the product of the fraction of inelastic interactions, \( \epsilon_{\text{opt}}/\epsilon_{\text{ elastic}} \sim 0.5 – 0.8 \) (where \( \epsilon_{\text{opt}} \) and \( \epsilon_{\text{ elastic}} \) are the elastic and inelastic cross sections, respectively; Kamae et al. 2006), and the fraction \( \approx 1/3 \) of inelastic events placed into the \( p + p \rightarrow \pi^0 \rightarrow \gamma + \gamma \) channel. We thus typically expect that \( \epsilon_{\gamma} \approx 0.2 -0.3 \), depending on the details of the accelerated proton spectrum, while a lower value could result if protons are not efficiently trapped, i.e. at times \( t > t_{\text{coll}} \).

To summarize, we expect \( \epsilon_{\text{th}}/\epsilon_p \approx 0.03 \) in hadronic scenarios if the magnetic field within the ejecta is perpendicular to the shock plane, but this value may be considerably lower if the field is instead parallel to the shock plane, as would be expected given a phase of quasi-steady outflow from the white dwarf surface.

### 2.3.2 Leptonic scenario

In supernova shocks, the ratio of the energy placed into relativistic protons to that in relativistic electrons is estimated to be \( \epsilon_{\text{p}} \sim 10^{-4} -10^{-2} \) based on observations of individual remnants (Völk et al. 2005; Morlino & Caprioli 2012) and synchrotron emission from Galactic cosmic rays (Beck & Krause 2005; Strong et al. 2010). This value is consistent with recent particle-in-cell simulations of non-relativistic shocks, which find \( \epsilon_{\text{p}} < 10^{-3} \) when extrapolated to shock velocities \( v_{sh}/c \ll 0.01 \) characteristic of those in novae (Park et al. 2014; Kato 2014). Assuming shocks accelerate protons with a maximum efficiency \( \epsilon_{\text{max}} \approx 0.1 \), the ratio \( \epsilon_{\text{p}}/\epsilon_{\text{e}} \) in the shocked ejecta is \( \epsilon_{\text{p}}/\epsilon_{\text{e}} \sim 10^{-2} \).

To Compton upscatter optical seed photons of energy \( E_{\text{opt}} \sim 10^{17} \text{ eV} \) to \( \sim 1 \times 10^{-2} \text{ GeV} \) requires electrons with energy \( E_e \approx \gamma_e m_e c^2 \) and Lorentz factor \( \gamma_e \sim 10^4 -10^5 \). The ratio of the Compton cooling timescale of the electron to the expansion timescale is given by

\[ t_{\text{cool}} \sim \frac{3 m_e c}{4 \sigma_T U_{\text{opt}} \gamma_e} \sim 0.07 \tau_{\text{opt}} \left( \frac{\gamma_e}{10^4} \right) \left( \frac{L_{\text{opt}}}{10^{38} \text{ erg s}^{-1}} \right)^{-1} \varepsilon_{\text{th}} \]  

(18)

where \( U_{\text{opt}} \sim L_{\text{opt}} \tau_{\text{opt}}/4\pi c R_{\text{sh}}^2 \) is the radiation energy density near the shock (for shocks below the photosphere, \( \tau_{\text{opt}} > 1 \)). Equation (18) shows that \( \gamma_e > 10^4 \) electrons are in the fast cooling regime \( t_{\text{cool}} \ll t_{\text{cool}} \) for luminosities \( L_{\text{opt}} \gtrsim 10^{38} \text{ erg s}^{-1} \) and timescales \( t \sim \text{weeks} \) of relevance. Thus, a high fraction of the shock energy placed into electrons with the necessary Lorentz factors to produce the observed gamma-ray emission will in fact radiate their energy.

Compton scattering of electrons with an accelerated spectrum \( dN_e/dE_e \sim E_e^{-\Gamma} \), where \( E_e \approx \gamma_e m_e c^2 \), produces a gamma-ray spectrum

\[ \frac{dN_e}{dE_e} \sim \begin{cases} E_e^{-(\rho+1)/2}, & E_e \ll E_c, \\ E_e^{-(\rho+2)/2}, & E_e \gg E_c. \end{cases} \]  

(19)

where \( E_c \sim E_{\text{opt}} \gamma_e^2 \) and \( \gamma_e \) is the minimum electron Lorentz factor obeying the fast-cooling condition \( t_{\text{cool}} \approx \tau_{\text{cool}} \) (eq. [18]). The fraction of the total energy in relativistic electrons radiated in the LAT bandpass is given by

\[ \epsilon_{\gamma} \sim \int_{10^{53}}^{10^{55}} \frac{dN_{\gamma}}{dE_{\gamma}} E_{\gamma} dE_{\gamma} \sim \begin{cases} 10^{-4}\rho - 10^{-5}\rho, & \rho > 2, \\ 0.2, & \rho = 2. \
\end{cases} \]  

(20)

Power-law fits to the gamma-ray spectra of three classical novae yield best-fit photon indices \( \Gamma \approx 2.5 \) (Ackermann et al. 2014) corresponding to \( \rho = 2(\Gamma - 1) \approx 2 - 3 \), for which we estimate \( \epsilon_{\gamma} \sim 0.2 - 0.4 \). To be conservative we hereafter assume \( \rho = 2 \), corresponding to \( \epsilon_{\gamma} = 0.2 \).

In summary, adopting \( \epsilon_{\text{th}} \sim 10^{-3} - 10^{-2} \) motivated by simulations and observations, we estimate that \( \epsilon_{\text{th}} \epsilon_{\gamma} \sim 10^{-6} - 10^{-4} \) in the leptonic scenario.

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\( ^4 \) Relativistic bremsstrahlung is subdominant compared to the Inverse Compton emission (Ackermann et al. 2014)
3 DATA

This section describes our analysis of the gamma-ray and optical light curves of V1324 Sco and V3229 Del, events chosen as currently being the only classical novae with both published gamma-ray data and contemporaneous optical coverage. We do not consider symbiotic novae in this analysis because the argument for the shocks being radiative is less secure given that the latter propagate into the extended wind of the giant companion star, instead of the potentially denser slow ejecta in the case of internal shocks. Our goal is to use the measured ratio of gamma-ray and optical fluxes to constrain the particle acceleration efficiency $\epsilon_{\text{nth}}$ and the fraction of the classical nova optical light curve powered by shocks $f_{\text{sh}}$, following the arguments outlined in [42,3].

3.1 Optical

For each nova we use photometric measurements at both optical and near-infrared (NIR) wavelengths. Optical photometry was taken from the database of the American Association of Variable Star Observers (Henden 2013). Both nova Sco and Del had BVR band measurements, but only Del included I band. To supplement the NIR for nova Sco, we used $JHK$ measurements from the Small and Moderate Aperture Research Telescope (F. M. Walter, private communication). Only photometric measurements coincident with the published Fermi gamma-ray light curve were used; specifically August 16–September 14 2013 and June 15–June 30 2012 for V339 Del and V1324 Sco, respectively. Photometric reddening corrections were applied to the datasets using an $E(B - V) = 0.2$ for V339 Del (Munari et al. 2013) and $E(B - V) = 1.0$ for V1324 Sco (Pinzella et al. 2015). Filter specific reddening corrections for the optical/NIR were taken from Schlafly et al. (2011) assuming $R_V = 3.1$.

The total optical/NIR flux each night, $F_{\text{opt}}$, was determined by approximating the SED as a blackbody, fitting a temperature and constant offset, and integrating over frequency. The early-time flux obtained for V339 Del are consistent with the values found in the more detailed analysis of Skopal et al. (2014), who also showed that the spectral shape was a blackbody all the way out to the mid-IR. We confirmed our blackbody assumption for V1324 Sco by extending our spectral fit up to the near-UV and comparing with flux values from Swift UVOT (Page et al. 2012). After applying reddening corrections from Brown et al. (2010) our blackbody fit predicts a near-UV flux of $2.1 \pm 0.4 \times 10^{-13}$ erg s$^{-1}$ cm$^{-2}$, a value in agreement with the observed flux, $2.4 \pm 0.16 \times 10^{-13}$ erg s$^{-1}$ cm$^{-2}$. As the near-UV flux is a factor $\sim 10^5$ times lower than in the optical, this excellent agreement validates the assumed blackbody spectral profile. Results for the total optical/NIR flux and best-fit temperature as a function of time are shown in Figure 2.

As only three nights have complete NIR + BVR color data in nova Sco, we also consider separately the flux ratio on other nights with LAT detections obtained using just the V magnitude to calculate the optical/NIR flux, assuming a bolometric correction identical to that measured with better color data on day 15. Uncertainty in the total flux is estimated by combining the uncertainty in the best-fit constant offset with an uncertainty in the temperature fit derived using the error in individual flux measurement.

3.2 Gamma-Ray

The Fermi LAT data for nova Sco and Del are taken from Ackermann et al. (2014). Spectral fits to the gamma-ray emission (over its entire duration) are provided by these authors for two profiles: Power Law (PL, $dN/dE \propto E^{-\Gamma}$) and Power Law with Exponential Cut-Off at energy $E = E_0$ (EPL, $dN/dE \propto E^{-\Gamma} \exp(-E/E_0)$). The best-fit values of the free parameters, $\Gamma$ (PL) or $\Gamma$ and $E_0$ (EPL), are provided along with their uncertainties.

Also provided for each day is the average photon number above 100 MeV, $N_\gamma \propto \int_{100 \text{ MeV}}^{\infty} E dN/dE dE$. The energy flux $F_\gamma \propto \int_{100 \text{ MeV}}^{\infty} E dN/dE dE$ is calculated from $F_\gamma$ according to

$$F_\gamma = N_\gamma \int_{100 \text{ MeV}}^{\infty} E dN/dE dE,$$

where uncertainties in $F_\gamma$ are derived from both the quoted uncertainties in $N_\gamma$ and in the best-fit parameters of the spectral fits.

Table 1. $\gamma$-ray/optical flux ratio in classical novae and implications.

| Nova   | $\langle \log(f_{\text{opt}}/F_{\text{opt}}) \rangle$ | $\log(\epsilon_{\text{nth}})_{\text{min,\,max}}$ | $\gamma_{\text{sh}}^{(a)}$ |
|--------|-----------------------------------------------|-----------------------------|--------------------------|
| V1324 Sco | -2.2±0.4                                        | -1.5±0.4                     | 0.13($f_{\text{sh}}/(0.1)^{-1}$) |
| V339 Del | -3.5±0.2                                        | -2.8±0.2                     | 0.016($f_{\text{sh}}/(0.1)^{-1}$) |

$^{(a)}$ Minimum fraction of shock powered placed into non-thermal particles (eq. [14]), calculated assuming radiative shocks and (conservatively) $\epsilon_{\gamma} = 0.2$ ($\gamma_{\text{sh}}^{(a)}$).

$^{(b)}$ Fraction of optical luminosity powered by shocks (eq. [13]), calculated assuming radiative shocks and (conservatively) $\epsilon_{\gamma} = 0.2$.}

4 DISCUSSION AND CONCLUSIONS

The high densities within the nova ejecta at early times coincide with the observed gamma-rays imply that the shocks responsible for this emission are likely to be radiative (§2.1). Though not observed directly, thermal X-rays from the shocks are absorbed by the ejecta and reprocessed to optical frequencies (§2.2). Based on nova Del did show X-ray emission at later times, after the epoch of peak gamma-ray and optical emission. This presumably occurred after the nova had expanded or became sufficiently photo-ionized to reduce the column of neutral absorbing gas (see below).
Gamma-ray novae as probes of particle acceleration

Figure 2. Total optical flux (top panel) and best-fit blackbody temperature (bottom panel) as a function of time since outburst for novae V1324 Sco (blue) and V39 Del (red). Time is measured starting on June 1, 2012 for Nova Sco (beginning of optical outburst), and starting on August 16, 2013 for Nova Del (epoch of first gamma-ray detection, within days of the optical rise).

the observed values of $L_\gamma/L_{\text{opt}}$ for the highest gamma-ray efficiencies $\epsilon_{\text{sh}} \sim 0.1$ expected in hadronic scenarios, we conclude that a significant fraction $f_{\text{sh}} \gtrsim 10$ per cent and $\gtrsim 1$ per cent of nova optical emission in nova Sco and nova Del, respectively, is powered by shocks. These limits approach unity for more realistic efficiencies $\epsilon_{\text{sh}} \lesssim 0.01$ expected from the difficulty of accelerating protons given the expected geometry of the magnetic field in the pre-shocked ejecta (see below).

This shock-powered emission mechanism is similar to that at work in some core collapse supernovae when the ejecta from the exploding star interacts with dense circumstellar matter (e.g. Chevalier & Fransson 1994, Larsson et al. 2011), which in extreme cases powers the most luminous supernovae yet discovered (Smith et al. 2007). Our finding that a large fraction of nova optical emission is powered by shocks requires a revision of standard models that instead assume nova thermal emission results from the direct outwards transport of thermal energy released in the white dwarf envelope by nuclear burning (e.g. Hillman et al. 2014). The ultimate energy source (the thermonuclear runaway) is not in question, but shock-powered emission taps into the differential kinetic energy of the outflow instead of its thermal content advected or diffused from small radii. As suggested by Metzger et al. (2014), shock-powered optical emission might also help explain some of the irregular light curve shapes, including plateaus and secondary maxima, observed in some novae (e.g., Strope et al. 2010). A similar mechanism of internal shocks within a time-variable outflow may power an optical transient from the remnant produced by the merger of binary white dwarfs (Beloborodov 2014).

If the majority of nova optical light is powered by radiative shocks, then the ratio of optical and gamma-ray fluxes should be relatively constant in time, assuming that the microphysical parameters of the shocks remain constant (and that the conditions $\tau_{\text{p-p}} > \gamma c$ and $\tau_{\text{cool}}/t < 1$ remain satisfied in hadronic and leptonic scenarios, respectively). Nova Sco is consistent with a temporally constant ratio $L_\gamma/L_{\text{opt}}$, while nova Del shows evidence for the gamma-ray emission becoming relatively stronger with time (Fig. 3). The apparent increase of $L_\gamma/L_{\text{opt}}$ in nova Del could result from the reverse shock thermal emission becoming less radiatively efficient with time, i.e. $\chi \propto t$ (eq. 7). Alternatively, it may reflect the early transition in this event to a nebular phase (as occurred on
Day 11, [Skopal et al. 2014], resulting in a loss of flux out of the measured optical band and into the UV. If nova shocks are indeed radiative, then we predict that $\epsilon_r/\epsilon_{\text{opt}}$ should not exceed ~0.06 at any time, since this represents the limit of shock-dominated emission $f_{28} = 1$ (eq. [13]) for realistic maximum efficiencies $\epsilon_t \lesssim 0.3$, $\epsilon_{\text{th}} \lesssim 0.2$.

The gamma-ray to optical flux ratio also places a lower limit on the acceleration efficiency $\epsilon_{\text{th}}$ of relativistic non-thermal particles at non-relativistic shocks. For nova Sco and nova Del we constrain $\epsilon_{\text{th}} \gtrsim 10^{-2}$ and $10^{-3}$, respectively (Table 1). Unlike traditional studies modeling particle acceleration in supernova remnants, nova shocks evolve in real-time, in principle allowing for the study of time-dependent effects. Models used to constrain particle acceleration in supernova remnants also sometimes require assumptions about the escape fraction of relativistic particles, e.g. if gamma-rays are produced by the collision of relativistic protons with nearby molecular clouds of irregular geometry. By contrast, nova ejecta serves as a relatively efficient hadronic “calorimeter” since relativistic protons are probably trapped when gamma-rays are observed (eq. [17]).

Neither hadronic nor leptonic scenarios for nova gamma-ray emission can be ruled out by modeling the $\gamma$-ray spectrum alone ([Ackermann et al. 2014]). However, the tension between the value $\epsilon_{\text{th}} \gtrsim 10^{-2}$ we find is in nova Sco and the much lower electron acceleration efficiency $\epsilon_{\text{th}} \lesssim 10^{-3}$ inferred from observations of supernova remnants (e.g. [Morlino & Caprioli 2012]) and theoretical modeling ([Kato 2014], [Park et al. 2014]) appears to disfavor the leptonic scenario. Leptonic models also require a flat injected spectrum $\epsilon_t = 2$ to be energetically feasible (eq. [20]). In contrast, ion acceleration efficiencies as high as $\epsilon_{\text{th}} \sim 0.1$ are found for non-relativistic shocks by hybrid PIC simulations, but only in cases when the upstream magnetic field is parallel to the shock normal ([Caprioli & Spitkovsky 2014]), for perpendicular fields essentially no ion acceleration is seen, also inconsistent with our lower limits on $\epsilon_{\text{th}}$. Hadronic scenarios involving ion acceleration at quasi-parallel shocks thus appear to be the most viable source of relativistic particles responsible for nova gamma-rays.

However, this conclusion is problematic. Most of the power dissipated in the system occurs at the reverse shock, making it the most natural location for particle acceleration. If the fast outflow is ejected over a timescale which is long compared to the dynamical timescale of the nova, it will therefore become dominated by its azimuthal component at the much larger radii where the reverse shock occurs (Fig. 1). Such a field, being perpendicular to the radial shock normal, is not conducive to ion acceleration.

This mystery could be resolved if particle acceleration occurs at localized regions of quasi-parallel shocks, creating special regions of efficient acceleration ($\epsilon_{\text{th}} \sim 0.1$) with $\epsilon_{\text{opt}} = 0$ across the larger bulk, thereby resulting in an average value of $\epsilon_{\text{th}} \sim 10^{-3} - 10^{-2}$ consistent with observations. These localized regions of conductive field geometry could be caused by global asymmetries in the ejecta, such as oblique shocks occurring where the edge of the slow ejecta torus meets the faster bipolar wind (Fig. 1). Alternatively, inhomogeneity of the nova ejecta (“clumpiness”) could also result in localized regions of oblique shocks between the clumps. Indeed, low ejecta filling factors are inferred observationally by emission line modeling of the nebular phase (e.g. [Shore et al. 2013]). Radial fields may also be produced near the shock as the result of Rayleigh-Taylor instabilities at the interface between the fast now outflow and the dense central shell (as may occur in supernova remnants; e.g. [Blondin & Ellison 2001]).

Leptonic versus hadronic models for nova gamma-ray emission may be further distinguished by their predictions for non-thermal X-ray emission. In nova Sco, Swift XRT observations placed an upper limit of $F_X < 9 \times 10^{-14}$ erg s$^{-1}$ cm$^{-2}$ on the 0.3–10 keV flux on days 22–41 ([Page et al. 2012]), i.e. three orders of magnitude less than the simultaneous LAT-detected flux of $F_\gamma \sim 2 \times 10^{-10}$ erg cm$^{-2}$ s$^{-1}$. Such a low value of $F_X/F_\gamma$ appears to be barely compatible with the leptonic scenario, even when considering the most conservative case of $\epsilon_t = 2$ and a slow cooling spectrum $vF_\gamma \propto v^{1/2}$ from keV to ~GeV energies (eq. [19]). The hadronic scenario fares even worse: secondary $e^\pm$ pairs from $\pi^0$ decay carry a similar total energy to that released in $\gamma$-rays from $\gamma$0 decay. These pairs’ energies, ~0.1 – 1 GeV, correspond to Lorentz factors $\gamma_\nu \sim 100 – 1000$ that will upscatter ~eV optical to energies ~10 – 1000 keV with reasonably high efficiency (eq. [18]), again overpredicting the Swift upper limits.

This dearth of early ~keV emission from nova Sco is consistent with the high column of neutral, X-ray absorbing gas ahead of the shock at times coincident with the gamma-ray emission ([Metzger et al. 2014]). However, higher energy X-rays $\gtrsim 10$ keV should escape without attenuation, motivating one to consider the prospects for nova detection with NuSTAR ([Harrison et al. 2013]), which has unprecedented X-ray sensitivity above 10 keV. If a fraction $f_X$ of the LAT luminosity is radiated as X-rays of energy $\epsilon_X \sim 30$ keV, this results in a number flux of $N_X \sim f_X F_\gamma / \epsilon_X \propto 4 \times 10^{-2} (f_X / 0.1)$ s$^{-1}$ for a typical value $F_\gamma \sim 10^{-10}$ erg cm$^{-2}$ s$^{-1}$ ([Ackermann et al. 2014] and the NuSTAR effective area $A_{\text{eff}} \sim 200$ cm$^2$ [Harrison et al. 2013], their Fig. 2). Given the NuSTAR estimated background of $N_X \sim 2$ counts s$^{-1}$ in this energy range ([Harrison et al. 2013], their Fig. 11), we estimate that the integration time required for a $3\sigma$ detection is approximately $N_{\text{crit}} / N_X^2 \sim (10 f_X / 0.1)$ s$^{-1}$. Reasonable constraints on $f_X$ could thus be achieved for a modest 10 ks exposure, while sensitivity to the minimum value of $f_X \sim 0.01$ predicted by the leptonic model could be achieved for an ambitious ~Ms observation. Earlier observations are more likely to yield a detection, as the accelerated electrons are more likely to be fast-cooling at early times (eq. [18]).

[Orio et al. 2014] reported NuSTAR observations of the fast symbiotic nova V745 Sco within 10 days of the outburst, placing an upper limit of $\lesssim 10^{-11}$ ergs cm$^{-2}$ s$^{-1}$ on the ~50 keV emission. Gamma-ray emission was detected by LAT at $2 \sim 3\sigma$ significance and a flux of $\sim 10^{-10}$ erg cm$^{-2}$ s$^{-1}$, but lasting only for 2 days near the onset of the outburst ([Cheung et al. 2014], approximately a week before the NuSTAR observations and when (perhaps not coincidentally in light of our results on shock-powered optical emission) the optical light curve was at its maximum. By the time of the NuSTAR observations, the optical flux had decreased by a factor of ~100 from its peak value, such that if $F_{\gamma}/F_{\text{opt}}$ had remained approximately constant in time (as in the nova in our sample; Fig. 3), then the NuSTAR upper limit corresponds to $f_X < 10$ in the notation above, i.e. not particularly constraining. We strongly encourage future NuSTAR observations of gamma-ray novae, ideally closer to the optical peak and coincident with LAT detections.

The sample of gamma-ray novae should expand rapidly in the next few years thanks to anticipated enhancements in the sensitivity of Fermi LAT resulting from improvements in its ability to perform low-level event reconstruction. The method for jointly analyzing optical and gamma-ray data employed in this paper will thus soon be applicable to a greater sample of events. Our results highlight the importance of obtaining broad (ideally bolometric) frequency coverage of nova light curves, including near infrared and ultraviolet wavelengths, at times coincident with the gamma-ray emission.
The utility of our model rests heavily on the assumption of radiative shocks, which depends sensitively on the quantity of mass and velocity of the nova ejecta, and how it is partitioned between the fast and slow components. Additional radio monitoring of gamma-ray novae is thus also key to better constraining the ejecta mass from these systems and its geometry, thereby strengthening or refuting the argument for radiative shocks.

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