Annihilation of positrons from $^{22}\text{Na}$ in novae

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Abstract

We explore for the first time effects of the magnetic field on the escape of $^{22}\text{Na}$ positrons and on the flux evolution of annihilation 511 keV line in novae. It is shown that for the white dwarf magnetic field of $\sim 10^6 \text{G}$ the field of the expanding nova shell is able to significantly impede positrons escape and increase the time of the nova emission in 511 keV up to hundreds days.

Keywords: stars — novae — nucleosynthesis
1 Introduction

Nova phenomenon is caused by a thermonuclear runaway (TNR) in a hydrogen shell accumulated on a white dwarf (WD) by an accretion in a binary system with a normal companion (usually red dwarf). Despite a general understanding of the nova outburst, widely accepted models for major processes involved in the nova phenomenon are lacking. Particularly, a mixing mechanism of the WD matter with the hydrogen shell that is crucial for the TNR is still debatable. Moreover, it is not clear, whether the nova ejects less or more material than accumulated between consecutive outbursts (Epelstain et al. 2007; Starrfield et al. 2020). Uncertainties remain also in the description of the shell ejection and shell morphology.

Given the theoretical status, a significant role in the theoretical progress belongs to observations. This is supported by two fundamental facts laid in the basis of the nova theory: (i) novae are binary system containing the WD (Kraft 1964); (ii) nova shell material is enriched by a factor of 100 with CNO elements (Mustel & Baranova 1966), which indicates a crucial role of mixing between accreting and WD matter for the nova energetics (Sparks et al. 1976). It should be emphasised that apart from novae related to the TNR on the CO WD there is an abundant class of novae related to the TNR on ONe dwarfs (Starrfield et al. 1986). This conclusion is prompted by the detection of Ne overabundance in spectra of some novae (Ferland & Shields 1978). Noteworthy, according to the interesting conjecture (Shara & Prialnik 1994) neon novae could arise from CO WD with outer layers significantly enriched by Ne and Mg due to the preceded hydrogen accretion with the rate of $\sim 10^{-6} M_\odot \text{yr}^{-1}$.

The nova TNR results in the synthesis of radioactive isotopes which decay with half life from hundreds of seconds ($^{13}\text{N}$) till several years ($^{22}\text{Na}$). As far back as half century Clayton & Hoyle (1974) emphasised the importance of observations of gamma-lines from novae. The isotope composition of nova shell depends on WD type (CO or ONe) and on some poorly fixed parameters, particularly, the WD mass, the accretion rate and the admixing of the WD material into the hydrogen shell. A detection of gamma-lines from novae would permit us to determine the mass of a certain isotope and thus to obtain useful constraints on outburst models.

Highly prospective for gamma-line observations is the isotope $^{22}\text{Na}$ with the half life of 2.6 years that provides a possibility to apply a long integration time. The outburst theory predicts a high $^{22}\text{Na}$ abundance in neon novae, $\sim 10^{-8}$ by mass (Denissenkov et al. 2014; Kudryashov 2019), which is by a three order higher compared to that of CO novae. The $^{22}\text{Na}$ decay is accompanied by the emission of the 1275 keV gamma-line, the 511 keV annihilation line, and three-photon annihilation continuum. Five neon novae have been observed with the gamma-ray telescope COMPTEL resulting in the upper limit of the $^{22}\text{Na}$ mass between $3 \cdot 10^{-8} M_\odot$ and $2 \cdot 10^{-6} M_\odot$ based on the 2$\sigma$ upper limit of the flux in the 1275 keV line (Iyudin et al. 1995). These upper limits do not contradict to theoretical $^{22}\text{Na}$ abundance in current models of...
neon novae provided a total ejecta mass is $< 3 \cdot 10^{-5} M_\odot$. Subsequently, Iyudin (2010) reported on the COMPTEL detection of 1275 keV line at the confidence level of 4$\sigma$ from the slow nova V723 Cas of 1975, which is rather unexpected result for these category of novae.

The model luminosity of novae in the annihilation 511 keV line of $^{22}$Na has been addressed by Gomez et al. (1998) who demonstrated that already a week after the TNR the nova shell becomes transparent for positrons that escape without annihilation. If this is the case, the possibility to detect 511 keV line from $^{22}$Na in novae becomes doubtful. On the other hand, the conclusion about the high transparency of the nova shell for $^{22}$Na positrons is made neglecting a possible magnetic field in the shell. It may well be that a moderate value of the magnetic field would be enough to essentially impede the positrons escape and thus to prolong the time of a significant nova luminosity in 511 keV line. In this case the detection possibility for 511 keV line from neon novae would become feasible.

This paper for the first time explores the issue, to which extent the inclusion of the magnetic field could affect the escape of $^{22}$Na positrons from the shell and the evolution of 511 keV line flux from the neon nova. The shell model with the frozen-in magnetic field and the prescription for the annihilation are considered in Section 2. The computations of the 511 keV flux evolution for different values of relevant parameters are presented in Section 3.

2 Model overview

2.1 Nova shell with frozen-in magnetic field

The issue of the shell structure and kinematics for different subclasses of novae currently is not completely clear. The most natural is the assumption that the ejecta forms in the regime of the optically thick wind (Ruggles & Bath 1979; Kato & Hachisu 2007). The idea is based on the fact that the major stage of the TNR is longer compared to the hydrodynamic time at the WD surface. Following this arguments we assume that the ejected shell is produced by the wind with the constant mass loss rate and constant velocity of $v \sim 1500$ km s$^{-1}$ that forms at the outer boundary of the expanded envelope with the radius $r_1 \sim 2GM/v^2 \sim 10^{10}$ cm, where $M \sim 1 M_\odot$ is the WD mass. The convenient wind parameter is the wind kinetic luminosity $L_w = 0.5wv^3$, where $w = M/v = 4\pi r^2\rho$ is the wind density parameter. The $L_w$ value should be of the order of the Eddington luminosity, $1.26 \cdot 10^{38} (M/M_\odot)$ erg s$^{-1}$. For the standard model we adopt $L_w = 10^{38}$ erg s$^{-1}$. It takes thus $t_w = M_s/\dot{M} \sim 26$ days for the wind to form ejecta with the mass of $M_s \sim 10^{-5} M_\odot$.

For the adopted wind velocity the ejecta perturbation from the binary components are insignificant. Indeed, for the typical orbital period of 3.5 h, the WD mass of $M_1 = 1 M_\odot$, and the red dwarf mass of $M_2 = 0.5 M_\odot$ one
gets the large semi-axis $a = 9 \cdot 10^{10}$ cm, and orbital velocities of components
$v_1 = 150 \text{ km s}^{-1}$ and $v_2 = 300 \text{ km s}^{-1}$. The maximal variation of the wind
velocity caused by the WD motion is about 10% close to the orbital plane.

The variation of the density and velocity caused by the gravitation from the
secondary is also small due to high wind velocity. Below we adopt that the
shell expands spherically with the constant velocity. Apropos, the neon nova
V382 Vel (Takeda & Diaz 2019) shows the spherical shell expanding with the
velocity of 1200 km s$^{-1}$. A possible deviation from the spherical symmetry
(cf. Chomiuk et al 2020) is discussed in the final section.

In our picture the hydrogen shell before the outburst resides in the mag-
netic field of the WD with the average strength of $B_0 = 10^6$ G. This value is
rather moderate compared to the field of polars and intermediate polars of
$3 \cdot 10^6 - 2 \cdot 10^8$ G. (Ferrario et al. 2020). The assumption on the magnetic field
value is also consistent with the magnetic momentum of the WD of classical
nova 1934 DQ Her, viz., $(1 - 3) \cdot 10^{32}$ G cm$^3$ (Patterson 1994).

For the WD radius of $\sim 4 \cdot 10^8$ cm this value corresponds to the magnetic field on the WD
surface of $(1 - 3) \times 10^6$ G.

The TNR is accompanied by a vigorous convection all over the hydrogen
shell (Casanova et al. 2018; Starrfield et al. 2020). The adopted mag-
netic field ($\sim 10^6$ G) is weak in the energetic and dynamic respect, so the
convection unavoidably results in the creation of the entangled field in the
convective shell with the thickness of $h \approx 250$ km in the case of a ONe WD
of $1.25 M_\odot$ (Casanova et al. 2018). We neglect both the amplification of the
magnetic field by the convection and the field dissipation at the small turbu-
lent scales. The average magnetic field in the convective shell is assumed
to be equal to the average field on the WD surface $B_0 = 10^6$ G.

Our scenario suggests that the nova shell forms in three stages: (i) the
TNR accompanied by the convection, (ii) the hydrogen shell expansion, and
(iii) the wind outflow. The envelope expansion at the stage (ii) from the
WD radius of $r_0 = 4 \cdot 10^8$ cm up to the radius of $r_1 = 10^{10}$ cm results in
the weakening of the frozen-in magnetic field down to $B_1 = B_0(r_0h/r_1^2)$ that
is adopted as the initial field of the wind at the level of its formation. The
subsequent evolution of the field in the wind proceeds differently for the radial
($B_r$) and tangential ($B_t$) components. The radial field decreases as $B_r =
B_1(r_1/r)^2$, whereas the tangential as $B_t = B_1r_1^2/(r\Delta r)$, where $\Delta r = vt_w$ is the
shell thickness that remains constant. At the late time $t \gg \max(r_1/v, \Delta r/v)$
the entangled field $B$ in the shell is dominated by the tangential components,
so we adopt $B = B_t$. We neglect the shell radial spread-out due to thermal
velocities and ignore a possible shell compression by the fast rarefied wind
at the nebular stage of the nova evolution ($t \gtrsim 100$ d).

2.2 Positron annihilation and escape

The $^{22}$Na isotope decays via the positron emission with the probability of 0.9
and the $e^-$ capture with the probability of 0.1. In both cases the daughter nu-
cleus $^{22}$Ne ends up in the excited state with a subsquent emission of 1275 keV
gamma-quantum (Firestone et al. 1999). Positron energies are distributed in the range of 0 - 546 keV. Our description of positron annihilation takes into account the numerical result that the probability of a free annihilation is \( p_f = 0.1 \) (Krannell et al. 1976; Leising and Clayton 1987), whereas 90\% of positrons annihilate via the positronium formation at thermal energies.

A positron slows down via ionization losses that depend on the matter composition and the ionization degree. We take into account only hydrogen (X = 0.7) and helium and adopt that the hydrogen ionization fraction is \( y_f = 0.35 \) and the fraction of bound electrons is \( y_b = 0.5 \). The positron slowdown is calculated according to the Bethe formula for ionization losses that is splitted into terms related to neutrals and free electrons (Ahlen 1980)

\[
\frac{dE}{d\mu} = -\frac{4\pi e^4 N_A}{mc^2 \beta^2} \left( y_b \ln \frac{2mc^2 \beta^2 \gamma^2}{I} + y_f \ln \frac{2mc^2 \beta^2 \gamma^2}{\hbar \omega_0} \right),
\]

where \( \mu \) is the surface density along the positron path (g cm\(^{-2}\)), \( e \) is the elementary charge, \( N_A \) is the Avogadro number, \( m \) is the electron mass, \( c \) is the speed of light, \( \beta = u/c \) (\( u \) is the positron velocity), \( \gamma \) is the positron Lorentz factor, \( I = 29 \) eV is the average ionization potential for the adopted mixture of hydrogen and helium and the ionization fraction, \( \omega_0 \) is the plasma frequency.

The positron range for a given energy is calculated taking into account the energy distribution (Xiao et al. 2018). As an illustration, for the initial energy of 200 keV and 400 keV the positron range at \( t = 100 \) d (\( \omega_0 \) depends on the shell age) is \( R = 0.012 \) g cm\(^{-2}\) and \( R = 0.035 \) g cm\(^{-2}\), respectively. These values differ from the average range of 0.1 g cm\(^{-2}\) estimated earlier (Leising & Clayton 1987). The difference presumably is related to the adopted ionization degree. Indeed, we find for the neutral gas the comparable value \( R \sim 0.1 \) g cm\(^{-2}\). The slowdown due to free electrons is efficient because of the significant energy loss in Coulomb collisions.

The escape probability is determined by the ratio of the total surface density along the random walk \( \mu \) (g cm\(^{-2}\)) in the shell and the range: \( p_{esc} = (1 - p_f) \exp \left( -\mu / R \right) \). The \( \mu \) value is assumed to be a product of the total geometrical length of the random walk \( l \) and the density at the radius of the positron creation. An alternative version, based on the average density of the shell, produces the similar result. With the probability of \( (1 - p_{esc}) \) the positron radiatively recombines with the branching ratio of 1/4 and 3/4 into the singlet and triplet states, respectively, in line with statistical weights. The subsequent kinetics depends on the balance between annihilation rates \( t(1S) = 1.2 \cdot 10^{-10} \) s, \( t(3S) = 1.38 \cdot 10^{-7} \) s (Karshenboim 2004), on the one hand, collisional transition rates \( 1S \leftrightarrow 3S \) and the photoionization rate, on the other hand; these processes are taken into account. The photoionization might affect the ratio between two- and three-photon annihilation at the early epoch. However, already the day after the TNR, the photoionization rate \( P \) in the wind turns out already small enough, \( P \tau(3S) < 0.1 \), so at
the subsequent stage the ratio between the two-photon and three-photon annihilation does not deviate from 1/3.

A frozen-in entangled magnetic field impedes the positrons escape and this should change the evolution of the annihilation line from the $^{22}\text{Na}$ decay compared to the case without magnetic field. The positron diffusion in the entangled field is unlike the case of a uniform field. In the latter case the diffusion across the magnetic field is strongly inhibited in the absence of collisions and field perturbations, whereas along the field the positrons can propagate freely. In the entangled field the diffusion across the field can occur because at some coherent length $l_B$ the field significantly changes that results in the positron drift across the field. The diffusion in the entangled field looks therefore like an excursion along the field on the scale $l_B$ and a drift across the field. Although there are approaches to the description of the diffusion in the entangled field based on the scale $l_B$ and Larmor radius (Narayan & Medvedev 2001; Chandran & Maron 2004) we do not see a possibility to apply these approaches without three-dimensional modelling of the magnetic field and the positron random walk.

Here we use the diffusion treatment in terms of the average Bohm diffusion coefficient $D_B = (1/3) R_L u$, where $R_L$ is the Larmor radius and $u$ is the positron velocity. For the diffusion coefficient we adopt parametric representation $D = \xi D_B$ with the coefficient $\xi \geq 1$. This approach with $\xi$ in the range of 1...10 is adopted, e.g., for the diffusion in the chaotic field when addressing the particle acceleration in supernova remnants (Marcowith et al. 2006).

The characteristic diffusion time for the positron in the expanding shell with the width $d$ is $t_{\text{dif}} \sim d^2 / D$ while the total surface density ignoring a possible annihilation is $\mu = p u t_{\text{dif}}$. This value combined with the positron range $R$ due to ionization losses determines the average escape probability of positrons $p_{\text{esc}}$ introduced above. In the adopted model of the nova shell and the magnetic field evolution for the fixed initial value at the WD surface the probability $p_{\text{esc}}$ is determined by the parameter $\xi$ only. The adopted form of the diffusion coefficient is equivalent to a random walk with the step $l$ being proportional to the shell radius $R_s$. Indeed, in our model $l \propto R_L \propto B^{-1} \propto R_s$. It is interesting that in this case the diffusion coefficient formally corresponds to the approximation $\lambda = l/R_s = \text{const}$ used earlier in the description of the positrons diffusion in the entangled field of SN Ia ejecta (Churazov & Khabibulin 2018).

3 Results

A standard model suggests the shell with the mass $M_s = 10^{-5} M_\odot$, created by the wind with the velocity $v = 1500 \text{km s}^{-1}$ and the kinetic luminosity $L_w = 10^{38} \text{erg s}^{-1}$. The mass fraction of $^{22}\text{Na}$ in the shell is $10^{-3}$ and this isotope is assumed to be mixed in the shell homogeneously with the constant fraction in any point. The wind presumably forms at the radius $r_1 = 10^{10} \text{cm}$;
the result is weakly sensitive to the $r_1$ value. The average magnetic field at the WD surface is $B_0 = 10^6$ G.

The flux of directly escaped 511 keV quanta from the $^{22}$Na decay in nova at the distance of 1 kpc is shown for $\xi = 1$, 10, and 100, as well as for the case without magnetic field (Fig. 1). The plot also shows the flux in the 1275 keV line of $^{22}$Na. As expected, the magnetic field essentially changes the evolution of the annihilation line, so its detection on the time scale of one year becomes feasible. Remarkably, in the absence of the magnetic field the flux in the 511 keV line does not reach the maximal possible value because of the early positrons escape. For models with the magnetic field and $\xi = 1, 10$, and 100 the fraction of escaped positrons is 0.41, 0.64, and 0.8, respectively.

To remind, we do not take into account the magnetic field amplification by the convection that can be significant. Indeed, the MHD 3D-modelling of the solar convection demonstrates that a weak magnetic field can be amplified to the value corresponding the Alfvén velocity $V_A \sim 0.45 u$ (Cataneo et al. 2003). This relation suggests that the convection during the nova outburst with the convective velocity $u \sim 100 \text{ km s}^{-1}$ (Casanova et al. 2018) in the layer with the density of $\sim 3 \cdot 10^2 \text{ g cm}^{-3}$ (Denissenkov et al. 2014) is able to amplify the field upto $\sim 10^8$ G. For $B_0 > 10^6$ G positrons trapping will be even more effective.

The effect of other parameters — the shell mass, the shell velocity, and the width of the convective shell — in the evolution of the annihilation line from $^{22}$Na decay is shown in Figure 2. As a fiducial model we take the standard model with $\xi = 10$. The twice as large shell mass increases the flux and the time of the flux maximum. The lower expansion velocity by a factor of 1.5 increases the positron trapping and also decreases the time of the flux maximum. The increase of the width of the convective shell results in the larger field at the late time and therefore more efficient trapping. We do not demonstrate the case with the kinetic wind luminosity of $2 \cdot 10^{38}$ erg s$^{-1}$ because the only effect of this variation is the decrease of the maximum age. Versions with the different magnetic field are not shown as well because effect of smaller/larger value of the initial field is similar to the effect of larger/smaller value of the $\xi$ parameter by the same factor. This suggests that the effect of the field and the diffusion coefficient is actually determined by the combined parameter $B_0/\xi$.

4 Discussion and conclusions

The paper has been aimed at the study of effects of the magnetic field in the nova shell on the escape of positrons from the $^{22}$Na decay and on the flux evolution in the annihilation line of 511 keV. We present the model for the positron annihilation in a nova shell that takes into account a magnetic field. The modelling shows that for the sensible value of the average magnetic field on the WD surface of $10^6$ G, the frozen-in magnetic field is able to essentially impede the positrons escape and thus to prolong the luminosity
of the annihilation line from $^{22}$Na up to hundreds days.

Results, however, depend not only on the strength of magnetic field, but on the diffusion coefficient as well. The dimensionless diffusion coefficient $\xi$ used above is in the range $1 \leq \xi \leq 100$, that includes the range of $1 \leq \xi \leq 10$ in models of the diffusive shock particle acceleration constrained by the non-thermal X-ray brightness at the limb of young supernova remnants (Marcowith et al. 2006).

Yet there remains an open question: to which extent an assumption of the homogeneous chaotic magnetic field is applicable to nova shells. It may well be that the turbulent convection results in the intermittent chaotic field with regions of strong and weak field. If this the case and the filling factor of regions with the weak field is significant, the diffusion and escape of positrons could be more efficient compared to the homogeneous case. Partially this situation is taken into account by the large value of $\xi = 100$. However, actually this does not resolve the issue of a possible intermittency of the chaotic field and the problem requires separate study. It is noteworthy that a simultaneous detection of $1275$ keV and $511$ keV lines will become a valuable tool for the diagnostics not only of the synthesised $^{22}$Na mass but also of the parameter $B_0/\xi$ that determines the positron escape.

Shells of historical novae have clumpy structure, e.g., nova 1970 FH Ser (Gill & O'Brien 2000), and this, at first glance, is in conflict with our model of the spherical shell. However, the effect of the clumpy structure on the $^{22}$Na positron escape depends on the stage at which this clumpiness does emerge. There is no reason for the clumpiness to appear at the stage of the optically thick wind. More naturally to admit that the shell experiences a fragmentation significantly later under the dynamical effect of a high velocity rarefied wind (Cassatella et al. 2004). The high velocity wind is formed by the remnant of the hot envelope of the WD due to the radiative acceleration. A signature of the wind with the velocity of $\sim 3000$ km s$^{-1}$, versus the main shell velocity of $\sim 1800$ km s$^{-1}$, has been observed in nova V1974 Cyg (Cassatella et al. 2004). Based on the WR-star wind one expects that the kinetic luminosity of the high-velocity wind ($L_{w,f}$) should be $\sim 10^{-2}$ of the radiative luminosity of the WD (Sander et al. 2020), i.e., $L_{w,f} \sim 10^{36}$ erg s$^{-1}$ in our case. One can estimate the fragmentation time assuming the fast wind velocity of $3000$ km s$^{-1}$. The major fragmentation mechanism is the Rayleigh-Taylor instability initiated by the shell acceleration driven by the fast wind. The fragmentation time scale exceeds the initial phase of the acceleration determined by the shell compression due to the shock propagation. For the shell mass of $M_s = 10^{-5} M_\odot$ and the thickness $d = M_s/w$ formed by the wind with the kinetic luminosity $L_w = 10^{38}$ erg s$^{-1}$ the compression by the fast wind takes the time $t_c \sim d/D$, where $D$ is the shock wave speed determined from the momentum conservation $D = (\rho_f/\rho)^{1/2}(v_f - v) = 53$ km s$^{-1}$, $\rho_f$ is the density of the fast wind, $v_f$ is its velocity, while the rest of values are related to the nova shell. In the considered case $t_c \sim 2$ years. This time refers to the initial stage of the nova shell fragmentation that requires some-
what longer time, of the order of $t_f \sim 4t_c$ (Klein et al. 1994), i.e., about 8 years. The found fragmentation time indicates that the approximation of the spherical shell on the time scale of one year generally does not contradict to the observed clumpy shells of historical novae.

Although novae have been considered as a possible source of Galactic positrons, in effect their contribution to the Galactic 511 keV gamma-line is unlikely significant even for 100% positron escape. With the Galactic nova rate of $50 \text{ yr}^{-1}$, neon novae fraction of 1/4, nova shell mass of $10^{-5} M_\odot$, and $^{22}\text{Na}$ mass fraction of $10^{-3}$ the Galactic production rate of positrons by novae is only of $2 \cdot 10^{41} \text{ s}^{-1}$ that is by a two order lower compared to the Galactic annihilation rate of $5 \cdot 10^{43} \text{ s}^{-1}$ (Siegert et al. 2016).

It remains unclear, whether WD of novae are always magnetic with the field $\gtrsim 10^6 \text{ G}$? The recent study of the cataclismic variables from the Gaia DR2 within 150 pc (Pala et al. 2020) shows that 36% of them possess mega-gauss fields. Given the fact that the field of $\sim 10^6 \text{ G}$ is a border value for the cataclismic variables to be classified as an intermediate polar with the accretion on the poles, a significant fraction of WD in cataclismic variables could have magnetic fields $<10^6 \text{ G}$ and thus not show signatures of the intermediate polar. Yet in this case the magnetic field can be nevertheless strong enough to impede $^{22}\text{Na}$ positrons escape. For example, for the field of $10^5 \text{ G}$ and $\xi = 10$ the situation with the trapping of $^{22}\text{Na}$ positrons in effect is equivalent to the case of $B_0 = 10^6 \text{ G} \dot{E} \xi = 100$ (Fig. 2) that demonstrates the efficient positron annihilation on the time scale of $\sim 100$ days.

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Figure 1: Flux of gamma-quanta in 1275 keV line (thin line) and in the annihilation line 511 keV from the $^{22}$Na decay in ONe nova shell at the distance of 1 kpc. Numbers next to lines indicate the value of $\xi$. Dotted line shows the model without magnetic field.
Figure 2: Flux in the annihilation line 511 keV from the $^{22}$Na decay in the standard model with $\xi = 10$ (thick line) and in models with the variation of a certain parameter (shell mass, velocity, and thickness of the convective shell on the WD surface).