LETTERS

Optical spectrum of the unusual supernova remnant G109.1 − 1.0

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The new supernova remnant G109.1 − 1.0 has important X-ray and radio parallels to the SS433–W50 pair. We have obtained a spectrum of the faint optical filaments associated with the large shell-like radio source. The spectrum shows strong [S II] λλ 6,717, 6,731 emission relative to Hα, which is characteristic of shock-heated gas and which confirms that the filaments are part of the supernova remnant, and is generally similar to spectra of the Cygnus Loop. By assuming pressure equilibrium between the optical filaments and the interior of the remnant, we find an initial energy of 3×10^{51} erg, which is higher than that found from X-ray measurements. When compared with W50, the G109.1 − 1.0 remnant has filaments of higher density, and weaker [N II] emission.

G109.1 − 1.0 is a newly discovered galactic supernova remnant (SNR) which has been found to be a source of both X-ray[1] and non thermal radio radiation[2]. The X-ray image obtained with the Einstein Observatory shows an incomplete shell of emission with a diameter of ~36 arc min. A strong, compact X-ray source with no known optical counterpart appears at the centre of the X-ray shell. The radio observations show a smaller but complete shell with a spectral index $\alpha \approx -0.5$ ($S_\nu \propto \nu^\alpha$), which is typical for SNRs. These observations also show enhancements in the radio shell along a NE–SW line which approximately line up with a jet-like feature seen both in radio and X ray. This morphology has strong parallels with the association[3–5] between the SNR W50 and the bizarre object SS433.

The distance of G109.9 − 1.0 has been estimated from its radio surface brightness, Σ, assuming a relationship between Σ and the diameter (D). However, because the angular size of the SNR is different in X-ray and radio maps, and because of differences in the various Σ–D relations, the distance is only roughly known, with estimates ranging from 3.6 to 5.2 kpc. Using the X-ray angular diameter of ~36 arc min, and parameterizing to distance $d = 4$ kpc, we find the diameter of the SNR is

$$D = 41.9 \left(\frac{d}{4\ \text{kpc}}\right) \text{pc}$$

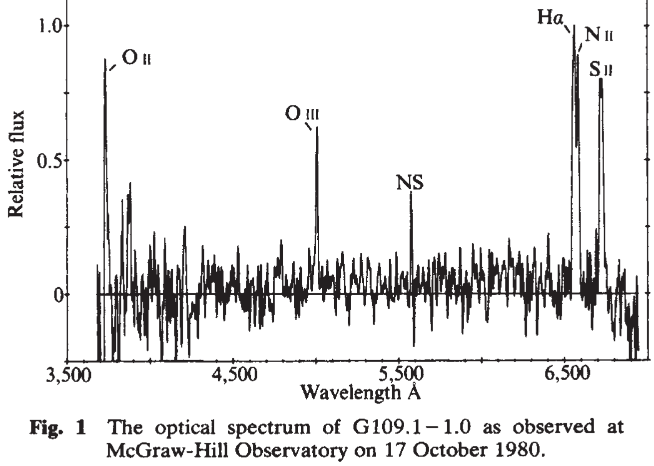

Fig. 1 The optical spectrum of G109.1 − 1.0 as observed at McGraw-Hill Observatory on 17 October 1980.

Table 1 Line strengths for G109.1 − 1.0

Line	λ	F(Hα = 100)	I(Hα = 100)†	Cygnus‡ Loop-Pos 2
[O II]	3,727	85	762	430
Hβ	4,861	< 15	< 48	34
[O III]	4,959	15	44	28
[O III]	5,007	52	151	85
[N II]	6,548	29*	29	29
Hα	6,563	100	100	100
[N II]	6,584	89	86	87
[S II]	6,717	73	66	72
[S II]	6,731	73	66	55

* Assumes [N II] λ 6,584/λ 6,548 = 3.0.
† Assumes $E(B-V) = 1.1$.
‡ From ref. 8.

Near the regions of enhanced radio emission are several faint optical filaments visible on the Palomar Observatory Sky Survey (POSS) red plate and on deep Hα and [S II] Schmidt plates[2]. We obtained a spectrum of one of these filaments on 17 October 1980 with the intensified Reticon scanner on the 1.3-m telescope at McGraw-Hill Observatory. This filament is located in the southern portion of the remnant at (1950) α = 22 h 59 min 45 s, δ = 58° 25′ 05″; it runs in a nearly radial direction and is about 2 arc min in length. A 4 × 40 arc s slit was placed along the filament and alternate 10-min integrations were obtained on the object and on a sky position ~10 arc min to the south—well outside the observed emission region. This technique along with the extreme faintness of the filament caused inexact sky subtraction in the resulting spectrum, which is shown reduced to fluxes by standard methods[6] (Fig. 1). The total integration time on the object was 3,000 s.

The observed line fluxes, $F(\lambda)$, relative to $F(\text{H}\alpha) = 100$ are shown in Table 1. We were not able to detect Hβ in our spectrum and hence cannot determine the reddening directly. However, the reddening has been estimated for two Sharpless H II regions[7] nearby which are thought to be at approximately the same distance as the remnant[1]. Using a median value of $E(B-V) = 1.1$, we show these reddening corrected line intensities, $I(\lambda)$, in Table 1. The relative line strengths are accurate to 30% except for [O II] λ 3,727, which is worse because the detector sensitivity is poor and rapidly changing in this region.

We have also listed in Table 1 the relative line strengths[8] of Miller's position 2 in the Cygnus Loop SNR. Comparing these with our observations of G109.1 − 1.0, we see broad similarities in the relative line strengths. The main difference is that the lines of [O II] λ 3,727 and [O III] λλ 4,959, 5,007 are stronger in G109.1 − 1.0 although the estimates of these lines are particularly vulnerable to the uncertain reddening. The strong lines of [S II] λλ 6,717, 6,731 and [N II] λ 6,584 relative to Hα are the usual indicators of shock heating[9] and this confirms that we are seeing the optical filaments of a SNR.

Our spectrum is not good enough to warrant a detailed analysis and comparison with published shock models, but we can use the relative strengths of the [S II] λλ 6,717, 6,731 lines to estimate the electron density in the S^+ zone of the emitting region. Using updated collisional cross-sections for [S II][10] and a five–level atom calculation, we estimate $N_e(\text{S II}) \approx 700\ \text{cm}^{-2}$.

As the remnant is not bright optically, it is probably still in the adiabatic expansion phase. The optical filaments are probably dense clouds which have been overtaken by the blast wave from the supernova as it travels through the interstellar gas. Using the formulation of McKee and Cowie[11], we can estimate the initial energy of the supernova by assuming rough pressure equilibrium between the clouds and intercloud gas (now emitting X rays), given by

$$E_0 = 5.6 \times 10^{43} (\beta')^{-1} N_e(\text{S II}) D(\text{pc})^3\ \text{erg}$$

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here β' is a factor of the order of unity which relates the pressure in the cloud and intercloud gas. Using values of N_e(S II) and D mentioned earler, we find

$$E_0 = 2.88\times 10^{51}\left(\frac{d}{4\ \text{kpc}}\right)^3 \text{erg}$$

This is considerably greater than the value of $\sim 7\times 10^{50}$ erg which has been found for the Cygnus Loop (which is about the same diameter) using the same method[6]. More importantly, it is more than an order of magnitude greater than the initial energy calculated from the X-ray luminosity. Using the X-ray parameters derived by Gregory and Fahlman[1], we find

$$E_0(\text{X ray}) = 1.84\times 10^{50}\left(\frac{d}{4\ \text{kpc}}\right)^3 \text{erg}$$

This type of discrepancy for SNR has been noticed previously[12] and may indicate either that the magnetic pressure in the filaments is non-negligible or that there are beams from the central object, as in W50.

The presence of a central radio and X-ray source in G109.1−1.0 and a possible jet-like structure[1] indicate a striking similarity to W50. The emission-line object SS433, which is believed to be associated with W50, also shows a double jet structure in both X rays[3] and radio[13]. The axis of this jet seems to be aligned with the enhanced 'radio ear' structure of W50 and it is in this region where faint optical filaments are visible on the POSS red print. Spectra of some of these filaments[5,14,15] indicate a qualitative similarity to the position we have observed in G109.1−1.0, except that in W50 the [N II] and [S II] lines are much stronger relative to Hα and the density as measured from the [S II] lines is lower in W50.

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Cosmological consequences of grand unified theories on density fluctuations

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Recent investigations into the cosmological consequences of grand unified theories (GUTs) of elementary particles have shown that the observed matter–antimatter asymmetry of the Universe can be explained without recourse to the hypothesis of specific initial conditions. It is shown here that the origin of inhomogeneities in the matter distribution, which are thought to be responsible for the later formation of galaxies, cannot be explained by a simple addition of density fluctuations to the standard model. The appearance of these fluctuations, after the epoch when baryon number is fixed, is almost purely adiabatic, any departure from adiabaticity falling off in inverse proportion to the mass of the perturbation.

Fry *et al.*[1] have calculated detailed models for the creation of baryon number in the early stages of the Universe. Their results are parametrized by a quantity $K = 2.9\times 10^{17}\alpha(m_X/\text{GeV})^{-1}$, where $\alpha \simeq 10^{-2}$ is the GUT coupling constant, and m_X is the mass of the superheavy X-boson whose decay is responsible for the generation of baryon number. If $K < 1$, the X-bosons freeze out of thermal equilibrium while they still have relativistic energies, and their number density relative to other particles is independent of the freeze-out temperature and of K. If, on the other hand, $K > 1$, the freeze-out temperature is less than m_X, and the number of X-bosons is reduced by a Boltzmann factor. The resulting ratio B of baryon to photon number then decreases with increasing K.

In models of slightly inhomogeneous cosmologies, one conventionally imagines the Universe to consist of a patchwork of pieces, each larger than the horizon at any given time, and each evolving like a separate Friedmann model with a particular value of curvature[2]. In one of the perturbed regions, the expansion rate will be slightly altered by the curvature term, and this will mean that the freeze-out epoch t_f and temperature T_f will differ from the corresponding values in the unperturbed region (which is taken to have zero curvature). If $K < 1$, this variation in T_f will have no effect on the final value of B, but if $K > 1$, B depends on T_f through the Boltzmann factor. (Strictly speaking, the dependence of B on T_f only vanishes in the limit $T_f/m_X \to \infty$, but it falls off rapidly as the X-bosons become relativistic at freeze-out.) Consider a perturbed region just entering the horizon at t_f. If the density variation is $\delta\rho/\rho = \kappa$, the expansion rate[2] will be

$$(\dot{R}/R) = \left(\frac{8\pi G}{3}\rho\right)^{1/2}(1-\kappa)^{1/2} \simeq \left(\frac{8\pi G}{3}\rho\right)^{1/2}\left(1-\frac{\kappa}{2}\right) \qquad (1)$$

Fry *et al.*[1] give a simple analytic expression for B which applies in the limit $K \gg 1$. Using a variable $z = m_X/T$, then $B = \varepsilon/z_f K$, where ε is the baryon number produced in the decay of an $X\bar{X}$ pair; z_f denotes the freeze-out temperature, and is the larger of the two solutions of

$$Kz^{7/2}\,\mathrm{e}^{-z} = 1 - \frac{\kappa}{2} \qquad (2)$$

This equation, although like that given by Fry *et al.*, incorporates the modified expansion rate of equation (1). Now, let z_f be the solution of equation (2) in the unperturbed ($\kappa = 0$) region, and let $z_f' = z_f + \mathrm{d}z_f$ be the solution for non-zero κ. It is easy to show that $\mathrm{d}z_f/z_f = \kappa/(2z_f - 7)$. The bad behaviour at $z_f = 7/2$ is because equation (2) is an approximation which fails for small z; we therefore require $z_f \gg 7/2$, and take $\mathrm{d}z_f/z_f \simeq \kappa/2z_f$. With the expression for B, then $\delta B/B = -\mathrm{d}z_f/z_f \simeq \kappa/2z_f$. The minus sign means that regions of higher density (positive κ) have smaller values of baryon-to-photon ratio, because freeze-out occurs later in the overdense regions, where the expansion rate is reduced. These conclusions were reached qualitatively by Turner and Schramm[3] and by Barrow[4].

It can be concluded that if K is small, density fluctuations do not alter B: that is, they are entirely adiabatic. If $K > 1$, then departures from adiabaticity can be achieved, but the effect decreases with increasing K (which implies increasing z_f). There must therefore be some pair of values of K and z_f (presumably near unity) for which the non-adiabaticity is maximized. The appearance of the fluctuations is, in a sense, the opposite of isothermal, because the variation in baryon density is smaller than that in radiation density. Note that these departures from adiabaticity have no dynamical significance at these early stages of cosmic evolution, because at these times baryon number is simply a quantum number attached to the various particles comprising the radiation gas. Variations in baryon density reflect variations in the locally averaged baryon number at different points. Only at a later stage ($t \simeq 10^{-4}$ s), when nucleons become massive (non-relativistic) particles carrying the baryon number, does the distinction between adiabatic and isothermal fluctuations have a dynamical role.

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