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Filaments in the southern giant lobe of Centaurus A: constraints on nature and origin from modelling and GMRT observations

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ABSTRACT
We present results from imaging of the radio filaments in the southern giant lobe of Centaurus A using data from Giant Metrewave Radio Telescope observations at 325 and 235 MHz, and outcomes from filament modelling. The observations reveal a rich filamentary structure, largely matching the morphology at 1.4 GHz. We find no clear connection of the filaments to the jet. We seek to constrain the nature and origin of the vertex and vortex filaments associated with the lobe and their role in high-energy particle acceleration. We deduce that these filaments are at most mildly overpressured with respect to the global lobe plasma showing no evidence of large-scale efficient Fermi I-type particle acceleration, and persist for \( \sim 2–3 \) Myr. We demonstrate that the dwarf galaxy KK 196 (AM 1318 \(-\)444) cannot account for the features, and that surface plasma instabilities, the internal sausage mode and radiative instabilities are highly unlikely. An internal tearing instability and the kink mode are allowed within the observational and growth time constraints and could develop in parallel on different physical scales. We interpret the origin of the vertex and vortex filaments in terms of weak shocks from transonic magnetohydrodynamical turbulence or from a moderately recent jet activity of the parent AGN, or an interplay of both.

Key words: instabilities – turbulence – techniques: image processing – galaxies: individual (Centaurus A) – galaxies: jets – radio continuum: galaxies.

1 INTRODUCTION
From studies of both high- and low-power radio galaxies over the past three decades, considerable observational evidence has emerged for inhomogeneous, filamentary lobes, e.g. Cygnus A (Perley, Dreher & Cowan 1984), 3C 310 (van Breugel & Fomalont 1984), Hercules A (Dreher & Feigelson 1984; Gizani & Leahy 2003), Fornax A (Fomalont et al. 1989), Pictor A (Perley, Röser & Meisenheimer 1997), 3C 353 (Swain, Bridle & Baum 1998), M87 (Owen, Eilek & Kassim 2000; Forman et al. 2007), NGC 193 (Laing et al. 2011), B2 0755+37 (Laing et al. 2011), with the bulk of the observations being conducted with the Very Large Array (VLA) in the GHz regime. The filamentarity has implications for the internal structure of the lobes. However, no consensus exists on whether the magnetic field in the lobes has a low filling factor and electrons are uniformly distributed, or the electron population tracks the magnetic field enhancements closely.
Positionally varying magnetic field strength was claimed for lobes of a number of Fanaroff–Riley class II (FR II; Fanaroff & Riley 1974) sources (e.g. Hardcastle & Croston 2005; Goodger et al. 2008), in contrast with the western giant lobe of the source Fornax A,\textsuperscript{1} for which a positionally varying electron energy spectrum is favoured (Seta, Tashiro & Isobe 2011).

Filamentary structure does not necessarily imply turbulence, but magnetohydrodynamical (MHD) turbulence implies filamentary structure in synchrotron emission (e.g. Eilek 1989; Hardcastle 2013; Wykes et al. 2013). MHD turbulence amplifies and transports magnetic fields which in turn control lobe viscosity, conductivity and resistivity, as well as the acceleration and propagation of cosmic rays (e.g. Lee et al. 2003; Jones et al. 2011). The presence of turbulence might in some cases beakin to a development of plasma instabilities. Various types of instabilities, promoting growth of filament-like features, could develop inside or on the surface of a radio lobe. Hydrodynamical (HD) instabilities, such as Kelvin–Helmholtz (KH), Rayleigh–Taylor (RT) and Richtmyer–Meshkov (RM) are relevant since they can lead to flow patterns that naturally filament and can amplify ambient magnetic fields (e.g. Jun & Norman 1995; Ryu, Jones & Frank 2000). MHD instabilities such as the resistive tearing instability, the sausage mode and the kink mode are also apposite, as are radiative instabilities.

Utilizing the Australia Telescope Compact Array (ATCA) and the 64 m Parkes telescope for imaging at 1.4 GHz with 49 arcsec angular resolution, Feain et al. (2011) have discovered intricate filamentary features associated with the northern and southern giant lobes of Centaurus A (Fig. 1). Centaurus A is the nearest (3.8 ± 0.1 Mpc; Harris, Rejkuba & Harris 2010)\textsuperscript{2} Fanaroff–Riley class I (FR I) radio galaxy, hosted by the massive elliptical galaxy NGC 5128. Due to its luminosity and proximity, Centaurus A is an outstanding testbed for models of jet energetics, particle acceleration and the evolution of low-power radio galaxies in general. Centaurus A’s northern jet (angular size ∼4.0 arcmin) and its immediate surroundings, the bright inner lobes (∼5.5 arcmin each) and the northern middle lobe (∼33 arcmin) have been extensively studied (e.g. Tingay et al. 1998; Morganti et al. 1999; Hardcastle et al. 2003; Kraft et al. 2003, 2009; Hardcastle, Kraft & Worrall 2006; Croston et al. 2009; Müller et al. 2011; Neff et al. 2014; Israel et al., in preparation). However, Centaurus A’s proximity to the Earth has hampered for a long time comprehensive investigations of its giant (i.e. outer) lobes (∼4.3 each), whose substructure has been mapped only recently in the aforementioned work by Feain et al. (2011).

Topics of great current interest are the ages of the giant lobes and the lobe particle content and pressure. Hardcastle et al. (2009) and Yang et al. (2012) have determined radiative ages of Centaurus A’s giant lobes: the former obtaining ∼30 Myr based on synchrotron ageing fitting the single-injection Jaffe–Perola model (Jaffe & Perola 1973), the latter <80 Myr reasoning that ages significantly larger than a few tens of Myr are not consistent with the observations of gamma-ray inverse-Compton emission. The above values would imply that the giant lobe front ends expand at, respectively, ∼0.030 and >0.011c, i.e. faster than Centaurus A’s inner lobes (∼0.009c; Croston et al. 2009), in discord with expectations. The dynamical age calculations by Wykes et al. (2013) give ∼560 Myr

\textsuperscript{1} Morphologically, Fornax A is FR II class by the original (Fanaroff & Riley 1974) definition; in terms of luminosity, it is on the boundary FR I/FR II.

\textsuperscript{2} At that distance, 1 arcmin corresponds to 1.1 kpc.
Origin of filaments in Centaurus A

2 OBSERVATIONS AND DATA REDUCTION

GMRT observations of the vertex and vortex filaments at 325, 235 and 150 MHz were carried out in 2012 May and 2013 February (project codes 22_038 and 23_060). Due to the low elevation of the target, observations were limited to <5 h per night, requiring a total of eight nights. A journal of these observations is given in Table 1. Visibilities for two polarizations (RR and LL) were recorded in spectral line mode to enable narrow-band RFI excision and prevent bandwidth smearing. For calibration purposes, we observed two standard calibrators: 3C 286 as the (primary) flux density and bandpass calibrator, observed for 20 min at the start and/or end of each run, and 3C 283 as the (secondary) phase calibrator, observed for 5 min every 50 min. During observations, the standard GMRT pointing error correction model was applied on both targets, as well as 3C 286.

The data\(^3\) reduction was conducted using aIPS (version 31DEC12; Greisen 2003), the PYTHON-based extension S\(\text{PAM}\) (Intema et al. 2009; Intema 2014), and the cl\(\text{ean}\) imaging task in CASA (version 4.1.0; McMullin et al. 2007) to create the final low-resolution maps. S\(\text{PAM}\) was used to correct for (direction-dependent) ionospheric phase corrections, which can be a dominant source of error at sub-GHz frequencies. Overall, ionospheric conditions were relatively quiet during all observing nights, as judged from the slowly varying gain phases. For all observations, except 2012 May 28, flux and bandpass calibration were derived from 3C 286, adopting the modified Perley–Taylor flux model as described by Intema et al. (2011). For this, we excluded the shortest (central square) baselines, and performed some manual flagging of obviously bad antennas/times/polarizations based on gain calibration tables. Calibration results were then applied to the target field data. In the 2012 May 28 observation, a combination of factors rendered the data on 3C 286 unusable. In this case, we have used 3C 283 as primary calibrator, adopting a 325 MHz flux density of 23.3 Jy (as

\(^3\)MHD simulations (e.g. Jones et al. 2011) indicate that this is also true for trans-Alfvénic and mildly super-Alfvénic turbulence. Filaments last roughly an eddy turnover time for the scale of the eddies that stretch them.
of Centaurus A, and also the background FR I radio galaxy
the GMRT primary beam strongly attenuated the inner lobes
446. Outside the main beam, we found that at 325 MHz
−
PKS 1320
most useful to reduce the DR-limiting effects of the point source
843 MHz catalogue (Mauch et al. 2003), followed by wide-field
dam 2004) of 10
spheric calibration and imaging. This includes peeling (Noor-
amphlet) have a DR (defined here as the ratio of peak flux to the
PKS B1318−434, so that they do not put any DR limitations on our
image quality. At 235 MHz, the inner lobes caused slight ripples
across our image, but these were reduced by peeling. At 150 MHz,
the inner lobes strongly affected the entire field of view, rendering
the image inutile for our purpose of detecting faint diffuse emission.
Many attempts (including peeling) to improve this failed. We will
disregard the 150 MHz observations in what follows. Table 2 lists
the properties of the final high-resolution images.

2.2 Imaging of large-scale emission
With gain calibration sorted out during the high-resolution imaging,
we have explored several ways of imaging the highly resolved
diffuse components that are embedded in the southern giant lobe
of Centaurus A. To maximize our signal-to-noise on the vertex and
vortex, we converged on a method in which we used the CASA
imager task c\texttt{lean} in mosaicking mode, simultaneously imaging and
\texttt{cleaning} all visibilities of the two pointings into one final image
per frequency. We were restrained in the use of multiscale decon-
volution, because it is as yet unsupported in mosaicking mode.
Before importing the visibilities into \texttt{CASA}, we have pre-
subtracted (in \texttt{AIPS / SPAM}) the \texttt{clean} components (of ≥5 times
the background noise level) of all point sources found in the compact
emission image from the visibility data sets, while temporarily ap-
plying the appropriate direction-dependent gain calibrations. Fur-
thermore, since \texttt{SPAM} functionality is not yet available in \texttt{CASA},
we have used a single gain table to calibrate each data set, namely
the one corresponding to the field centre. This lack of direction-
dependent ionospheric calibration seems to have limited effect while
imaging using solely the shortest baselines.
For imaging the diffuse emission, we used a Gaussian weight ta-
er, suppressing visibilities from baselines longer than 2.5 kλ. The
few visibilities from baselines shorter than 150 kλ were excluded
during imaging to (i) remove strong ripples coming from the very
few high-amplitude visibilities that sense (but completely under-
sample) the largest scale structures, and (ii) roughly match our data
to the ATCA interferometric data that is part of the study by Feain
et al. (2011), and which is also used in our analysis.
The final GMRT low-resolution images (Fig. 2, middle and right-
hand panel) have a DR (defined here as the ratio of peak flux to the

2.1 Calibration and imaging of compact emission
The target field data were initially calibrated and imaged at ‘high
resolution only’ by excluding visibilities from the inner 1 kλ of the
uv-plane and choosing a weighting scheme between robust and uni-
form (robust = −1 in the \texttt{AIPS} convention). With this approach,
we can obtain good gain calibrations for all antennas without hav-
ning to deal with the large-scale emission immediately. Indeed, the
same calibration can be used later to image the large-scale emission
(see Section 2.2). Each target field was initially phase-calibrated
against a simple ∼10 point source model derived from the SUMSS
843 MHz catalogue (Mauch et al. 2003), followed by wide-field
(facet-based) imaging and \texttt{clean} deconvolution. In all imaging, we
automatically put tight \texttt{clean} boxes encompassing emission peaks
(above five times the central noise level) to guide the deconvolution
and suppress \texttt{clean} bias. Within the \texttt{clean} boxes, emission was
\texttt{cleaned} down to twice the central noise level. Extra facets were
added at the locations of known bright outlier sources within four
primary beam radii. This was followed by two rounds of phase-
only self-calibration, and one round of amplitude and phase self-
calibration. In between rounds, bad data were removed by flagging
spurious points in the gain calibration tables and residual visibility
amplitudes.
Self-calibration was followed by two rounds of (\texttt{SPAM}) iso-
npheric calibration and imaging. This includes peeling (Noord-
dam 2004) of 10−20 of the brightest sources. The latter was
most useful to reduce the DR-limiting effects of the point source
PKS 1320−446. Outside the main beam, we found that at 325 MHz
the GMRT primary beam strongly attenuated the inner lobes
of Centaurus A, and also the background FR I radio galaxy

<table>
<thead>
<tr>
<th>Table 1. Journal of GMRT observations.</th>
</tr>
</thead>
<tbody>
<tr>
<td>Observing dates</td>
</tr>
<tr>
<td>-----------------</td>
</tr>
<tr>
<td>2012 May 27−28</td>
</tr>
<tr>
<td>2013 Feb 24−27</td>
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<tr>
<td>2013 Feb 22−23</td>
</tr>
</tbody>
</table>

Note. rms values corrected for \texttt{T}\textsubscript{sys.}

<table>
<thead>
<tr>
<th>Table 2. Properties of GMRT maps per region and frequency.</th>
</tr>
</thead>
<tbody>
<tr>
<td>Image</td>
</tr>
<tr>
<td>-------</td>
</tr>
<tr>
<td>Vertex high resolution</td>
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<tr>
<td>Vertex high resolution</td>
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<td>Vertex high resolution</td>
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<td>Vortex high resolution</td>
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<td>Vortex/Vortex high resolution</td>
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<td>Vortex/Vortex low resolution</td>
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<td>Vortex/Vortex low resolution</td>
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Table 3. Flux scale corrections, based on $T_{\text{sys}}$ estimates in the context of the model described by Sirothia (2009).

<table>
<thead>
<tr>
<th>Pointing</th>
<th>$T_{\text{sys}}$ (K)</th>
<th>Flux correction factor</th>
</tr>
</thead>
<tbody>
<tr>
<td>3C 286 at 323 MHz</td>
<td>123.57</td>
<td>1.0</td>
</tr>
<tr>
<td>3C 286 at 234 MHz</td>
<td>222.96</td>
<td>1.0</td>
</tr>
<tr>
<td>3C 283 at 323 MHz</td>
<td>129.28</td>
<td>1.05</td>
</tr>
<tr>
<td>3C 283 at 234 MHz</td>
<td>235.17</td>
<td>1.05</td>
</tr>
<tr>
<td>Vertex at 323 MHz</td>
<td>248.78</td>
<td>2.01</td>
</tr>
<tr>
<td>Vertex at 234 MHz</td>
<td>469.93</td>
<td>2.11</td>
</tr>
<tr>
<td>Vertex at 323 MHz</td>
<td>227.07</td>
<td>1.84</td>
</tr>
<tr>
<td>Vertex at 234 MHz</td>
<td>437.15</td>
<td>1.96</td>
</tr>
<tr>
<td>Vertex/vortex at 325 MHz</td>
<td>239.68</td>
<td>1.93</td>
</tr>
<tr>
<td>Vertex/vortex at 235 MHz</td>
<td>455.56</td>
<td>2.04</td>
</tr>
</tbody>
</table>

Note. The $T_{\text{sys}}$ values are based on a model for our duration of the observations.

off-source rms away from the image centre) of $10^2$: 1. We determine the rms noise level (1σ) in those images as $\sim 3.8$ mJy beam$^{-1}$ (325 MHz) and $\sim 10.8$ mJy beam$^{-1}$ (235 MHz), see Table 2.

2.3 Flux density levels

Large-scale radio emission that is resolved out by the GMRT interferometer, but is still registered by individual GMRT antennas, enters the system as sky noise, routinely expressed as sky temperature $T_{\text{sky}}$. At low radio frequencies, the sky noise can make a significant contribution to the total system temperature $T_{\text{sys}}$, which also includes receiver noise, ground radiation and potentially other, smaller caches of noise. GMRT observations are by default not corrected for variations in $T_{\text{sys}}$ when moving the telescope across the sky (e.g. Tasse et al. 2007; Sirothia 2009; Intema et al. 2011). This causes a target field flux scale error $T_{\text{sys, target}}/T_{\text{sys, fullcal}}$ when transferring the gain calibration from flux calibrator to target field.

In our case, the large-scale emission of the southern giant lobe of Centaurus A is filling most (or all) of our target fields of view at all observing frequencies, resulting in considerably higher sky temperatures than in the field of our calibrators 3C 286 and 3C 283, leading to large flux scale errors. For correcting the resulting flux scale error of our target fields, we rely on the model discussed essentially by Sirothia (2009), with frequency dependence also incorporated. The resultant $T_{\text{sys}}$ values, which are based on the duration of our observations, are codified in Table 3. Instantaneous values of $T_{\text{sky}}$ and $T_{\text{receive}}$ independently are not meaningful for our purpose and were not computed. All gain ratios and correction factors were calculated with regard to our flux density calibrator 3C 286. The errors on the flux correction factors amount to $\sim 4$ per cent. Note that systematics in calibration while imaging will dominate the flux density calibration at these frequencies with the GMRT which are in the range $\sim 10$–12 per cent. Thus, the total error on the flux density is $\sim 15$ per cent.

Since the final 325 and 235 MHz low-resolution images have been produced with visibility data from both pointings combined, we assign single flux correction factors per frequency. Based on Table 3, we adopt, respectively, 1.93 at 325 MHz and 2.04 at 235 MHz.

2.4 1.4 GHz image processing

An overview of the ATCA and Parkes observations and imaging of the giant lobes at 1.4 GHz can be found in Feain et al. (2009, 2011). We have used a vertex/vortex ATCA data subset from the above work and subtracted the background sources from the image plane using the source extraction tool PYBDSM. This map is used later in the article for spectral index extraction (Sections 3.1 and 3.2) and turbulence modelling (Section 4.2.2).

3 RESULTS

3.1 Radio morphology and filament flux densities

As is apparent from Fig. 2, the vertex, vortex and other bright filaments maintain their coherence over the frequency range 1.4 GHz–235 MHz, i.e. the 1.4 GHz emission appears spatially closely associated with the emission at 325–235 MHz. Given the better sampling of large-scale structures at low frequencies, auxiliary filamentary structure could have been revealed by our observations, if present. We find some additional filamentary features but no clear indication in the GMRT images for the vertex/vortex and other features being connected to the core/jet. The topology of the filamentary mesh is unexplored; however, the vertex filament seems twisted, bright at what may be intersections, and its northern part appears to be built up of two segments which themselves potentially contain substructure in form of approximately parallel threads (most clearly visible in Fig. 3). At all available frequencies, the vortex morphology resembles a fleshy fungus; the eastern ‘cap of the mushroom’s head’ is displaced by about 4 arcmin to the south as we go to lower frequencies. Pronounced in the 325 MHz image (Fig. 2, middle panel) is a short bow-like filament north-east of the vertex, which is ‘behind’ the interferometric rings caused by the PKS 1320–446 source in the 1.4 GHz image (Fig. 2, left-hand panel). Fainter, but still well discernible (in Fig. 2 left and middle, and Fig. 3) is a pipe-like filament close to and west of the vertex.

We have measured the vertex and vortex flux densities from FITS files making use of ds9\(^6\) and the FUNTOOLS\(^7\) library and defining rectangular regions encompassing the filaments (see Figs 3 and 4): two regions of, respectively, 5217 and 5828 pixels for the vertex, and nine regions (to account for its curvature, and for a slightly different position of its south-east part as a function of frequency) for, respectively, 2310, 1519, 504, 912, 819, 800, 3956, 4268 and 8320 pixels for the vortex (on a synthesized beam area of 28.65 pixels at 325 MHz and 28.20 pixels at 235 MHz). Equally sized rectangular regions were used elsewhere on the maps, well away from obvious outliers (positive and negative) and from edge-effects, for background determination for which we effected 15 trials. We handled the standard deviation of the background values to estimate the error on the source regions.

Table 4 lists the final, background-corrected vertex and vortex flux densities. The 1.4 GHz value is an order of magnitude less than the flux density of these filaments quoted in Feain et al. (2011); however, there they represent the combined ATCA+Parkes flux densities at 1.4 GHz and are not background subtracted.

3.2 Filament spectral index

We have utilized our GMRT 325–235 MHz radio continuum images and the existing ATCA 1.4 GHz continuum image to determine the spectral index. To this end, we have matched the shortest baselines, and have disregarded visibilities from baselines shorter than

\(^6\) https://hea-www.harvard.edu/RD/ds9/site/Home.html
\(^7\) https://www.cfa.harvard.edu/~john/funtools

Origin of filaments in Centaurus A 2871
Figure 2. From left to right: ATCA 1.4 GHz, GMRT 325 MHz and GMRT 235 MHz continuum maps of the large-scale emission in the vertex and vortex fields. The beam size is indicated by the ellipse in the bottom left-hand corner. The rms noise of the GMRT images is given in Table 2.

Figure 3. GMRT 325 MHz map of the vertex filament with spectral extraction regions used for Table 4 overlaid. The scale is in Jy beam$^{-1}$.

Figure 4. GMRT 325 MHz map of the vortex filament with spectral extraction regions used for Table 4 overlaid. The scale is in Jy beam$^{-1}$.

0.15 kλ. We have fitted power laws in frequency to the flux densities and errors in Table 4 augmented with the errors on the calibration in the GMRT data, minimizing $\chi^2$. From our analysis, the resultant best-fitting spectral indices are $\alpha = 0.81 \pm 0.10$ (vertex) and $\alpha = 0.83 \pm 0.16$ (vortex).

Given the large errors on the spectral index for the entirety of the individual filaments, we cannot determine variations of the spectral index along them (i.e. a spectral index analysis for filament subregions would be meaningless).

4 INTERPRETATION

4.1 Vertex and vortex: the nature

A detailed radio spectral energy distribution (SED) analysis can provide lower limits on the ages of the filaments and clues about the filament’s nature and origin. The radio SED of distinct filaments in a handful of lobes where such measurements are accessible shows a slightly flatter spectrum compared to that of the general lobe plasma, e.g. $\alpha_{1.4 \text{GHz}} = 4.8 \pm 0.3$ GHz $\sim 1.1$ versus $\alpha_{1.3 \text{GHz}} = 1.6$ in the FR I/II source Hercules A (Gizani & Leahy 1999), which might indicate either more recent particle acceleration or magnetic field enhancement at the filament (Tribble 1994). More specifically, the spectral indices of the ring-like filaments in Hercules A obtained by spectral tomography (Gizani & Leahy 2003) are in the range $\alpha_{1.4 \text{GHz}} = 0.70$–1.15, and at lower GHz frequencies in the range $\alpha_{1.3 \text{GHz}} = 0.70$–0.90, with a general trend of flatter spectral indices at lower frequencies. The arc filament in the western lobe of Hercules A, comparable in morphology and orientation to the vertex and vortex, shows a significantly steeper index than any other filamentary feature in Hercules A’s lobes: $\alpha_{1.4 \text{GHz}} = 0.90 \pm 0.05$ and $\alpha_{1.3 \text{GHz}} = 1.14 \pm 0.04$ (see their table 5): this is surprising if it is associated with a recently shocked region. Given that Hercules A is probably in a driving phase, more useful comparisons with Centaurus A might include FR I sources.
such as Hydra A and 3C 310, or the source Fornax A. However, no spectral index measurements of individual filaments within the lobe volume of these sources (and any other FR I or FR II radio galaxy) are yet at hand.

Centaurus A’s integrated southern giant lobe diffuse emission shows a spectral index $\alpha_{\nu_{\text{4 GHz}}} = 0.55 \pm 0.02$ and $\alpha_{\nu_{\text{10 MHz}}} = 0.47 \pm 0.06$ (Hardcastle et al. 2009) and $\alpha_{\nu_{\text{18 MHz}}} = 0.63 \pm 0.01$ (McKinley et al. 2013) in the region of the vertex. Around the location of the vertex, the measured spectral indices are $\alpha_{\nu_{\text{4 GHz}}} = 0.71 \pm 0.05$ and $\alpha_{\nu_{\text{10 MHz}}} = 0.62 \pm 0.12$ (Hardcastle et al. 2009) and $\alpha_{\nu_{\text{18 MHz}}} = 0.65 \pm 0.01$ (McKinley et al. 2013) although this region does not encompass the vortex in its entirety. Various authors (Feain et al. 2011; Stawarz et al. 2013; Stefan et al. 2013; Wykes et al. 2013; Eilek 2014) have modelled (parts of) the giant lobes as recently undergoing particle re-acceleration; this idea is consistent with the rather flat spectral indices of the global lobe plasma. An important point is that the spectral indices that we measure for the filaments (Section 3.2) are not significantly flatter than these lobe spectral indices; instead, they are steeper (90 per cent confidence level) than or similar to (given the relatively large errors on the GMRT flux densities) the global lobe plasma values measured by McKinley et al. (2013). This places constraints on the possible character of the filaments that we discuss in more detail in Sections 4.1.2 and 4.1.3.

A number of theoretical works have suggested that a correlation between B-field and particle density may be weak, essentially because B-field fluctuations relate to Alfvénic turbulence which does not compress the plasma, and also on grounds of the so-called reconnection diffusion which violates the flux-freezing condition (see for more in-depth discussions, e.g. Cho & Lazarian 2003; Passot & Vázquez-Semadeni 2003; Santos-Lima et al. 2010). This is supported by X-ray observations of radio lobes (e.g. Pictor A, Hardcastle & Croston 2005; 3C 353, Goodger et al. 2008) where the inverse-Compton emission does not follow the pattern seen in radio-synchrotron emission from the lobes. We discuss the possible behaviour of the magnetic field in the filaments in Sections 4.1.1 and 4.1.3.

It has thus far been unclear whether the vertex and vortex filaments are overpressured as argued for M87 (Hines, Owen & Eilek 1989; Forman et al. 2007) and for Pictor A (Perley et al. 1997), or in pressure equilibrium with the medium within which they are embedded. The large pressure jump factors associated with the vertex and vortex ($p_{2}/p_{1} \sim 30$ for the vertex and $\sim 240$ for the vortex) for the heat capacity ratio of $\gamma = 5/3$, and $p_{2}/p_{1} \sim 15$ for the vertex and $\sim 80$ for the vortex for $\gamma = 4/3$) presented by Feain et al. (2011) on grounds of an idealized case of an initially spherical cocoons collapsing into a torus after a passage of a strong shock, would point towards large filament overpressure, and therefore to fairly flat filament spectral indices if these are set by particle acceleration. For $p_{2}/p_{1} \sim 30$ and $p_{2}/p_{1} \sim 240$ pressure ratios and $\gamma = 5/3$, the Mach number ($\mathcal{M}$) from the Rankine–Hugoniot jump relations is, respectively, $\mathcal{M} \sim 4.9$ and $\mathcal{M} \sim 13.9$; for $p_{2}/p_{1} \sim 15$ and $p_{2}/p_{1} \sim 80$ pressure ratios and $\gamma = 4/3$, it is, respectively, $\mathcal{M} \sim 3.6$ and $\mathcal{M} \sim 8.4$. However, it is not clear whether such large overpressure factors are consistent with the radio observations. We will elaborate on pressure considerations in Sections 4.1.1 and 4.2.1.

### 4.1.1 Filament pressure

In this section, we consider the pressure in the filaments and its implications for their dynamics.

Since the radio emission from the vertex and vortex is synchrotron radiation, we can calculate the pressure of these filaments from their flux densities and estimated volumes if we make some assumptions about the relative energy densities in B-field, electrons and non-radiating particles. Ideally, we would use the lowest radio frequency at our disposal for the flux density measurement, i.e. 150 MHz, to minimize the effect of any age-related steepening in the electron spectrum. However, because the 150 MHz data have defined imaging attempts (see Section 2.1) and the 235 MHz data possess a relatively large error on the flux, we resort to the flux densities at 325 MHz. We estimate the volumes of the vertex and vortex, treating them as cylinders with length $l = 48$ kpc and radius $r = 3.2$ kpc (vertex) and $l = 105$ kpc and $r = 1.8$ kpc (vortex), to be, respectively, $V \sim 4.5 \times 10^{7}$ and $V \sim 3.1 \times 10^{6}$ cm$^{3}$.

One possible approach is to determine the minimum energy in the lobes and filaments. To do this, we adopt the Myers & Spangler (1985) minimum-electron-energy model and perform the minimum-energy calculations numerically using the code of Hardcastle, Birkinshaw & Worrall (1998). This yields for the vertex an equipartition B-field strength of $B \sim 2.4$ $\mu$G and a minimum pressure of $p_{\text{min}} = 2U_{\text{th}}/3 \sim 1.5 \times 10^{-12}$ dyn cm$^{-2}$ while for the vortex, we find an equipartition field $2.2$ $\mu$G and $p_{\text{min}} \sim 1.3 \times 10^{-13}$ dyn cm$^{-2}$. Compared to the global minimum pressure of the lobes of $p_{\text{min}} \sim 4.5 \times 10^{-14}$ dyn cm$^{-2}$ (based on results by Hardcastle et al. 2009 building as well on the Myers & Spangler 1985 formalism), these are a factor $\sim 3$ higher. Such an overpressure could be identified with weak shocks, $\mathcal{M} \sim 1.7$ (vertex) and $\mathcal{M} \sim 1.6$ (vortex), from the Rankine–Hugoniot jump relations.

However, in the model by Wykes et al. (2013), the giant lobes are clearly not in minimum pressure; the dominant pressure component is provided by thermal material, with $p_{\text{th}} \sim 1.5 \times 10^{-12}$ dyn cm$^{-2}$, with the pressures due to relativistic electrons and B-field being much lower. In this situation, it is not clear what we should assume for the pressures in the various components in the filaments. We consider one possible limiting case, in which the thermal and relativistic electron pressures remain constant in the filaments while the B-field strength increases to provide

---

Table 4. Measured flux density, LAS and diameter (over the minor axis) of the vertex and vortex filaments at 1.4 GHz and at 325 and 235 MHz.

<table>
<thead>
<tr>
<th>Filament</th>
<th>$S_{\nu_{\text{1.4 GHz}}}$ (Jy)</th>
<th>$S_{\nu_{\text{325 MHz}}}$ (Jy)</th>
<th>$S_{\nu_{\text{235 MHz}}}$ (Jy)</th>
<th>LAS (arcmin)</th>
<th>Diameter (arcmin)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Vertex</td>
<td>1.25 ± 0.15</td>
<td>4.90 ± 0.24</td>
<td>4.61 ± 0.73</td>
<td>31</td>
<td>7</td>
</tr>
<tr>
<td>Vortex</td>
<td>0.72 ± 0.14</td>
<td>2.60 ± 0.40</td>
<td>3.05 ± 0.60</td>
<td>58</td>
<td>4</td>
</tr>
</tbody>
</table>

Note. Flux density corrected for background.
the observed increase in synchrotron emissivity. We find that, by keeping the relativistic electron number density \( n_{e,\text{rel}} \) fixed, the \( B \)-field has to increase by a factor \( \sim 6.5 \) (vertex) and factor \( \sim 6.0 \) (vortex) with respect to the ambient \( B \)-field (for which we adopted 0.9 \( \mu \)G from Abdo et al. 2010), to produce the radio-bright filaments. We obtain \( U_{B,e} \sim 1.7 \times 10^{-12} \text{ dyn cm}^{-2} \) and thus \( p_{B,e} \sim 5.8 \times 10^{-13} \text{ dyn cm}^{-2} \) (vertex), and \( U_{B,e} \sim 1.2 \times 10^{-12} \text{ dyn cm}^{-2} \) and thus \( p_{B,e} \sim 3.9 \times 10^{-13} \text{ dyn cm}^{-2} \) (vortex). In this case, the filaments would be only mildly overpressured (a factor \( \sim 1.3 \)) with respect to the medium in the giant lobes, which translates to \( \mathcal{M} \sim 1.0 \). (We note that as this is a limiting case, the overpressure factor could be even lower than this if both electron and \( B \)-field energy densities are varied to produce the filament emissivity.)

The \( B \)-field estimated above, although strong compared to the field estimated from inverse-Compton observations in the giant lobes as a whole (Abdo et al. 2010), is not energetically dominant: with a thermal pressure of \( p_{th} \sim 1.5 \times 10^{-12} \text{ dyn cm}^{-2} \) (Wykes et al. 2013), the filaments would be \( B \)-field dominated for a filament strength \( B \gtrsim 6 \mu \)G. However, this conclusion depends strongly on our assumption about the thermal pressure in the giant lobes. With \( p_{th} \sim 3.2 \times 10^{-11} \text{ dyn cm}^{-2} \) (Eilek 2014), they would already be \( B \)-field dominated for \( B \gtrsim 3 \mu \)G. In the less likely case of giant lobe pressure close to the minimum pressure, the \( B \)-field limit would be lower. If the \( B \)-fields are amplified through turbulence, a hard upper bound to the strength of the filament \( B \)-field could be placed in terms of a balance between Maxwell and Reynolds stresses (i.e. balance between the magnetic tension and turbulence stresses, which is the saturation regime), leading to \( B^2/(8\pi) = 0.5\rho v_\text{turb}^2 \), with \( \rho \) the mass density and \( v_\text{turb} \) the turbulent speed. We have no means of independently calculating \( v_\text{turb} \) (see also the remarks in Section 4.2.2); hence, we cannot obtain a maximum attainable \( B \)-field of the filament via this route.

The above results are inconsistent with the cocoon collapsing scenario suggested by Feain et al. (2011), which implied \( \mathcal{M} \gg 1.7 \) (see the foregoing section); furthermore, the results make models relying on high-energy particle acceleration by the Fermi I mechanism in the giant lobes (as per e.g. Pe’er & Loeb 2012 and Fraija 2014) unlikely (see also Section 4.1.2).

### 4.1.2 Particle confinement and ageing

As noted in Section 4.1, we find a somewhat steeper spectrum in the filaments than in the global lobe plasma. The steeper index is hard to explain in any model; but if it is real, it rules out models in which the excess emissivity in the filament region is related to (any type of) particle acceleration. Instead, we need to consider models in which the electrons in the filaments have cooled more rapidly than those in the giant lobes. Probably, the only viable excess-loss scenario involves particles confined in the filaments and so ageing faster (Tribble 1993); hence in this section, we consider whether confinement in the filaments can take place for relevant time-scales.

To see whether electrons of the required energies (in the middle of the available observing frequencies of 235 MHz to 2.2 GHz (Shimwell, private communication, for the latter), i.e. \( \sim 1.2 \text{ GHz} \)) stay confined in the vertex filament, we estimate the diffusion time \( \tau_{\text{diff}} \simeq \xi r^2 / r_{\nu, e} c \),

\[
\tau_{\text{diff}} \simeq \xi r^2 / r_{\nu, e} c ,
\]

where \( \xi \) is a fudge factor for which we use \( \xi \sim 10 \) (see Wykes et al. 2013 and references therein). Assuming the \( B \)-field is parallel to the long axis of the filament (see Section 4.2.3), and taking the vertex radius \( r \sim 3.2 \text{ kpc} \), \( B \sim 5 \mu \text{G} \) and electron gyroradius \( r_{\nu, e} \) corresponding to the above energy of \( \gamma \sim 9 \times 10^5 \), we get \( \tau_{\text{diff}} \sim 330 \text{ Myr} \). So for cross-field diffusion,\(^9\) the electrons could remain in the filaments for a long time relative to the loss timescale. Note, however, that we assume that the filaments have aligned along their length everywhere so as to imply cross-field diffusion for motion perpendicular to the filament. If we also assume effective pitch angle scattering, then the electrons can stream along the length of the filaments (and, presumably, off the end) on a timescale \( t_{\text{age}} \sim 0.1 \text{ Myr} \), which is much shorter than the loss timescale; moreover, unaged electrons must stream in at the same rate.\(^10\) If there is some component of the field perpendicular to the filament long axis, then the time-scale to move along the filament can become longer, but the time-scale to move across it becomes shorter. So the confinement time-scale is likely to be in the range \( \sim 0.1 \text{ Myr–300 Myr} \), and for a substantial fraction of that range \textit{in situ} losses may be significant, thus allowing the electrons in the filaments to be more aged than those around them. However, we emphasize that the lifetime of the filaments themselves places a hard upper bound on the possible excess ageing of electrons trapped in them. We return to this point in the following section.

### 4.1.3 Spectral ages

We now attempt to put constraints on the spectral (radiative) ages of the electrons in the filaments. Previously, Hardcastle et al. (2009) derived a synchrotron age of 24 ± 1 Myr for Centaurus A’s southern giant lobe region around the vertex and 29 ± 1 Myr for the region around the vertex. It is clear that, naively, the steep spectra of the filaments would imply spectral ages older than those of the lobes, which is hard to understand, as noted in the previous section. In this section, we explore whether excess-loss models of the sort described above can quantitatively explain the observed filament spectra. We use an injection index of 0.5 to be consistent with Hardcastle et al. (2009).\(^11\)

In this type of model, we need to consider two \( B \)-fields: \( B_{\text{present}} \), i.e. the field that the electrons are currently in, and \( B_{\text{age}} \), i.e. the one that they have undergone most of their ageing in (see in this context the preceding section). \( B_{\text{age}} \) is an emission-weighted average of all the \( B \)-fields the electron has been in since acceleration; it is possible for it to be lower, equal to or higher than \( B_{\text{present}} \), depending on the stage of the filament evolution. For a fixed break frequency, the synchrotron age is

\[
\tau_{\text{sync}} \propto \frac{B_{\text{present}}^2}{B_{\text{age}}^2 + B_{\text{CMB}}^2} ,
\]

where \( B_{\text{CMB}} \) is the equivalent magnetic field strength giving the cosmic microwave background (CMB) energy density, and represents inverse-Compton losses. If \( B_{\text{present}} > B_{\text{age}} \), e.g. if the electrons had spent most of their radiative lifetimes in a low-field region, this

\(^9\) There is potential for field line wandering and other mechanisms enhancing somewhat particle diffusion; our estimate might be considered a hard upper bound on the diffusion time.

\(^10\) In the presence of MHD turbulence, isolation of energetic particles and thus streaming of electrons at the speed of light over kpc-scale distances is unlikely; therefore, the quoted value should be considered a hard lower limit.

\(^11\) It is plausible that the injection index is slightly steeper (as indicated for FR I sources by, e.g. Young et al. 2005; Laing & Bridle 2013); the exact value of the injection index does not have a strong effect on our conclusions.
would make the filament spectral age exceed the spectral age of the lobe by an even larger factor. If we increase $B_{\text{age}}$ (relative to the rest of the lobe), we can in principle obtain a filament spectral age below that of the lobe; however, $B_{\text{age}}$ has to increase substantially because its equipartition value is significantly below $B_{\text{CMB}}$ (which is about 3.3 $\mu$G). In practice, though, a high value of $B_{\text{age}}$ could only be maintained on a time-scale comparable to the dynamical lifetime of the lobe, which we estimate (see Section 4.2.2) to be at most 3 Myr. Forcing the spectral age to be as low as this requires extreme values of $B_{\text{age}}$: fixing $B_{\text{present}}$ to its maximum plausible value of 5.5 $\mu$G (see Section 4.1.1), we obtain $B_{\text{age}} \sim 30$ $\mu$G, and even with $B_{\text{present}} \sim 0.9$ $\mu$G as in the giant lobes at present (i.e. allowing an increase in the electron number density to produce the observed increased emissivity of the filaments), we require $B_{\text{age}} \sim 18$ $\mu$G. These $B$-field strengths seem very high. Assuming that electrons are confined in the filament as discussed above, one can envisage a ‘squeezed state’ followed by a ‘relaxed state’ (i.e. recoil); then during the squeeze, the electrons in the filament are ageing faster; so when the system relaxes back to the ambient $B$-field strength, the electrons in the filament are ‘older’ than the electrons around them, and exhibit a steeper spectrum. However, in this picture very high magnetic fields must have been present at the peak of the squeezing.

We conclude that there is no very convincing explanation of the steep spectra of the filaments, if it is real, in the range of models we have considered. We emphasize that, because of the lack of a broad frequency range, our spectral age fit is not very robust, and more data are needed to confirm our picture. In addition, we have assumed a single electron population in the lobes. There could, in fact, be a second electron population (not to be confused with ‘secondary electrons’ from proton–proton or proton–photon collisions), from two or more episodes of jet activity with some time in between for radiative/adiabatic losses; this would elevate the emission at the low-frequency end of the range and might be made visible in the filaments by elevated $B$-field strengths there. However, there is hardly any spectral evidence for this second population either in Centaurus A itself or in other FR I and FR II sources, and so we regard this explanation as speculative without confirmation from low-frequency observations of the giant lobes.

### 4.1.4 Magnetic twist?

Both the ATCA 1.4 GHz and the low-frequency GMRT images seem to indicate a twisting of the vertex filament. In a flux tube with a quasi-cylindrical geometry and with axial field $B_z$ and azimuthal field $B_\varphi$, the amount of twist of a field line follows from the relation $\delta \phi = r \varphi / B_\varphi$ along a given field line. If the radius of the tube is $r$ and its length is $l$, the total amount of twist of a field line at the edge of the tube is obtained from (e.g. Priest 1994)

$$\Delta \phi = \frac{1}{r} \frac{B_\varphi(r)}{B_z}. \tag{3}$$

We have observational constraints on the filament length and width (see Table 4 and Section 4.1.1). If the twist is $\sim 2 \pi$, as seems from the available radio images, and hence the length over which the twist exists is $\sim 31$ kpc, then the ratio $B_\varphi/B_z \sim 0.5$. We will return to this in Section 4.2.3.

In the simple case, where the cylindrical filament carries a constant current density over its cross-section and for constant $B_z$, the static MHD equilibrium gives the gas pressure on the tube’s axis as $p(r = 0) = p(r) + B_\varphi^2(r)/4\pi$. From this and equation (3) follows that the observed amount of twist $\Delta \phi$ constrains $B_\varphi$, and thereby the pressure difference: using $B_\varphi \sim (r/\Delta \phi/l)B_z$ we can write

$$\frac{\Delta p}{p} = \frac{(r/\Delta \phi/l)^2}{1 + (r/\Delta \phi/l)^2} \frac{B_z^2}{4\pi p}. \tag{4}$$

Thus, if the twisting scenario is correct, for plasma $\beta = p/(B^2/8\pi) \sim 1$ (as is likely for the brightest filaments, see Section 4.1.1) and $r/\Delta \phi/l \sim 0.5$, and assuming that $B_z$ is constant inside and at the edge of the filament, the vertex cannot be compressed by a large factor and hence not greatly overpressured (i.e. compatible with strong shocks) with respect to the interfilament plasma.

### 4.2 Vertex and vortex: the origin

In this section, we put forward as the most attractive origins of the vertex and vortex (i) temporary enhancements of the activity of the extant or extinct jet leading to squeezing of the plasma in the southern giant lobe, and (ii) internal MHD turbulence. Feain et al. (2011) have likewise suggested jet activity, and additionally (iii) surface instabilities related to turbulence and (iv) the passage of the dwarf galaxy KK 196 through the lobe. We perform a ‘sanity check’ of the turbulent properties of the giant lobes, and we investigate whether the lobe plasma is prone to a development of HD and MHD instabilities.

#### 4.2.1 Jet activity

The idea that the vertex and vortex filaments might originate from an episode (or episodes) of (increased) jet activity and possibly represent weak shocks arises from, jointly, their 1.4 GHz and 325–235 MHz morphology and their orientation (see also Feain et al. 2011). The vertex and vortex may connect to the jet/core (as appears in the radio continuum images of e.g. M87; but note that weak, propagating shocks from the jet/core do not inevitably connect to it) whereas turbulence-induced filaments are not expected to be anchored to one another or anything else. At 1.4 GHz, 235 MHz, and 235 MHz, no connection to the current jet has been observed. For a filament age of about 2–3 Myr (see Section 4.2.2), it is not obvious that the filaments could originate from direct activity of the current jet if we believe its physical age (discussed by Wykes et al. 2013) is $\sim 2$ Myr. If the filaments result from variations in the core/jet activity and represent subsequent outbursts, the brighter of these, the vertex, ought to be spectrally and dynamically younger than the vortex, since it is smaller in size and positioned closer to the core. Arguments involving filament diameter in cases where the filaments appear as (half) rings have been used to compare filament ages in Hercules A and 3C 310 (Morrison & Sadun 1996). In the scenario in which the vertex and vortex filaments originate from a relatively recent (enhanced) direct jet activity, it would seem natural for them to exhibit a ‘leading edge’ and a ‘trailing edge’, but this is not seen at either 1.4 GHz or the GMRT frequencies. In this scenario, and if the cooling time of the electrons were short enough, we would also expect to see a spectral gradient along a path from upstream to downstream, which is not seen either. In any case, it is not plausible that Centaurus A’s jet (current jet power $\sim 1 \times 10^{43}$ erg s$^{-1}$, pre-existing jet power probably $\sim 1-5 \times 10^{43}$ erg s$^{-1}$; see Wykes et al. 2013) would produce these filaments as either terminal shocks or

\[12\] The morphology alone can easily be generated in simulations of isotropic turbulence.
as a bow shock, in this source. We additionally note that angular momentum from the jet rotation on small scales may add left-over rotational component to the lobes on large scales and possibly help to explain a twisting of the vortex; however, the spectral index gradient would then unlikely be as clear cut for this filament.

If the filaments are slightly higher pressure than the surroundings and are energetically plasma dominated (rather than B-field dominated), they expand with a velocity \( v_{\exp} \approx (\Delta p/\rho)^{1/2} \), where \( \Delta p = p_{\text{flow}} - p_{\text{fil}} \) is the difference between the filament and external pressure and \( \rho \) is the mass density (taken to be similar inside and outside the filament). For a filament width \( d \), the expansion time-scale is \( t_{\exp} = d/v_{\exp} \approx (p/\Delta p)^{1/2} (d/c_s) \), with \( c_s = \sqrt{\gamma kT/\mu m_\text{H}} \) the adiabatic sound speed. In what follows, we assume \( \Delta p/p = O(1) \), so that the expansion time-scale satisfies \( t_{\exp} \approx t_c \), the sound crossing time, and we estimate the time-scale for expansion considering the time-scale for sound to cross (the shortest dimension of) the structures

\[
t_c = d/c_s. \tag{5}
\]

Using the heat capacity ratio \( \gamma = 5/3 \), the mean particle mass \( \mu = 0.62 \) and the lower limit on the lobe temperature of \( 1.6 \times 10^8 \) K (Wykes et al. 2013), we obtain \( c_s \gtrsim 1.9 \times 10^8 \) cm s\(^{-1} \). Using the lobe temperature of \( 2 \times 10^12 \) K, derived by Wykes et al. (2013) based on entrainment modelling and pressure constraints for the lobes,\(^\text{13} \) we can set an upper limit on the sound speed using the relativistic expression \( c_s = c/\sqrt{3} \) which yields \( c_s \lesssim 1.7 \times 10^8 \) cm s\(^{-1} \). This gives us, taking the projected width of the vortex filament of 6.4 kpc, a time-scale for its expansion in the range \( t_c \sim 37 \) kyr—3.3 Myr. For the vortex filament projected width of 3.6 kpc, we obtain a time-scale for expansion in the range \( t_c \sim 21 \) kyr—1.9 Myr. The derived time-scales are possible to reconcile with the notion that these filaments, if plasma dominated and mildly overpressured, are driven directly by a pre-existing jet given the limit on the age of the current jet of \( \sim 2 \) Myr (see Wykes et al. 2013 and references therein).

If the filaments were more strongly overpressured (for example, if they are magnetically dominated as discussed in Section 4.1.1), then the expansion speeds would be faster than the external sound speed and the time-scales estimated above would be correspondingly reduced. However, our overall conclusion is that it is marginally possible for weakly overpressured filaments to have been produced directly by the last-existing jet, opening the giant lobes.

### 4.2.2 Internal MHD turbulence

Simulations (e.g. Clarke 1993; Lee et al. 2003; Tregillis, Jones & Ryu 2004; Schekochihin et al. 2004; Falceta-Gonçalves et al. 2010; Jones et al. 2011; Hardcastle 2013; TenBarge & Howes 2013) have shown that magnetic filaments develop naturally as a consequence of MHD turbulence. The vortex, vortex and also the fainter, ribbon-like filaments in the southern giant lobe of Centaurus A are reminiscent of the outcomes of such simulations, and we conjecture that their origin could be due to the turbulent motions in the lobe, presumably driven by current or recently terminated\(^\text{14} \) jet activity (Wykes et al. 2013), Turbulence folds and stretches the B-field, leading to its amplification (dynamo action) and a spatially intermittent B-field distribution. Filaments develop with lengths of the order of the largest eddies and transverse dimensions possibly as small as the viscous dissipation scale. B-fields coming via the turbulent dynamo concentrate into such structures. From the Jones et al. MHD simulations, the filaments are likely representing flux features stretched around the larger eddies in the turbulence. Even if the global helicity vanishes, some local ‘helicity’ might be present that could instigate filament twisting. Weak shocks with a Mach number \( M \sim 2 \) are evident in density images from two-dimensional HD simulations (e.g. Lee et al. 2003), and Eilek (2014) alluded to the existence of transonic flows, and hence weak shocks, augmenting locally the Alfvénic turbulence, in Centaurus A’s giant lobes.

Filaments persist about an eddy turnover time for the scale of the eddies that stretch them; filament longevity goes as \( \lambda^{13} \), which builds upon the Kolmogorov kinetic energy scaling (\( \langle E(k) \rangle \propto k^{-5/3} \)), where \( k = 2\pi/\lambda \) is the wavenumber). Velocity \( v_{\lambda} \) associated with scale \( \lambda \) compares to that on the driving scale (i.e. outer turbulence scale) \( v_{\lambda_{\max}} \) as

\[
v_{\lambda} \approx v_{\lambda_{\max}} (\lambda/\lambda_{\max})^{3/2}, \tag{6}
\]

with \( \lambda_{\max} \) the eddy size on the driving scale. However, the Kolmogorov scaling is not exactly in play when targeting the driving scale as processes on that scale take longer than on the inertial range, and so the filament longevity invoking \( \lambda_{\max} \) might be considered a lower limit. To establish the eddy turnover time on the driving scale, we use the turbulent speed \( v_{\lambda_{\max}} \) (i.e. \( v_{\lambda_{\max}} \) from Wykes et al. (2013), \( v_{\lambda_{\max}} \sim 1.9 \times 10^8 \) cm s\(^{-1} \), and a driving scale of 60 kpc (i.e. about a mid-value from their range 30—100 kpc); this gives \( \sim 3.1 \) Myr. Taking the vortex and vortex LAS (respectively, 34 and 53 kpc, see Section 1 and Table 4) as a scale below \( \lambda_{\max} \) and employing equation (6), then the speed at these scales is \( v_{\lambda} = 34 \) kpc \( \sim 1.6 \times 10^8 \) cm s\(^{-1} \) and \( v_{\lambda} = 53 \) kpc \( \sim 1.8 \times 10^8 \) cm s\(^{-1} \) and consequently their longevity, or ‘turbulent age’, is, respectively, \( \sim 2.1 \) and \( \sim 2.8 \) Myr.

We can test the turbulence-generated model by considering the power spectrum of structure in the lobe, since a given magnetic field power spectrum will give rise to a characteristic power spectrum of projected synchrotron emissivity (e.g. Eilek 1989) if we assume that the variations in electron energy density are small. Taking the two-dimensional Fourier transform of the 1.4 GHz synchrotron emission (as the 1.4 GHz image is the most sensitive available to us) and averaging its amplitude in radial bins to find the power on each spatial scale in the manner described by Hardcastle (2013), we find that the power spectrum on the scales of the lobe that can be reliably measured (i.e. smaller than the largest scale of the lobe itself and larger than the resolution) is flatter than expected from a Kolmogorov spectrum for the B-field but there is no particular a priori reason to expect such a spectrum in real MHD turbulence; on the other hand, despite the subtraction of point sources from the image (see Section 2.4), there is still a large amount of spurious small-scale structure which will affect the measured power spectrum. Overall, we can say that the turbulence model passes this consistency check, in the sense that the structure in the lobes does appear to have a power spectrum consistent with a power law; in particular, there is no evidence that the vortex and vortex are formed in a different

\(^{13} \) We have confirmed by considering the thermal bremsstrahlung emissivity of the giant lobes in the model of Wykes et al. that the INTEGRAL upper limits on photon counts from the giant lobes presented by Beckmann et al. (2011) are at least two orders of magnitude above the predictions of the model for any lobe temperature.

\(^{14} \) For a giant lobe disconnected from the energy supply, the time-scale for decay of turbulence, with the parameters derived by Wykes et al. (2013), is of the order of 6 Myr.
way from the other filaments in the lobes, as would be provided by, e.g. a deviation from a power law on scales comparable to those structures.

It is difficult to put (independent) quantitative limits on the fundamental parameters of fully developed turbulence, such as the turbulent speed $v_t$ at the driving scale, or the kinematic viscosity $v$ and the corresponding viscous dissipation scale $\lambda_r$, in an environment such as Centaurus A’s lobes. The rarefaction and high temperature of the lobe plasma (in the model described by Wykes et al. 2013) imply that binary (Coulomb) collisions between ions (mainly protons) and/or electrons are too infrequent to be physically relevant. This makes the conventional collisional estimate for the kinematic viscosity of the form $v \simeq v_{th,i}^3 \tau_{\text{scat}}/3 \simeq v_{th,i}^3/(3 \varepsilon_{\text{scat}})$, with $v_{th,i} = (3kT_i/m_i)^{1/2}$ the ion thermal velocity, $\tau_{\text{scat}}$ the ion–ion Coulomb collision time and $\varepsilon_{\text{scat}}$ the collisional frequency,16 equally inappropriate. To illustrate: the collisional mean free path for proton–proton collisions in a Maxwellian plasma of temperature $T$ is of the order of $\lambda_{\text{mfp,i}} \simeq v_{th,i} \tau_{\text{scat}}$. With a lower limit on the lobe temperature of $1.6 \times 10^8 \ K$ and a thermal proton number density $n_{\text{th,i}} \sim 5.4 \times 10^{-9} \ cm^{-3}$ (Wykes et al. 2013), and thus $\tau_{\text{scat}} \sim 1.4 \times 10^{20} \ s$, $\lambda_{\text{mfp,i}} \gtrsim 9 \ Gpc$, i.e. larger than the source size.

However, collective processes, in particular gyro-resonant interactions between protons and MHD waves or low-level MHD turbulence, likely produce an effective mean free path much smaller than the above estimate (see also e.g. Lazarian & Beresnyak 2006; Schekochihin & Cowley 2006; Santos-Lima et al. 2014). In the simplest model, with a low level of isotropic MHD turbulence with magnetic amplitude $\delta B$, pitch angle scattering by MHD waves (e.g. Wenzel 1974) restricts the mean free path along a large-scale magnetic field $B_0$ to $\lambda_{\text{mfp,i}} \simeq (\delta B/B_0)^2$ with $v_{\text{scat,i}} = v_{th,i}/\Omega_i$, the gyro-radius of thermal ions and $\Omega_i = ZeB_0/(\mu_0 c)$ the ion gyrofrequency.

There is some evidence that the power spectrum of MHD turbulence retains the Kolmogorov slope (e.g. Goldreich & Sridhar 1995; Goldstein, Roberts & Matthaeus 1995; Cho & Lazarian 2003, 2014; Kim & Ryu 2005; Kowal, Lazarian & Beresnyak 2007; Gaspari & Churazov 2013), with a flattening towards the driving scale. However, the ‘knee’ (i.e. the transition from the flat-spectrum range to the Kolmogorov slope) is not well constrained. Observations provide a turbulence level $\delta B/B_0$ on the largest expanse, for which one generally adopts the fiducial value $\delta B/B_0 \sim 1$ (e.g. O’Sullivan et al. 2009). A crude estimate of turbulence levels on small scales can be obtained by scaling down from the driving scale, to a dimension that we determine invoking the relations $\delta B/B_0 \propto (\lambda/\lambda_{\text{max}})^{1/3}$ and $\lambda_{\text{mfp,i}} \propto g_{\text{scat}}/(\delta B/B_0)^{2}$, and assuming $B_0$ similar on both scales. The mean free path derived in this way gives an upper limit on the true value since the magnetic field power spectrum may flatten towards the largest scales. Working from the driving scale of $\lambda_{\text{max}} = 60 \ kpc$ as above, we arrive at $\lambda_{\text{mfp,i}} \lesssim r_{\text{scat}}^{2/3} \lambda_{\text{max}}^{1/3} \lesssim 3.4 \times 10^{15} \ cm (\lesssim 1 \times 10^{-3} \ pc)$; such a proton effective mean free path is fairly below the size of the resolved fine structure in Centaurus A’s giant lobes. At this scale, the derived turbulence level is $\delta B/B_0 \lesssim 2.6 \times 10^{-3}$, which leads to an effective viscosity where large-scale motions (on scales significantly larger than $\lambda_{\text{mfp,i}}$) with velocity $v$ are dissipated with a dissipation rate per unit volume $\dot{v}$. If $\varepsilon_{\text{scat}}$ is weak ($\varepsilon_{\text{scat}} \ll \Omega_i$), the dissipation

\begin{equation}
\dot{v} = n_i m_i \varepsilon_{\text{eff}} \left( \frac{\partial v}{\partial \Omega} \right) = \frac{1}{3} \nabla \cdot v^2,
\end{equation}

see e.g. Kaufman (1960) and Braginskii (1965). The above assumes the effective collision time $\tau_{\text{scat}} \simeq \lambda_{\text{mfp,i}}/v_{\text{th,i}} \sim 1/\Omega_i (\delta B/B_0)^2$. Note that the dissipation rate scales as $\dot{v} \propto n_i m_i \varepsilon_{\text{eff}} v_i^2(\lambda/\lambda_r)^2$ with $\lambda$ the turbulent scale and $v_i(\lambda)$ the turbulent speed at that scale; this scaling with $v_i$ and $\lambda$ is the same as for ordinary (collisional) viscosity.

The purpose of the next three steps is to compute the scales at which the turbulent cascade dissipates: the viscous cutoff that marks the end of the HD structure evolution, and the resistive cutoff that marks the end of the MHD architecture, to finally obtain a crude estimate of the current sheet dimension likely associated with the resistive scale. This allows us to constrain the growth time of the tearing mode.

In a simple Krook collision model (Bhatnagar, Gross & Krook 1954) with parallel scattering mean free path $\lambda_{\text{mfp}}$ and thermal velocity $v_{\text{th,i}}$, the effective kinematic viscosity entering equation (7) is $\nu_{\text{eff}} \simeq v_{\text{th,i}} \lambda_{\text{mfp}} / 15$. Adopting $T = 1.6 \times 10^8 \ K$, $B = 0.9 \ \mu G$ and $\delta B/B_0 = 2.6 \times 10^{-4}$, we get $v_{\text{eff}} \sim 4.5 \times 10^{22} \ cm^2 \ s^{-1}$; this is considerably smaller than the kinematic viscosity based on Coulomb collisions. Using such effective viscosity, the Reynolds number is, adopting the driving scale and turbulent speed figures as above, $Re = \lambda_{\text{max}} v_{\text{th}}/v_{\text{eff}} \sim 7.7 \times 10^8$, which implies that a Kolmogorov cascade breaks off at $\lambda \simeq \lambda_{\text{max}}/Re^{1/4} \sim 7.1 \times 10^3 \ cm (\sim 2 \times 10^{-3} \ pc)$. Thus, in principle, a turbulent cascade could be set up from kpc to sub-pc scales comparable to $\lambda_{\text{mfp,i}}$, assuming $T$, $B$ and $\delta B/B_0$ levels as above.

The magnetic diffusivity due to electron scattering with scattering time $\tau_{\text{scat,e}}$, $\eta_{\text{m}} = c^2/4\pi \sigma$, with $\sigma = n_ec^2\tau_{\text{scat,e}}/m_e$ the electrical conductivity, is expected to be $\ll \nu_{\text{eff}}$ for typical parameters. In this situation, magnetic dynamo action is allowed at scales $\ll$ the dissipation scale of the large-scale kinetic turbulence (e.g. Tobias, Cattaneo & Boldyrev 2011). An actual computation of $\eta_{\text{m}}$ is hampered by the same problem as calculating the effective viscosity: the theory of Coulomb collisions, in this case for the current-carrying electrons, does not apply as the electron mean free path is comparable to the source size. We parametrize this uncertainty by expressing all parameters in terms of the (unknown) electron scattering mean free path $\lambda_{\text{mfp,e}}$ in the turbulent medium. For thermal electrons, the scattering time is $\tau_{\text{scat,e}} \simeq \lambda_{\text{mfp,e}}/(kT/m_e)^{1/2}$ and the magnetic diffusivity becomes $\eta_{\text{m}} = (c/\omega_{\text{pe}})^2 \tau_{\text{scat,e}}$, where $c/\omega_{\text{pe}}$ is the electron skin depth and $\omega_{\text{pe}} = (4\pi n_e e^2/m_e)^{1/2}$ the electron plasma frequency.

The magnetic Reynolds number $Re_{\text{m}}$ for large-scale fields is based on the magnetic diffusivity (as above) $\eta_{\text{m}} = c^2/4\pi \sigma$, with here $\sigma = \omega_{\text{pe}}^2 \tau_{\text{scat,e}}/4\pi r$ the conductivity of the plasma. Using for the scattering time the anomalous (i.e. turbulence-induced) value $\tau_{\text{scat}} \simeq \lambda_{\text{mfp,i}}/v_{\text{th,i}}$, with the bulk of the electrons (since they carry the current), we have $r_{\text{scat}} = g_{\text{scat}}(m_i/m_e)^{1/2} \sim 5.4 \times 10^8 \ cm (\sim 2 \times 10^{-10} \ pc)$ and $\tau_{\text{scat}} = \tau_{\text{scat,e}}(m_i/m_e)^{1/2} \sim 9.4 \times 10^8$. For $B = 0.9 \ \mu G$ and $\delta B/B_0 = 2.6 \times 10^{-3}$, this yields an anomalous conductivity $\sigma \sim 1.3 \times 10^8 \ s^{-1}$. The specific magnetic diffusivity is then $\eta_{\text{m}} \sim 5.6 \times 10^{23} \ cm^2 \ s^{-1}$, which is smaller than the

15 The Wykes et al. (2013) estimate of $v_t$ is based on the condition that a turbulent dynamo can reach energy equipartition, $U_p \sim U_i$.

16 The ion–ion collisional frequency is (e.g. Braginskii 1965) $\varepsilon_{\text{scat}} = 4\pi^2/2Z^2 n_in_i e^2 \ln \Lambda/(3m_i^{1/2}kT_i^{1/2})$, where $\ln \Lambda$ is the Coulomb logarithm in which $\Lambda = 4\pi n_i \lambda_D^3/3$, and $\lambda_D = (kT/4\pi n_i e^2)^{1/2}$ is the Debye length.
effective kinematic (ion) viscosity. This means that the magnetic Prandtl number \( \text{Pr}_m \), defined as \( \text{Pr}_m \equiv v_{\text{el}}/\nu > 8.1 \times 10^6 \), is large; therefore, magnetic turbulence can be maintained well below the viscous dissipation scale \( \lambda_\nu \). However, the character of the turbulence is different on those scales, since (both fast and slow) magnetosonic fluctuations are damped, velocity fluctuations are supposedly far sub-Alfvénic and the B-field lines are mostly ‘shuffled’. Relating \( \lambda_\nu \) to the magnetic Prandtl number, one finds the resistive dissipation scale, \( \lambda_\nu = \text{Pr}_m^{1/2} \lambda_\nu \sim 2.5 \times 10^2 \, \text{cm} \) \((\sim 8 \times 10^{-7} \, \text{pc})\). This is well below any observable scale, but we will re-appeal to the resistive cutoff in Section 4.2.3 in the context of the tearing mode.

### 4.2.3 (Surface) instabilities

We do not as yet have a tight observational constraint on whether the vertex and/or vortex filaments are surface features or are fully embedded in the lobe. An origin as magnetic flux ropes created by surface instabilities, such as KH or RT, is therefore somewhat appealing. Feain et al. (2011) have suggested the KH instability to be at the origin of some of the filamentary features in the Southern lobe, albeit without detailed arguments or support from modelling. Wykes et al. (2013) and Eilek (2014) have argued that the giant lobes are in an approximate pressure balance with their surroundings (a prerequisite for the KH instability to operate) and we further verify whether the KH instability could develop by considering its growth time (e.g. Chandrasekhar 1961):

\[
T_{KH} = \left[ \frac{(\rho_1 + \rho_2)^2}{k^2 \rho_1 \rho_2 v_{\text{el}}^2} \right]^{1/2},
\]

where \( k \) is the wavenumber, \( \rho_1 \) and \( \rho_2 \) the giant lobe and intragroup plasma mass densities and \( v_{\text{el}} \) the velocity of the lobe flow relative to the intragroup flow. Taking 5 kpc for the scale (that is, the mean filament thickness, see Table 4) and thus \( \sim 4.1 \times 10^{-22} \, \text{cm}^{-1} \) for the wavenumber, \( 1.1 \times 10^{-8} \, \text{cm}^{-3} \) (Wykes et al. 2013) and \( 1 \times 10^{-5} \, \text{cm}^{-3} \) (O’Sullivan et al. 2013; Eilek 2014) for, respectively, the total lobe and intragroup densities, and the buoyancy velocity of \( 4.9 \times 10^3 \, \text{cm} \, \text{s}^{-1} \) (\( \sim 0.002c \)); Wykes et al. (2013) for the lobe flow speed (while setting the intragroup speed to zero), we obtain a growth rate of \( \sim 0.1 \times 10^{-7} \, \text{yr}^{-1} \). From this it follows that the growth time is of the order of \( T_{KH} \sim 240 \, \text{Myr} \), which makes the origin of the vertex/vortex as a KH instability improbable. If, however, the interior plasma is moving faster than the buoyancy speed at which we propose the giant lobes are rising and/or the thermal vortex (projected) edges of the lobes with sizes comparable to the observationally given the fact that we do not see perturbations at the vertex or vortex.

The RT instability is driven by buoyancy in a stratified plasma and requires heavy material on top of lighter matter. It is not clear that the effect of RT (‘fingers’ of external medium intruding into the lobes) could give rise to structures resembling the vertex or vortex. The RM instability is analogous to the RT mode. It requires low-density material intruding impulsively into a higher density matter, or vice versa, producing voids in the former case and spike-like features in the latter. The RM instability is probably more relevant for very young lobes; we have no clear indication at such RM-like features in the available radio images of the filaments, and we do not expect strong shocks at the distance from the core of the vertex/vortex filaments in the giant lobes (see e.g. Sections 4.1.1 and 4.2.1).

The tearing instability develops due to small non-zero resistivity in high-current regions with B-field reversals. The primary requirement for this instability to be triggered is a thin field-reversal layer (the geometry as in magnetic reconnection), which could occur at the lobe/intragroup medium interface and probably also inside the lobes. We can legitimately consider anomalous resistivity, i.e. resistivity arising from particle-wave interactions (see e.g. Laval, Pellat & Vuillemin 1966; Melrose 1994 and references therein; Section 4.2.2). We can adopt an extreme upper limit on the anomalous collision rate (following Hines et al. 1989 in their treatment of the tearing mode possibly associated with lobe filaments in M87) which is the electron plasma frequency \( \omega_{pe} = (4\pi n_e e^2/m_e)^{1/2} \), and so the collision rate is \( v_{\text{coll}} \lesssim \omega_{pe}/2\pi \). This gives us an upper limit on the resistivity: \( \eta \lesssim 2\pi \omega_{pe} \). A lower bound on the growth time for the tearing mode then reads (Hines et al. 1989)

\[
T_{\text{TM}} \gtrsim \left( \frac{2 \pi \omega_{pe}}{v_{\text{A}}} \right)^{3/5} \left( \frac{a}{v_{\text{A}}} \right)^{8/5},
\]

in which \( a \) denotes the width of the current sheet and \( v_{\text{A}} \) the Alfvén speed. We use \( v_{\text{A}} = 2.4 \times 10^3 \, \text{cm} \, \text{s}^{-1} \) \((\sim 0.081c)\) and the thermal electron content \( n_{e,ib} = 5.4 \times 10^{-9} \, \text{cm}^{-3} \) (Wykes et al. 2013) as an approximation to the total electron number density (it is the full plasma which carries the waves). The width of the current sheet is not well known, yet an obvious hard lower limit is the electron gyroradius which is for our adopted conditions in the giant lobes \( r_{\text{e}} \sim 5.4 \times 10^4 \, \text{cm} \) \((\sim 2 \times 10^{-13} \, \text{pc})\). Probably, the most realistic estimate of the current sheet width\(^1\) is the resistive dissipation scale of the turbulence that we have estimated in Section 4.2.2 to be \( \lambda_\nu \sim 2.5 \times 10^2 \, \text{cm} \) \((\sim 8 \times 10^{-7} \, \text{pc})\). This gives \( T_{\text{TM}} \gtrsim 6.3 \, \text{h} \), which obviously falls well within both the dynamical and spectral ages of the lobes, and also within the turbulent ages of the individual vertex and vortex filaments (derived in Section 4.2.2). The resistive tearing mode is generally suppressed in high Re number media; however, the eddy turnover time at the tearing scale \((\sim 0.2 \, \text{yr})\) is longer than the lower limit on the tearing instability growth time; hence, the tearing is presumably not suppressed by this route. The topological consequences of the tearing instability are tiny closed B-field loops in 2D simulations and long magnetized filaments in 3D. Thus, while the tearing instability may be associated with the vertex and vortex filaments, it would be expected to show a non-vertex/vortex morphology and to appear on unobservable scales within them.

Various authors considered radiative instabilities, which arise as a result of runaway cooling, in radio galaxies’ lobes: synchrotron cooling, if relativistic electrons dominate the pressure (de Gouveia dal Pino & Opher 1989; Hines et al. 1989; Bodo et al. 1990; Rossi et al. 1993), or thermal bremsstrahlung cooling, if thermal plasma dominates (Hines et al. 1989; Bodo et al. 1990). To deal with the synchrotron instability, we consider relativistic electrons with a distribution of energies \( n_{\text{e}}(E_e) = N_0 E_e^{-\beta} \) between \( E_{e,\text{min}} \) and \( E_{e,\text{max}} \), where \( N_0 \) is the normalization of the electron energy spectrum and \( \beta \) is the electron spectral index. The growth time of the synchrotron

\(^1\)In the presence of MHD turbulence, a single stochastic magnetic reconnection region is associated with a multitude of such current sheets (e.g. Lazarian & Vishniac 1999). Magnetic reconnection and weak shocks associated with the filaments could locally boost the Alfvénic particle acceleration likely operating throughout the lobes (Eilek 2014).
instability, with also the inverse-Compton component included, can be written as the total electron energy divided by the total electron energy loss rate:

$$t_{\text{sync}} = \frac{\int E_{\text{e, min}} n_{\text{e,rel}}(E) \ E_d E_d}{C \int E_{\text{e, min}} n_{\text{e,rel}}(E) E_d^2 \ dt},$$

(12)

where $C$ is a constant such that for an electron $dE_e/dt = -CE_e^2$, and it represents both synchrotron and inverse-Compton losses: $C = 4 \pi \sigma T (U_{\text{B}} + U_{\text{CMB}})/(2m_e c^4)$, where $\sigma T$ denotes the Thomson cross-section, $U_{\text{B}}$ the energy density in aged $B$-field (as in Section 4.1.3) and $U_{\text{CMB}}$ the energy density in CMB photons (see also Hardcastle 2013). With $B = 0.9 \mu G$, $p = 2$, and $E_{\text{e, min}} = 5 \text{ MeV}$ and $E_{\text{e, max}} = 100 \text{ GeV}$ (which we have invoked in calculations of the filaments and the giant lobes) as energy cutoffs, the instability growth time becomes $\approx 110 \text{ Myr}$. This exceeds the turbulent ages of the individual filaments as well the spectral ages of the lobes. The synchrotron instability may be also excluded on grounds of the likely dominance of non-radiating particle pressure in the lobes (Wykes et al. 2013 for Centaurus A, e.g. Croston & Hardcastle 2014 for other FR I sources). The growth time of the thermal bremsstrahlung instability goes as (e.g. Karzas & Latter 1961)

$$t_{\text{th-b}} = 1.8 \times 10^{11} \frac{T^{1/2}}{n_{e, h} \bar{T}_H} \text{yr},$$

(13)

where $T$ is the lobe temperature and $\bar{T}_H$ the temperature-averaged Gaunt factor. Other thermal energy and emissivity parameters are absorbed into the numerical pre-factor. Using the range of temperatures $1.6 \times 10^9$ to $2.0 \times 10^{12} \text{ K}$ and $n_{e, h} = 5.4 \times 10^{-9} \text{ cm}^{-3}$ as above, and setting $\bar{T}_H = 10$ (Karzas & Latter 1961), we obtain a growth time range $\approx 1.3$–$99.6 \text{ Myr}$ which is much larger than the Hubble time. Even if the thermal density were as high as $1 \times 10^{-3} \text{ cm}^{-3}$ and temperature as low as $5.8 \times 10^9 \text{ K}$ (Stawarz et al. 2013), the thermal bremsstrahlung instability would still exceed the Hubble time, as expected. This makes a development of radiative instabilities in the giant lobes unrealistic.

The current-driven instabilities do not as much address the question of the origin of the filaments but rather the question what happens to them once formed; we will treat two such instabilities below.

The sausage instability, formally an axisymmetric $m = 0$ mode in a cylindrical magnetized plasma column (pinch), occurs when the azimuthal field $B_\phi$ becomes too strong. In that case, a sausage-like series of bulges and compressions in the tube radius grow; the exact growth rate of this mode depends on the field geometry both inside the pinch and in the surrounding medium. Filaments in other radio lobes, where measured, usually show a relatively high degree of linear polarization, implying the presence of significant axial field $B_\phi$. (e.g. Pictor A, Perley et al. 1997; 3C 310, van Breugel & Fomalont 1984; MG 0248+0641, Conner et al. 1998; Hercules A, Dreher & Feigelson 1984; Saxton, Bicknell & Sutherland 2002; Gizani & Leahy 2003). Also, the Junkes et al. (1993) polarization images of Centaurus A seem to show $E$-field vectors perpendicular to the $z$-axis of the filaments (albeit the resolution of these maps is not very high). This makes the sausage instability as a means of affecting the vertex/$vortex$ shape rather unlikely; moreover, inspecting in more detail the radio images, the vertex and vortex morphology does not well match the morphology as expected for the sausage mode.

The kink instability ($m = 1$ mode) eventuates when a plasma column moves more or less as a whole, and its axis is deformed into a

sinusoidal or helical shape. Analogously to the sausage mode, an exact determination of growth rates and the stability criterion depends on the detailed properties of the magnetic field inside and outside the pinch; the presence of an azimuthal field is destabilising.

The instability develops when the magnetic twist $\Delta \phi$ exceeds a critical value\(^{18}\) (Lundquist 1951; Hood & Priest 1979, 1981; Kliem, Titov & Török 2004). The evolution of the kink instability has been investigated by, e.g. Begelman (1998) and Haynes, Arber & Verwichte (2008), and Srivastava et al. (2010) have provided observational evidence for the $m = 1$ mode in the solar corona. The growth time of the kink instability can be approximated as (e.g. Bateman 1978)

$$t_{\text{m, k}} = \frac{r}{v_A} \left[ \frac{1}{k^2 r^2 - (B_\phi(r)/B_0)^2} \right]^{1/2},$$

(14)

where $k$ is again the wavenumber, $r$ is the flux tube radius and $B_\phi(r)$ the azimuthal field at the pinch surface (i.e. the edge of the filament). In the available radio images, the kink mode is mostly matched by the vertex morphology; hence, we will calculate the growth time for this particular filament. We consider the northern part of the vertex which seems to show a turn over $> 1$ wavelength (see Figs 2 and 3). The vertex radius is (at the widest point) $3.6 \text{ kpc}$; as before, we adopt $k = 4.1 \times 10^{-22} \text{ cm}^{-1}$ and $v_A = 2.4 \times 10^8 \text{ cm}^{-1}$, and $B_0/B_\phi \sim 0.5$ (from Section 4.1.4). This results in a growth time of $\approx 32 \text{ kyr}$. This is, as for the tearing mode, within both the dynamical and spectral ages of the lobes and also within the turbulent ages of the vertex and vortex filaments. A cotemporary development of the kink and the tearing mode, on disparate scales, is possible. Thus, the kink instability is likely to be operating: not destroying the vertex filament, just making it appear kinked.

4.2.4 Dwarf galaxy KK 196 (AM 1318–444)

We consider explanation (iv) in Section 4.2.4 the least likely given the low mass of the dwarf irregular galaxy KK 196, its low relative velocity to NGC 5128, a combination of high temperature and low particle density of the giant lobe plasma, and the vertex and vortex morphology; moreover, we would expect a single wake. We first test this assertion by performing a simple check for a galaxy to have a Bondi–Hoyle wake: the accretion radius is (based on its size) $r_{\text{acc}} = 2 G M_{\text{gal}}/(v_0^2 + c_s^2)$, where $M_{\text{gal}}$ and $v_0$ are the mass and the velocity of the traversing galaxy, and $c_s$ the local speed of sound (Bondi 1952; Sakelliou 2000). Adopting 15 arcsec as KK 196’s angular radius (Jerjen et al. 2000), we infer $\approx 1.8 \times 10^{21} \text{ cm}$ for its physical size.

The total baryonic mass of KK 196 is $\approx 5.9 \times 10^9 \text{ M}_\odot$ (based on the results from Jerjen et al. 2000; Warren, Jerjen & Koribalski 2007), and its relative radial velocity to NGC 5128 is approximately $189 \text{ km s}^{-1}$, so we estimate a relative 3D velocity of $v_\parallel \approx 189 \times \sqrt{3} \approx 327 \text{ km s}^{-1}$ ($\approx 0.001 \text{ c}$). Note that Crnojević, Grebel & Cole (2012) derive KK 196’s total baryonic mass as $5.1$–$7.2 \times 10^8 \text{ M}_\odot$ based on its $B$-band luminosity, a stellar mass-to-light ratio $L < M/L < 2$ and the measured H I mass, and as $6.9$–$8.2 \times 10^8 \text{ M}_\odot$ based on KK 196’s star formation history and the measured H I mass. The baryonic mass is thus well constrained and within $5$–$8 \times 10^8 \text{ M}_\odot$.

Assessing the non-baryonic dark matter (DM) mass fraction in dwarf irregular galaxies is difficult because the rotation curves are generally hard to measure, the H I rotation velocities increase only
very slowly as a function of galactocentric distance and peak typically at an amplitude not much larger than the turbulent motion of the gas. From a few known cases (see Côté, Carignan & Freeman 2000) follows that the DM mass is of the same order as the stellar mass. For KK 196, this translates into a total mass $M_{\text{tot}}$ of $4.5 \times 10^7 < M_{\text{stellar}} + M_{\text{gas}} + M_{\text{DM}} < 7.3 \times 10^7 M_\odot$. Note that the dwarf irregular ESO 444-G084 in the Côté et al. sample has a luminosity comparable to KK 196 and its $M_{\text{DM}}/M_{\text{lum}}$ value at $R_{25}$ (radius at the 25th mag arcsec$^{-2}$ surface brightness) is $0.3 \pm 0.5$. Thus, with the estimate given above we are likely to slightly overestimate KK 196’s total mass. However, they also quote $M_{\text{DM}}/M_{\text{lum}} = 10 \pm 12.9$ for ESO 444-G084 at $R_{\text{max}}$. If the comparison between KK 196 and ESO 444-G084 holds out to that larger radius then the total mass of KK 196 is of the order of $M_{\text{tot}} \sim 5 \times 10^8 M_\odot$. This value is also consistent with $M_{\text{tot}} = M_{\text{dyn}} = 1 - 4 \times 10^8 M_\odot$, where $M_{\text{dyn}}$ is the dynamical mass, inferred from Eder & Schombert (2000) adopting $M_B = -12.16$ and a colour $B - I = 1.5$ (i.e. total I-band luminosity of $-13.66$) for KK 196.

In the non-relativistic limit and assuming $T \geq 1.6 \times 10^5 K$ (Wykes et al. 2013), the local sound speed in the giant lobes reaches $\sim 1.9 \times 10^3 \text{ cm s}^{-1}$ ($\sim 0.006c$; see also Section 4.2.1) and consequently $R_{\text{acc}} \lesssim 3.6 \times 10^{18}$ cm. Hence, the accretion radius is at least three orders of magnitude smaller than KK196’s size and no Bondi–Hoyle wake is expected. If the internal sound speed in the lobes is relativistic (as is required in any scenario where there is stochastic acceleration of ultra-high energy cosmic rays), then the dwarf galaxy has even less effect.

Another possible manifestation of the galaxy–lobe interaction is a bow shock in front of the traversing galaxy and associated stripping of the galaxy ISM, then the filaments may be associated with a ram-pressure-induced wake. This however requires KK 196’s velocity to be in excess of the intralobe sound speed which is not the case (see above). Still a moderate gas stripping could occur from KK 196 in the absence of a bow shock: the non-detection of H I (Banks et al. above). Still a moderate gas stripping could occur from KK 196 in the absence of a bow shock: the non-detection of H I (Banks et al. above). Still a moderate gas stripping could occur from KK 196 in the absence of a bow shock: the non-detection of H I (Banks et al. above). Still a moderate gas stripping could occur from KK 196 in the absence of a bow shock: the non-detection of H I (Banks et al. above). Still a moderate gas stripping could occur from KK 196 in the absence of a bow shock: the non-detection of H I (Banks et al. above). Still a moderate gas stripping could occur from KK 196 in the absence of a bow shock: the non-detection of H I (Banks et al. above). Still a moderate gas stripping could occur from KK 196 in the absence of a bow shock: the non-detection of H I (Banks et al. above). Still a moderate gas stripping could occur from KK 196 in the absence of a bow shock: the non-detection of H I (Banks et al. above). Still a moderate gas stripping could occur from KK 196 in the absence of a bow shock: the non-detection of H I (Banks et al. above). Still a moderate gas stripping could occur from KK 196 in the absence of a bow shock: the non-detection of H I (Banks et al. above).

5 SUMMARY AND CONCLUSIONS

We have presented new, high-DR GMRT observations at 325 and 235 MHz of parts of the southern giant lobe of Centaurus A, and have modelled the origin of the filamentary structure associated with the lobe. The key results of this paper are as follows.

19 For dwarf galaxies, a significant increase in non-baryonic DM is common. The $R_{25}$ radius is roughly the size of the stellar component of the galaxy. However, the H I-based rotation curves tend to flatten out only further out (e.g. Warren, Jerjen & Koribalski 2006); hence, there might be a vast amount of non-baryonic DM in these systems beyond their optical radii (see also e.g. Simon & Geha 2007).
intralobe MHD turbulence or from last stages of the activity of the pre-existing jet, or an interplay of both.

There are several aspects that remain to be explored. The magnetic field direction and a possible variability of the spectral index along the filaments may be revealed by observations with the ATCA. ALMA will be the ideal facility for further investigations of the origin of the filaments, including the faint edge-like features, the wisps. In a future paper, we will report on XMM–Newton observations designed to constrain the magnetic field strength of the individual filaments. Reproducing the complex morphology of the vertex will be a challenge for future MHD simulations.

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