Broadband spectral modelling of bent jets of Active Galactic Nuclei

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In memory of
my beloved grandma
Josephine...
Abstract

The understanding of the physics of relativistic jets from active galactic nuclei (AGN) is still incomplete. A way to understand the different features of the AGN jets is to study it broadband spectra. In general, within the limits of present observations, AGN jets are observed in radio-to-X-ray energy band and they exhibit various intrinsic features such as knots. Particularly, the blazar jets which are pointed towards the observer, are observed in radio-to-\(\gamma\)-ray and their radio maps exhibit internal jet structures. Moreover, the high energy emission from blazars show rapid variability.

In this thesis, models have been developed to study the radiation emission processes from the knots of AGN jets as well as for blazar jets. A continuous injection plasma model is developed to study the X-ray emission from the knots of sources 1136-135, 1150+497, 1354+195 and 3C 371. The knot dynamics is then studied within the framework of internal shock model. In such a scenario, knots are formed due to the collision of two successive matter blobs emitted sporadically from the central engine of AGN. Shocks, generated in such collisions, accelerate electrons to relativistic energies. These electrons subsequently emit radiation via synchrotron and/or inverse Compton processes in the radio-to-X-ray energy range. The study of M87 knots involves a two zone model where the electrons with a power-law distribution are further accelerated. The synchrotron emission from these energetic electrons is then used to explain the observed spectrum.

Regarding the blazar jets, the limb-brightening feature observed in the radio maps of the BL Lac object MKN501 is studied considering shear acceleration of electrons at the boundary of the jet. This interpretation does not require a large viewing angle of the jet as demanded by the earlier models and is consistent with the constraints obtained from very high energy studies. In case of MKN421, the dependence of temporal behaviour of radiation emission on the particle acceleration mechanism has been studied within the framework of two zone model.
List of Publications relevant to this Thesis

(1) A Continuous Injection Plasma Model for the X-Ray/Radio Knots in Kiloparsec-Scale Jets of Active Galactic Nuclei
   Sahayanathan, S.; Misra, R.; Kembhavi, A. K.; Kaul, C. L.
   Astrophysical Journal Letters (2003), 588, L77-L80

(2) Interpretation of the Radio/X-Ray Knots of AGN Jets within the Internal Shock Model Framework
   Sahayanathan, S.; Misra, R.
   Astrophysical Journal (2005), 628, 611-616

(3) Particle acceleration process and temporal behaviour of non-thermal emission from blazar
   Bhattacharyya, S.; Sahayanathan, S.; Bhatt, N.
   New Astronomy (2005), 11, 17-26

(4) A two-zone synchrotron model for the knots in the M87 jet
   Sahayanathan, S.
   Monthly Notices of the Royal Astronomical Society Letters (2008), 388, L49-L53

(5) Boundary shear acceleration in the jet of MKN501
   Sahayanathan, S.
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## Contents

Abstract i

List of Publications relevant to this Thesis ii

Acknowledgements iii

Acronyms xi

1 Introduction 1

1.1 Active Galactic Nuclei 1

1.1.1 Historical Background 1

1.1.2 Morphology 5

1.1.3 Classification 6

1.1.4 Unification 8

1.2 Jets of Active Galactic Nuclei 9

1.2.1 Evidence for Relativistic flow 9

1.2.2 AGN Jet Features 13

2 Particle Acceleration and Radiative Processes 19

2.1 Fermi Acceleration 19

2.1.1 Stochastic Acceleration - Second Order Fermi Mechanism 22
| Section | Title | Page |
|---------|-------|------|
| 2.1.2  | Shock Acceleration - First Order Fermi Mechanism | 23 |
| 2.1.3  | Acceleration at Shear Layers | 24 |
| 2.1.4  | Accelerated particle distribution | 26 |
| 2.2    | Radiation Emission Mechanisms | 27 |
| 2.2.1  | Blackbody Radiation | 27 |
| 2.2.2  | Bremsstrahlung Radiation | 27 |
| 2.2.3  | Synchrotron Radiation | 28 |
| 2.2.4  | Inverse Compton Radiation | 32 |
| 2.3    | Equipartition Magnetic Field | 36 |
| 2.4    | Hadronic Processes | 38 |
| 2.5    | Emission Models | 40 |
| 2.5.1  | Leptonic Models | 40 |
| 2.5.2  | Hadronic Models | 44 |
| 2.6    | Aim of the Thesis | 46 |
| 3      | Emission Models and the Dynamics of AGN knots | 50 |
| 3.1    | A Continuous Injection Plasma Model | 52 |
| 3.1.1  | Results and Discussion | 54 |
| 3.2    | Internal Shock Interpretation | 60 |
| 3.2.1  | *Chandra* X-ray Data Analysis | 60 |
| 3.2.2  | The Model | 62 |
| 3.2.3  | Results and Discussion | 65 |
| 3.3    | A Two Zone Model for the knots of M87 | 70 |
| 3.3.1  | The Model | 70 |
| 3.3.2  | Results and Discussion | 74 |
## CONTENTS

4 Boundary Shear Acceleration in the Jet of MKN 501

4.1 Shear acceleration at MKN 501 jet boundary ........................................... 80
4.2 Particle Diffusion at the jet boundary and Limb-brightening .................... 84
4.3 Spectral index ................................................................................................. 85
4.4 Discussion ....................................................................................................... 86

5 A Two zone Model for Blazar Emission Mechanism

5.1 The Model ....................................................................................................... 89
5.2 Results and Discussion .................................................................................. 92
   5.2.1 Spectral evolution ..................................................................................... 92
   5.2.2 Flare ......................................................................................................... 96

6 Summary and Conclusions................................................................................ 103

Bibliography ....................................................................................................... 107
List of Figures

1.1 VLA image of the AGN 3C 175 ........................................... 5
1.2 AGN Classification .......................................................... 6
1.3 Schematic diagram of an AGN ............................................. 9
1.4 Superluminal motion - Representation ............................... 10
1.5 Superluminal motion - Plot of apparent velocity versus viewing angle .... 11
1.6 Multi wavelength image of 3C 273 .................................... 15
1.7 VLA image of 3C 308 ....................................................... 16
1.8 VLBA images of PKS 2136+141 ......................................... 17
1.9 VLA image of M87 with limb-brightened feature .................. 18
2.1 Illustration of Fermi acceleration mechanism .......................... 20
2.2 Particle acceleration at sheared flow .................................. 25
2.3 Intensity spectrum of bremsstrahlung radiation ...................... 29
2.4 Synchrotron power function ............................................. 30
2.5 Spectral fits of 3C279 using leptonic and hadronic models .......... 45
3.1 Multi wavelength image of 1136-135 .................................. 55
3.2 Multi wavelength image of 1150+497 and 1354+195 ................. 56
3.3 Multi wavelength image of 3C371 ...................................... 57
3.4 Spectral fit for the AGN knots using continuous injection plasma model .. 58
3.5 Spectral fits for the knots using internal shock model ........................... 66
3.6 Multi wavelength image of M87 ......................................................... 71
3.7 Spectral fit for the knots of M87 using two zone model ......................... 75
4.1 Radio images of MKN 501 jet ............................................................. 79
4.2 Spectral index map of MKN 501 jet ..................................................... 86
5.1 Spectral fits of MKN 421 using two zone model .................................... 94
5.2 Evolution of the model spectra ......................................................... 95
5.3 Synchrotron and SSC light curves (energy independent acceleration time-scale) 97
5.4 Hysteresis curves during a flare (energy independent acceleration time-scale) 98
5.5 Synchrotron and SSC light curves (energy dependent acceleration time-scale) 100
5.6 Hysteresis curves during a flare (energy dependent acceleration time-scale) 101
5.7 Variation of lag time-scale ............................................................ 102
List of Tables

3.1 Continuous Injection Plasma Model Parameters .......................... 59
3.2 Details of Chandra Observations of AGN jet ............................. 61
3.3 Observed Knot Features ...................................................... 62
3.4 Internal Shock Model Parameters ......................................... 67
3.5 Knot/Jet Properties ............................................................ 67
3.6 Two Zone Model Parameters for the Knots of M87 ................. 76
5.1 Two Zone Model Parameters used for fitting MKN 421 ............... 93
5.2 Values of $\eta$ for different frequencies ............................... 99
## Acronyms

| Acronym  | Description                                                                 |
|----------|-----------------------------------------------------------------------------|
| 3C       | Third Cambridge Catalogue of Radio Sources                                  |
| ACIS     | Advanced CCD Imaging Spectrometer                                          |
| ACT      | Atmospheric Cherenkov Telescope                                             |
| AGN      | Active Galactic Nuclei                                                     |
| ASCA     | Advanced Satellite for Cosmology and Astrophysics                          |
| BLRG     | Broad Line Radio Galaxies                                                  |
| CANGAROO | Collaboration of Australia and Nippon (Japan) for a GAmma Ray Observatory in the Outback |
| CAT      | Cherenkov Array at Themis                                                   |
| CCD      | Charge Coupled Device                                                       |
| CELESTE  | Cerenkov Low Energy Sampling and Timing Experiment                           |
| CIAO     | *Chandra* Interactive Analysis of Observations software                     |
| CMB      | Cosmic Microwave Background                                                 |
| CTA      | Cherenkov Telescope Array                                                   |
| EGRET    | Energetic Gamma-Ray Experiment Telescope                                    |
| FR I     | Fanaroff-Riley I                                                            |
| FR II    | Fanaroff-Riley II                                                           |
| FSRQ     | Flat spectrum Radio Quasars                                                 |
| HEGRA    | High-Energy-Gamma-Ray Astronomy                                             |
| HESS     | High Energy Stereoscopic System                                             |
| HST      | Hubble Space Telescope                                                      |
| IACT     | Imaging Atmospheric Cherenkov Telescope                                     |
| IC/CMB   | Inverse Compton scattering of Cosmic Microwave Background                   |
| M87      | Messeir 87                                                                  |
| MACE     | Major Atmospheric Cherenkov Experiment                                      |
| MAGIC    | Major Atmospheric Gamma-ray Imaging Cherenkov Telescope                     |
| MERLIN   | Multi-Element Radio Linked Interferometer Network                           |
| Acronyms | Description |
|----------|-------------|
| MKN      | Sources from Markarian Catalogue |
| NASA     | National Aeronautics and Space Administration |
| NGC      | New General Catalogue |
| NLRG     | Narrow Line Radio Galaxies |
| PACT     | Pachmarhi Array of Cherenkov Telescopes |
| PKS      | Parkes Catalogue of Radio Sources |
| QSO      | Quasi Stellar Objects |
| SED      | Spectral Energy Distribution |
| SSC      | Synchrotron Self Compton |
| SSRQ     | Steep Spectrum Radio Quasars |
| STACSE   | Solar Tower Atmospheric Cherenkov Effect Experiment |
| TACTIC   | TeV Atmospheric Cherenkov Telescope with Imaging Camera |
| VERITAS  | The Very Energetic Radiation Imaging Telescope Array System |
| VLA      | Very Large Array |
| VLBA     | Very Long Baseline Array |
| VHE      | Very High Energy |
| XJET     | X-ray emission from Extragalactic Radio Jets |
| (http://hea-www.harvard.edu/XJET/) |
| XSPEC    | X-Ray Spectral Fitting Package |
Chapter 1

Introduction

1.1 Active Galactic Nuclei (AGN)

The nucleus of a galaxy with luminosity \( \gtrsim 10^{44} \text{ ergs s}^{-1} \) and exceeding the overall luminosity of the entire host galaxy is known as Active Galactic Nucleus. Galaxies which host an AGN are called as active galaxies.

1.1.1 Historical Background

AGN was first observed in optical by Fath \([1]\) in 1908 at Lick Observatory. His aim was to test the claim of that time that the spectra of the spiral nebulae are continuous, consistent with a collection of stars. The continuous spectrum and absorption lines he observed for most of the sources suggested the presence of an unresolved collection of stars. However, for the nebula NGC 1068, the spectrum was composite showing emission and absorption lines. Later a higher quality and better resolution spectrum of NGC 1068 obtained by Slipher \([2]\) at Lowell Observatory confirmed the presence of emission and absorption lines. Slipher also observed the emission lines are broad and spread over a substantial range of frequencies. Seyfert \([3]\) was the first to do a systematic study of such galaxies with strong nuclear emission and broad lines. He attributed the broadness of emission lines to Doppler shifts and obtained a maximum velocity spread of 8500 km s\(^{-1}\) for the hydrogen lines of NGC 3516 and NGC 7469. This class of objects with broad emission lines and a luminous core was later called as Seyfert galaxies (the most numerous type of AGN known).

Though the observation of AGN began as early as the beginning of the twentieth century,
a major study of these objects started only after the advent of radio astronomy initiated by Karl Jansky \[4\], a radio engineer working at Bell Telephone Laboratories. Later Grote Reber \[5\], a radio engineer working on radio astronomy during his spare time, published a map of the radio sky at 160 MHz using his 31 feet reflector in his backyard. The map included the bright source located in the constellation Cygnus which we know at present as an AGN and commonly called as Cygnus A. In 1951, Smith \[6\] used radio interferometry and obtained accurate positions of the radio sources Taurus A, Virgo A, Cygnus A and Cassiopeia A, which led Bade and Minkowski \[7\] to identify the latter two in the optical band. The observed emission lines of these sources were very similar to that of Seyfert galaxies and radio sources with this characteristic feature in their spectrum are now identified as radio galaxies. There was also a considerable progress in the study of the radio source structures during this period. Jennison and Das Gupta \[8\] used radio interferometry to study the structure of Cygnus A which showed two bright components separated by \(\sim 1.5\) arcmin. This morphology was observed to be common for extragalactic radio sources. Subsequent observations at GHz frequencies provided better angular resolution and revealed the presence of a compact nucleus often associated with extended radio sources. In addition to these radio galaxies, some star-like radio sources were also detected. Their spectra were continuous with broad emission lines and no absorption lines. These peculiar stars were then called as “radio stars”. Many attempts were made to understand them as an old novae or white dwarfs (\[9\] and references therein) until Schmidt \[10\] identified the lines of 3C 273 as nebular emission lines with redshift \(z = 0.158\) and Greenstein and Matthews \[11\] identified emission lines in 3C 48 with redshift \(z = 0.367\) (the largest redshift known at that time). It was clear then that radio stars are highly luminous sources situated at large distances. These radio stars are today called as “quasi-stellar radio sources” or simply “quasars” (a term coined by Chinese-born U.S. astrophysicist Hong-Yee Chiu \[12\]). Seyfert galaxies, radio galaxies, quasars and objects of similar type are all collectively termed as \textit{Active Galactic Nuclei (AGN)}. Quasars are the distant AGN for which the host galaxy cannot be resolved.

Janksy \[13\] suggested that the origin of the observed radio emission may be (a) from the stars or (b) secondary emission from the atmosphere due to the interaction of high energy particles emitted by the stars or (c) thermal emission from the interstellar dust. Former two options suggested maximum radio intensity from the direction of the sun since it being the closest star. However such an excess from the direction of the sun was not observed and hence these options are not viable. Meanwhile, Whipple and Greenstein \[14\] calculated the interstellar dust temperatures and it was found that they are too low to produce the radio
intensity observed by Jansky. Reber [15] suggested the free-free emission by ionised gas in the interstellar medium as a plausible mechanism for the observed radio emission. However, Henyey and Keenan [16] and Townes [17] showed that even this cannot reconcile the required Jansky’s brightness temperature. These studies thereby ruled out the possibility of the radio emission from the interstellar dust. Finally in 1950, Kiepenheuer [18] explained the galactic radio background in terms of synchrotron radiation by cosmic rays in the galactic magnetic field. By the end of 1950’s, the synchrotron theory got accepted for explaining the radio emission from the extragalactic radio sources. Later the power-law\(^1\) nature of the spectra and high degree of polarisation observed also supported this theory.

In 1959, Burbidge [19] used synchrotron theory and found that the minimum energy content of the radio sources is extremely large \(\sim 10^{60}\) ergs. Burbidge attributed this large energy to a chain of supernovae explosions occurring in a tight packed set of stars at the nuclear region of the galaxy. Such a situation is plausible since the nuclear region is expected to be much denser than the average density of the galaxy. In 1962, Hoyle and Fowler [20] suggested that such close packed stars can as well form a single super massive object. Accordingly they proposed the existence of a super massive star at the nuclear region of the galaxy and conjectured that it would evolve into a ‘super-supernovae’ providing the required energy. However the thermonuclear process is less efficient in converting mass into energy compared with the gravitational one for the masses of this order. Again, Hoyle and Fowler [21] proposed in their pioneering paper in Nature in early 1963, that the energy of the radio sources was of gravitational origin derived from the slow contraction of a super massive object under its own strong gravitational field. This theory of a collapsed super massive object powering the AGN is widely accepted at present. We refer to this object as Central Engine hereafter.

Today AGN are detected up to \(\gamma\)-ray energies and their structures are studied even at X-ray energies (see Chapter 3). AGN are first detected in \(\gamma\)-rays by satellite based experiments followed by ground based experiments using atmospheric Cherenkov techniques. 3C 273 was the first AGN detected in \(\gamma\)-ray (MeV) by Cos-B \(^2\) in 1978 [22]. Later EGRET \(^3\) operating at MeV-GeV \(\gamma\)-ray energy range detected around 60 AGN during 1991-2000 [23].

---

\(^1\)We define a power-law flux as \(F_\nu \propto \nu^{-\alpha}\) (ergs cm\(^{-2}\) s\(^{-1}\) Hz\(^{-1}\)) where \(\nu\) is the observed photon frequency and \(\alpha\) the power-law spectral index.

\(^2\)Cos-B was an European Space Research Organisation satellite mission to study \(\gamma\)-ray sources launched by NASA.

\(^3\)Energetic Gamma-Ray Experiment Telescope (EGRET) is a satellite borne \(\gamma\)-ray telescope onboard Compton Gamma-Ray Observatory (CGRO) launched by NASA.
tional with the detection of MKN 421 in TeV $\gamma$-ray by Whipple$^4$ in 1992 [24]. Later imaging atmospheric Cherenkov telescopes (IACT) (e.g. HEGRA$^5$, TACTIC$^6$, CAT$^7$, CANGAROO$^8$ etc.) and wavefront sampling Cherenkov telescopes (e.g. CELESTE$^9$, PACT$^{10}$, STACEE$^{11}$ etc.) detected 6 AGN at TeV $\gamma$-ray energies [25]. All AGN detected in $\gamma$-ray energies during this period are blazars (a class of AGN with a jet of matter flowing at relativistic speed towards the observer ([1.1.4] and [1.1.3])) except the radio galaxy Cen A which was detected by EGRET. Presently the satellite based experiment Fermi$^{12}$ operating at GeV-TeV $\gamma$-ray energies detected around 600 AGN till now which include blazars and radio galaxies [26]. The second generation IACT namely MAGIC$^{13}$, HESS$^{14}$ and VERITAS$^{15}$ operating at GeV-TeV $\gamma$-ray energies detected around 40 AGN which again include blazars and radio galaxies$^{16}$. The number of AGN detected in $\gamma$-ray energy range are likely to increase with the help of these experiments and the future experiments (MACE$^{17}$ and CTA$^{18}$). Many of the AGN detected by Fermi and second generation IACT along with simultaneous observation at radio, optical and X-ray energies cannot be explained with our present understanding and the problems regarding AGN physics are still open [27].

---

$^4$Whipple is located at the Fred Lawrence Whipple Observatory in Southern Arizona, USA.
$^5$High-Energy-Gamma-Ray Astronomy (HEGRA) telescope is located at Roque de los Muchachos Observatory on La Palma.
$^6$TeV Atmospheric Cherenkov Telescope with Imaging Camera (TACTIC) is located at Mt. Abu, Rajasthan, INDIA.
$^7$Cherenkov Array at Themis (CAT) is located at Themis, France
$^8$Collaboration of Australia and Nippon (Japan) for a GAmma Ray Observatory in the Outback (CANGAROO) is located at Woomera, Australia.
$^9$CErenkov Low Energy Sampling and Timing Experiment (CELESTE) is located at Themis, France
$^{10}$Pachmarhi Array of Cherenkov Telescopes (PACT) is located at Pachmarhi, Madhya Pradesh, INDIA
$^{11}$Solar Tower Atmospheric Cherenkov Effect Experiment (STACEE) is located near Albuquerque, New Mexico.
$^{12}$Fermi Gamma-ray Space Telescope is a satellite based $\gamma$-ray telescope. The mission is a joint venture of NASA, the United States Department of Energy, and government agencies in France, Germany, Italy, Japan, and Sweden.
$^{13}$Major Atmospheric Gamma-ray Imaging Cherenkov Telescope (MAGIC) is a system of two IACT situated at the Roque de los Muchachos Observatory, La Palma
$^{14}$High Energy Stereoscopic System (HESS) is a stereoscopic IACT located at Namibia.
$^{15}$Very Energetic Radiation Imaging Telescope Array System (VERITAS) is an array of four IACT located at Fred Lawrence Whipple Observatory in southern Arizona, USA.
$^{16}$http://www.mppmu.mpg.de/ rwagner/sources/
$^{17}$Major Atmospheric Cerenkov Telescope Experiment (MACE) is an upcoming stereoscopic IACT at Hanle, India.
$^{18}$Cherenkov Telescope Array (CTA) is a proposed open observatory and will consist of two arrays of IACT with one array at northern hemisphere and a second array at southern hemisphere.
1.1.2 Morphology

AGN comes in a variety of morphological structures and sizes. At one extreme we find an unresolved compact luminous core and in the other end we have, complex structures extending up to hundreds of kiloparsec (kpc, $1\text{pc} = 3.0857 \times 10^{18} \text{ cm}$). However one can develop a primary morphological feature of an AGN based on the commonly observed characteristics (Figure 1.1).

- **Core**: These are compact unresolved luminous radio components coinciding with the nucleus of the associated galaxy (if resolved). They mostly have a spectrum which is a flat power-law with an index $\alpha < 0.5$. Cores are found in almost all quasars and in nearly 80 per cent of all radio galaxies.

- **Lobes**: These are two extended regions of radio emission located on opposite sides of the galaxy or the nuclei. Lobes can be separated by several hundred kpc or in some extreme cases even up to a few megaparsec (Mpc) (e.g. 3C 236 has an overall size of 4 Mpc). They have a steep power-law spectrum with index $\alpha > 0.5$. Within the lobes one often finds local intensity maxima commonly called as **hotspots**.
AGN Classification

Fig. 1.2: AGN Classification

- **Jets**: These are narrow features that connect the compact core to the outer regions commonly referred to radio jets. They extend from pc scale to kpc scales. Radio jets often have bright regions along its length which are called as radio knots or simply knots. We shall consider the jets in detail in §1.2.

1.1.3 Classification

AGN are broadly classified into two groups, namely radio loud and radio quiet based on the ratio of their radio luminosity at 5 GHz to the optical luminosity at the B-band. Conventionally the demarcating value of this ratio is taken around \( \frac{L_{5\text{GHz}}}{L_B} \approx 10 \) and roughly 15-20% of AGN are radio loud.

**Radio Quiet AGN**

The commonly observed radio quiet AGN are the Seyfert galaxies which have a morphology of spiral galaxies. They are classified based on their optical/UV properties.
Chapter 1. Introduction

- **Seyfert Type 1**: These are the Seyfert galaxies having a bright star-like nucleus and emit a strong continuum from far infrared to the X-ray band. Their spectrum contain broad emission lines with a width of the order of few thousand kilometers per second.

- **Seyfert Type 2**: These are the Seyfert galaxies with weak continuum and narrow emission lines with a line width of the order of few hundred kilometers per second.

Seyfert galaxies with properties intermediate between type 1 and type 2 are classified as types 1.2, 1.5 and so on. Although Seyfert galaxies are categorized as radio quiet sources there also exists radio loud Seyfert galaxies detected even up to $\gamma$-ray energies [29].

**Radio Loud AGN**

Radio loud AGN are classified based on their morphology and optical/UV properties.

- **Fanaroff-Riley I (FR I) radio galaxies**: These are extended sources often having morphology with core, jet and lobes. Their jets are often symmetric and become fainter as one approaches the outer extreme. Due to this feature they are called as edge-darkened sources. Their spectrum contain narrow emission lines features.

- **Fanaroff-Riley II (FR II) radio galaxies**: These are more luminous than FR I radio galaxies with knotty jet and lobes with hot spots. In contrast to FR I type, their intensity falls off towards the nucleus and hence they are called as edge-brightened sources. Jets of these types of sources are well collimated and often one sided. Their spectrum contain narrow emission lines features.

  *FR I* and *FR II* galaxies are often elliptical and together are referred to as *Narrow Line Radio Galaxies (NLRG)*.

- **Broad Line Radio Galaxies (BLRG)**: These sources have a continuum and emission lines resembling those from Seyfert 1 galaxies.

- **Radio Quasars**: These sources are distinguished from BLRG based on the luminosity and have a luminous nucleus which outshines the light from the host galaxy. These sources are often observed to have one sided jets with superluminal knots (see §1.2.1). They are further classified into *Steep Spectrum Radio Quasars (SSRQ)* and *Flat Spectrum Radio Quasars (FSRQ)* depending on their radio spectral index is either $\alpha \gtrsim 0.5$ or $\alpha \lesssim 0.5$. 
• **BL Lacs:** These sources have strong nuclear continuum with high polarisation and rapid flux variability. Their continuum extends from radio to γ-ray energies with few of them detected at TeV energies [25]. Their emission lines are absent or weak.

BL Lacs and FSRQ are collectively called as *Blazars* since they share many common properties like, rapid variability, high and variable polarisation, superluminal motion etc.

### 1.1.4 Unification

The idea of unifying the different classes of AGN as a variant of objects belonging to a single population emanated from the fact that they share many common observational features. The most successful theory in AGN unification is based on the *Orientation hypothesis*, which assumes that the observed differences between the different classes of AGN are due to different orientations of similar objects with respect to the line of sight [30, 31] as shown in Figure 1.3.

The unification of *Seyfert 1* and *Seyfert 2* is based on the fact that broad emission lines are present in the polarised spectra of the latter. Hence the unified picture derived for these sources assume that the broad line emitting regions are located closer to the central engine and a “dusty torus” is wrapped around the central region. Consequently the broad emission lines are hidden by the torus when the source is viewed edge-on (i.e. when the angle between the line of sight and the axis of the torus is large). On the other hand, narrow line emitting regions are farther from the central engine and hence are not obscured. Also, the unified picture assumes the presence of hot electrons far from the central region which are believed to reflect the broad lines towards the observer when the source is viewed edge-on. These reflected lines will therefore be visible in the polarised spectrum. Considering the above mentioned picture of the source, the edge-on view will resemble as *Seyfert 2* while the face-on view as *Seyfert 1* [30]. Also the unified picture of AGN contains twin relativistic jets emanating from the central engine, normal to the obscuring torus. High luminosity sources, *Quasars* and *FR II* are considered to be similar objects with the jets of the former aligned close to the line of sight (angle < 14°) and thereby causing relativistic beaming. Similarly, *BL Lacs* are the aligned jet version of *FR I* (low luminosity sources) [31].
1.2 Jets of Active Galactic Nuclei

The bridge of radiation connecting the central compact source with the extended lobes of the AGN is called as a “Jet”. They are the common feature observed in most of the AGN and are either two-sided or one-sided. They may extend from pc to kpc scales and are visible in radio, optical and X-rays. Jets are interpreted as conduits for the transport of energetic particles from the nucleus to the extended radio structures at relativistic velocities.

1.2.1 Evidence for Relativistic flow

Superluminal Motion

For many AGN, jet components are observed to move with velocities greater than the velocity of light and this phenomenon is referred to superluminal motion \cite{32}. Rees \cite{33} interpreted this phenomenon as a result of bulk relativistic flow at an angle close to the line of sight. To understand this, let us consider a blob ejected from a stationary source with a velocity $v$ at an angle $\psi$ with respect to the line of sight of the observer as shown in the
Figure 1.4. Let the blob emit signals for a duration $dt_e$, measured from the frame of the stationary source. The distance travelled by the blob during this time is $v dt_e$. Due to the inclined motion, the signal emitted by the blob at the end of the duration $dt_e$, travel lesser distance compared to the one emitted at the beginning of $dt_e$ for a distant observer. Hence the observer will measure this interval as

$$dt_o = dt_e \left( 1 - \frac{v}{c} \cos \psi \right) \quad (1.1)$$

where $c$ is the velocity of light. Since the projected distance travelled by the blob in the sky plane is $v dt_e \sin \psi$, the apparent velocity measured by the observer will be

$$v_a = \frac{v dt_e \sin \psi}{dt_o} = \frac{v \sin \psi}{1 - \frac{v}{c} \cos \psi} \quad (1.2)$$

or

$$\beta_a = \frac{\beta \sin \psi}{1 - \beta \cos \psi} \quad (1.3)$$

where $\beta_a = v_a/c$ and $\beta = v/c$ are dimensionless velocities. In Figure 1.5, we plot the dependence of $\beta_a$ with the viewing angle $\psi$ for different relativistic motion of the blob. It is evident from the plot that the apparent speed can exceed $c$ for the blobs with relativistic velocities and moving closer to the line of sight.
Chapter 1. Introduction

Gamma-ray transparency

High energy $\gamma$-rays can interact with low energy photons to produce electron-positron pairs. The cross-section for this process is maximum when

$$\epsilon_\gamma \epsilon_x \sim 2 (m_e c^2)^2$$

where $\epsilon_\gamma$ and $\epsilon_x$ are the energy of the high energy and the low energy photons respectively and $m_e$ is the electron mass. For example, a 100 MeV photon can pair produce with 5 keV X-ray photon. In general, the transparency of a medium to a radiation is described by its optical depth. Optimal depth of a medium corresponding to a radiation of frequency $\nu$ is defined as

$$\tau_\nu(s) = \int_{s_0}^{s} \alpha_\nu(s') \, ds'$$

where $s_0 \rightarrow s$ is the path of light travel and $\alpha_\nu$ is the absorption coefficient. The medium is said to be optically thick (or opaque) when $\tau > 1$ and optically thin (or transparent) when $\tau < 1$. The optical depth $\tau$ to pair production for a $\gamma$-ray of energy $\epsilon_\gamma$ in a homogeneous
region of size $R$ can be approximated as

$$\tau(\epsilon_\gamma) \simeq \frac{0.2 \sigma_T L(2\epsilon_x)}{4\pi Rc} \quad (1.6)$$

where $L(2\epsilon_x)$ is the luminosity of the target photon at energy $2\epsilon_x$ and $\sigma_T$ is the Thompson cross-section. If the power-law spectra of AGN and quasars in the X-ray region extend to much higher energies, then the source may be opaque to $\gamma$-rays if $\tau(\epsilon_\gamma) > 1$. This situation can be conveniently stated in terms of compactness parameter ($l$) defined as

$$l = \frac{L}{R} \frac{\sigma_T}{m_e c^3} \quad (1.7)$$

where $L$ is the total luminosity produced in the region of size $R$. The opacity condition, $\tau > 1$, can now be translated in terms of compactness parameter as $l > 60$. Estimation of compactness parameter requires the knowledge of X-ray luminosity and the size of the emission region. The latter can be estimated from the observed time variability of the flux as

$$R \sim ct_{\text{var}} \quad (1.8)$$

where $t_{\text{var}}$ is the variability time-scale. The inferred values of $l$ for blazars from the observed X-ray luminosities are much larger than 60. For 3C 279, $l \sim 5000$ and for PKS 0528+134, $l \sim 15000$. However, EGRET mission detected many blazars at energies greater than 100 MeV. In order to observe the $\gamma$-rays of these energies, the actual luminosity $L$ of the source in its proper frame must be much smaller than the observed value and the actual size should be larger than the inferred value in order to fulfill the condition $l < 60$. Indeed, this can happen if the source is moving at relativistic speed towards the observer. The relativistic beaming effects will enhance the luminosity as

$$L_{\text{obs}} = \delta^4 L_{\text{int}} \quad (1.9)$$

and the size of the emission region will be reduced as

$$R \sim \frac{ct_{\text{var}}}{\delta} \quad (1.10)$$
where $L_{\text{obs}}$ and $L_{\text{int}}$ are the observed and the intrinsic luminosity and the relativistic Doppler factor $\delta$ is given by

$$
\delta = \frac{1}{\Gamma(1 - \beta \cos \psi)}
$$

(1.11)

Here, $\Gamma = (1 - \beta^2)^{-1/2}$ is the bulk Lorentz factor, $\beta = v/c$ is the dimensionless velocity and $\psi$ is the angle between the line of sight of the observer and the direction of motion of the source. The compactness parameter $l$ in such a case will reduce to

$$
l = \delta^{-5} \frac{L_{\text{obs}}}{\rho_{\text{var}} \sigma_T m_e c^4}
$$

(1.12)

Dondi and Ghisellini [37] estimated the minimum value of the Doppler factor, corresponding to the optical depth $\tau = 1$, for the $\gamma$-ray bright blazars detected by EGRET. Their inferred minimum values of $\delta$ were spread within 1.3 to 11.3.

1.2.2 AGN Jet Features

Lobes

Lobes are the regions where the jets terminate and release their energy and momentum into the ambient intergalactic medium. The radio spectrum of the lobes are steep with index $\alpha > 0.5$ and the emission is usually polarised. This suggests the lobes to be an optically thin synchrotron sources driven by efficient cooling of relativistic particles. They contain enhanced emission regions known as hot spots which are often collinear with the central core (see Figure 1.1). In low resolution radio maps of the lobes, the presence of hot spots will make it appear as an edge-brightened source. The radio spectrum of the hot spots are flatter than the lobe spectrum with the index in the range $\alpha \sim 0.5 - 1$. This suggests the hot spots as the location where the jet hit the ambient medium and the bulk kinetic energy of the beam is converted into the random energy through the shock formed at the collision (see §2.1.2). The shock accelerated energetic particles diffuse from hotspot to the lobes thereby providing a continuous supply of energy.
Knots

These are the bright regions observed along the length of the AGN jet (Figure 1.6). The emission from these regions are strongly polarised and the polarisation angle is often perpendicular to the jet axis. Knots are seen in radio, optical and X-ray images of AGN jets and many show superluminal motion. The non-thermal (power-law) nature of the knot spectrum along with the observed polarisation suggests that the low-energy (radio-to-optical) emission must be synchrotron radiation emitted by relativistic distribution of charged particles cooling in a magnetic field. However, the X-ray emission from the knots can be an extension of synchrotron spectrum itself or the inverse Compton radiation emitted via scattering of soft target photons by relativistic electrons (see chapter 3). The target photons for the inverse Compton scattering can be either synchrotron photons itself or the photons external to the jet. The enhanced brightness of the knots is often interpreted as a result of efficient acceleration of charged particles probably by a shock present in the jet.

Jet Asymmetry

One of the intriguing characteristics of the AGN jets is that they are often observed to be one sided (Figure 1.7). The straightforward way to interpret this asymmetry is to relate it to relativistic beaming, since there are enough evidences that the jets are relativistic. If we assume that the central engine produces two similar jets ejected in opposite directions, then the brightness of the jet moving towards the observer will be enhanced due to relativistic beaming. Whereas the counter-jet will be dimmed (or invisible) since it moves away from the observer. If \( F_{\text{adv}} \) is the flux of the advancing jet and \( F_{\text{rec}} \) is the flux of the receding one, then the jet/counter-jet flux ratio \( J \) can be written as

\[
J = \frac{F_{\text{adv}}}{F_{\text{rec}}} = \left( \frac{1 + \beta \cos \psi}{1 - \beta \cos \psi} \right)^p
\]

where \( \beta \) is the bulk velocity of the jet in units of \( c \) and \( \psi \) is the angle between the jet direction and the line of sight of the observer. If \( \alpha \) is the power-law spectral index of the intrinsic jet flux, then the index \( p \) can be either \( 2 + \alpha \) in case of a continuous jet flow or \( 3 + \alpha \) for a moving isotropic source. Using equation (1.3), \( J \) can be expressed in terms of superluminal velocity \( (\beta_a) \) as

\[
J = (\beta_a^2 + \delta^2)^p
\]
Fig. 1.6: Image of 3C 273 with knots. Radio image by MERLIN (left), optical image by Hubble Space Telescope (middle) and X-ray image by Chandra with optical contour overlaid. Figure reproduced from Marshall et al. [38]
Here $\delta$ is the relativistic Doppler factor given by equation (1.11). The fact that all quasars are detected with one sided jet suggests that their jets are directed close to the line of sight of the observer. For example, the predicted value of $J$ for the quasar 3C 273 is $9 \times 10^6$ [32].

The interpretation of the one sided jet as an outcome of relativistic beaming effects can be tested by measuring the difference in polarisation between the lobe in jet and counter-jet side. The lobe situated at the counter-jet side should be less polarised than the lobe at the jet side since the emission from the farther lobe face more depolarising medium along the line of sight. Indeed, such a difference was observed by Garrington and Conway [39] for 49 sources from a sample of 69, supporting the above interpretation.

**Bent Jets**

Sky maps of most AGN jets show curvature and in particular this feature is prominent in the radio trails or head-tail

$^{19}$sources. In high resolution radio maps of these sources, one can see jets emerging from the nucleus and bent through a large angle [40]. For example, the sky map of the radio loud quasar PKS 2136+141 show a jet which is bent by an angle $\sim 210^\circ$ and this is the largest bending angle seen till today [41] (Figure 1.8). The bending of the radio jet can be interpreted in many ways. A distortion in the jet shape can occur as a result of ram pressure associated with the motion of the host galaxy through a dense intracluster medium. Bent jets with reflection symmetry can be explained if the host galaxy is associated

$^{19}$These are sources with a radio morphology consisting of a bright nucleus and a single faint jet streaming on one side.
with a companion. The jet curvature, in such a case is the result of the acceleration of the parent galaxy introduced by its companion [42]. Jets of the radio source can also appear as an inversion symmetry if the jet precesses. In this case though the jet matter follows a linear path, it will appear curved in the sky plane due to precession. Alternatively, the jet can also be deflected in spite of the kinetic motion of the host galaxy or by geometrical effects as explained above. As the jet moves through the surrounding medium, it gets disrupted when the medium is an extremely dense intracluster gas. Pressure stratification of the intracluster gas can also cause the jet to curve [43] [44].

**Limb-brightened Jets**

Another feature seen in some AGN jets is *limb-brightening* at pc scales. High resolution radio maps of these jets show the edges are brighter than the central spine of the jet (Figure 1.9). Such features are commonly observed in few FR I sources and blazars [45]. This feature is usually explained by the “spine-sheath” model where the velocity at the jet spine is larger compared to the velocity at the boundary. Such a radial stratification of velocity across the jet arises when jet moves through the ambient medium and the viscosity involved causes a shear at the boundary. The difference in the flow velocity of the jet between the
spine and the boundary can cause differential Doppler boosting and this in turn can produce a limb-brightened jet. Alternatively, velocity stratification at the jet boundary can accelerate the particles via shear acceleration \[46, 47\] (see chapter 4). These particles can then emit radiation via synchrotron process giving rise to a limb-brightened feature. Particles at the jet boundary can also be accelerated by turbulent waves initiated by the instabilities \[48\].
Chapter 2

Particle Acceleration and Radiative Processes

The observed very high energy $\gamma$-ray emission (up to GeV-TeV energies) from blazars suggests the presence of extremely relativistic particles in AGN jets. Also the power-law spectra observed over a broadband starting from radio-to-$\gamma$-ray indicates the jet emission to be dominated by non-thermal processes. The promising mechanism by which particles can be accelerated to very high energies in AGN jets is “Fermi mechanism” suggested by Enrico Fermi [50] to explain the power-law nature of the cosmic ray spectrum.

2.1 Fermi Acceleration

In 1949, Fermi [50] proposed that particles can be accelerated to high energies when they are scattered by magnetic irregularities (magnetic mirrors) associated with moving clouds in the interstellar medium. Consider a cloud moving with a velocity $V$ along the x-axis in the observer’s frame (Figure 2.1). Let us assume the cloud to be massive in comparison with the scattered particle and hence the centre of mass frame (CM-frame) is the frame of the cloud itself. The energy and the momentum of the particle before scattering in the CM-frame can be obtained using Lorentz transformation as [51]

$$E' = \Gamma_c (E - V p_x)$$

$$p'_x = \Gamma_c \left( p_x - \frac{V E}{c^2} \right)$$

(2.1)

(2.2)
Fig. 2.1: Illustration of Fermi acceleration mechanism: Collision between a particle with velocity $v$ and a massive cloud moving with velocity $V$. (a) Head-on collision; (b) Follow-on collision.
where \( E \) and \( E' \) are the energy of the particle in the observer’s frame and CM-frame respectively and \( p_x \) and \( p'_x \) are the corresponding momenta along x-axis (we represent the quantities in CM-frame with prime). The Lorentz factor of the moving cloud \( \Gamma_c \) is given by

\[
\Gamma_c = \left( 1 - \frac{V^2}{c^2} \right)^{-1/2}
\]

(2.3)

Since the CM-frame is the frame of scatterer itself, the energy and the x-component of momentum after collision will be

\[
E'_s = E'
\]

(2.4)

\[
p'_{x,s} = -p'_x
\]

(2.5)

Transforming the scattered energy back into the observer frame we get

\[
E_s = \Gamma_c \left( E' - Vp'_x \right)
\]

\[
= \Gamma_c^2 \left[ E \left( 1 + \frac{V^2}{c^2} \right) - 2 Vp_x \right]
\]

(2.6)

where we have used equations (2.1) and (2.2). If we express the momentum along x-axis in terms of energy

\[
p_x = \frac{E v_x}{c^2}
\]

(2.7)

The scattered energy can then be written as

\[
E_s = \Gamma_c^2 E \left( 1 + \frac{V^2}{c^2} - \frac{2 Vv_x}{c^2} \right)
\]

(2.8)

Here \( v_x \) is the velocity of the particle along x-axis. For \( V \ll c \), the change in the particle energy due to collision can be obtained using equations (2.8) as

\[
\Delta E = E_s - E
\]

\[
\approx E \left[ 2 \left( \frac{V}{c} \right)^2 - \frac{2 Vv_x}{c^2} \right]
\]

(2.9)

where we have retained the terms only up to second order in \( V/c \). If the particle arrives at an angle \( \theta \) with respect to the velocity of the scatterer as shown in Figure 2.1 then \( v_x = -v \cos \theta \)
and we get

$$\frac{\Delta E}{E} = 2 \frac{V}{c^2} \left[ v \cos \theta + V \right]$$

(2.10)

From equation (2.10) it is evident that the particle gains energy in head-on collisions. On the other hand, for follow-on collisions $\cos \theta$ is negative, and the particle loses energy. Since in the scatterer’s frame the velocity of the approaching particle is more than the receding one due to addition of velocities, the probability for head-on collision is more than that of follow-on collision. Hence there will be a net gain in particle energy.

### 2.1.1 Stochastic Acceleration - Second Order Fermi Mechanism

Let us consider a situation where the scattering centers move randomly, then the net gain is obtained by averaging equation (2.10) over the angle $\theta$. The acceleration process is then referred to stochastic acceleration. An example of stochastic acceleration process is turbulent acceleration, where particles are energized via scattering by moving magnetic inhomogeneities associated with turbulence in a flowing fluid. Since the collision probability between the particle and the scatterer is anisotropic, we need to know the probability of a collision happening at an angle $\theta$ to perform the averaging. If we assume $v \approx c$, then the rate of collision in the scatterer’s frame will be greater by a factor $\Gamma_c [1 + (V/c) \cos \theta]$ [52]. Hence the probability of collision at angle $\theta$ will be proportional to this factor and the average will be

$$\left\langle \frac{2V \cos \theta}{c} \right\rangle \approx \left( \frac{2V}{c} \right)^2 \frac{\int_{-1}^{+1} [1 + (V/c) \mu] d\mu}{\int_{-1}^{+1} [1 + (V/c) \mu] d\mu}$$

$$= \frac{2}{3} \left( \frac{V}{c} \right)^2$$

(2.11)

where $\mu = \cos \theta$. From equation (2.10), the average energy gain will then be

$$\left\langle \frac{\Delta E}{E} \right\rangle \approx \frac{8}{3} \left( \frac{V}{c} \right)^2$$

(2.12)

Since the average increase in energy is of the order $(V/c)^2$, this process is called as second order Fermi acceleration mechanism.
Chapter 2. Particle Acceleration and Radiative Processes

2.1.2 Shock Acceleration - First Order Fermi Mechanism

If we consider only head-on collisions in Fermi acceleration, then the dominant term in equation (2.10) will be of the order of \((V/c)\) (for \(v \approx c\)) and the acceleration process is called as first order Fermi mechanism. Acceleration of charged particles at a shock front is an example of this mechanism. A shock is a discontinuity initiated by perturbation in a supersonic fluid flow. Fluid on either side of the shock will be in different state of equilibrium and are connected by the conservation equations. The strength of a shock is determined by its Mach number \(M\) defined as

\[
M = \frac{v_1}{c_s}
\]

where \(v_1\) is the velocity of the upstream fluid with respect to the shock front and \(c_s\) is the local speed of sound in the upstream region. Here, we denote upstream as the fluid ahead of the shock front and downstream as the one behind. Strong shocks are the one with \(M \gg 1\). Particles in a magneto hydrodynamic fluid (plasma) are scattered by magnetic inhomogeneities associated with turbulence in the flowing fluid. In presence of a shock in the fluid they get energised by crossing the shock front from upstream to downstream or vice versa.

Let us consider the case of non-relativistic strong shock where the shock velocity \(U \gg c_s\) and \(U \ll c\). In the shock frame the plasma will pass through it with an upstream velocity \(v_1(=U)\) and the downstream velocity \(v_2\). From the mass conservation we get

\[
\rho_1 v_1 = \rho_2 v_2
\]

where \(\rho_1\) and \(\rho_2\) are the mass density of the upstream and downstream plasma. For a strong shock in the limit of \(M \to \infty\) we can write

\[
\frac{\rho_1}{\rho_2} = \frac{\gamma_s + 1}{\gamma_s - 1}
\]

where \(\gamma_s\) is the ratio of the specific heats at constant pressure and volume. For a fully ionised gas we have \(\gamma_s = \frac{5}{3}\) and hence \(\frac{\rho_2}{\rho_1} = 4\) and \(v_2 = \frac{1}{4}U\). Hence the plasma on either side of the shock (upstream or downstream) will see the plasma approaching from the other side of the shock (downstream or upstream) with velocity \(\frac{3}{4}U\). In the proper frame of the plasma (upstream or downstream), the particle distribution is isotropic due to scattering. Let us consider a particle of energy \(E\) with \(x\)-component momentum \(p_x\) in the upstream plasma. We have chosen a coordinate system where \(x\) axis normal to the shock front. The energy of
Chapter 2. Particle Acceleration and Radiative Processes

the particle $E'$ in the frame of downstream plasma will be

$$E' = E + p_x V$$

(2.16)

where $V = \frac{3}{4}U$ is the velocity of the downstream plasma with respect to the upstream plasma. Since the particles are relativistic, we can write $p \approx E/c$ and $p_x \approx \frac{E}{c} \cos \theta$. Here $\theta$ is the angle between the particle momentum and the shock normal. Hence the particle enters the downstream with an energy increment

$$\Delta E = E \left( \frac{V}{c} \right) \cos \theta$$

(2.17)

The probability of the particle crossing the shock front and entering into the downstream within the angle interval $\theta$ to $\theta + d\theta$ is proportional to $\sin \theta \cos \theta d\theta$. Hence the average increase in energy in crossing the shock once is

$$\langle \frac{\Delta E}{E} \rangle = \frac{2}{3} \left( \frac{V}{c} \right)$$

(2.18)

A similar situation happens for the particle crossing the shock front from downstream to upstream plasma. Hence the average fractional energy gain in making one round trip is

$$\langle \frac{\Delta E}{E} \rangle = \frac{4}{3} \left( \frac{V}{c} \right)$$

(2.19)

Thus a particle in the vicinity of a shock is scattered by magnetic inhomogeneities and gets accelerated to higher energy by crossing the shock front multiple number of times. Also, since the average energy gain is of the order of $(V/c)$, shock acceleration is efficient than the stochastic acceleration.

2.1.3 Acceleration at Shear Layers

A velocity shear can arise in a plasma flow if it passes through a viscous medium. Particles can then get accelerated when they are scattered between different velocity layers by magnetic irregularities [46](Figure 2.2). Let us consider the case of a gradual shear, where the mean free path of the particle is much smaller than the transverse width of the sheared velocity layer. Suppose the flow is non-relativistic and we choose a reference frame in which the local fluid is at rest and the flow velocity is along z-axis, $U = U_z(x)\hat{e}_z$. If $\tau$ is the
mean scattering time, then the distance travelled by the particle along x-axis before getting scattered will be

\[ \delta x = \frac{p_1}{m} \cos \theta \tau \]  

(2.20)

where \( p_1 \) is the momentum of the particle, \( m \) is its mass and \( \theta \) is the angle between the particle momentum and x-axis. The change in the fluid velocity due to this displacement will be

\[ \delta u = \left( \frac{\partial U_z}{\partial x} \right) \delta x \hat{e}_z \]  

(2.21)

where \( \left( \frac{\partial U_z}{\partial x} \right) \) is the shear velocity gradient. The momentum of the particle \( p_2 \) relative to the local fluid will then be [54]

\[ p_2^2 = p_1^2 \left( 1 + 2 \frac{m}{p_1} \sin \theta \cos \phi + \frac{m^2 \delta u^2}{p_1^2} \right) \]  

(2.22)
where $\phi$ is the angle between y-axis and the projection of the particle momentum in yz-plane. Since $E = \frac{p^2}{2m}$, the average fractional energy gain due to shear acceleration will be

$$\left\langle \frac{\Delta E}{E_i} \right\rangle \propto \left( \frac{\partial U_z}{\partial x} \right)^2 \tau^2$$  \hspace{1cm} (2.23)

The first order term in equation (2.22) will get cancelled because for every forward scattering of particle between two velocity layers there will be an equivalent reverse scattering with negative $\delta u$.

### 2.1.4 Accelerated particle distribution

The Fermi acceleration process leads to a power-law distribution of particles [55]. Consider after each encounter the particle energy increases by a factor $\xi$, i.e. $\Delta E = \xi E$. Let us assume a particle with initial energy $E_o$ enters the acceleration region. After $n$ encounters the particle energy will be

$$E = E_o (1 + \xi)^n$$  \hspace{1cm} (2.24)

Also, let the probability of escape from the acceleration region after an encounter be $P_{esc}$. Hence, after $n$ encounters the probability of the particle to remain in the acceleration region will be $(1 - P_{esc})^n$. Thus the number of particles in the acceleration region with energies greater than $E$ will be

$$N(\geq E) \propto \sum_{j=n}^{\infty} (1 - P_{esc})^j$$

$$= \frac{(1 - P_{esc})^n}{P_{esc}}$$  \hspace{1cm} (2.25)

From equation (2.24), $n$ can be written as

$$n = \frac{\ln(E/E_o)}{\ln(1 + \xi)}$$  \hspace{1cm} (2.26)

Substituting equation (2.26) in equation (2.25), we get

$$N(\geq E) \propto \frac{1}{P_{esc}} \left( \frac{E}{E_o} \right)^{-p}$$  \hspace{1cm} (2.27)
where the index $p$ is given by

$$p = \frac{\ln\left(\frac{1}{1 - \nu_{\text{esc}}}\right)}{\ln(1 + \xi)}$$  \hspace{1cm} (2.28)

### 2.2 Radiation Emission Mechanisms

#### 2.2.1 Blackbody Radiation

The radiation emitted by a distribution of charged particle which is in thermal equilibrium with itself as well as the emitted radiation is called as *blackbody radiation*. The specific energy density of a blackbody spectrum is given by Planck spectrum

$$u_{\nu}(\Omega) = \frac{2h\nu^3}{c^3} \exp\left(\frac{h\nu}{k_B T}\right) - 1 \text{ ergs cm}^{-3} \text{ Hz}^{-1} \text{ sr}^{-1}$$  \hspace{1cm} (2.29)

where $\nu$ is the observed photon frequency, $\Omega$ the solid angle, $h$ the Planck constant, $k_B$ the Boltzmann constant and $T$ is the blackbody temperature. The peak photon frequency at which the energy density is maximum for a given temperature $T$ is given by Wien’s law, $\nu_{\text{peak}} = 2.82(k_B/h)T$. Integrating equation (2.29) over the entire photon frequency and the solid angle we get the *Stefan-Boltzmann* law for the energy density of a blackbody spectrum

$$u(T) = a T^4$$  \hspace{1cm} (2.30)

where $a = 4\sigma_{SB}/c$ and $\sigma_{SB}$ is the *Stefan-Boltzmann* constant.

#### 2.2.2 Bremsstrahlung Radiation

When a charged particle moves in a Coulomb field it gets accelerated and emits electromagnetic spectrum. This radiation is known as *Bremsstrahlung* radiation. Let us consider the motion of an electron in the Coulomb field of an ion. We will assume small angle scattering where the deviation of the electron path from a straight line is negligible. The variation in the dipole moment of the electron-ion system gives rise to dipole radiation and it can be shown
that the total energy emitted per frequency by the electron is \[51\]

\[
\frac{dW(b)}{d\omega} = \begin{cases} 
\frac{8 Z^2 e^6}{3\pi c^2 m_e^2 v^2 b^2}, & b \ll v/\omega \\
0, & b \ll v/\omega 
\end{cases}
\tag{2.31}
\]

where \(\omega = 2\pi \nu\) is the angular frequency of the emitted photon, \(b\) is the impact parameter, \(Ze\) is the charge of ion, \(v\) the electron velocity and \(m_e\) and \(e\) are the electron mass and its charge respectively. The total power emitted per frequency per volume for a medium with ion density \(n_i\) and electron density \(n_e\) assuming a fixed electron speed \(v\) will be \[51\]

\[
\frac{dW}{d\omega dV dt} = \frac{16\pi e^6}{3\sqrt{3} c^3 m_e^2 v} n_e n_i Z^2 g_{ff}(v, \omega)
\tag{2.32}
\]

where \(g_{ff}\) is known as Gaunt factor which is a function of energy of the electron and the frequency of the emitted photon. For a thermal distribution, the particles follow a Maxwellian velocity distribution and we obtain the power per frequency per volume as \[51\]

\[
\frac{dW}{d\omega dV dt} = \frac{2^5 \pi^2 e^6}{3 m_e c^3} \left( \frac{2\pi}{3 k m_e} \right)^{1/2} T^{-1} Z^2 n_e n_i e^{-h\nu/kT} \bar{g}_{ff}
\tag{2.33}
\]

where \(\bar{g}_{ff}\) is the velocity averaged Gaunt factor. From equation (2.33) we find that the bremsstrahlung emission from a thermal particle distribution gives rise to a flat spectrum with an exponential cutoff at about \(h\nu \sim kT\) (Figure 2.3).

2.2.3 Synchrotron Radiation

The radiation emitted due to the helical motion of a relativistic charged particle along a magnetic field is known as synchrotron radiation \[51, 56\]. The radiative power emitted by a relativistic electron with velocity \(\beta(= v/c)\) moving at an angle \(\alpha\) (pitch angle) with respect to an uniform magnetic field \(B\) is given by \[51\]

\[
P_{\text{syn}} = \frac{2}{3} r_e^2 c \gamma^2 \beta^2 B^2 \sin^2 \alpha
\tag{2.34}
\]

where \(r_e = e^2/m_e c^2\) is the classical electron radius and \(\gamma\) is the Lorentz factor of the electron, \(\gamma = (1 - \beta^2)^{-1}\). For an isotropic distribution of mono energetic electrons we need to
average equation (2.34) over all angles and we get

$$P_{\text{syn}} = 4 \frac{\beta^2 \gamma^2 e}{3} \sigma_T U_B$$

(2.35)

where $\sigma_T$ is the Thomson cross section and $U_B = B^2 / 8\pi$ is the magnetic field energy density. Equation (2.35) is also equal to the average energy lost by an electron via synchrotron process.

The synchrotron spectrum emitted by an electron of energy $\gamma m_e c^2$ moving with pitch angle $\alpha$ can be written as [57]

$$P_{\text{syn}}(\gamma, \nu) = \sqrt{3} \frac{e^3 B \sin \alpha}{m_e c^2} F \left( \nu \frac{\nu}{\nu_c} \right)$$

(2.36)

where $\nu$ is the frequency of the emitted photon and

$$\nu_c = \frac{3 e B \gamma^2}{4\pi m_e c} \sin \alpha$$

(2.37)
Chapter 2. Particle Acceleration and Radiative Processes

Fig. 2.4: Synchrotron power function

is the critical frequency. The synchrotron power function $F(x)$ is defined as

$$F(x) = x \int_x^\infty K_{5/3}(\xi) \, d\xi$$

(2.38)

where $K_{5/3}$ is the modified Bessel function of order $5/3$. The shape of the spectrum is determined by $F(x)$ with a peak located at $\approx 0.29(\nu/\nu_c)$ (Figure 2.4). Alternatively one can write the single particle emission spectrum using equation (2.35) as

$$P_{syn}(\gamma, \nu) = \frac{4}{3} \beta^2 \gamma^2 c \sigma_T U_B \phi(\gamma)$$

(2.39)

where $\phi$ is a function of $\nu$ and $\gamma$ satisfying the relation

$$\int_0^\infty \phi(\gamma) \, d\nu = 1$$

(2.40)

Considering the shape of the spectrum (Figure 2.4), one can approximate the function $\phi$ as a
Dirac delta function (to the orders of unity)

\[ \phi_\nu(\gamma) \rightarrow \delta(\nu - \gamma^2 \nu_L) \quad (2.41) \]

where \( \nu_L = eB/2\pi m_e c \) is the Larmor frequency.

For an isotropic power-law electron distribution given by

\[ N(\gamma) = k \gamma^{-p} \quad \gamma_{\text{min}} < \gamma < \gamma_{\text{max}} \quad (2.42) \]

the radiation energy emitted per second per frequency for \( \gamma_{\text{max}} \gg \gamma_{\text{min}} \) can be shown as [57]

\[ \frac{dW}{d\nu dt} \approx \frac{4\pi k e^3 B^{(p+1)/2}}{m_e c^2} \left( \frac{3e}{4\pi m_e c} \right)^{(p-1)/2} a(p) \nu^{-(p-1)/2} \quad (2.43) \]

where \( a(p) \) is a function of particle spectral index. Hence the emitted synchrotron spectrum is a power-law with index \( (p - 1)/2 \).

Moreover, a charged particle in a magnetic field besides emitting synchrotron radiation, can absorb a photon and get energized. This absorption process is called as Synchrotron Self Absorption. Also a photon can induce a charged particle to emit in a direction and at a frequency of the photon itself (stimulated emission or negative absorption). The absorption coefficient for synchrotron self absorption process is given by [58]

\[ \kappa_\nu = -\frac{1}{8\pi m_e \nu^2} \int \frac{N(\gamma)}{\gamma(\gamma^2 - 1)^{1/2}} \frac{d}{d\gamma} \left[ \gamma(\gamma^2 - 1)^{1/2} P_{\text{syn}}(\gamma, \nu) \right] \quad (2.44) \]

The specific intensity \( (I_\nu) \) of the synchrotron radiation can then be found using the radiative transfer equation as [56]

\[ I_\nu = S_\nu (1 - e^{-\tau_\nu}) \quad (2.45) \]

where we have assumed a source with uniform properties and no background illumination. Here \( S_\nu (= j_\nu / \kappa_\nu) \) is the synchrotron source function and \( \tau_\nu \) is the optical depth defined over a distance \( s \) as

\[ \tau_\nu = \int_s \kappa_\nu \, ds' \quad (2.46) \]
Chapter 2. Particle Acceleration and Radiative Processes

The synchrotron emissivity $j_\nu$ for an isotropic emission is given by

$$j_\nu = \frac{1}{4\pi} \int_1^\infty P_{\text{syn}}(\gamma, \nu) N(\gamma) \, d\gamma$$ \hspace{1cm} (2.47)

For a power-law distribution of particle (equation (2.42)), the synchrotron self absorption coefficient will be

$$\kappa_\nu \propto B^{(p+2)/2} \nu^{-(p+4)/2}$$ \hspace{1cm} (2.48)

Hence it may be possible for a source to be optically thick ($\tau_\nu > 1$) at low frequencies but optically thin ($\tau_\nu < 1$) at high frequencies. From equations (2.39), (2.47) and (2.48) one can find that the source function $S_\nu$ for the optically thick region will be a power-law of the form

$$S_\nu \propto \nu^{5/2}$$ \hspace{1cm} (2.49)

The total synchrotron spectrum will then be a broken power-law with spectrum changing from $\nu^{5/2}$ at lower frequencies to $\nu^{-(p-1)/2}$ at high frequencies. The frequency at which the index changes is called as synchrotron self absorption frequency. For AGN jet emission this frequency is observed to be within a range of few gigahertz.

### 2.2.4 Inverse Compton Radiation

Scattering of low energy electrons by high energy (or hard) photons is called as Compton scattering and the reverse process where high energy electrons scatter off low energy (or soft) photons is called as inverse Compton scattering. The spectrum obtained due to the scattering of soft photons by relativistic electrons is called as inverse Compton spectrum. If the energy of the incident photon in electron’s rest frame is much smaller than the electron rest mass energy, then one can ignore the recoil of electron and the scattering process leave the photon energy unchanged in electron’s rest frame. In such case the scattering process is described by Thomson cross section with differential cross section given by

$$\frac{d\sigma_T}{d\Omega} = \frac{1}{2} r_e^2 \left(1 + \cos^2 \theta\right)$$ \hspace{1cm} (2.50)

Here $r_e$ is the classical electron radius and $\theta$ is the angle between the incident and the scattered photon directions. On the other hand if the recoil of the electron becomes considerable,
then the scattering cross section is described by Klein-Nishina cross section. The differential cross section in this case will include the quantum effects and is given by

$$\frac{d\sigma_T}{d\Omega} = \frac{r_e^2}{2} \frac{\epsilon_s^2}{\epsilon^2} \left( \frac{\epsilon}{\epsilon_s} + \frac{\epsilon_s}{\epsilon} - \sin^2 \theta \right)$$

(2.51)

where $\epsilon$ and $\epsilon_s$ are the energies of the incident and the scattered photon. These conditions can be expressed in terms of $\epsilon$ and the Lorentz factor of the relativistic electron $\gamma$ as, $\gamma \epsilon \ll m_e c^2$ for scattering in Thomson regime and $\gamma \epsilon \gg m_e c^2$ for scattering in Klein-Nishina regime [51, 57]. Also, in the Thomson limit the scattered photon energy ($\epsilon_s$) can be shown as [57]

$$\epsilon_s \approx \gamma^2 \epsilon$$

(2.52)

Hence for large $\gamma$ the photon energy gain is very large. Nevertheless this gain is quite small compared with the electron energy and the electron loses only a small fraction of energy in each scattering. However in Klein-Nishina regime the electron loses almost its entire energy to the photon in a single scattering. Hence the scattered photon energy in this case will be

$$\epsilon_s \approx \gamma m_e c^2$$

(2.53)

The radiative power emitted (or the power lost by an electron) due to inverse Compton scattering of an isotropic soft photon distribution can be shown as [51]

$$P_{com} = \frac{4}{3} \beta^2 \gamma^2 \epsilon \sigma_T U_{ph}$$

(2.54)

where $U_{ph}$ is the energy density of the soft target photon distribution and the scattering is assumed to be in Thomson regime. Comparison of equation (2.54) with the power emitted due to synchrotron emission (equation (2.35)) one finds

$$\frac{P_{syn}}{P_{com}} = \frac{U_B}{U_{ph}}$$

(2.55)

When the scattering happens in Klein-Nishina regime the power lost by an electron can be computed using

$$P_{com, kn} = \int \int (\epsilon_s - \epsilon) \frac{dN}{dlde_s} d\epsilon_s$$

(2.56)
where \( \frac{dN}{dtd\epsilon_s} \) is the emission rate of the scattered photon per frequency given by [57]

\[
\frac{dN}{dtd\epsilon_s} = \frac{2\pi r_e^2 c}{\gamma^2} \frac{n(\epsilon)\,d\epsilon}{\epsilon} \times \left[ 2q \ln q + (1 + 2q)(1 - q) + \frac{1}{2} \frac{(\Gamma_{\epsilon} q)^2}{(1 + \Gamma_{\epsilon} q)} (1 - q) \right]
\]  \((2.57)\)

Here \( n(\epsilon) \) is the number density of the soft target photons and the quantities \( \Gamma_{\epsilon} \) and \( q \) are defined as

\[
\Gamma_{\epsilon} = \frac{4\epsilon \gamma}{m_e c^2} \quad \text{and} \quad q = \frac{\epsilon_s}{\Gamma_{\epsilon} (\gamma m_e c^2 - \epsilon_s)}
\]  \((2.58)\) \((2.59)\)

For a power-law distribution of particle given by equation \((2.42)\) the emitted photon spectrum when the scattering happens in Thomson regime will be [57]

\[
\frac{dW}{dtd\epsilon_s} = \pi r_e^2 c k 2^{p+3} \frac{p^2 + 4p + 11}{(p + 3)^2 (p + 1)(p + 5)} \epsilon_s^{-(p-1)/2} \times \int e^{(p-1)/2} n(\epsilon) \, d\epsilon
\]  \((2.60)\)

From equations \((2.43)\) and \((2.60)\) we find that the spectrum emitted by both synchrotron and inverse Compton process can be represented by a power-law with same index \((p - 1)/2\). In case of extreme Klein-Nishina limit the emitted spectrum will be

\[
\frac{dW}{dtd\epsilon_s} = \pi r_e^2 c k (m_e c^2)^{p+1} \epsilon_s^{-p} \times \int \frac{d\epsilon}{\epsilon} n(\epsilon) \left( \ln \frac{\epsilon \epsilon_s}{m_e^2 c^4} + C(p) \right)
\]  \((2.61)\)

where \( C(p) \) is a parameter of order unity [57]. Thus we see the inverse Compton spectrum in extreme Klein-Nishina limit is much steeper than that of Thomson limit (equation \((2.60)\)).

**Synchrotron Self Compton**

In many astrophysical systems, the high energy emission is explained as a result of the inverse Compton scattering of synchrotron photons. Here, the same electron population which is responsible for the synchrotron emission will scatter off these photons to higher energies.
This process is commonly referred as synchrotron self-Compton (or SSC) mechanism. For example, the ratio of the X-ray to TeV $\gamma$-ray fluxes obtained during simultaneous observation of BL Lac objects, both in quiescent and flaring state, can be explained in the context of SSC process \[59\].

Let us assume that the radiation emitting plasma of AGN jet be confined in a spherical region with tangled magnetic field. This plasma moves down the jet at relativistic speed. The distribution of electrons in the emission region is assumed to be a power-law described by equation (2.42). Therefore, the optically thin synchrotron spectrum will be a power-law with index $\alpha = (p - 1)/2$ (equation (2.43)). The SSC flux at the photon energy $\epsilon_s$ can then be predicted from the observed synchrotron flux as \[60, 61\]

$$
F_{SSC}(\epsilon_s) \approx d(\alpha) \theta_d^{-2(2\alpha+3)} \nu_m^{-(3\alpha+5)} (F_m^{Syn})^{2(\alpha+2)} \epsilon_s^{-\alpha} \ln \left( \frac{\nu_{max}}{\nu_m} \right) \left( \frac{1 + z}{\delta} \right)^{2(\alpha+2)}
$$

where $\theta_d$ is the angular size of the source, $\nu_m$ is the synchrotron self absorption frequency, $F_m^{Syn}$ is the synchrotron flux at frequency $\nu_m$, $\nu_{max}$ is synchrotron high frequency cutoff corresponding to the high energy cutoff in the particle spectrum, $z$ is the redshift of the source and $\delta$ is the Doppler factor of the jet. Here $d(\alpha)$ is a function depending only on $\alpha$ and has values $d(0.25) = 130, d(0.50) = 43, d(0.75) = 18$ and $d(1.00) = 9.1$. If we assume the X-ray emission from AGN jets as due to SSC process, then by comparing the predicted SSC flux at X-ray energy with the observed flux at that energy one can estimate the Doppler factor of the jet \[52\].

**External Compton**

In external Compton mechanism, the photons which are produced outside the emission region are scattered off to high energies by inverse Compton process. A general case assuming the scattering of an isotropic distribution of soft target photons to high energies by a distribution of relativistic electrons can be studied using equations (2.60) and/or (2.61). However in case of AGN jet these equations are invalid. Here as the emission region moves down the jet at relativistic speed, the plasma see an anisotropic distribution of target photons due to Doppler boosting. Consider the case of a spherical emission region moving down the jet with bulk Lorentz factor $\Gamma$. Let the inclination angle of the jet to the line of sight of the observer be $\theta_o$. For simplicity let us assume the external target photon distribution be monochromatic and isotropic in the frame of the central source. The energy density of these
target photons in the emission region frame will then be \( \approx \Gamma^2 u_{iso}^* \), where \( u_{iso}^* \) is the energy density of the isotropic external radiation field. The target photon distribution in the frame of the emission region will peak at energy \( \epsilon' \approx \Gamma \epsilon^* \), where \( \epsilon^* \) is the energy of the external photon distribution. The resultant inverse Compton emissivity \( j_c(\epsilon, \Omega) \) at photon energy \( \epsilon \) emitted in a direction \( \Omega \), due to a power-law distribution of particles described by equation (2.42), is given by [62]

\[
j_c(\epsilon, \Omega) \approx \frac{c \sigma_T u_{iso}^* K}{8\pi \epsilon^*} \left[ \Gamma (1 + \mu) \right]^{1+\alpha} \left( \frac{\epsilon}{\epsilon^*} \right)^{-\alpha} \text{ergs cm}^{-3} \text{s}^{-1} \text{sr}^{-1}
\]

where \( \alpha = (p - 1)/2 \) and \( \mu = \cos \theta_z \), with \( \theta_z \) being the angle between the emitted photon and the jet axis. Since the scattering is assumed to happen in Thomson regime, equation (2.63) is valid only for the scattered photon energies satisfying the relation

\[
\gamma_{min}^2 \leq \epsilon / \Gamma \epsilon^* (1 + \mu) \leq \gamma_{max}^2
\]

The quantity \( \Gamma (1 + \mu) \) can be written in terms of the viewing angle \( \theta_o \) and the Doppler factor of the jet \( \delta \) as

\[
\Gamma (1 + \mu) = \delta \frac{1 + \cos \theta_o}{1 + v/c}
\]

where \( v \) is the velocity of the emission region along the jet.

### 2.3 Equipartition Magnetic Field

Consider a spherical source of volume \( V \) with a power-law distribution of electrons described by equation (2.42) cooling in a magnetic field \( B \). The total energy of the electrons will be

\[
U_e = V m_e c^2 \int_{\gamma_{min}}^{\gamma_{max}} \gamma N(\gamma) \, d\gamma
\]

\[
= V m_e k c^2 \frac{2 - p}{2 - p} \left( \gamma_{max}^{2-p} - \gamma_{min}^{2-p} \right)
\]
The total synchrotron luminosity of the source using equation (2.35) (assuming $\beta \approx 1$) will be

$$L = V \int_{\gamma_{\text{min}}}^{\gamma_{\text{max}}} P_{\text{syn}}(\gamma) N(\gamma) \, d\gamma$$

$$= 4 \sigma_T c U_B kV \frac{3}{3 - p} \left( \frac{\gamma_{\text{max}}^{3 - p} - \gamma_{\text{min}}^{3 - p}}{\gamma_{\text{max}}^{3 - p} - \gamma_{\text{min}}^{3 - p}} \right)$$

(2.67)

(2.68)

We can write $\gamma_{\text{min}}$ and $\gamma_{\text{max}}$ in terms of characteristic synchrotron photon frequency $\nu_{\text{min}}$ and $\nu_{\text{max}}$ using equation (2.41) as $\gamma_{\text{min}} = (\nu_{\text{min}}/\nu_K)^{1/2}$ and $\gamma_{\text{max}} = (\nu_{\text{max}}/\nu_K)^{1/2}$. Then the ratio of total energy of the electrons to the synchrotron luminosity will be

$$\frac{U_e}{L} = \frac{A}{B^{3/2}}$$

(2.69)

where $A$ is a constant that depends only on the particle spectral index $p$. If the source had other particles like hadrons along with electrons then the total particle energy will be $U_p = a U_e$, where $a > 1$. Then the total energy of the source will be

$$U_{\text{tot}} = U_p + U_B$$

$$= a AL \frac{V}{8\pi} + \frac{B^2}{2B^{3/2}}$$

(2.70)

(2.71)

The magnetic field $B_{\text{min}}$ for which the total energy of the system is minimum can be obtained by solving

$$\left( \frac{\partial U_{\text{tot}}}{\partial B} \right)_{B = B_{\text{min}}} = 0$$

(2.72)

and we get

$$B_{\text{min}} = \left( \frac{6\pi a AL}{V} \right)^{2/7}$$

(2.73)

On the other hand, the equipartition magnetic field $B_{eq}$ obtained from the relation $U_p = U_B$ will be

$$B_{eq} = \left( \frac{8\pi a AL}{V} \right)^{2/7}$$

(2.74)
Equations (2.73) and (2.74) differ by an factor less than 10 percent and the total energy corresponding to $B_{eq}$ and $B_{min}$ will be

$$U_{tot}(B_{eq}) = 2V \left( \frac{B_{eq}^2}{8\pi} \right) \quad \text{and} \quad (2.75)$$

$$U_{tot}(B_{min}) = \frac{7}{3}V \left( \frac{B_{min}^2}{8\pi} \right) \quad (2.76)$$

or

$$U_{tot}(B_{eq}) = \frac{6}{7} \left( \frac{B_{eq}}{B_{min}} \right)^2 U_{tot}(B_{min}) \approx 1.01 U_{tot}(B_{min}) \quad (2.77)$$

Hence it is customary to use equipartition value for the magnetic field while modelling the sources to ensure a minimum energy condition [19, 34].

### 2.4 Hadronic Processes

If protons are accelerated to high energies they can lose their energy through synchrotron emission and hadronic interactions. The main hadronic interactions by which an energetic proton can lose its energy are the following:

- Bethe-Heitler process:

  $$p + h\nu \rightarrow p + e^+ + e^- \quad (2.78)$$

  where $h\nu$ is a photon. The photon threshold energy in the rest frame of proton for the Bethe-Heitler process is the sum of rest mass energies of the electron and positron (i.e. 1.022 MeV). The cross section for this process in case of an ultra relativistic proton ($\beta \approx 1$) can be expressed in terms of the photon momentum $k'$ when $2 \leq k' \leq 4$ as [63]

  $$\sigma_{BH}(k') \approx \frac{2\pi}{3} \alpha_{fs} r_e^2 Z^2 \left( \frac{k' - 2}{k'} \right)^3 \left( 1 + \frac{1}{2} \eta + \frac{23}{40} \eta^2 + \frac{37}{120} \eta^3 + \frac{61}{192} \eta^4 \right) \quad (2.79)$$

  where $\alpha_{fs}$ is the fine structure constant, $r_e$ is the classical electron radius, $Z$ is the charge of the ion in units of $e$ and $\eta = (k' - 2)/(k' + 2)$. For $k' > 4$ the cross section...
can be approximated as [63]

\[
\sigma_{BH}(k') \approx \alpha_f s r_e^2 Z^2 \left\{ \frac{28}{9} \ln 2k' - \frac{218}{27} + \left( \frac{2}{k'} \right)^2 \left[ 6 \ln k' - \frac{7}{2} + \frac{2}{3} \ln^2 2k' - \frac{1}{3} \pi^2 \ln 2k' + 2\zeta(3) + \frac{\pi^2}{6} \right] \\
- \left( \frac{2}{k'} \right)^4 \left( \frac{3}{16} \ln 2k' + \frac{1}{8} \right) \\
- \left( \frac{2}{k'} \right)^6 \left( \frac{29}{9.256} \ln 2k' - \frac{77}{27.512} \right) \right\} \quad (2.80)
\]

• proton-proton collision:

\[ p + p \rightarrow X + \sum_{i=1}^{m} \pi_i \quad \text{where} \quad X \rightarrow \text{hadrons} \quad (2.81) \]

Here \( m \) is the multiplicity of secondary pions. The threshold energy for this reaction is \( E_{th} = 2m_\pi c^2(1 + m_\pi/4m_p) \approx 280 \text{ MeV} \), where \( m_\pi \) and \( m_p \) are the masses of the \( \pi^0 \)-meson and the proton. When the incident proton is in the GeV to TeV energy region, the total cross section can be approximated by [64]

\[
\sigma_{pp}(E_p) \approx 30 \left[ 0.95 + 0.06 \ln \left( \frac{E_{kin}}{1\text{GeV}} \right) \right] \text{ mbarn} \quad (2.82)
\]

where \( E_p \) is the initial proton energy and \( E_{kin} = E_p - m_p c^2 \), and \( E_{kin} \geq 1 \text{ GeV} \). Here it is assumed that \( \sigma_{pp} = 0 \) at lower energies.

• photo-meson process:

\[ p + h\nu \rightarrow X + \sum_{i=1}^{m} \pi_i \quad \text{where} \quad X \rightarrow \text{hadrons} \quad (2.83) \]

The cross section for the photo-meson process increases starting from the threshold energy of the photons \( E_{th} = 150 \text{ MeV} \) (in the rest frame of protons) reaching their maximum value \( \sim 3 \times 10^{-28} \text{ cm}^2 \) at \( E \sim 300 - 400 \text{ MeV} \), and then decrease [65].
Decay modes for the pions produced in these reactions are as follows

\[ \pi^0 \rightarrow 2\gamma \]  \hspace{1cm} (2.84)

\[ \pi^+(\pi^-) \rightarrow \mu^+(\mu^-) + \nu_\mu(\bar{\nu}_\mu) \]  \hspace{1cm} (2.85)

\[ \mu^+(\mu^-) \rightarrow e^+(e^-) + \nu_e(\bar{\nu}_e) + \bar{\nu}_\mu(\nu_\mu) \]  \hspace{1cm} (2.86)

The decay products, \( \gamma \)-rays and pairs, can then initiate an electromagnetic cascade by causing the production of further pairs and \( \gamma \)-rays.

### 2.5 Emission Models

The radio-to-UV/X-ray radiation from AGN jets are generally attributed to synchrotron emission due to cooling of relativistic non-thermal electrons in a magnetic field. However there are two different approaches concerning the high energy emission viz. leptonic model and hadronic model. High energy radiation ranging from MeV to TeV energies are generally observed from blazars. Hence these models are discussed in the context of these sources. In leptonic models, the high energy radiation will be dominated by the inverse Compton emission from the same ultra relativistic electrons producing synchrotron radiation \[66, 67, 68, 69\]. On the other hand, in hadronic models, the high energy radiation is mainly due to pair cascades initiated by the interaction of relativistic protons with photons and proton synchrotron radiation \[70, 71\].

#### 2.5.1 Leptonic Models

Leptonic models assume the electrons in the jets are accelerated to ultra relativistic velocities via Fermi acceleration process. Whereas protons are not sufficiently accelerated and their energies remain lower than the threshold energy required to initiate the hadronic interactions. Some models assume the hadrons are cold and mainly provide the inertia required for the jet to reach up to kpc/Mpc scales \[72, 73\]. The accelerated electrons beside emitting synchrotron radiation also scatter off soft target photons to hard X-ray and \( \gamma \)-ray energies by inverse Compton process. The possible target photons for this process are the synchrotron photons produced within the jet (SSC) (\[2.2.4\]) and/or the external photons entering into the jet. The sources of external photons which can play an important role in explaining the high energy radiation in AGN jet are
• the radiation from the accretion disk around the central massive object
• the reprocessed accretion disk radiation from the broad line emitting region
• the infra-red photons from the dusty torus.

Some parameters of the leptonic model can be constrained from the relativistic Doppler boosting required for the high energy radiation to be transparent against the pair production opacity with soft photons (§1.2.1). However this effect may be non-negligible at very high energies and the resultant spectrum will be hard. Also, the synchrotron radiation from the secondary electrons may become important [74]. Moreover, the detection of subluminal velocities ($\beta_{\text{app}} < 1$) in the sub-pc scale jets of few TeV blazars suggest that the relativistic jets of these sources decelerate. The varying Doppler factor due to this deceleration will have a significant impact on the observed properties of the blazars [75].

In simplistic approaches, the underlying electron distribution is either a single or broken power-law with index/indices inferred from the observed photon spectral index/indices. A reasonable estimate of the parameters can then be obtained from the observations. For example, the underlying magnetic field can be estimated considering equipartition between the electrons and the magnetic field energy densities (§2.3). As discussed earlier, the Doppler factor can be estimated from the observed superluminal motion of the knots or using $\gamma$-ray transparency (§1.2.1). Also, from the measured variability time-scale, the size of the emission region can be constrained. While these simplistic models have been successful in reproducing the spectrum of AGN jets, they lack a self-consistent basis for the shape of the electron distribution.

A more realistic approach consists of the solution of a kinetic equation involving acceleration of the particles and radiative as well as non-radiative cooling mechanisms [76]. One zone models assume the observed emission as a result of efficient cooling of non-thermal electrons from a region with tangled magnetic field. The evolution of the particle distribution $N(\gamma, t)$ in this region can be described in its simplest form by the kinetic equation as

$$\frac{\partial N(\gamma, t)}{\partial t} + \frac{\partial}{\partial \gamma} [P(\gamma, t)N(\gamma, t)] = Q(\gamma, t) \quad (2.87)$$

where $\gamma$ is the Lorentz factor of the electron, $P(\gamma, t)$ is the energy loss rate and $Q(\gamma, t)$ is the injection rate of non-thermal particles. The injection can be a single burst of non-thermal particles injected at time $t = 0$ (one-time injection models) or a continuous injection of non-thermal particles. If the losses are mainly due to synchrotron and inverse Compton processes,
then for a power-law distribution of particles similar to the one given in equation (2.42), the resultant photon spectrum will be a power-law with index \((p - 1)/2\). However since the energy loss rate is proportional to \(\gamma^2\) for these processes (§§2.2.3 and 2.2.4), the high energy particles cool more efficiently than the low energy ones. This leads to a depletion of high energy particles in case of one time injection models and gives rise to a time-dependent high energy cut off in the non-thermal particle distribution. The emitted photon spectrum will then exponentially decrease at high energies corresponding to this cut off energy in the particle spectrum. Alternatively one can infer the age of the emission region by translating this exponentially decreasing feature in the photon spectrum to the high energy cut off in the particle distribution. On the other hand, in the continuous injection models, the depleted high energy electrons are continuously replenished. This gives rise to a broken power-law particle distribution with a break at energy for which the cooling time-scale is equal to the age of the emission region. The resultant photon spectrum will then be a broken power-law instead of one with an exponential cut off.

Two zone models are more involved than one zone models where acceleration of particles are also considered along with the cooling mechanisms. According to these models particles are accelerated to relativistic energies in an acceleration region. These high energy particles are then injected into a cooling region where they lose most of their energies by radiative and non-radiative processes. The evolution of the particles are governed by the kinetic equations corresponding to acceleration region and cooling region. These equations can be written in their simplest form as

\[
\frac{\partial n(\gamma, t)}{\partial t} + \frac{\partial}{\partial \gamma} [(P_{AR}(\gamma, t) + \dot{\gamma}_{\text{acc}}) n(\gamma, t)] + \frac{n(\gamma, t)}{t_{\text{esc}}} = Q(\gamma, t) \quad (2.88)
\]

\[
\frac{\partial N(\gamma, t)}{\partial t} + \frac{\partial}{\partial \gamma} [P_{CR}(\gamma, t) N(\gamma, t)] = \frac{n(\gamma, t)}{t_{\text{esc}}} \quad (2.89)
\]

where the equation (2.88) governs the evolution in the acceleration region and the equation (2.89) in the cooling region. Here \(n(\gamma, t)\) and \(N(\gamma, t)\) are the particle distribution in the acceleration region and the cooling region, \(P_{AR}(\gamma, t)\) and \(P_{CR}(\gamma, t)\) are the respective energy loss rates, \(\dot{\gamma}_{\text{acc}}\) is the particle acceleration rate, \(t_{\text{esc}}\) is the particle escape time-scale in acceleration region and \(Q(\gamma, t)\) is the particle injection rate. The injection into the acceleration region can be mono energetic electrons or a residual particle distribution of an earlier acceleration process. If the acceleration happens at a shock front then the acceleration rate \(\dot{\gamma}_{\text{acc}}\)
can be approximated as

$$\dot{\gamma}_{\text{acc}} \approx \frac{\gamma}{t_{\text{acc}}} \tag{2.90}$$

where $t_{\text{acc}}$ is the acceleration time-scale and can be estimated from the theory of diffusive shock acceleration [77]. The index of the particle spectrum in acceleration region is governed by the acceleration and escape time-scales. Whereas the maximum energy to which the particles can be accelerated is determined by the acceleration and cooling time-scales. One zone and two zone models involve many parameters even in their simplistic form and are often cannot be constrained. However these models can throw light on the underlying physics of the source and its predictions compared with future observations will help us to understand these sources better.

Sambruna et al. [78] used a one-time injection model to explain the broadband emission from the knots of several AGN detected in radio, optical and X-ray. The spectra modelled by them predicts an exponential cut off at optical/UV or X-ray energies. Whereas Liu & Shen [79] proposed a two zone model to explain the X-ray emission from the knots of M87 (a nearby FRI radio galaxy). One zone models failed to reproduce the observed flux and/or the spectral index from the knots of this source.

In case of blazars the picture is different. One-zone models are used by various authors to explain the emission from blazars [80, 81]. These models assume a homogeneous distribution of particles and magnetic field throughout the emission region. The broadband spectra of the blazars are successfully explained by this model. However the time lags observed between the flares at different frequencies cannot be explained under this model. Also this model requires the light travel time to be shorter than the synchrotron cooling time-scales in order to satisfy the homogeneity of the particle distribution. This introduces a constraint on the size of the emission region. Two zone models offers more insight into these sources [82, 83]. Kirk et al. [82] solved the kinetic equation considering the temporal as well as the spatial variation of the particle distribution in the cooling region. They showed that the time lags between the flares at different frequencies can be an outcome of the difference between cooling and acceleration time-scales.
2.5.2 Hadronic Models

In hadronic models, the protons are accelerated along with electrons to ultra-relativistic energies and cool off mainly through proton synchrotron and photo-meson interactions (synchrotron-proton blazar model) [71]. The dominant channels for the photo-meson process are [84]

\[ p + \gamma \rightarrow \pi^0 + p \]
\[ p + \gamma \rightarrow \pi^+ + n \]
\[ p + \gamma \rightarrow \pi^+ + \pi^- + p \]

The pions decay as shown in equations (2.84), (2.85) and (2.86) and decay products initiate an electromagnetic cascade. Hadronic models explain the high energy emission as a result of these cascades and proton synchrotron emission. The type of the resulting cascade spectrum depends upon the compactness parameter \( l \) which measures the optical depth with respect to pair creation (equation (1.7)). For emission region with small compactness parameter \( l \ll 60 \) the electromagnetic cascade terminates after few generations and hence the photon luminosity is concentrated at high energies. On the other hand, as \( l \) increases, more and more generations shift power towards lower energies.

The target photons for photo-meson process can be the electron synchrotron radiation and/or the external photons. Mücke & Protheroe [85] using Monte Carlo technique simulated the proton interactions and the subsequent cascades. They considered the co-acceleration of protons along with electrons while the synchrotron emission from the latter is responsible for the low energy emission from blazars. These photons serve as target photons for the \( p\gamma \) interactions. They showed that the cascades initiated by the \( \pi^0 \) decay and \( \pi^\pm \) decay generate a featureless \( \gamma \)-ray spectra. In contrast, the proton synchrotron cascades and \( \mu^\pm \) synchrotron cascades produce a two-component \( \gamma \)-ray spectrum commonly observed in flaring blazars. In general, direct proton and \( \mu^\pm \) synchrotron radiation is mainly responsible for the high energy bump in blazars, whereas the low energy bump is dominated by synchrotron radiation from the primary electrons, with a contribution from the secondary electrons [86].

Unlike leptonic models, hadronic blazar models result in neutrino emission through the production and decay of charged mesons (equations (2.85) and (2.86)). Another important source of high energy neutrinos is the production and decay of charged kaons. In the case of \( p\gamma \) interactions positively charged kaons are produced [87, 88]. They decay into muons and direct high energy muon-neutrinos. These muon-neutrinos will not have suffered energy
losses through $\pi^\pm$ and $\mu^\pm$ synchrotron radiation unlike the ones originating from $\pi^\pm$ and $\mu^\pm$ decay. Therefore they appear as an excess in comparison to the remaining neutrino flavors at the high energy end of the emerging neutrino spectrum. Detection of predicted neutrino spectrum play an important role in validating the hadronic models.

Investigation of time-dependent hadronic models is very difficult because of time consuming Monte-Carlo cascade simulations. Also it is difficult to reconcile their rapid variability observed in blazars ($< 1$ hour) with the radiative cooling time-scales of protons \cite{89}.

Fig. 2.5: Spectral fits to the spectral energy distribution of 3C279 using a leptonic external-Compton model (solid (red)); leptonic SSC model (short-dashed (red)); hadronic model with electron synchrotron photons as targets for the photo-meson process (dot-dashed (maroon)) and hadronic model with electron synchrotron + external photons as targets for photo-meson process (long-dashed (maroon)). Figure reproduced from Böttcher et al. \cite{27}

Böttcher et al.\cite{27} studied the simultaneous multi wavelength observation of the very high energy (VHE) blazar 3C279 using the leptonic and hadronic models. Leptonic one zone model requires unrealistic parameters to explain the observed spectrum. Whereas, the hadronic synchrotron-proton blazar model is able to fit the broadband spectrum successfully. They also considered the contribution from the external target photons for hadronic interactions
in order to reduce the energy loss time-scale. In Figure 2.5 we show their fit to the spectral energy distribution of 3C279 based on leptonic and hadronic models.

2.6 Aim of the Thesis

Despite the availability of enormous amount of information about AGN by virtue of high resolution and high sensitivity experiments at present, there exist a large amount of uncertainties regarding the physics of various observed features [90, 91, 92, 93]. In the work presented in this thesis we shall attempt to understand certain features of AGN jets in view of the recent observations.

- In Chapter 3, we describe our works to interpret the underlying physics of the AGN knots. The on-board X-ray satellite Chandra\(^1\) studied the knots of several AGN for which radio and optical informations are already available [94, 95, 78, 96]. The X-ray emission from these knots can be either due to synchrotron emission or inverse Compton emission [94, 95, 96, 97, 78, 98]. If the X-ray flux lies below the extrapolation of radio-to-optical flux then the synchrotron origin is plausible else the emission may be due to inverse Compton process. Sambruna et al. [78] used one zone model to explain the X-ray emission from these knots. This model gives rise to a spectrum with a time-dependent exponential high-frequency cutoff. Hence for the knots with synchrotron origin of X-ray, their model predicted an exponentially decreasing steep spectrum at this energy. However the photon spectral index measured from the short duration observations of the knots at X-ray energies contradicts this prediction [98]. On the other hand, the acceleration process may exist for a longer duration such that a continuous injection of non-thermal particles may be viable. We model the observed radio-optical-X-ray spectra of the knots of the AGN 1136-135, 1150+497, 1354+195 and 3C 371 by using a continuous injection plasma model. We assume the knot to be a uniform expanding sphere with continuous injection of non-thermal particles. The electron distribution and resultant radiation spectrum is computed by taking into account synchrotron cooling, inverse Compton scattering of cosmic microwave background and adiabatic cooling due to the expansion of the sphere. The continuous injection of particles will generate a break in the electron distribution at energy where the cooling time-scale is equal to the age of the emission region (cooling break).

\(^1\) Chandra is a satellite borne X-ray telescope launched by NASA.
energy at which the index changes is time dependent, and synchrotron/adiabatic cool-
ing is dominant for the particles with energy greater than this break energy. We show
that this model can successfully reproduce the observed spectrum from the knots of
these sources.

We also interpret the knots as an outcome of an internal shock which we discuss in
§3.2. An internal shock interpretation for the knots was first suggested by Rees [99]
to explain the knots of M87. We study the jet dynamics and the viability of the above
mentioned continuous injection model using a simple internal shock model. The cen-
tral engine of the AGN is assumed to emits blobs of matter sporadically at relativistic
speeds. In the process, the fast moving blobs will collide with the previously ejected
slow moving ones thereby forming a shock. We implement the model for the knots
of AGN and compute the time-evolution of the non-thermal particles produced. Also
we compare the results obtained with the broadband fluxes from knots of several AGN
jets and their observed positions. The motivation here is to find quantitative values of
the model parameters by demanding that the model can self consistently explain the
observation.

It is also noted that the observed X-ray flux from the knots in the jet of the nearby
galaxy M87 cannot be explained by considering simple one zone models involving
continuous injection or one-time injection of non-thermal particles [100]. Perlman &
Wilson [100] proposed a modified CI model where the volume within which particle
acceleration occurs is energy dependent. Using this phenomenological model they ob-
erved that the particle acceleration takes place in a larger fraction of the jet volume in
the inner jet than the outer jet. Also particle acceleration region occupy a smaller frac-
tion of the jet volume at higher energies. Liu & Shen [79] proposed a two-zone model
with the acceleration region and cooling region spatially separated. The advection of
particles from the acceleration region to the cooling region introduces a break in the
particle spectrum which along with the cooling break in the cooling region produces a
double broken power law. The synchrotron emission from such a particle distribution
is used to fit the observed spectra. We propose an alternate two-zone model to explain
the observed X-ray flux from these knots and which we discuss in §3.3. Here we con-
sider the acceleration of power-law distribution of electrons in an acceleration region
which are then injected into a cooling region and cool via radiative processes. We show
that one will obtain a broken power-law particle spectrum when the initial spectrum
injected into the acceleration region is flatter than the characteristic spectrum of the
acceleration region. This particle distribution will then develop an additional cooling break depending upon the age of the knot while cooling in the cooling region. The synchrotron emission from the resultant particle distribution in the cooling region is used to reproduce the broadband spectrum of the knots of M87 jet.

- In Chapter 4 we discuss a model to explain the observed limb-brightened structure of the blazar MKN 501. The high resolution image of the nearby BL Lac object MKN 501 in radio show a transverse jet structure with the edges being brighter than the central spine. Such a feature is commonly referred as “limb-brightened” structure. This feature is usually explained by the “spine-sheath” model where the velocity at the spine is larger than the velocity at the boundary. The limb-brightened structure can then be explained by considering the differential Doppler boosting for a proper combination of the bulk Lorentz factors and the viewing angle. The viewing angle deduced for MKN 501 based on this model is $\gtrsim 15^\circ$ [101]. However the high-energy studies of MKN 501 demand the viewing angle to be $\sim 5^\circ$. Since the high-energy emission is originated from the inner part of the jet close to the nucleus, a possible bending of the jet was suggested by the earlier work [101]. We explain the observed limb-brightened structure of MKN 501 jet as an outcome of efficient particle acceleration process at the jet boundary. Here the particles can be accelerated by shear acceleration or turbulent acceleration. We deduce the required condition for the shear acceleration to be dominant over turbulent acceleration and discuss the diffusion of particles accelerated at the boundary into the jet medium. Also we derive the spectral index of the particle distribution accelerated via shear acceleration process and turbulent acceleration process and show the observed index at the boundary of MKN 501 jet supports the former. The bending of the jet is not a requirement for this interpretation unlike the explanation based on differential Doppler boosting.

- In Chapter 5 we study the temporal and spectral behaviour of the non-thermal emission from blazars. The flux variability observed in blazars have been studied by several authors using one zone and two zone models [80, 82, 102, 83]. Kirk et al. [82] and Kusunose et al. [83] assumed a two zone model where particles are accelerated in a region presumably by a shock and escape into emission region where they lose their energy by radiative processes. Chiaberge & Ghisellini [102] explained the short time variability observed in MKN 421 by dividing the emission region into thin slices. We study the spectral and temporal behaviour of blazars by considering a two zone model under two different scenarios of acceleration process. Mono energetic particles are
assumed to be accelerated in the acceleration region and then injected into the cooling region where they lose most of their energy via radiative processes. The behaviour of the resultant spectrum is then studied for two cases of particle acceleration process in the acceleration region. In the first case, the rate of particle acceleration is assumed to be energy dependent and in the second case, it is independent of the same. The model is then applied on the BL Lac object MKN 421 and its flare characteristics are studied under these two cases. It is found that, detailed information about the temporal behaviour of blazars can throw light on the underlying particle acceleration mechanism.

- Finally in Chapter 6, we summarize the work presented in the thesis and discuss the possible future work.
Knots in kpc scale jets of several AGN are recently studied in X-ray by the on-board X-ray satellite *Chandra* [94, 95, 78, 96]. *Chandra* due to its excellent spatial resolution is able to resolve bright X-ray knots and in most of the cases it coincides with their radio/optical counterparts [95, 78, 98]. The radio-to-optical emission from these knots are generally accepted to be of synchrotron origin, whereas the X-ray emission could be due to synchrotron [97, 78, 98] or inverse Compton processes depending on its radio-to-optical ($\alpha_{RO}$) and optical-to-X-ray index ($\alpha_{OX}$) [95, 78, 98, 96, 94]. If $\alpha_{RO} > \alpha_{OX}$, then the X-ray flux lies above the extrapolation of radio-to-optical flux and hence a single emission mechanism may not explain the observed fluxes. In such a case, the X-ray emission may be due to inverse Compton process or it may arise from a different electron population other than the one responsible for the radio/optical emission. However, it should be noted here that such a spectrum can still be a resultant synchrotron radiation from an electron distribution modified due to inverse Compton scattering happening at extreme Klein-Nishina regime [103]. The synchrotron spectrum in this case can satisfy the spectral indices requirement, $\alpha_{RO} > \alpha_{OX}$. On the other hand, if $\alpha_{RO} < \alpha_{OX}$, the X-ray flux lies below the extrapolation of radio-to-optical flux and synchrotron origin of X-ray is acceptable [96, 78, 98]. The synchrotron origin of X-rays for knots was strengthened in case of the knots of 3C’271 since the alternate inverse Compton model would require exceptionally large Doppler factors [98].

When the X-ray emission can be attributed to the inverse Compton process, the possible choices of target photons are radio/optical synchrotron photons (SSC) [104] or external photons. The source of external photons which can be dominant at kpc scale jet is the cosmic
microwave background\(^1\) (IC/CMB) [96, 95, 78, 98]. The SSC interpretation of X-ray emission would require large jet powers and magnetic fields much lower than the equipartition values whereas IC/CMB requires relatively low jet power and near equipartition magnetic fields [96].

These possible radiative process identifications have to be associated with (and confirmed by) dynamical models regarding the origin and subsequent evolution of the radiating non-thermal particles. In many models, these non-thermal particles are assumed to be generated by a short duration acceleration process and the particle distribution is determined by radiative losses (one-time injection) [95, 78, 96, 105, 76, 106]. The high energy particles cool more efficiently (\(\S 2.2.3\) and \(\S 2.2.4\)) and hence get depleted in time. This give rise to a time-dependent high energy cut off in the non-thermal particle distribution. If the X-ray emission is attributed to synchrotron emission by these particles, then these models predict an exponentially decreasing X-ray spectrum [95, 78, 96, 98]. This can be translated to a high energy cutoff in the electron distribution which in turn gives an estimate of the age of the knot. These non-thermal electrons move with a bulk speed \(v \approx c\) along the jet. Thus from the age of the knot, one can determine the location in the jet where the short duration acceleration process occurred. The distance of the knot from the central object and the short duration of acceleration (much less than the age of the knot) may naturally put strong constraints on any models of the acceleration process. Also the photon spectral slope measured during the short duration observations of 3C 371 is \(\alpha_X \approx 1.7 \pm 0.4\) which is in apparent contradiction to the predicted exponential X-ray spectrum by this model [98]. Moreover, the model requires the coincidence that the age of the knot be equal to the time required for X-ray emitting electrons to cool. A larger survey of X-ray jets have to be sampled to confirm whether this is statistically plausible.

On the other hand, it may also be possible that the acceleration process exists for a duration longer than the age of the knot, and hence, there is a continuous injection of non-thermal particles. The acceleration process may be due to internal shocks formed as a result of collision between sporadically ejected relativistic blobs of matter from the central engine [107]. In such a case the duration of acceleration can be roughly the cross-over time of the collided blobs. If we assume a typical blob size \(\sim \) kpc (which is roughly the size of the knots seen in radio) and its velocity \(\sim c\), the acceleration duration can be as large as \(\sim 10^{11}\) s. In the next section we study the case where the knots of AGN are modelled as a spherical expanding

\(^1\)According to the Big Bang theory, cosmic microwave background (CMB) radiation is the relic radiation left over from the formation of the universe.
emission region with a continuous injection of non-thermal particles and discuss the results obtained. In §3.2 we interpret the same considering an internal shock scenario and study its dynamic properties. Though one-time injection and continuous injection models can successfully reproduce the observed spectrum of knots of many AGN, they failed to explain the same for the knots of a nearby AGN, M87. The X-ray flux of M87 knots when compared with radio-to-optical flux suggests a synchrotron origin but its flux and/or spectral index cannot be reproduced by these simple models. In §3.3 we propose a modified synchrotron model to understand the broadband emission from the knots of M87.

3.1 A Continuous Injection Plasma Model

We consider the knot of the AGN as a plasma moving relativistically along the jet with a bulk Lorentz factor $\Gamma$. In the rest frame, it is assumed that the plasma uniformly occupies an expanding sphere with radius $R(t) = R_o + \beta_{\text{exp}}ct$, where $R_o$ is the initial size of the sphere and $\beta_{\text{exp}}c$ is the expansion velocity. Initially at $t = 0$ there are no non-thermal particles in the system. A continuous and constant particle injection rate for $t > 0$ is assumed, with a power-law distribution of energy,

$$Q(\gamma)\, d\gamma = q_o \gamma^{-p} \, d\gamma \quad \text{for} \quad \gamma > \gamma_{\text{min}} \tag{3.1}$$

where $\gamma$ is the Lorentz factor of the electrons. The evolution of the total number of non-thermal particles in the system, $N(\gamma, t)$, can be conveniently described by a kinetic equation of form

$$\frac{\partial N(\gamma, t)}{\partial t} + \frac{\partial}{\partial \gamma}[P(\gamma, t)N(\gamma, t)] = Q(\gamma) \tag{3.2}$$

Here $P(\gamma, t)$ is the particle energy loss rate given by

$$P(\gamma, t) = -(\dot{\gamma}_S(t) + \dot{\gamma}_{\text{IC}}(t) + \dot{\gamma}_A(t)) \tag{3.3}$$

where $\dot{\gamma}_S(t)$, $\dot{\gamma}_{\text{IC}}(t)$ and $\dot{\gamma}_A(t)$ are the cooling rates due to synchrotron, inverse Compton scattering of the cosmic microwave background (IC/CMB) radiation and adiabatic expansion.
respectively. These cooling rates are given by

\[ \dot{\gamma}_S(t) = \frac{4}{3} \frac{\sigma_T}{m_e c} \frac{B^2(t)}{8\pi} \gamma^2 \]  
(3.4)

\[ \dot{\gamma}_{IC}(t) = \frac{16}{3} \frac{\sigma_T}{m_e c^2} \Gamma^2 \sigma_{SB} T_{\text{cmb}}^4(z) \gamma^2 \]  
(3.5)

\[ \dot{\gamma}_A(t) = \frac{\beta_{\text{exp}} c \gamma}{R(t)} \]  
(3.6)

The former two cooling rates are associated with radiative losses and the latter is a non-radiative energy loss associated with the loss in internal energy used up for expansion [76]. Here the evolving magnetic field is parameterized to be 

\[ B(t) = B_0 \left( \frac{R(t)}{R_0} \right)^m \]

and 

\[ T_{\text{cmb}}(z) = 2.73(1 + z) \]

is the temperature of the CMB radiation at the redshift \( z \) of the source. Note that the time \( t \) and other quantities in the above equations are in the rest frame of the plasma.

We solved the equation (3.2) numerically for \( N(\gamma, t) \) using the finite difference scheme described by Chang & Cooper [108] and the resultant synchrotron and inverse Compton spectra are computed at an observing time \( t = t_o \). Finally, the flux at the Earth is computed taking into account the Doppler boosting [36], characterized by the Doppler factor 

\[ \delta \equiv [\Gamma(1 - \beta \cos \theta)]^{-1} \]

where \( \beta c \) is the bulk velocity of the plasma moving down the jet and \( \theta \) is the angle between the jet and the line of sight of the observer\(^2\).

While the total non-thermal particle distribution has to be computed numerically, a qualitative description is possible by comparing cooling time-scales with the observation time \( t_o \). The cooling time-scale due to synchrotron and inverse Compton cooling at a given time \( t \) and Lorentz factor \( \gamma \) is 

\[ t_c(t, \gamma) \approx \gamma / (\dot{\gamma}_S + \dot{\gamma}_{IC}) \]

Then, \( \gamma_c \), defined as the \( \gamma \) for which this cooling time-scale is equal to the observation time, 

\[ t_c(t_o, \gamma_c) \approx t_o \]

becomes,

\[ \gamma_c \approx \frac{m_e c}{\sigma_T} \left[ \frac{B^2(t_o)}{8\pi} + 4 \Gamma^2 c \sigma_{SB} T_{\text{cmb}}^4(z) \right] \]  
(3.7)

The adiabatic cooling time-scale \( t_a \) also turns out to be \( \approx t_o \), since \( t_a \approx R(t) / \beta_{\text{exp}} c \approx t_o \) for 

\[ R_o \ll R(t_o) \]

Thus, the non-thermal particle distribution at time \( t = t_o \) can be divided into three distinct regions [76]:

1. In the regime \( \gamma \ll \gamma_c \), radiative cooling is not important and \( N(\gamma, t_o) \approx q_0 \gamma^{-p} t_o \). The

\(^2\)Here and everywhere else in this thesis \( H_0 = 75 \text{ km s}^{-1} \text{ Mpc}^{-1} \) and \( q_0 = 0.5 \) are adopted
corresponding spectral index for both synchrotron and inverse Compton emission are 
\[ \alpha = \frac{p - 1}{2}. \]

2. In the regime \( \gamma \gg \gamma_c \), either synchrotron or inverse Compton cooling is dominant and \( N(\gamma, t_o) \propto \gamma^{-(p+1)} \). The corresponding spectral index for both synchrotron and inverse Compton emission are \( \alpha = p/2 \).

3. In the regime \( \gamma \approx \gamma_c \), either synchrotron or inverse Compton cooling as well as adiabatic cooling are important and the spectral slope is in the range \((p - 1)/2 > \alpha > p/2\).

The computed spectrum depends on the following ten parameters: the observation time \( t_o \), the magnetic field at the time of observation \( B_f = B(t = t_o) \), the magnetic field variation index \( m \), the radius of the knot at the time of observation, \( R_f = R(t = t_o) \), the index \( p \), the minimum Lorentz factor \( \gamma_{min} \), the Doppler factor \( \delta \), the bulk Lorentz factor \( \Gamma \), expansion velocity \( \beta_{exp}c \) and the normalization of the injection rate \( q_o \). On the other hand, there are only three observational points, namely, the radio, optical and X-ray fluxes. Clearly, the parameters are under constrained, and it is not possible to extract meaningful quantitative estimates. However, the motivation here is to show that this model can explain the observed data with reasonable values of the above parameters.

### 3.1.1 Results and Discussion

We applied the above model to the knots of the AGN observed by Chandra along with the information available in radio and optical energies (Figure 3.1, 3.2, 3.3). In Figure 3.4, the computed spectra are compared with the data for different knots for four sources namely 1136-135, 1150+497, 1354+195 and 3C 371. The values of the parameters used are tabulated in Table 3.1. The injected power in non-thermal particles in the rest frame of the knot can be written as

\[
P_{\text{inj}} = \int_{\gamma_{\text{min}}}^{\infty} (\gamma m_e c^2) Q(\gamma) \, d\gamma = q_o \frac{m_e c^2}{p - 2} \gamma_{\text{min}}^{-(p-2)} \quad (3.8)
\]

while the total jet power can be approximated as

\[
P_{\text{jet}} = \pi R^2 \Gamma^2 \beta c (U_p + U_e + U_B) \quad (3.9)
\]
where $U_p$, $U_e$ and $U_B$ are the energy densities of the protons, electrons and the magnetic field respectively. Here it has been assumed that the protons are cold and the number of protons is equal to the number of electrons. We found the jet power ranges from $10^{46}$ to $2 \times 10^{48}$ ergs s$^{-1}$, while the injected power is generally three orders of magnitude lower. This means that the non-thermal acceleration process is inefficient and most of the jet power is expected to be carried to the lobes. The magnetic field $B_f$ is nearly equal to the equipartition values.

The X-ray emission for the knots in 3C371 and Knot A of 1136-135 are identified as being due to synchrotron emission which is consistent with earlier works [78, 98]. However, in this case the predicted X-ray spectral index is $\alpha_X = \alpha_R + 1/2$ instead of being exponential. Note that this relation between the spectral indices is independent of the parameters used to fit the data. For the rest of the sources the X-ray emission is attributed to IC/CMB which is again consistent with the results obtained from the earlier works [78]. However, for some of the sources the optical spectral index is now $\alpha_O = \alpha_R + 1/2$ instead of being exponential.

The present model can be confirmed (or ruled out) vis-à-vis one-time injection models, by future measurements of the radio $\alpha_R$, optical $\alpha_O$ and X-ray $\alpha_X$ spectral indices. In particular, the following cases are possible:
Fig. 3.2: Chandra image of 1150+497 and 1354+195 overlaid with VLA radio contours. Figure reproduced from Sambruna et al. [78]
Fig. 3.3: Multi wavelength image of 3C371 in X-ray by Chandra (top), in optical by HST (middle) and in radio by Merlin (radio). Figure reproduced from Pesce et al. [98]
Fig. 3.4: The observed fluxes in radio, optical and X-ray compared with model spectrum using parameters given in table 3.1. The data for 3C 371 is taken from Pesce et al. [98], while the rest are taken from Sambruna et al. [78]. Triangles correspond to knot A and circles to knot B.
Chapter 3. Emission Models and the Dynamics of AGN knots

Table 3.1: Parameters for model fittings

| Source/knot   | $B_f$ ($\times 10^{-5}$ G) | $\gamma_{\text{min}}$ | $p$ | $t_o$ ($\times 10^{11}$ s) | $\delta$ | $\Gamma$ | $\beta_{\text{exp}}$ | $P_{\text{inj}}$ (ergs s$^{-1}$) | $P_{\text{jet}}$ (ergs s$^{-1}$) | $B_f/B_{\text{equ}}$ |
|---------------|-----------------------------|------------------------|-----|--------------------------|---------|----------|-----------------------|-------------------------------|-------------------------------|---------------------|
| 1136-135A     | 0.9                         | 2.0                    | 2.4 | 0.2                      | 5       | 5        | 0.8                   | 44.7                          | 47.5                          | 0.4                 |
| 1136-135B     | 4.0                         | 20.0                   | 2.9 | 9                        | 5       | 5        | 0.1                   | 44.2                          | 47.9                          | 0.5                 |
| 1150+497A     | 2.5                         | 30.0                   | 2.85| 9                        | 5       | 3.5      | 0.1                   | 44.2                          | 47.3                          | 0.4                 |
| 1150+497B     | 4.3                         | 30.0                   | 3.3 | 9                        | 5       | 3.5      | 0.1                   | 44.1                          | 47.3                          | 0.75                |
| 1354+195A     | 1.7                         | 40.0                   | 3.0 | 9                        | 3.5     | 2        | 0.1                   | 45.8                          | 48.4                          | 0.04                |
| 1354+195B     | 8.0                         | 25.0                   | 3.2 | 9                        | 3.5     | 2        | 0.1                   | 44.7                          | 47.5                          | 0.63                |
| 3C 371A       | 1.3                         | 10.0                   | 2.4 | 12                       | 3.5     | 3.5      | 0.1                   | 42.8                          | 46.4                          | 0.8                 |
| 3C 371B       | 1.0                         | 10.0                   | 2.4 | 1                        | 3.5     | 3.5      | 0.5                   | 43.5                          | 46.0                          | 0.9                 |

Columns: (1) Source and knot name taken from Pesce et al. [98] for 3C371 and the rest from Sambruna et al. [78]. (2) Magnetic field at the observation time, $B = B(t = t_o)$; (3) Minimum Lorentz factor $\gamma_{\text{min}}$; (4) Power-law index of the injected non-thermal particle $p$; (5) Observation time $t_o$; (6) Doppler factor $\delta$; (7) Bulk Lorentz factor $\Gamma$; (8) Velocity of expansion $\beta_{\text{exp}}$ (in units of $c$); (9) Log of the injected power $P_{\text{inj}}$; (10) Log of the total jet power $P_{\text{jet}}$; (10) Ratio of the magnetic field to the equipartition value. For all cases, the magnetic field variation index $m$ and the size of the source at $t = t_o$ is fixed at $1$ and $5 \times 10^{21}$ cm, respectively.

1. In the case $\alpha_R \approx \alpha_X$, the X-ray emission is probably due to IC/CMB. Both the continuous injection and one-time injection models are equally viable.

2. In the case $\alpha_X \approx \alpha_R + 1/2$, the X-ray emission would be due to synchrotron emission from electrons in the cooling dominated region. The continuous injection scenario will be favored in such case.

3. In the case $\alpha_X > \alpha_R + 1/2$, when the X-ray emission is exponentially decreasing, it should be attributed to the high energy cutoff in the electron distribution. The one-time injection scenario will be favored.

4. In the case $\alpha_X < \alpha_R$, when the X-ray emission is exponentially increasing, it should be attributed to the low energy cutoff ($\gamma_{\text{min}}$) in the electron distribution and the X-ray emission should be due to IC/CMB. Both the continuous injection and one-time injection models are equally viable.

Similar arguments can be put forth for the optical spectral index $\alpha_O$ as compared to the radio. It should be noted that for some older systems the one-time injection would be the natural scenario, while for younger systems the continuous injection would be more probable. The technique described above will be able to differentiate between the two, and a generic constraint on the acceleration time-scales and typical age of the knots may be obtained. A generic model where the injection rate decays in time may then be used to fit the observations. The measurement of spectral indices in different wave-lengths will also reduce
the number of unconstrained parameters in the model fitting, leading to reliable estimates of
the system parameters.

3.2 Internal Shock Interpretation

The possibility of a continuous injection of non-thermal particles into the knots of AGN as
well as its dynamical properties can be studied using a specific injection model. While there
is no consensus on the origin of these non-thermal particles, one of the standard model is
the internal shock scenario \cite{99,107}. Here the particles are energized by Fermi acceleration
in shocks produced during the interaction of relativistically moving blobs ejected from
the central engines with different speeds. A detailed description of the shock formation and
subsequent electron acceleration is complicated and would require numerically difficult mag-
neto hydrodynamic simulations. Moreover from the limited number of observables, which
can be obtained from the featureless spectrum in two or three different energy bands, one
may not be able to constrain the various assumptions and/or the initial conditions of such
a detailed study. Nevertheless, a qualitative idea as to whether the internal shock model is
consistent with the present observations (and if so, qualitative estimates of the model param-
eters) would be desirable. Such an estimate would provide insight into the temporal behavior
of the central engine.

We implement an internal shock model with simplifying assumptions and compute the time
evolution of the non-thermal particles produced. We also compare the results obtained with
the broadband fluxes from knots of several AGN jets and their observed positions. The mo-
tivation here is to find a consistent set of model parameters that can explain the observations
and thereby make qualitative estimates of their values. Apart from the fluxes at different
energy bands, the spectral indices in each band can also provide important diagnostic infor-
mation about the nature of these sources. Hence, we have analyzed long (> 40 ks) Chandra
observations of three AGN and present the constrains that were obtained on the X-ray spec-
tral indices of the individual knots.

3.2.1 Chandra X-ray Data Analysis

Long-exposure Chandra observations of the sources PKS 1136-135, PKS 1150+497, and
3C 371 were performed with the Advanced CCD Imaging Spectrometer (ACIS-S) with the
Table 3.2: Chandra Observations

| Source name | Obs Id | Exposure (ks) | Knots | $F_{0.3-3.0}$ (ergs cm$^{-2}$ s$^{-1}$) | $\alpha_X$ |
|-------------|-------|---------------|-------|--------------------------------------|------------|
| 1136 − 135  | 3973  | 77.37         | A     | 0.63                                 | $1.24^{+1.51}_{-0.66}$ |
|             |       |               | B     | 1.56                                 | $0.65^{+0.69}_{-0.28}$ |
| 1150 + 497  | 3974  | 68.50         | A     | 3.15                                 | $0.66^{+0.28}_{-0.26}$ |
|             |       |               | B     | 0.61                                 | $0.90^{+1.02}_{-1.02}$ |
| 3C371       | 2959  | 40.86         | A     | 3.32                                 | $1.43^{+0.85}_{-0.76}$ |
|             |       |               | B     | 7.84                                 | $1.07^{+0.26}_{-0.23}$ |

**Columns:** (1) Source name; (2) Chandra Observation Id; (3) Exposure time; (4) Knots prominent in X-ray; (5) Flux in 0.3 – 3.0 keV energy band; (6) X-ray energy spectral index.

source at the aim point of the S3 chip. The Observation ID (ObsID) and the exposure time of the observation are given in Table 3.2. Earlier shorter duration observations of these sources revealed two bright knots for each source, whose positions from the nucleus are given in Table 3.3. These longer duration observations allow for better constraint on the X-ray spectral indices of these knots.

The data from Chandra X-ray observatory$^3$ were analyzed using the Chandra data analysis software CIAO$^4$ and the latest calibration files were used to produce the spectrum. The X-ray counts from each individual knot was extracted using a circular region centered at the knot. The background was estimated from the counts obtained from same size regions located at the same distance from the nucleus but at different azimuth angles. The radius of the circular region was chosen to be 0.74$''$ for the sources 1136 − 135 and 1150 + 497, while for 3C371 a smaller radius of 0.6$''$ was used. These sizes were chosen to minimize any possible contamination from the nucleus and/or the other knot.

Spectral fits were undertaken on the data using the XSPEC package$^5$ in C statistic mode, which is the appropriate statistic when the total counts are low. The flux and the energy spectral indices obtained are tabulated in Table 3.2. The quoted spectral indices are not very meaningful due to large errors.

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3http://cxc.harvard.edu/cda/
4http://cxc.harvard.edu/ciao/
5http://heasarc.gsfc.nasa.gov/docs/xanadu/xspec/
Table 3.3: Observed Knot Features

| Source name | Type   | z       | Knot | Position (arcsec) | $\alpha_{RO}$ | $\alpha_{OX}$ | Ref |
|-------------|--------|---------|------|-------------------|---------------|---------------|-----|
| 1136 – 135  | FSRQ   | 0.554   | A    | 4.5               | 0.73          | 0.83          | 78  |
|             |        |         | B    | 6.7               | 1.04          | 0.68          |     |
| 1150 + 497  | FSRQ   | 0.334   | A    | 2.1               | 0.99          | 0.83          | 78  |
|             |        |         | B    | 4.3               | 1.24          | 0.44          |     |
| 1354 + 195  | FSRQ   | 0.720   | A    | 1.7               | 1.05          | 0.6           | 78  |
|             |        |         | B    | 3.6               | 1.15          | 0.68          |     |
| 3C273       | QSO    | 0.158   | A    | 13                | 0.86          | 0.61          | 95  |
|             |        |         | B    | 15                | 0.9           | 0.73          |     |
| 3C371       | Bl Lac | 0.051   | A    | 1.7               | 0.9           | 1.28          | 98  |
|             |        |         | B    | 3.1               | 0.76          | 1.14          |     |

**Columns:** (1) Source name; (2) Type of the source; (3) Redshift; (4) Knots prominent in X-ray (The nomenclature is same for all the knots as they are in the literature except for 3C371 where A and B are reversed); (5) Position of the knot; (6) Radio-to-Optical index; (7) Optical-to-X-ray index; (8) References: Sambruna et al. [78], Sambruna et al. [95], Pesce et al. [98].

### 3.2.2 Model

In the internal shock model framework, temporal variations of the ejection process produces density fluctuations (moving with different velocities), which collide at some distance from the source to produce an observable knot. This distance will depend on the time-scale over which the variation takes place. In general the system will exhibit variations over a wide range of time-scales and knots like features would be produced at different distance scales. Here, we consider large-scale jets (with deprojected distances $\approx 100$ kpc) which are expected to arise from variability occurring on a corresponding large time-scale. Variations on smaller time-scales would produce knot structures on smaller distance scales, for example, pc scale or even smaller jets, which would be unresolved for these sources. These smaller time-scale variabilities will be smoothed out at large distances, and hence one expects the jet structure of these sources to be determined by variations over a single characteristic time-scale. To further simplify the model, we approximate the density and velocity fluctuations as two discrete blobs with equal masses, $M_1 = M_2 = M$, having Lorentz factors, $\Gamma_1$ and $\Gamma_2$ that are ejected one after the other, from the central engine with a time delay of $\Delta t_{12}$. The collision of the blobs is considered to be completely inelastic; i.e. the blobs coalesce and move as a single cloud, which is identified with the observed knot. From conservation of
momentum, the Lorentz factor of the knot is

$$\Gamma = \sqrt{\left(\frac{\Gamma_1 \beta_1 + \Gamma_2 \beta_2}{2}\right)^2 + 1} \quad (3.10)$$

where $\beta_{1,2} = v_{1,2}/c$ are the velocities of the blobs. Since the collision is inelastic, a fraction of the bulk kinetic energy is dissipated. Denoting all quantities in the rest frame of the knot by subscript $K$, this dissipated energy $\Delta E_K$ can be estimated as

$$\Delta E_K = [\Gamma_{1K} + (\Gamma_{2K} - 2)]Mc^2 \quad (3.11)$$

The Lorentz factors of the blobs in the knot’s rest frame $\Gamma_{1K,2K} = (1 - \beta_{1K,2K}^2)^{-1/2}$ are computed using

$$\beta_{1K,2K} = \frac{\beta_{1,2} - \beta}{1 - \beta_{1,2} \beta} \quad (3.12)$$

where $\beta = (1 - 1/\Gamma^2)^{1/2}$. The time-scale on which this energy will be dissipated can be approximated to be the crossing-over time of the two blobs,

$$T_{ON,K} \approx \frac{2 \Delta x_K}{c(\beta_{2K} - \beta_{1K})} \quad (3.13)$$

where, $\Delta x_K$ is the average size of the two blobs in the rest frame of the knot. It is assumed that this dissipated bulk kinetic energy, $\Delta E_K$, gets converted efficiently to the energy of the non-thermal particles produced during the collision. The number of non-thermal particles injected per unit time into the knot is taken to be,

$$Q_K(\gamma) \, d\gamma = A \gamma^{-p} \, d\gamma \quad \text{for} \quad \gamma > \gamma_{\text{min}}, \quad (3.14)$$

where $\gamma$ is the Lorentz factor of the electrons, $p$ is the particle index and $A$ is the normalization constant given by,

$$A = \frac{\Delta E_K}{T_{ON,K}} \frac{(p - 2)}{m_e c^2} \gamma_{\text{min}}^{(p - 2)} \quad (3.15)$$

Here the injection is assumed to be uniformly occurring for a time $T_{ON,K}$. The cloud is assumed to be permeated with a tangled magnetic field, $B_K$.

We used the kinetic equation (3.2) to study the evolution of the total number of non-thermal...
particles in the system. However for the total loss rate we considered only the synchrotron and IC/CMB (equations (3.4) and (3.5)). Using equation (3.6), the adiabatic cooling time-scale can be written as $t_a \approx \gamma/\dot{\gamma}_A = R(t)/(\beta_{\exp} c)$. And for any given time of observation, $t_{O,K}$, the size of the blob will be $R(t_{O,K}) = R_0 + \beta_{\exp} c t_{O,K}$. From the above one finds that $t_a$ will be always larger than $t_{O,K}$ and hence the adiabatic cooling can be neglected. For $t_{O,K} < T_{ON,K}$, the resultant particle distribution will be a broken power-law giving rise to a composite synchrotron spectrum with a spectral break ($\ §3.1$).

The predicted spectrum and size of a knot depends on nine parameters, which are the mass of the blobs $M$, their average size $\Delta x_K$, the Lorentz factors $\Gamma_1$ and $\Gamma_2$, the particle injection index $p$, the minimum Lorentz factor $\gamma_{\min}$, the magnetic field $B_0$, the inclination angle of the jet $\theta$, and the observation time $t_{O,K}$.

From these parameters and the location of the knot in the sky plane, one can infer the time delay $\Delta t_{12}$ between the ejection of the two blobs. The projected distance of the knot from the source $S$ can be written as

$$S = c(\beta_1 t_c + \beta t_O) \sin \theta$$  \hspace{1cm} (3.16)

where $t_O = \Gamma t_{O,K}$ is the time of the observation after the formation of the knot in the source frame. The time elapsed $t_c$, after the ejection of the first blob and the start of the collision, is given by,

$$t_c = \frac{v_2 \Delta t_{12} - \Delta x_1}{v_2 - v_1}$$  \hspace{1cm} (3.17)

where $\Delta x_1$ is the size of the first blob $\approx \Gamma_{1,K} \Delta x_K / \Gamma$. Thus $\Delta t_{12}$ can be estimated using the above equation, where $t_c$ is given by equation (3.16), and it essentially depends on four parameters, $\theta, \Gamma_1, \Gamma_2$ and $t_{O,K}$.

A total time, $t_{tot}$, can be defined to be the time that has elapsed between the ejection of the first blob and the observation, $t_{tot} = t_c + t_O$. For two knots, A and B, the time difference between the ejection of their first blobs, $t^{AB}$ is then

$$t^{AB} = t_A^{tot} - t_B^{tot} - t_{LT}$$  \hspace{1cm} (3.18)

where $t_{LT}$ is the light travel time difference between the two knots, which is approximated
to be

$$t_{LT} \approx \frac{S^A - S^B}{c \tan \theta}$$  \hspace{1cm} (3.19)

The power of the jet, can be defined in two different ways. The instantaneous power, which is the power when the system is active, can be defined to be the average energy of the blobs divided by the time-scale on which the blobs are ejected. This power can be estimated for each knot to be

$$P_{ins} \approx \frac{M c^2 (\Gamma_1 + \Gamma_2)/2}{\Delta t_{12}}$$  \hspace{1cm} (3.20)

On the other hand, the time averaged power of the jet can be defined as the typical energy ejected during active periods divided by the time-scale on which such activity occurs. For two knots, A and B, this can be approximated to be

$$P_{ave} \approx \frac{[M^A c^2 (\Gamma_1^A + \Gamma_2^A) + M^B c^2 (\Gamma_1^B + \Gamma_2^B)]/2}{\Delta t_{AB}}$$  \hspace{1cm} (3.21)

### 3.2.3 Results and Discussion

The model has been applied to those knots of kpc scale jets, which have been detected by Chandra and for which radio and optical data are available. This criterion was satisfied by the two brightest knots of the AGN: 1136 − 135 (Figure 3.1), 1150 + 497 (Figure 3.2), 1354 + 195 (Figure 3.2), 3C273 (Figure 1.6) and 3C371 (Figure 3.3). In this work, the knot closer to the nucleus is referred to Knot A and the further one as Knot B. For these sources this nomenclature is same as in the literature [78, 95] except for 3C371, for which the farther one has been referred to Knot A [98]. For three of these sources, the X-ray spectral indices were constrained using long exposure observations as described in §3.2.1. The observed properties of the sources and the knots are tabulated in Table 3.2 and 3.3.

Figure 3.5 shows the observed radio, optical and X-ray fluxes of these knots along with the computed spectrum corresponding to model parameters that are given in Table 3.4. Since the number of parameters is large as compared to the observables, a unique set of parameter values cannot be obtained. Two consistency checks have been imposed on the parameter values: that the number of non-thermal electrons which will be injected into the system, \( N_{nth} \), is smaller than the total number of protons, \( N_k \), and that the magnetic field, \( B \) should
Fig. 3.5: The observed fluxes in radio, optical and X-ray compared with model spectrum using parameters given in Table 3.4. Knot A fluxes are represented by filled triangles, and knot B fluxes are represented by filled circles.
Table 3.4: Internal Shock Model Parameters

| Source     | Knot | $\theta$ (deg) | $\Gamma_1$ | $\Gamma_2$ | $\Gamma_\ast$ | $t_{O,K}$ (10^{11}s) | $t_{\ast}^c$ (10^{12}s) | $t_{\ast}^{tot}$ (10^{12}s) | $\log M$ (g) | $\gamma_{\min}$ | $p$ | $B$ (10^{-5}G) |
|------------|------|---------------|------------|------------|----------------|-----------------------|------------------------|----------------------|--------------|---------------|-----|----------------|
| 1136-135   | A    | 11.5          | 4.6        | 5.4        | 5.0           | 0.25                  | 11.8                   | 11.6                 | 38.0         | 2             | 2.4 | 1.1           |
|            | B    | ...           | 4.1        | 5.9        | 5.0           | 8.5                   | 13.1                   | 17.3                 | 36.4         | 20            | 2.9 | 4.5           |
| 1150+497   | A    | 10.2          | 2.5        | 3.0        | 2.8           | 16.0                  | 1.2                    | 5.1                  | 37.5         | 30            | 2.8 | 1.4           |
|            | B    | ...           | 2.6        | 3.5        | 3.0           | 6.2                   | 8.8                    | 10.0                 | 36.8         | 30            | 3.3 | 4.0           |
| 1354+195   | A    | 8.21          | 1.7        | 2.3        | 2.0           | 9.0                   | 8.1                    | 8.1                  | 38.1         | 37            | 3.0 | 1.7           |
|            | B    | ...           | 1.7        | 2.3        | 2.0           | 8.7                   | 15.0                   | 17.0                 | 37.1         | 20            | 3.2 | 8.0           |
| 3C 273     | A    | 8.23          | 2.2        | 4.0        | 3.1           | 2.1                   | 27.0                   | 24.0                 | 36.8         | 50            | 2.7 | 0.4           |
|            | B    | ...           | 1.7        | 2.3        | 2.0           | 8.0                   | 30.0                   | 32.0                 | 37.8         | 80            | 2.8 | 0.6           |
| 3C 371     | A    | 15.8          | 1.6        | 2.4        | 2.0           | 5.9                   | 0.18                   | 1.4                  | 35.5         | 20            | 2.5 | 1.0           |
|            | B    | ...           | 2.0        | 2.4        | 2.2           | 2.3                   | 0.39                   | 0.75                  | 36.8         | 10            | 2.4 | 0.6           |

Eight model parameters and derived quantities. The ninth parameter is $\Delta x_K = 5.0 \times 10^{21}$ cm for all sources. Columns marked with an asterisk are derived quantities and not parameters.

Columns: (1) Source name; (2) Knot; (3) Viewing angle; (4) Lorentz factor of the first blob; (5) Lorentz factor of the second blob; (6) Lorentz factor of the Knot; (7) Observation time; (8) Collision time; (9) Total time; (10) Mass of the blobs; (11) Minimum Lorentz factor of the particle injected into the knot; (12) Injected particle spectral index; (13) Magnetic field.

Table 3.5: Knot/Jet Properties

| Source     | Knot | $\Delta t_{12}$ (10^{11}s) | $t_{AB}$ (10^{11}s) | $B_{eq}$ | $N_{nth}/N_K$ | $\log P_{ins}$ (ergs s^{-1}) | $\log P_{ave}$ (ergs s^{-1}) | $T_{ON,K}$ (10^{11}s) | $\frac{t_{O,K}}{T_{ON,K}}$ | D (kpc) |
|------------|------|-----------------------------|----------------------|----------|----------------|-------------------------------|-------------------------------|------------------------|--------------------------|---------|
| 1136-135   | A    | 1.1                         | 3.1                  | 0.55     | 0.88           | 48.60                         | 48.17                         | 20.4                   | 0.01                     | 112.8    |
|            | B    | 2.5                         | 0.74                 | 0.74     | 46.65          | 9.1                           | 0.93                         | 163.4                  | 0.01                     | 112.8    |
| 1150+497   | A    | 0.9                         | 5.0                  | 0.12     | 0.13           | 47.98                         | 47.31                         | 17.0                   | 0.94                     | 50.5     |
|            | B    | 3.9                         | 0.65                 | 0.43     | 46.63          | 10.7                          | 0.58                         | 96.1                   | 0.94                     | 50.5     |
| 1354+195   | A    | 7.5                         | 18.0                 | 0.04     | 0.38           | 47.49                         | 47.15                         | 9.6                    | 0.94                     | 78.7     |
|            | B    | 17.0                        | 0.64                 | 0.78     | 46.10          | 9.6                           | 0.90                         | 135.5                  | 0.90                     | 135.5    |
| 3C 273     | A    | 20.0                        | 34.0                 | 0.03     | 0.80           | 46.00                         | 46.63                         | 5.4                    | 0.39                     | 240.7    |
|            | B    | 31.6                        | 0.02                 | 0.16     | 46.60          | 9.6                           | 0.84                         | 248.5                  | 0.84                     | 248.5    |
| 3C 371     | A    | 1.4                         | 1.5                  | 0.39     | 0.88           | 45.59                         | 46.98                         | 7.2                    | 0.83                     | 11.2     |
|            | B    | 1.0                         | 0.28                 | 0.29     | 47.14          | 16.3                          | 0.14                         | 7.7                    | 0.14                     | 7.7      |

Columns: (1) Source name; (2) Knot; (3) Time delay between the ejection of the blobs; (4) Time delay between the ejection of the first and the third blob; (5) Ratio of the magnetic field to equipartition magnetic field; (6) Ratio of non-thermal electrons to the total number of protons; (7) The instantaneous power; (8) Time-averaged power; (9) Time-scale over which non-thermal particles are injected; (10) Ratio of the observation time to particle injection time-scale; (11) Deprojected distance.
not deviate far from the equipartition value, $B_{eq}$. Both of these conditions are satisfied by the parameter sets as shown in Table 3.5 where the ratios $B/B_{eq}$ and $N_{nth}/N_K$ are given.

For each source, the time delay $\Delta t_{12}$ between the ejection of the two blobs that form the knots are nearly equal to the time difference between the ejection of the first blobs of Knot A and Knot B, $t^{AB}$. This gives an overall single time-scale of activity for each source which ranges from $10^{11} - 10^{12}$ s and can reproduce the knot properties, as had been assumed in the development of the simple internal shock model. This result is important since if it had not been true, a more complex temporal behavior would have to be proposed, wherein the jet structure is due to variability of the source at two different time-scales, the first being the time difference between the ejection of two blobs that form a knot, and the second being the time difference between the activities that produced the two knots.

The cross-over time, $T_{ON,K}$, is determined here by equation (3.13). Since, $\beta_{2K} - \beta_{1K} \approx (\Gamma_1 - \Gamma_2)/\Gamma \approx 0.3$, $T_{ON,K} \approx 10^{12}(\Delta x_K/5 \times 10^{21}\text{cm})$ s. During this time, i.e., the time when there is injection of particles into the system, the knot would travel a distance $\approx cT_{ON} \approx c\Gamma T_{ON,K} \approx 50(\Gamma/5)(\Delta x_K/5 \times 10^{21}\text{cm})$ kpc. This is a significant fraction of the total observable distance traveled by the knot, from formation $\approx c t_c \approx 100$ kpc (see Table 3.5) to the termination of the jet in the radio lobe, $\approx 200$ kpc. Hence, as shown in Table 3.5 it is possible to fit the spectrum of all the knots with an observation time $t_{O,K}$, which is less than the crossing-over time $T_{ON,K}$, implying that there is continuous injection of particles into the system. In this scenario, the synchrotron and IC spectra will have a break corresponding to a Lorentz factor $\gamma_c$ where the cooling time-scale equals the observation time $\tau_{3.1}$.

The knots of 3C371 and Knot A of 1136-135 are unique in this sample, since their the X-ray flux lies below the extrapolation of the radio-to-optical spectrum to X-ray wavelengths. This allows for the interpretation that the X-ray flux is due to synchrotron emission [78, 98]. For Knot A of 1136-135 and Knot B of 3C371, this implies that the spectral break for the synchrotron emission occurs at the X-ray regime (Figure 3.5), which in turn indicates that these sources are relatively younger. Indeed, the ratio of the observation time to the injection time, $t_{O,K}/T_{ON,K}$, for these sources are smallest (Table 3.5). On the other hand, for Knot A of 3C371, the spectral break can occur at the optical band even if the X-ray flux is interpreted as being due to synchrotron emission (Figure 3.5) and hence this source need not be relatively young. However, this is only possible if $t_{O,K} < T_{ON,K}$ and there is continuous injection of particles. Otherwise, a sharp cutoff in the spectrum at the optical band would have occurred and the X-ray emission would not be due to synchrotron emission.
Figure 3.5 shows that the radio, optical and X-ray spectral indices for different knots may vary and highlights the need for more spectral measurements in all bands. A definite prediction of this model is that for most knots, the X-ray spectral index should be equal to the radio spectral index, indicating that the X-ray flux is due IC/CMB. Such spectral constrains would be particularly important since, although it has been demonstrated here that the internal shock model can explain the broadband spectrum of these sources, there could be other models that may be physically and observationally more favorable. Analysis of the Very Large array (VLA) and Hubble Space Telescope (HST) images of 3C273 have shown that the optical and radio spectral indices are different, indicating the presence of an additional emission mechanism for the source \(^{109}\). Also, the X-ray flux may be due to a second population of non-thermal electrons, rather than being the IC/CMB spectrum of the same distribution that produces the radio and optical emission \(^{110}\). An argument in support of this is that since in the IC/CMB model the X-ray emission is due to electrons that are only a factor ten more energetic than those which produce the radio, a source in which the X-ray flux falls rapidly from the center should also exhibit a similar decrease in radio emission. However this feature is not observed (e.g. 3C273 and 1354+195). Moreover, the jet power required in the IC/CMB model can be very large, \(\approx 10^{48}\) ergs/s, which may be larger than the power inferred from the giant radio lobes \((\lesssim 10^{47}\) ergs/s). While the former argument may not strictly be applicable to the internal shock model (since each knot is a separate entity and the distance from the source is not a measure of the age of the source), the power requirement for some sources may indeed be very large, for e.g. Knot A of 1136-135 requires \(P_{\text{ave}} \approx 2 \times 10^{48}\) ergs/s (Table 3.5). However, the energy requirement may be decreased if the magnetic field is sub-equipartition e.g. 3C273 (Table 3.5). Thus it is desirable to obtain direct observational signatures, such as spectral indices, to discriminate between models.

A realistic description of the knots is more complicated than the simple model considered here. For example, the forward and reverse shocks that should form when the blobs collide, may provide different injection rates and at different locations within the Knot. However, the physics of these shock formations and the subsequent acceleration of particles is complicated and unclear, especially if they are mediated by magnetic fields. In the future, results from sophisticated numerical simulations could be compared with higher resolution data (which can resolve the internal structure of the knots) to prove (or disprove) the internal shock model.
3.3 A Two Zone Model for the knots of M87

The models discussed in earlier sections can explain the broadband spectrum from the knots of many AGN jets. However, they fail in the case of M87, a nearby giant elliptical galaxy, at a distance = 16 Mpc [111], possessing an one-sided jet with projected distance \( \approx 2 \) kpc. The jet is bright in radio, optical, and X-ray energies. The jet structure is very well studied in radio, infrared, optical, and X-ray energies [49, 112, 113, 114, 115, 100]. The flux and the spectral indices at X-ray energies indicate a possible continuation of synchrotron emission of the radio-to-optical spectrum with the change in the spectral index beyond optical energies [116, 117]. Simple theoretical models, namely the continuous injection model [76, 118, 119, 120] (§3.1) and the one-time injection model [105, 76, 106] were unable to explain the observed X-ray flux and/or the spectral index. The X-ray flux predicted by the continuous injected model is more than the observed flux whereas the one-time injection model with pitch angle scattering under predicts the X-ray flux and the one without pitch angle scattering fails to predict the observed X-ray spectral index [100].

We propose a two zone model to explain the non-thermal emission from the knots of M87 jet. In a two zone model, particles are accelerated in a region, namely acceleration region (AR) (probably around a shock front), and subsequently cool off through radiative processes in an associated cooling region (CR).

3.3.1 The Model

We consider the acceleration of a power-law distribution of particles (which may be a relic of a past acceleration process) at a shock front and cooling via synchrotron radiation in a homogeneous magnetic field. We treat the present scenario as two zones: one around the shock front where the particles are accelerated (AR) and the downstream region where they lose most of their energy through the synchrotron process (CR). This model is then used to explain the radio-optical-X-ray spectrum of the knots in the M87 jet. We assume the CR to be a spherical blob of radius \( R \) with tangled magnetic field \( B_{CR} \) and the AR is assumed to be a very thin region with magnetic field \( B_{AR} \). A power-law distribution of electrons is continuously injected into the AR characterized by an acceleration time-scale \( t_{acc} \). Particles are then accelerated at a rate \( 1/t_{acc} \) to a maximum energy determined by the loss processes. The AR is assumed to be compact and the emission from the CR mainly contributes the overall photon spectrum.
Fig. 3.6: Grey-scale image of M87 in radio (top), optical (middle) and X-ray (bottom). Figure reproduced from Wilson & Yang [121].
The kinetic equation governing the evolution of electrons in AR is given by

$$\frac{\partial n(\gamma, t)}{\partial t} = \frac{\partial}{\partial \gamma} \left[ \left( \zeta_{AR} \gamma^2 - \frac{\gamma}{t_{acc}} \right) n(\gamma, t) \right] - \frac{n(\gamma, t)}{t_{esc}} + Q(\gamma) \quad (3.22)$$

where

$$Q(\gamma) d\gamma = q_o \gamma^{-p} d\gamma \quad \text{for} \quad \gamma_{min} < \gamma < \gamma_b \quad (3.23)$$

Here $\gamma$ is the Lorentz factor of the electron, $t_{esc}$ is the escape time-scale and $\zeta_{AR} = \frac{\sigma_T}{6\pi m_e c} B_{AR}^2$.

Equation (3.22) can be solved analytically using Green’s function and the electron distribution for an energy-independent $t_{acc}$ and $t_{esc}$ at time $t$ is given by \[122\]

$$n(\gamma, t) = t_{acc} \gamma^{-(\alpha+1)} \left( 1 - \frac{\gamma}{\gamma_{max}} \right)^{\alpha-1}$$

$$\times \int_{x_o}^{\gamma} Q(x) \left[ \frac{1}{x} - \frac{1}{\gamma_{max}} \right]^{-\alpha} dx \quad (3.24)$$

where $\alpha = t_{acc}/t_{esc}$ and the lower limit of integration $x_o$ is given by

$$x_o = \left[ \frac{1}{\gamma_{max}} + \left( \frac{1}{\gamma} - \frac{1}{\gamma_{max}} \right) \exp(t/t_{acc}) \right]^{-1} \quad (3.25)$$

$\gamma_{max} = 1/(\zeta_{AR} t_{acc}) > \gamma_b$ is the maximum Lorentz factor an electron can attain in AR. For $t >> t_{acc}$ equation (3.25) can be approximated to be $\gamma_{min}$ as the injection term in equation (3.24) vanishes for $x < \gamma_{min}$.

The evolution of the electrons in CR is governed by the equation

$$\frac{\partial N(\gamma, t)}{\partial t} = \frac{\partial}{\partial \gamma} \left[ \zeta_{CR} \gamma^2 N(\gamma, t) \right] + Q_{AR}(\gamma) \quad (3.26)$$

where $\zeta_{CR} = \frac{\sigma_T}{6\pi m_e c} B_{CR}^2$ and the last term is the injection from AR, $Q_{AR}(\gamma) = n(\gamma)/t_{esc}$. 
Chapter 3. Emission Models and the Dynamics of AGN knots

For $t >> t_{acc}$

$$Q_{AR}(\gamma) \approx q_o \alpha \gamma^{-(\alpha+1)} \left(1 - \frac{\gamma}{\gamma_{max}}\right)^{\alpha-1}$$

$$\times MIN(\gamma, \gamma_b) \int_{\gamma_{min}}^{\gamma_{max}} x^{-p} \left[1 - \frac{1}{x} \frac{1}{\gamma_{max}}\right]^{-\alpha} dx$$

(3.27)

The distribution of electron at time $t$ in CR from equation (3.26) is given by

$$N(\gamma, t) = \frac{1}{\zeta_{CR} \gamma^2} \int_{\gamma}^{\Gamma_o} Q_{AR}(x) \, dx$$

(3.28)

where $\Gamma_o = \gamma/(1 - \gamma \zeta_{CR} t)$.

From equation (3.24), it can be shown that the injection into CR (for $\alpha + 1 > p$) is a broken power-law with index $-p$ for $\gamma < \gamma_b$ and $-(\alpha + 1)$ for $\gamma > \gamma_b$. The synchrotron losses in CR introduces an additional break $\gamma_c$ in the electron spectrum depending upon the age of CR ($t_{obs}$) and the $B_{CR}$ (§3.1).

$$\gamma_c = \frac{1}{\zeta_{CR} t_{obs}}$$

(3.29)

The electron spectrum in CR at $t_{obs}$ can then have two different spectral shapes depending on the location of $\gamma_c$ with respect to $\gamma_b$.

(i) $\gamma_c > \gamma_b$: The final spectrum will have two breaks with indices

$$N(\gamma, t_{obs}) \propto \begin{cases} \gamma^{-p}, & \gamma_{min} < \gamma < \gamma_b \\ \gamma^{-(\alpha+1)}, & \gamma_b < \gamma < \gamma_c \\ \gamma^{-(\alpha+2)}, & \gamma_c < \gamma < \gamma_{max} \end{cases}$$

(3.30)
(ii) $\gamma_c < \gamma_b$: In this case the indices are

$$N(\gamma, t_{\text{obs}}) \propto \begin{cases} 
\gamma^{-p}, & \gamma_{\text{min}} < \gamma < \gamma_c \\
\gamma^{-(p+1)}, & \gamma_c < \gamma < \gamma_b \\
\gamma^{-(\alpha+2)}, & \gamma_c < \gamma < \gamma_{\text{max}} 
\end{cases}$$  \hspace{1cm} (3.31)

The resultant synchrotron emissivity can then be calculated by convolving $N(\gamma, t)$ with single particle emissivity averaged over an isotropic distribution of pitch angles (equation (2.47)).

The predicted spectrum depends on nine parameters, which are $q_o, \alpha, \gamma_{\text{min}}, \gamma_b, \gamma_{\text{max}}, p$, $B_{CR}$, $R$ and $t_{\text{obs}}$. Here $\alpha$ and $p$ are estimated from the radio-to-optical and optical-to-X-ray spectral indices; $q_o$ and $B_{CR}$ are constrained using the observed luminosity and equipartition magnetic field. For $R$ we assume the physical sizes measured in radio [123]. The age of the knot $t_{\text{obs}}$ is chosen to introduce a break in the observed spectrum at optical band and $\gamma_b$ is fitted to reproduce the observed X-ray flux. $\gamma_{\text{min}}$ and $\gamma_{\text{max}}$ are used as free parameters and are fixed at 5 and $10^8$.

### 3.3.2 Results and Discussion

We applied the above model to explain the knots D, E, A and B of M87 jet (Figure 3.6). The results of the fitting are shown in Figure 3.7 and the parameters used for the fit are given in Table 3.6. The spectrum of knot E can be explained by simple continuous injection model and the parameters we quote corresponds to continuous injection model. We did not model knot C due to significant differences in X-ray-optical properties. The optical image of knot C is a diffuse region with a single maximum whereas in the X-ray image there exists two distinct maxima coincident with the diffuse optical knot [100].

For all the fits shown in Figure 3.7, $\gamma_c < \gamma_b$. However, one can fit the spectrum with $\gamma_c > \gamma_b$ with proper choices of the parameters $\alpha$, $\gamma_b$ and $\gamma_c$. This degeneracy arises due to the unavailability of the ultraviolet (UV) spectral index since the present model predicts the corresponding particle spectral index as $-(\alpha + 1)$ or $-(p + 1)$ (equations (3.30) and (3.31)) depending upon the above two conditions. Future observations at these photon energies may help in validating the present model and also will remove this degeneracy. Also, to obtain a precise values for $p$ and $\alpha$, spectral indices at radio, optical and X-ray energies should be known accurately. It should be noted here that in general $t_{\text{acc}}$ and $t_{\text{esc}}$ can be energy
Fig. 3.7: The spectral fit for the knots of M87 with the observed fluxes in radio, optical (Perlman et al. [124]) and X-ray (Perlman & Wilson [100]). Knot labels are the same as used in the literature [121].
Table 3.6: Model Parameters for the Knots of M87

| Knot | $q_o$ | $\alpha$ | $p$ | $\gamma_b$ | $B_{CR}$ | $t_{obs}$ | $\gamma_c$ | $R$ |
|------|-------|---------|-----|-----------|---------|---------|----------|-----|
| D    | 7.5   | 1.75    | 2.35 | 1.8       | 9.3     | 1.6     | 5.7      | 12  |
| E    | 4.7   | ...     | 2.36 | ...       | 5.9     | 2.5     | 9.0      | 17  |
| F    | 0.6   | 2.3     | 2.3  | 2.1       | 4.5     | 4.7     | 8.2      | 29  |
| A    | 1.0   | 2.45    | 2.29 | 2.0       | 4.7     | 3.5     | 10.1     | 55  |
| B    | 0.9   | 2.75    | 2.3  | 1.4       | 4.7     | 3.9     | 9.0      | 50  |

Columns: (1) Knot name; (2) Normalisation of the power-law injected into AR (for knot E it is the normalisation of the power-law injected into CR (see text)); (3) Ratio between the acceleration time-scale and escape time-scale in AR; (4) Index of the power-law spectrum injected into AR; (5) Maximum energy of the electron Lorentz factor injected into AR; (6) CR magnetic field; (7) Time of observation; (8) Break Lorentz factor of the electrons due to synchrotron cooling in CR; (9) Size of CR measured in radio [123].

For all cases, the minimum Lorentz factor ($\gamma_{min}$) injected into AR is 5 and maximum Lorentz factor ($\gamma_{max}$) attained in AR is $10^8$.

dependent. In such a situation the solution (equation (3.24)) may differ from its form and the index beyond $\gamma_b$ may not be the one discussed above.

A possible scenario of the present model is where the AR is a region around an internal shock following an external shock. The electrons injected into the AR can be those that are already accelerated by an external shock and are advected downstream to be accelerated further by the internal shock [125, 126]. Alternatively, re-acceleration of a power-law electron distribution by turbulence at boundary shear layers can also be another possible scenario [127, 128]. Inclusion of these scenarios in their exact form into the present model will make it more complex and is beyond the scope of the present work.

Perlman & Wilson (2005) proposed a modified continuous injection model where the volume within which particle acceleration occurs is energy dependent [100]. This is expressed in terms of a filling factor $f_{acc}$ which is the ratio between the observed flux and the flux predicted by the simple continuous injection model. They found $f_{acc}$ declining with increasing distance from the nucleus suggesting particle acceleration taking place in a larger fraction of the jet volume in the inner jet than the outer jet. The energy dependence of $f_{acc}$ also indicates that particle acceleration regions occupy a smaller fraction of the jet volume at higher energies. Even though the model is phenomenological, it indicates that the process of high-energy emission from the knots is as complicated as their physical region. However, the mechanism responsible for the filling factor is not explained.

Liu & Shen (2007) proposed a two zone model to explain the observed spectrum of the knots of the M87 jet. In their model electrons are accelerated to relativistic energies in the
acceleration region (AR) and lose most of their energies in the cooling region (CR) through the synchrotron process. They considered that the AR and CR are spatially separated and introduced a break in the particle spectrum injected in the CR through the advection of particles from the AR to the CR. This along with the cooling break in the CR produces a double broken power-law with indices $-p$, $-(p + 1)$ and $-(p + 2)$ which is then used to fit the observed spectra. However, the present model assumes that the AR and CR are cospatial, supporting a more physical scenario where electrons accelerated by the shock cool in its vicinity.

We explored the possibility of the present model to reproduce the X-ray flux of other FR I galaxies (detected by Chandra) which are observed to have lower radio luminosity and relatively smaller jets when compared with FR II galaxies. The X-ray emission from the knots and/or the jets of the FR I galaxies, namely 3C 66B [129], 3C 346 [130], Cen A [131] and 3C 296 [132], listed in the online catalogue of extragalactic X-ray jets XJET6, which are not explained by synchrotron emission from simple one zone models, can be reproduced by the present model.

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6http://hea-www.harvard.edu/XJET/
Chapter 4

Boundary Shear Acceleration in the Jet of MKN 501

MKN 501 is a nearby BL Lac object (z = 0.034) and also the second extra galactic source detected in TeV photon energies by ground-based Cherenkov Telescopes [133]. It was later detected in MeV photon energies by the satellite-based experiment EGRET [134]. Radio images of MKN 501 show a jet emerging from a bright nucleus [135, 136, 137, 101]. The high-resolution (milliarcsecond) radio images show a transverse jet structure with the edges being brighter than the central spine commonly referred as “limb-brightened” structure [135, 136, 101] (Figure 4.1). This feature is usually explained by the “spine-sheath” model where the velocity at the jet spine is larger compared to the velocity at the boundary. Such a radial stratification of velocity across the jet arises when jet moves through the ambient medium and the viscosity involved will cause a shear at the boundary. Three-dimensional hydrodynamic simulations of relativistic jets [138] and two-dimensional simulations of relativistic magnetized jets [139] also support the presence of jet velocity stratification due to its interaction with the ambient medium. The existence of velocity shear at the jet boundary was first suggested by Owen, Hardee & Cornwell [49] to explain the morphology of M87 jet. Perlman et al. [140] later confirmed it through the polarization studies of M87 jet. Limb-brightening can occur in a misaligned jet with velocity stratification. For a proper combination of flow velocities, one will see a Doppler-boosted image of the boundary compared to the less boosted spine due to relativistic effects [141, 142]. A possible consequence of the velocity shear is the alignment of the magnetic field at the boundary parallel to the flow velocity due to stretching of the frozen-in field lines of the plasma [143]. The polarization angle observed at the jet boundary of MKN 501 is perpendicular to the jet axis indicating
a parallel magnetic field \[144, 137\]. However, it should be noted here that the polarization angle at the jet spine indicates a perpendicular magnetic field, and this along with the parallel magnetic field at the jet boundary can be an outcome of a dynamically dominant toroidal magnetic field structure \[144, 145, 146\]. The radial velocity stratification of the jet can introduce Kelvin-Helmholtz instability, and the stability of jets against this instability was studied by various authors \[147, 148, 149, 150, 151\].

Giroletti et al. \[101\] have studied the limb-brightened structure of MKN 501 jet considering the differential Doppler boosting at the jet spine and the boundary \[142, 141\], and concluded that the viewing angle (angle between the jet and the line of sight of the observer) of the radio jet should be more than \(15^\circ\). However, high-energy studies of MKN 501 demand that the viewing angle of the jet should be \(\approx 5^\circ\) in order to explain the observed rapid variability.
Chapter 4. Boundary Shear Acceleration in the Jet of MKN 501

and the high-energy emission \[152, 153\]. Considering the fact that the γ-ray emission is originated from the inner part of the jet close to nucleus, Giroletti et al. \[101\] suggested that a bending of the jet may happen immediately after the γ-ray zone to explain the required large viewing angle of the radio jet. However, the mechanism required to bend the jets is still not well understood (jets deflected due to the pressure gradient in external medium are studied by Canto & Raga \[43\]; Raga & Canto \[44\]; Mendoza & Longair \[154\]), and moreover the observed large bending of the jet in the radio maps can be the apparent one because of projection effects. This projection effects are even amplified when the jet is close to the line of sight. However it needs to be noted here that jets with large bending angle are indeed observed \[41\].

The limb-brightened structure can also be explained if we consider the synchrotron emission from the particles accelerated at the boundary, and this inference does not require large viewing angle. Eilek \[155, 48\] considered the acceleration of particles due to turbulence initiated by Kelvin-Helmholtz and Rayleigh-Taylor instabilities at the jet boundary. Particles at the boundary can also be accelerated via shear acceleration \[46, 47\], and this case is considered in the present work. The acceleration of particles in a shear flow or by turbulence is well studied by various authors for both relativistic and non-relativistic cases \[156, 157, 158, 159, 160, 161, 162\].

4.1 Shear acceleration at MKN 501 jet boundary

The particle acceleration process at the jet boundary can be described by the diffusion equation in momentum space. The evolution of an isotropic phase space distribution is given by \[163\]

\[
\frac{\partial f(p)}{\partial t} = \frac{1}{p^2} \frac{\partial}{\partial p} \left( p^2 D(p) \frac{\partial f(p)}{\partial p} \right)
\]

(4.1)

where \(D(p)\) is the momentum diffusion coefficient. The characteristic acceleration timescale can be written as

\[
t_{\text{acc}} = p^3 \left[ \frac{\partial}{\partial p} \left( p^2 D(p) \right) \right]^{-1}
\]

(4.2)

If we consider a sheared flow, the electrons are scattered across different velocity layers by turbulent structures which are embedded in the shear flow. Berezhko \[46\] showed in such
case that there will be a net gain of energy in the electrons getting scattered, and this process is referred to as shear acceleration ([2,1.3]). The momentum diffusion coefficient in case of a shear flow can be written as [160, 164]

\[ D_s(p) = \chi p^2 \tau \]  

(4.3)

where \( \tau \) is the mean scattering time given by \( \tau \simeq \lambda/c \) with \( \lambda \) being the mean free path and \( \chi \) is the shear coefficient given for a relativistic flow as [165]

\[ \chi = \frac{c^2}{15 \left( \Gamma(r)^2 - 1 \right)} \left( \frac{\partial \Gamma}{\partial r} \right)^2 \]  

(4.4)

where \( \Gamma(r) \) is the bulk Lorentz factor of the flow and \( r \) is the radial coordinate of the jet cross section. Using equation (4.3), the shear acceleration time-scale \( t_{acc,s} \) for \( \tau = \tau_o \xi \) will be

\[ t_{acc,s} = \frac{1}{(4 + \xi)\chi \tau} \]  

(4.5)

In case of turbulent acceleration (stochastic), the particles are scattered off by randomly moving scattering centres and gets energized by a second-order Fermi acceleration. The momentum diffusion coefficient in this case can be approximated as [164]

\[ D_t(p) \simeq \frac{p^2}{3 \tau} \left( \frac{V_A}{c} \right)^2 \]  

(4.6)

where the Alfvén velocity \( V_A \) is given by

\[ V_A = \frac{B}{\sqrt{4\pi \rho}} \]  

(4.7)

Here \( B \) is the magnetic field and \( \rho \) the mass density of the jet. Hence, the turbulent acceleration time-scale \( t_{acc,t} \) will be

\[ t_{acc,t} = \frac{3 \tau}{(4 - \xi) \left( \frac{c}{V_A} \right)^2} \]  

(4.8)

For shear acceleration to be dominant over turbulent acceleration \( t_{acc,s} < t_{acc,t} \). If we consider Bohm diffusion (\( \xi = 1 \)) then the mean free path of the electron aligned to the magnetic field \( (\lambda_\parallel) \) scales as the gyro radius \( (r_g) \) [166], \( \lambda_\parallel \simeq \eta \frac{m_e c^2}{e B} \), where \( \eta \) is a numerical factor (\( \eta > 1 \) for magnetized particles) and \( \gamma(\gg 1) \) is the Lorentz factor of the scattered ele-
tron. Since the magnetic field at the jet boundary of MKN 501 is parallel to the jet axis (or toroidal) \cite{137, 144, 145}, we consider \( \tau \simeq \lambda_{\parallel}/c \). Also if we consider
\[
\frac{\partial \Gamma}{\partial r} \simeq \frac{\Delta \Gamma}{\Delta r}
\]
(4.9)
where \( \Delta \Gamma \) is the difference between the bulk Lorentz factor at the jet spine and the jet boundary and \( \Delta r \) is the thickness of the shear layer, then the condition for shear acceleration to be dominant over turbulent acceleration will be
\[
\Delta r < \frac{\eta \gamma m_e c^3 (\Delta \Gamma)}{e B^2} \left[ \frac{4\pi \rho}{3 (\Gamma(r)^2 - 1)} \right]^{\frac{1}{2}}
\]
(4.10)
If we consider the mass density of the jet is dominated by cold protons and if the number of protons is equal to the number of non-thermal electrons, then the jet mass density can be written in terms of equipartition magnetic field \( B_{eq} \) as
\[
\rho \simeq \frac{m_p B_{eq}^2 (2 \alpha - 1)}{16\pi m_e c^2 \alpha \gamma_{\text{min}}}
\]
(4.11)
and equation (4.10) will be
\[
\Delta r < 0.29 \frac{\eta \gamma c^2 (\Delta \Gamma)}{e B_{eq}} \left[ \frac{m_e m_p (2 \alpha - 1)}{\alpha \gamma_{\text{min}} (\Gamma(r)^2 - 1)} \right]^{\frac{1}{2}}
\]
(4.12)
where \( \alpha \) is the observed photon spectral index, \( m_p \) is the proton mass and \( \gamma_{\text{min}} \) is the Lorentz factor of electron responsible for the minimum observed photon frequency \( \nu_{\text{min}} \). The equipartition magnetic field can be expressed in terms of observed quantities as
\[
B_{eq} \simeq 9.62 \frac{1}{\Gamma(r)} (m_e c e \nu_{\text{min}})^{\frac{1}{2}} \left[ \frac{d_L^2 F(\nu_{\text{min}})}{V \sigma_T (2 \alpha - 1)} \right]^{\frac{1}{2}} G
\]
(4.13)
where \( F(\nu_{\text{min}}) \) is the flux at the minimum observed frequency \( \nu_{\text{min}} \), \( d_L \) is the luminosity distance, \( V \) is the volume of the emission region and \( \sigma_T \) is Thomson cross section. Hence, for \( \Gamma(r)^2 \gg 1 \) and \( \alpha \simeq 0.7 \), shear acceleration will dominate the particle spectrum at the jet...
boundary of MKN 501 if the thickness of the shear layer

\[ \Delta r < 7.22 \times 10^{-9} \times \left( \eta \right) \left( \frac{\Delta \Gamma}{10} \right) \left( \frac{\nu_{\text{obs}}}{1.6 \, \text{GHz}} \right)^{\frac{1}{2}} \left( \frac{\nu_{\text{min}}}{10 \, \text{MHz}} \right)^{-\frac{6}{11}} \times \left( \frac{F(10 \, \text{MHz})}{910 \, \text{mJy}} \right)^{\frac{3}{11}} \left( \frac{R}{1.5 \, \text{pc}} \right)^{\frac{3}{11}} \, \text{pc} \]  

(4.14)

where \( R \) is the radius of the spherical region considered. (We assume 10 MHz as minimum observed frequency, and the flux at 10 MHz is obtained from the flux at 1.6 GHz considering the same spectral index. The flux at 1.6 GHz and \( R \) in equation (4.14) are obtained from a region around R.A. 10 mas and declination −10 mas from Figure 4.1 (bottom left). The corresponding equipartition magnetic field \( B_{eq} \) for \( \Gamma = 5 \) is \( 1.2 \times 10^{-3} \, \text{G} \).

The electrons accelerated by shear acceleration cool via synchrotron radiation. The cooling time for synchrotron loss is given by

\[ t_{\text{cool}} = \frac{6 \pi m_e c}{\gamma \sigma_T B_{eq}^2} \]  

(4.15)

Using equations (4.5) and (4.15), we find

\[ \frac{t_{\text{acc,s}}}{t_{\text{cool}}} \simeq 1.5 \times 10^{-12} \left( \frac{B}{1.2 \times 10^{-3} \, \text{G}} \right)^3 \left( \frac{\eta}{10} \right)^{-1} \left( \frac{\Gamma(r)}{5} \right)^2 \times \left( \frac{\Delta r}{10^{-9} \, \text{pc}} \right)^2 \left( \frac{\Delta \Gamma}{10} \right)^{-2} \]  

(4.16)

and since \( t_{\text{acc,s}} \ll t_{\text{cool}} \), shear acceleration dominates over synchrotron cooling. It can be noted that equation (4.16) is independent of the electron energy and hence the maximum energy of the electron will be decided by the loss processes other than synchrotron loss (which are not considered in this simplistic treatment).

If we maintain the general form of mean scattering time \( \tau = \tau_0 p^\xi \), then for shear acceleration to dominate over turbulent acceleration the thickness of the shear layer (\( \Delta r \)) should be

\[ \Delta r < 1.7 \times 10^6 \frac{\tau_0 p^\xi (\Delta \Gamma)}{\Gamma(r)} \left[ \frac{(4 + \xi)(2 \alpha - 1)}{\alpha (4 - \xi) \gamma_{\text{min}}} \right]^{\frac{1}{2}} \, \text{cm} \]  

(4.17)

It can be noted that equations (4.10) and (4.17) are equal, if we set in the latter \( \xi = 1 \) and \( \tau_0 p^\xi = \eta \, \gamma_{\text{g}} / c \).
4.2 Particle Diffusion at the jet boundary and Limb-brightening

Particles accelerated at the shear layer of the jet boundary diffuse into the jet medium before getting cooled off via synchrotron radiation. As the magnetic field at the jet boundary is parallel to the jet axis (or toroidal) \[137, 144, 145\], the radial diffusion of the electron into the jet medium is determined by cross-field diffusion. The cross-field diffusion coefficient can be approximated as \[167, 168, 166\]

\[
\kappa \approx \frac{1}{3 \eta} r_g c \quad (4.18)
\]

where \(\eta(>1)\) is the scaling factor determining the field-aligned mean free path (see §4.1). The radial distance \(R_{diff}\) that the electron diffuse before getting cooled can then be approximated as

\[
R_{diff} \approx \sqrt{\kappa \perp t_{cool}} \quad (4.19)
\]

Using equations (4.15) and (4.18) and considering the equipartition magnetic field, we get

\[
R_{diff} \simeq 2.9 \times 10^{-4} \left( \frac{\eta}{10} \right)^{-\frac{1}{2}} \left( \frac{B}{1.2 \times 10^{-3} \text{ G}} \right)^{-\frac{3}{2}} \text{ pc} \quad (4.20)
\]

Since the thickness of the shear layer \(\Delta r \ll R_{diff}\) (refer equations (4.14) and (4.20)), the thickness of the limb-brightened structure will be \(\approx R_{diff}\). This corresponds to an angular distance of \(4.7 \times 10^{-4}\) mas which is beyond the resolution of present-day telescopes.

For \(\tau = \tau_0 p^\xi\), the cross-field diffusion coefficient will be

\[
\kappa \perp \approx \frac{1}{3 \tau_0} r_g^2 p^{-\xi} \quad (4.21)
\]

Using equations (4.15) and (4.19) we get

\[
R_{diff} \simeq 5.2 \times 10^{15} B^{-2} \tau_0^{-\frac{1}{2}} p^{\frac{1-\xi}{2}} \text{ cm} \quad (4.22)
\]

and hence the thickness of the limb brightened structure will be energy dependent for \(\xi \neq 1\).
4.3 Spectral index

If we add mono-energetic particle injection term \( \delta(p-p_o) \) and particle escape term \(-1/t_{esc}\) in equation \((4.1)\), then the steady state equation in case of shear acceleration for \( p > p_o \) and \( \xi = 1 \) can be written as

\[
p^3 \frac{d^2 f_s}{dp^2} + 5p^2 \frac{df_s}{dp} - \frac{f_s}{\chi \tau_o t_{esc}} = 0 \tag{4.23}
\]

and in case of turbulent acceleration it will be

\[
p \frac{d^2 f_t}{dp^2} + 3 \frac{df_t}{dp} - \frac{f_t}{\psi t_{esc}} = 0 \tag{4.24}
\]

where \( \psi = \frac{V^2}{S c^2 \tau_o} \). If we substitute \( p = 1/x \) in equation \((4.23)\) we get

\[
x \frac{d^2 f_s}{dx^2} - 3 \frac{df_s}{dx} - \frac{f_s}{\chi \tau_o t_{esc}} = 0 \tag{4.25}
\]

Equations \((4.24)\) and \((4.25)\) can be solved analytically \([169]\). The solutions are complex and are given by

\[
f_s = \left( \frac{1}{\chi \tau_o p t_{esc}} \right)^2 \times \left[ a_s J_4 \left( 2i \sqrt{\frac{1}{\chi \tau_o p t_{esc}}} \right) + b_s Y_4 \left( 2i \sqrt{\frac{1}{\chi \tau_o p t_{esc}}} \right) \right] \tag{4.26}
\]

and

\[
f_t = \left( \frac{\psi t_{esc}}{p} \right) \left[ a_t J_2 \left( 2i \sqrt{\frac{p}{\psi t_{esc}}} \right) + b_t Y_2 \left( 2i \sqrt{\frac{p}{\psi t_{esc}}} \right) \right] \tag{4.27}
\]

where \( J_n(z) \) and \( Y_n(z) \) are the Bessel functions of first and second kind and \( a_s, b_s, a_t \) and \( b_t \) are constants. For negligible escape \( (t_{esc} \to \infty) \), using the limiting forms of Bessel functions \([170]\), the solutions, equations \((4.26)\) and \((4.27)\), approach a power-law \( f_s \propto p^{-4} \) and \( f_t \propto p^{-2} \). The shear-accelerated particle number density will then be \( n_s(p) \propto p^{-4} \) and the corresponding synchrotron photon flux will be \( S_{\nu,\text{shear}} \propto \nu^{-1/2} \). For turbulent acceleration, the number density will be independent of \( p \) \( (n_t(p) \propto p^0) \) and hence the observed synchrotron photon flux will be a flat one \( S_{\nu,\text{turb}} \propto \nu^{1/3} \) \([106]\). The spectral index map of MKN 501 jet indicates a steep photon spectrum at the boundary and a flat spectrum at the spine.
Chapter 4. Boundary Shear Acceleration in the Jet of MKN 501

Fig. 4.2: Low resolution 1.6 GHz-4.8 GHz spectral index map of MKN 501 jet. Figure reproduced from Giroletti et al. [101].

Hence, it can be argued that the shear acceleration may be dominant at the jet boundary of MKN 501 and turbulent acceleration at the jet spine. However, $\xi$ is usually related to the turbulent spectral index [171] which may be different at the jet boundary and jet spine.

4.4 Discussion

As the AGN jet moves through the ambient medium, the viscosity involved will cause a shear at the jet boundary, and hence acceleration of particles in these shear layers is unavoidable. If the shear gradient $\partial \Gamma / \partial r$ is very steep or if the shear layer is very thin (equation (4.14)), then shear acceleration can dominate over the turbulent acceleration initiated by the instabilities at the jet boundary [48]. Turbulent acceleration may play an important role at the interior regions of the jet [162] and can provide an alternative to explain the emission from the inter knot regions of AGN jets [172, 173]. The observed hard spectrum at the jet spine [101] also support this inference since turbulent acceleration can produce a hard particle spectra [162] (also shown in §4.3). The electrons accelerated by the turbulence can be reaccelerated by shocks and can form a broken power-law electron spectrum. This can possibly explain the break in the radio-to-X-ray spectrum of the knots of FRI jets (§3.3).
Giroletti et al. [101] calculated the jet viewing angle ($\theta$) using the correlation between the core power and the total power [174]. They estimated the jet viewing angle to be within $10^\circ < \theta < 27^\circ$ by comparing the observed core radio power and the expected intrinsic core power derived from the correlation. However, this estimation may vary if the core flux density variability is more than a factor of 2. Also considering the variation of the parameter values in the correlation with increased number of samples, this may not provide a strong constrain on the jet viewing angle. The estimate of $\theta$ based on the adiabatically expanding relativistic jet model [175] may not be a strong constraint as it considers a simplified situation. Also, the constrain is less severe in case of perpendicular magnetic fields, and observed polarization studies have indicated the presence of perpendicular magnetic fields at jet spine [144, 137]. Stawarz & Ostrowski [159] proposed a model similar to the present one; however, their aim was to show the observational implications of the two-component particle spectrum (power-law distribution with high-energy pile-up) formed at the boundary shear layer and the complex beaming pattern.
Chapter 5

A Two zone Model for Blazar Emission Mechanism

The spectral energy distribution (SED) of blazars are characterized by a typical double-hump feature. The first component peaks at IR/optical energies for low energy peaked blazars and at UV/soft X-ray energies for high energy peaked blazars. Whereas, the second component peaks at hard X-ray/\(\gamma\)-ray region. The first component is generally modelled as the synchrotron emission due to the cooling of relativistic electrons in a magnetic field while the second component is considered to be produced due to the inverse Compton scattering of soft photons by the relativistic electrons. The target photons for the inverse Compton process can be either synchrotron photons themselves (SSC) \[176,177,102,83\] or external radiation (EC) (\[2.5.1\]) \[178,69,179\].

Besides their non-thermal continuum, blazars also exhibit rapid and strong variability (for example MKN 421 \[180,181,182\], MKN 501 \[183,184\], PKS 2155-304 \[185,186,187\], etc.). In some cases the variability detected at different frequencies are non simultaneous and the flaring activity is referred as hard lag or soft lag depending upon the observed flare pattern. In a hard lag, the low frequency flare leads the high frequency one and for the soft lag it is the other way. Takahashi et al. \[176\] reported a soft lag in ASCA observation of MKN 421, while Fossati et al. \[180,181\] reported a hard lag in the BeppoSAX observation of MKN 421 in 1997-1998. Similarly for PKS 2155-304, a soft lag was reported by Chiappetti et al. \[185\] and Kataoka et al. \[186\] in BeppoSAX observation of 1997 and ASCA observation of 1994 respectively. However, Takahashi et al. \[182\] contradicted the conclusions of Fossati et al. \[180,181\] by reporting no spectral lag for the same data set. Also, Edelson et al. \[187\] reported no spectral lag for PKS 2155-304 from XMM-Newton observation in 2000.
Nevertheless, the study of short time-scale variability at different wavelength region and their interrelations is very important to understand the geometrical and causal connections between emission regions and emission processes.

The flux variability observed in blazars have been studied by several authors using one zone and two zone models \cite{80, 82, 102, 83}. Kirk et al. \cite{82} and Kusunose et al. \cite{83} assumed a two zone model where particles are accelerated in a region presumably by a shock and escape into a cooling region where they lose their energy by radiative processes. On the contrary, one zone models assume instantaneous particle acceleration and do not consider the particle acceleration process separately. Chiaberge & Ghisellini \cite{102} explained the short time variability observed in MKN 421 by dividing the emission region into thin slices. Moderski et al. \cite{188} modelled the variability detected in 3C279 considering SSC and EC processes.

We consider a two zone model, namely acceleration region and cooling region, to study the spectral behaviour of blazars. This model is similar to the one discussed in §3.3 used to model the knots of M87 jet. However here we assume mono energetic particles are injected into the acceleration region which are then accelerated. We studied the effect of particle acceleration time-scale on the variability pattern as well as spectral evolution under two different scenarios: (a) acceleration time-scale is independent of particle energy and (b) acceleration time-scale depends on particle energy.

## 5.1 Model

We model the broadband emission of blazar from a spherical blob (cooling region) of radius $R$ permeated by a tangled magnetic field $B$, moving down the jet with a bulk Lorentz factor $\Gamma$. The Doppler factor of the blob is given by $\delta = [\Gamma(1 - \beta \cos \theta)]^{-1}$, where $\theta$ is the angle made by the jet with the line of sight.

Relativistic electrons are injected into the cooling region from an acceleration region, around a shock front moving with velocity $v_s$. The acceleration region is continuously fed with mono-energetic electrons with Lorentz factor $\gamma_0$ at a rate $Q_0$ electrons per unit volume. It is assumed that electrons are accelerated by diffusive shock acceleration process \cite{77} and their evolution $Q(\gamma, t)$ is described by,

$$
\frac{\partial Q(\gamma, t)}{\partial t} + \frac{\partial}{\partial \gamma} \left[ \left( \frac{d\gamma}{dt} \right)_{\text{acc}} - \left( \frac{d\gamma}{dt} \right)_{\text{syn}} \right] Q(\gamma, t) + \frac{Q(\gamma, t)}{t_{\text{esc}}} = Q_0 \delta(\gamma - \gamma_0) \tag{5.1}
$$
where

\[
\left( \frac{d\gamma}{dt} \right)_{\text{acc}} = \gamma / t_{\text{acc}} \tag{5.2}
\]

and

\[
\left( \frac{d\gamma}{dt} \right)_{\text{syn}} = \frac{4}{3} \frac{\sigma T}{m_e c^2} \gamma^2 \frac{B_{\text{acc}}^2}{\pi} \tag{5.3}
\]

are the particle acceleration rate and synchrotron loss rate of the particles in the acceleration region. Here \( t_{\text{acc}} \) and \( t_{\text{esc}} \) are the acceleration and escape time-scale of the electrons respectively and \( B_{\text{acc}} \) is the magnetic field in the acceleration region. In the frame work of diffusive shock acceleration, acceleration time-scale can be written as

\[
t_{\text{acc}} = \left( \frac{20 c}{3 v_s^2} \right) \left( \frac{m_e c^2}{e B_{\text{acc}}} \right) \xi \gamma \tag{5.4}
\]

where

\[
t_{a,0} = \left( \frac{20 c}{3 v_s^2} \right) \left( \frac{m_e c^2}{e B_{\text{acc}}} \right) \xi \tag{5.5}
\]

The escape time-scale \( t_{\text{esc}} \) is assumed to be \( \alpha t_{\text{acc}} \) where \( \alpha \) is a constant factor. For energy independent scenario, \( \gamma \) in equation (5.4) is replaced by a constant value, \( \gamma_{\text{eff}} \) and for energy dependent scenario, equation (5.4) is maintained as such. The dependence of \( t_{\text{esc}} \) on \( t_{\text{acc}} \) is kept unchanged in both the cases. The analytical solutions of equation (5.1), under these two physical conditions, have already been discussed by [82].

Accelerated particles are injected into the cooling region from the acceleration region at a rate \( Q(\gamma, t) / t_{\text{esc}} \) where they lose energy via synchrotron and SSC processes. The evolution of the particle spectrum \( N(\gamma, t) \) in the cooling region is governed by

\[
\frac{\partial N(\gamma, t)}{\partial t} - \frac{\partial}{\partial \gamma} \left[ \left( \frac{d\gamma}{dt} \right)_{\text{loss}} N(\gamma, t) \right] + \frac{N(\gamma, t)}{\tau} = \frac{Q(\gamma, t)}{t_{\text{esc}}} \tag{5.6}
\]

where total energy loss rate of an electron,

\[
\left( \frac{d\gamma}{dt} \right)_{\text{loss}} = \left( \frac{d\gamma}{dt} \right)_{\text{syn}} + \left( \frac{d\gamma}{dt} \right)_{\text{SSC}} \tag{5.7}
\]
The synchrotron loss rate is given by

\[
\left( \frac{d\gamma}{dt} \right)_{\text{syn}} = \frac{4}{3} \frac{\sigma_T}{m_e c} \gamma^2 \frac{B_0^2}{8\pi} \tag{5.8}
\]

and the energy loss rate due to SSC process is given by

\[
\left( \frac{d\gamma}{dt} \right)_{\text{SSC}} = \frac{4}{3} \frac{\sigma_T}{m_e c} \gamma^2 U_{ph} \tag{5.9}
\]

\(U_{ph}\) is the energy density of synchrotron photons in the cooling region and can be calculated from the synchrotron specific intensity. Here \(B_0\) is the magnetic field in cooling region and \(\sigma_T\) is the Thomson cross-section.

The equation (5.6) is solved numerically using finite difference scheme [108, 102]. The specific intensity of synchrotron radiation at frequency \(\nu\) is given by

\[
I_s(\nu, t) = \frac{j_\nu(t)}{\kappa_\nu(t)} (1 - e^{-\kappa_\nu(t) R}) \tag{5.10}
\]

where, synchrotron emissivity \(j_\nu\) and synchrotron self-absorption coefficient \(\kappa_\nu\) are given by (equations (2.47) and (2.44))

\[
j_\nu(t) = \frac{1}{4\pi} \int_{\gamma_{\text{min}}}^{\gamma_{\text{max}}} N(\gamma, t) P(\gamma, \nu) \, d\gamma \tag{5.11}
\]

and

\[
\kappa_\nu(t) = -\frac{1}{8\pi m_e \nu^2} \int_{\gamma_{\text{min}}}^{\gamma_{\text{max}}} \frac{N(\gamma, t)}{\gamma (\gamma^2 - 1)^{3/2}} \frac{d}{dt} \left[ \gamma (\gamma^2 - 1)^{1/2} P(\gamma, \nu) \right] \tag{5.12}
\]

respectively. Here, \(P(\gamma, \nu)\) is single particle emissivity (equation (2.36)). The synchrotron photon energy density is then calculated using

\[
U_{ph}(\gamma, t) = \frac{4\pi}{c} \int_{\nu_{s,\text{min}}}^{\nu_{\text{max}}(\gamma)} I_s(\nu, t) \, d\nu \tag{5.13}
\]
where \[ \nu_{\text{max}}(\gamma) = \min[\nu_{s,\text{min}}, 3 m_e c^2 / 4h\gamma] \] (5.14)

Here \( \nu_{s,\text{min}} \) and \( \nu_{s,\text{max}} \) are the minimum and maximum frequency of the synchrotron photons. The specific intensity of inverse Compton scattering in the Thomson regime is calculated by using the standard formulation \([51]\).

Different possible mechanisms for the generation of flare have been discussed in the literature. For example, Kusunose et al. \([83]\) simulated flares by changing relation between acceleration time-scale\( (t_{\text{acc}}) \) and escape time-scale\( (t_{\text{esc}}) \) in the acceleration region for certain duration. Mastichiadis & Kirk \([81]\) discussed flares due to sudden changes in, (i) electron injection rate \( Q_0 \); (ii) maximum attainable energy of electrons \( \gamma_{\text{max}} \) and (iii) cooling region magnetic field \( B \). We simulated flares by increasing the injection into the acceleration region for a very short duration over and above the steady state injection. Finally, the observed flux is computed taking into account the Doppler boosting \([36]\) and cosmological effects.

### 5.2 Results and Discussion

#### 5.2.1 Spectral evolution

We considered the quiescent spectrum of MKN 421 as a test case and deduce the best fit parameter set for both energy dependent \( t_{\text{acc}} \) and energy independent \( t_{\text{acc}} \). Values of the parameters are tabulated in Table \([5.1]\) which are taken as the standard values and all the results presented in this work are based on these values. With these set of parameters \( t_{\text{acc}} \) turned out to be \( 1.4 \times 10^5 \) s (in source frame) for energy independent case, while \( 0.5\gamma \) s (in source frame) for energy dependent case. Hence the particle acceleration rate \( (\gamma / t_{\text{acc}}) \) becomes \( 0.70 \times 10^{-5}\gamma \) s\(^{-1}\) and \( 2 \) s\(^{-1}\) for energy independent and energy dependent cases respectively.

The fitted spectra of MKN 421 for both the cases of \( t_{\text{acc}} \) are shown in Figure \([5.1]\). The slope of the spectrum before synchrotron peak is decided by \( \alpha \), which relates the escape and the acceleration time-scale of electrons in the acceleration region. The escape of the particles from the cooling region produces a break at \( \nu \approx 10^{14} \) Hz. As the second hump of the spectrum is produced due to the self-synchrotron Compton process, the break in the
Table 5.1: Parameters used for fitting MKN 421

| Region           | Parameters             | Energy independent $t_{acc}$ | Energy dependent $t_{acc}$ |
|------------------|------------------------|------------------------------|----------------------------|
| Acceleration     | $\gamma_0$            | 1.00                         | 1.00                       |
|                  | $Q_0 \left(10^{-5} \# \text{ cm}^{-3} \text{ s}^{-1}\right)$ | 2.1                         | 2.1                        |
|                  | $B_{acc}$ (G)          | 0.137                        | 0.132                      |
|                  | $v_s$ (in units of c)  | 1                            | 0.17                       |
|                  | $\gamma_{eff}$        | $10^7$                       | ...                        |
|                  | $\xi$                 | $5 \times 10^3$             | $5 \times 10^3$            |
|                  | $\alpha$              | 1.70                         | 1.70                       |
| Cooling region   | $B_o$ (G)             | 0.18                         | 0.18                       |
|                  | $R$ (cm)              | $2.4 \times 10^{16}$         | $2.4 \times 10^{16}$       |
|                  | $\tau$ (s)            | $\frac{2.5 B}{\delta}$      | $\frac{2.5 B}{\delta}$    |
|                  | $\delta$              | 12                           | 12                         |

The synchrotron component of the spectrum is also reflected there.

As shown in Figure 5.1, the steady state spectrum of MKN 421 can be reproduced successfully by the two acceleration scenarios though rate of acceleration of particles in these scenarios are different. Hence the steady state spectrum of blazar cannot uniquely reflect the possible underlying electron acceleration mechanism and its dependencies on the electron energy. In case of energy independent $t_{acc}$, the rate of electron acceleration depends on the electron energy and is much less than the electron acceleration rate for energy dependent $t_{acc}$ (which is a constant in this case) for a wide range of $\gamma$. Hence the evolution of the spectrum for energy independent $t_{acc}$ will be much slower compared to the energy dependent $t_{acc}$ case as shown in Figure 5.2. This feature in the spectral evolution will also be reflected in the flare phenomena, if the flare is produced due to a sudden increase in the rate of injection of particles $Q_0$ in the acceleration region for a small duration.

For energy independent case the system attains the steady state on the 8th day (in the observer’s frame) of evolution. Although, in the energy dependent scenario, the spectrum attains the steady state faster but to make a comparative study of flares, we considered the steady state spectrum in both the scenarios only after the 8th day of evolution. We discuss below the characteristics of the flare phenomena under the two different scenarios of particle acceleration process discussed above.
Fig. 5.1: Two zone model fitting of MKN 421 archival data considering energy independent acceleration time-scale (top) and energy dependent acceleration time-scale (bottom) using the parameters given in Table 5.1. The observed data are taken from Kino et al. [190].
Fig. 5.2: Evolution of the model spectra in case of energy independent acceleration time-scale (top) and energy dependent acceleration time-scale (bottom). The spectra are plotted for every 0.1 days up to 1.5 days and there after at 2nd, 3rd, 6th and 8th days (in observer’s frame). The steady state is arrived after 8 days in both the cases.
5.2.2 Flare

As mentioned above, we simulate a flare by increasing the electron injection rate into the acceleration region by a factor of 50 for a duration of 0.2 days (in the observer’s frame) after the system reaches the steady state.

**Energy independent** $t_{acc}$

The simulated light curves for energy-independent $t_{acc}$ are shown in Figure 5.3. It is evident from top panel of Figure 5.3 that the high energy flares lag the low energy ones in time ($\sim$ kilo s) and it is also seen that the light curves are more asymmetric at lower energies. The lag in the light curves arise since the rate of acceleration of electrons in the acceleration region is proportional to energy of the electrons and the high energy electrons which are responsible for the high energy synchrotron emission, take longer time to attain the required energy. The asymmetry in the low energy light curves can be attributed to the longer synchrotron cooling time-scales of low energy electrons compared to the high energy ones. It can also be seen that the variability amplitude increases with frequency of emission. These features are in qualitative agreement with the *BeppoSAX* observations of MKN 421 [180, 181]. The lower panel of Figure 5.3 describes the simulated flare patterns for SSC component of emission. Since, the synchrotron photons are Compton boosted by the same population of electrons, the variability features in SSC component are qualitatively same as in the synchrotron component.

The variation of the spectral index with respect to the flux (*hysteresis loops*) have also been studied at different energies and are shown in Figure 5.4. The sense of the loops is anti-clockwise representing the hard lag in the system.

**Energy dependent** $t_{acc}$

In this case the acceleration time-scale is proportional to the energy of the electrons and hence the rate of acceleration is independent of electron energies ($1/t_{a,0}$). For the parameters given in table 5.1, we find $t_{a,0} = 0.5$ s (in the source frame). As the rate of particle acceleration is high and same for all electron energies, the flares are initiated almost immediately after the extra injection into the acceleration region occurs. This fact is evident from near-simultaneous flare patterns shown in Figure 5.5. The light curves at optical/UV
Fig. 5.3: Light curves at different frequencies for energy independent acceleration time-scale during a simulated flare created by the enhanced injection into the acceleration region (see text). The vertical dashed line represents the duration of the enhanced particle injection in the acceleration region. The top panel corresponds to the emission due to synchrotron process and bottom panel corresponds to the SSC process.
Fig. 5.4: The hysteresis curves showing the variation of spectral index with respect to flux during a simulated flare for energy independent acceleration time-scale. The top panel corresponds to the emission due to synchrotron process and bottom panel corresponds to the SSC process. The sense of rotation of the loops is anti clockwise.
frequencies in the synchrotron component show prominent breaks in their rising part. Since, the rate of acceleration is high and same for all energies, electrons in acceleration region attain higher energy in shorter time-scale compared to the escape time-scale, \( t_{\text{esc}} \). These high energy particles enter cooling region and get cooled giving rise to high energy emission. Then they join the freshly injected lower energy electrons and contribute to the low energy emission as well. The difference in the rates of these contributions give rise to breaks in the light curves at low energy emissions. This phenomena gives rise to a tendency that the light curves at optical and UV bands peak later than the light curves in the soft and hard X-ray bands. Although, this is not very prominent in the simulated light curves, it is reflected in the clockwise sense of the hysteresis loop shown in Figure 5.6. The absence of such breaks at very low energy emissions (namely far-IR band and below) can be justified by the fact that the escape time-scale from the cooling region is smaller than the time required for the high energy electrons to cool to these energies.

Flares follow the same features of asymmetry as in the case of energy independent \( t_{\text{acc}} \). Similar features are reflected in the flares of SSC component also.

Such variations in the photon spectral index during blazar flare (clockwise and anti-clockwise) are observed for many sources [176, 185, 180, 191].

**Table 5.2:** Values of \( \eta \) for different frequencies

| Frequency (Hz) | \( \eta \)   |
|---------------|-------------|
| \( 10^{13} \) | 2.05 ± 0.02 |
| \( 10^{15} \) | 1.71 ± 0.09 |
| \( 10^{17} \) | 1.99 ± 0.16 |
| \( 10^{24} \) | 1.67 ± 0.04 |

**Dependence of lag time-scale on \( t_{\text{acc}} \)**

As shown in the previous section, the energy independent acceleration time-scale gives rise to a hard lag in the flare patterns for different radiation frequencies. The variation of lag time-scale for different shock velocities \( v_s \) is shown in Figure 5.7. The lag time-scales are calculated with respect to the flare at \( 10^{11} \) Hz. It is found that the lag time-scale varies with \( v_s \) following a power-law \( (\propto v_s^{-\eta}) \). The fitted value of \( \eta \) for different frequencies are given in Table 5.2. As \( t_{\text{acc}} \) depends on the shock velocity \( v_s \) \( (t_{\text{acc}} \sim v_s^{-2}) \), so the lag time-scale is
Fig. 5.5: Light curves at different frequencies for energy dependent acceleration time-scale during a simulated flare created by the enhanced injection into the acceleration region (see text). The vertical dashed line represents the duration of the enhanced particle injection in the acceleration region. The top panel corresponds to the emission due to synchrotron process and bottom panel corresponds to the SSC process.
Fig. 5.6: The hysteresis curves showing the variation of spectral index with respect to flux during a simulated flare for energy dependent acceleration time-scale. The top panel corresponds to the emission due to synchrotron process and bottom panel corresponds to the SSC process. The sense of rotation of the loops is clockwise.
Fig. 5.7: Variation of lag time-scale with shock velocity in case of energy independent acceleration. Filled circles corresponds to $10^{13}$ Hz, open circles corresponds to $10^{15}$ Hz, filled triangles corresponds to $10^{17}$ Hz and open triangles corresponds to $10^{24}$ Hz.

proportional to $t_{\text{acc}}^{\eta/2}$. 
Chapter 6

Summary and Conclusions

The complete and comprehensive understanding of the physics of AGN jet requires the study of different aspects of jet in different wavelength ranges. Such studies may give rise to a global picture of jets describing the dynamics and radiation emission processes self-consistently. Although this is a mammoth task and requires a long time, an attempt has been made in this thesis to study certain aspects of jet physics phenomenologically and to understand the possible inter-connections between them. Here, we studied the dynamics and emission processes from knots in jets observed in radio-optical-X-ray bands. Similarly, we considered blazars where we studied the limb brightening effect in Mkn 501 observed in radio wave band. This feature has intimate connection with physics of jet flow and the particle acceleration process at jet boundary. Our work has been further extended to include the study of flares in blazars in infra-red-to-TeV energy band.

To model the emission from the knots we used the archival radio, optical and X-ray data. A continuous injection plasma model, where non-thermal relativistic electrons are injected into an expanding spherical region with a tangled magnetic field, is used to study the broadband emission from the knots. Injected relativistic electrons lose energy by synchrotron process in the tangled magnetic field of the knot and by up-scattering the cosmic microwave background photons. The expansion also introduces an adiabatic loss. For a given observation time, the particle distribution in the emission region will be a broken power-law with a break at an energy where the cooling time-scale is equal to the observation time. The resultant spectrum is fitted to the observed data to yield the parameters of the model. The parameters obtained from the spectral fitting are physically reasonable and they are used to obtain the kinetic powers of the jets.
Above work has been further extended to include the dynamics of knots. Assuming that the knots are produced due to the collision of matter shells ejected randomly in time from the central engine, the complete kinematics of the shells are used to study their location of collision and the energetics. Non-thermal relativistic electrons are produced in the shock generated due the collision of shells. These electrons emit radiation by synchrotron and inverse Compton process as described above. Apart from fitting the observed spectrum, the other main conclusion of this work is that the location of the knots can be reproduced from the physically acceptable choice of parameters in the kinematics of shell collision. Therefore internal shocks can be considered as one of the viable mechanism of knot generation. It is to be noted that the timescales obtained from the kinematics of shells and internal shock, convincingly support the continuous injection scenario.

A two zone model was proposed to explain the emission from the knots of the M87 jet since simple models involving continuous injection/one-time injection of non-thermal particles failed to reproduce the observed X-ray flux and the spectral index. In the proposed model, we consider the injection of a power-law distribution of particles into an acceleration region where they are accelerated further. The particles then escape from the acceleration region into a cooling region where it lose energy mostly via synchrotron radiation. The particle distribution in the cooling region will be a double broken power-law with one break at energy corresponding to the cutoff energy of the initial injected power-law into the acceleration region and the next break at energy for which the cooling time-scale equals to the age of the knot. The observed radio-optical-X-ray spectrum from the knots in M87 jet are reproduced by the resultant synchrotron emission. In its simplest form, the model does not consider any specific acceleration process but assumes an energy independent acceleration time-scale. The model can successfully reproduce the broadband spectrum from the knots/jets of other FRI galaxies, namely 3C 66B, 3C 346 and 3C 296, which are not explained by synchrotron emission from simple one zone models.

High-resolution radio maps of the jet of the BL Lac object MKN 501 shows a limb-brightened feature and an explanation of this feature based on the differential Doppler boosting of a stratified jet, requires large viewing angle ($> 15^\circ$). The viewing angle constraints inferred from the high-energy $\gamma$-ray studies of the source is very small ($\sim 5^\circ$). Since the $\gamma$-ray emission originates from the inner region jet, close to the central engine, this model requires the jet to be bent to accommodate the viewing angle conflict. However, the observed limb-brightened structure of the MKN 501 jet can be explained if we consider the shear acceleration of particles at the boundary due to velocity stratification and their diffusion into the jet medium.
This inference does not require a large viewing angle as demanded by the explanation based on differential Doppler boosting of the jet spine and boundary. We have shown that shear acceleration dominates over turbulent acceleration at the boundary if we consider thin shear layer or a sharp velocity gradient. Also for the estimated set of parameters, shear acceleration time-scale is much smaller than synchrotron cooling time-scale allowing acceleration of electrons to be possible. The thickness of the limb-brightened structure will be decided by the distance electrons have diffused into the jet medium before loosing its energy via synchrotron radiation. However the estimated thickness is beyond the resolution of present day telescopes. Simple analytical solution of the steady state diffusion equation considering mono-energetic injection and particle escape indicates a steep particle spectra for the electrons accelerated at the shear layer in comparison with turbulent acceleration. The radio spectral index map of MKN 501 jet is also observed to have steep spectrum at the boundary supporting the presence of shear acceleration.

The temporal behaviour of the blazar emission is studied under the framework of a two zone model. The spectral evolution has been examined for two different physical conditions of diffusive shock acceleration mechanism, namely energy independent acceleration time-scale and energy dependent acceleration time-scale. The model is applied on the BL Lac object, MKN 421, to study the implications on the flare characteristics for the above mentioned conditions. We found that in case of energy independent particle acceleration, the photon spectrum evolves at a slower rate compared to the energy dependent case though their steady state spectra are not differentiable. The flare patterns at different frequencies show a hard lag in the energy independent acceleration scenario while they are near simultaneous in the energy dependent scenario. Hence, the presence/absence of time lags in the flare pattern has direct bearing on the underlying particle acceleration mechanism in a blazar jet. Also the presence of a break in the rising part of high energy light curves in the case of near-simultaneous flares suggests that the acceleration time-scale may depend on particle energy. In the case of energy independent $t_{\text{acc}}$, it is also shown that the time-lag between two given frequencies has a power-law dependence on the shock velocity. Hence, a simultaneous multi wavelength study of blazar variability with good time resolution may be useful to constrain the physical parameters of the blazar jet and may also reveal the nature of underlying particle acceleration process which is crucial in understanding the dynamics of the jet.

The models described here to study the different aspects of jets from different AGN can be improved further to determine and constrain the model parameters unambiguously if better quality data are available. This requires truly simultaneous long-term observation of AGN.
using both ground-based and space-based telescopes in different wavelength bands. Particularly, radio observations with high spatial resolution are necessary to study the structure of knots as well as the boundary layers of the jets. This will shed further light on the jet dynamics, instabilities and the possible particle acceleration sites along the jet.

To study the flux variability of blazars it is important to have simultaneous measurements of flaring events at different wavelengths. The high sensitivity timing study is particularly important at GeV-TeV energies where blazars show the fastest flux variations. More observations of blazars with presently operating ground-based Cherenkov telescopes like MAGIC, HESS, VERITAS telescopes are extremely important in this respect. The upcoming Cherenkov telescopes with higher sensitivity like MACE and CTA will of course enhance the quality of data. This will in turn help us to improve the theoretical models to have better understanding of AGN jets.

In fact the present work of blazar can further be extended to study the effects of extragalactic background light (EBL) on the blazar spectrum and possible estimation of EBL. This is an important issue in blazar research at high energies and this will be pursued in future.
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